

LECTURE NOTES IN PHYSICS

J. Al-Khalili
E. Roeckl
(Eds.)

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Dr. Christian Caron
Springer Heidelberg
Physics Editorial Department I
Tiergartenstrasse 17
69121 Heidelberg/Germany
christian.caron@springer.com

Jim Al-Khalili Ernst Roeckl (Eds.)

Editors

Jim Al-Khalili
School of Electronics & Physical
Sciences
University of Surrey
GU2 7XH, U.K.
E-mail: j.al-khalili@surrey.ac.uk

Ernst Roeckl
GSI
Atomic Physics Group
Planckstraße 1
64291 Darmstadt, Germany
E-mail: e.roeckl@gsi.de

J. Al-Khalili and E. Roeckl, *The Euroschool Lectures on Physics with Exotic Beams, Vol. II*, Lect. Notes Phys. 700 (Springer, Berlin Heidelberg 2006), DOI 10.1007/b11743651

Library of Congress Control Number: 2004108216

ISSN 0075-8450

ISBN-10 3-540-33786-5 Springer Berlin Heidelberg New York

ISBN-13 978-3-540-33786-7 Springer Berlin Heidelberg New York

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Printed in The Netherlands

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Typesetting: by the authors and techbooks using a Springer L^AT_EX macro package

Cover design: *design & production* GmbH, Heidelberg

Printed on acid-free paper SPIN: 11743651 54/techbooks 5 4 3 2 1 0

Preface

This is the second volume in a series of lecture notes based on the highly successful Euro Summer School on Exotic Beams that has been running yearly since 1993 (apart from 1999) and is planned to continue to do so. It is the aim of the School and these lecture notes to provide an introduction to radioactive ion beam (RIB) physics at the level of graduate students and young postdocs starting out in the field. Each volume will contain lectures covering a range of topics from nuclear theory to experiment to applications.

Our understanding of atomic nuclei has undergone a major re-orientation over the past two decades and seen the emergence of an exciting field of research: the study of exotic nuclei. The availability of energetic beams of short-lived nuclei, referred to as radioactive ion beams (RIBs), has opened the way to the study of the structure and dynamics of thousands of nuclear species never before observed in the laboratory. In its 2004 report “Perspectives for Nuclear Physics Research in Europe in the Coming Decade and Beyond”, the Nuclear Physics European Collaboration Committee (NuPECC) states that the field of RIB physics is one of the most important directions for the future science programme in Europe. In 2005 it published its “Roadmap for Construction of Nuclear Physics Research Infrastructures in Europe”. In addition, the NuPECC report Nuclear Science in Europe: Impact, Applications, Interactions (June 2002) highlighted just how widely RIB physics impacts other areas, from energy and the environment to medicine and materials science. There is little doubt that RIB physics has transformed not only nuclear physics itself but many other areas of science and technology too, and will continue to do so in the years to come.

While the field of RIB physics is linked mainly to the study of nuclear structure under extreme conditions of isospin, mass, spin and temperature, it also addresses problems in nuclear astrophysics, solid-state physics and the study of fundamental interactions. Furthermore important applications and spin-offs also originate from this basic research. The development of new production, acceleration and ion storing techniques and the construction of new detectors adapted to work in the special environment of energetic radioactive beams is also an important part of the science. And, due to the fact that one is no longer limited to the proton/neutron ratio of stable-isotope beams, virtually the whole chart of the nuclei opens up for research, so theoretical

models can be tested and verified all the way up to the limits of nuclear existence: the proton and neutron drip lines.

The beams of rare and exotic nuclei being produced are via two complementary techniques: in-flight separation and post-acceleration of low-energy radioactive beams. Both methods have been developed in a number of European large scale facilities such as ISOLDE (CERN, Switzerland), GANIL (Caen, France), GSI (Darmstadt, Germany), the Accelerator Laboratory of the University of Jyväskylä (Finland), INFN Laboratori Nazionali di Legnaro (Italy) and the Cyclotron Research Centre (Louvain-la-Neuve, Belgium). Indeed, so important is the continued running and success of the School that a number of these European facilities have committed to providing financial support over the coming years.

Volume I of this series has proved to be highly successful and popular with many researchers reaching for it for information or providing it for their PhD students as an introduction to a particular topic. We stress that the contributions in these volumes are not review articles and so are not meant to contain all the latest results or to provide an exhaustive coverage of the field but written in the pedagogical style of graduate lectures and thus have a reasonably long ‘shelf life’. As with the first volume, the contributions here are by leading scientists in the field who have lectured at the School. They were chosen by the editors to provide a range of topics within the field and will have updated their material delivered at the School (sometimes several years ago) to incorporate recent advances and results.

Finally, we wish to thank the lecturers who have contributed to this volume for their hard work and diligence, and indeed for their patience. It is difficult to find the time to lay out a subject in such a careful, thorough and readable style. We also wish to thank Dr Chris Caron and his colleagues at Springer-Verlag for the help, fruitful collaboration and continued support on this project.

March 2006

Jim Al-Khalili
Ernst Roeckl

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List of Contributors

Jose Benlliure

Universidade de Santiago
de Compostela
15782 Santiago de Compostela
Spain
j.benlliure@usc.es

Piet Van Duppen

Instituut voor Kern- en
Stralingsfysica
K.U. Leuven
3001 Leuven, Belgium
piet.vanduppen@fys.
kuleuven.be

Jürgen Eberth

Institut für Kernphysik
Universität zu Köln
50937 Köln, Germany
J.Eberth@ikp.uni-koeln.de

William B. Gelletly

Physics Department
School of Electronic and
Physical Sciences
University of Surrey
Surrey GU2 7XH, U.K.
W.Gelletly@surrey.ac.uk

Rainer Neugart

Institut für Physik
Universität Mainz
55099 Mainz, Germany
neugart@uni-mainz.de

Gerda Neyens

Instituut voor Kern- en
Stralingsfysica
K.U. Leuven
3001 Leuven, Belgium
gerda.neyens@fys.kuleuven.ac.be

Karsten Riisager

Institut for Fysik og Astronomi
Århus Universitet
8000 Århus C, Denmark
kvr@phys.au.dk

Nuclear Halos and Experiments to Probe Them

K. Riisager

Institut for Fysik og Astronomi, Aarhus Universitet, 8000 Århus C, Denmark and
PH Department, CERN, 1211 Geneve 23, Switzerland

Abstract. Halo nuclei are characterized by an approximate decoupling – spatially and dynamically – of one or two nucleons from the remaining nucleus. This lecture identifies the key features of halo systems and reviews scaling relations that enable halos in different systems to be compared. Regions where halos have been seen or are likely to be found are outlined. A general overview is given of how halos are studied experimentally followed by a more detailed discussion on nuclear reaction experiments at energies above 20 MeV/u and on halo effects in beta decay.

1 Introduction

After two decades of halo physics, many reviews now exist [1, 2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13] that will guide the specialist through the literature. There are also more accessible introductions, e.g., two lectures in the earlier LNP volume in this series [14, 15] and many overviews given at conferences and schools, some examples can be found in [16, 17, 18, 19, 20]. The basic aim of this lecture is that it together with Al-Khalili's lecture [14] should give newcomers a good starting point in the field. To allow the reader to build up some intuition for halo systems the lecture will focus on key concepts and on many aspects only give a brief overview. When working with halos in practice one always need to worry about details; references are given for those who need to do this worrying.

The historical development of halo physics has been covered briefly in [4, 12]. Here it suffices to mention that halos had been anticipated theoretically quite early on [21] and that several aspects of typical halo behaviour were seen experimentally and understood as being due to the loose binding, e.g., the enhancement in (p, γ) reactions [22] and the unusually strong E1 transition within ^{11}Be [23]. However, the halo concept only emerged clearly when the first reaction cross section measurements by Tanihata's group [24, 25] led Hansen and Jonson [26] to suggest the halo structure in order to explain the large sizes of nuclei such as ^{11}Li .

I shall start in Sect. 2 by discussing how a halo state differs from normal nuclear states and how they can be classified. Section 3 then looks at where halos have been found and where they are likely to show up in the future. This is followed by a general discussion in Sect. 4 on how halo properties

can be probed quantitatively in experiments. Section 5 introduces the most important probe of halo structure, nuclear reactions, and Sect. 6 discusses the relevance of (and complications in) beta decays.

2 What is a Halo?

The most obvious feature of a halo system is that it is spatially large. To be more precise, it is unusually large for a system “of its type”. This is not sufficient to distinguish halos from other large systems: we furthermore require that the outer parts of the system are in a classically forbidden region and therefore appear due to quantum mechanical tunnelling. This requirement now ensures that halos in different physical systems are sufficiently alike that we can describe many of their properties in general terms.

The more exact criteria [2, 12] I adhere to can be summed up in two points. The first is that the total many-body wave function must have a cluster structure (or at least a large probability for finding a cluster component), i.e. we can talk about a core component and one or more halo particles. The second is that a large part of the wave function for the halo particles must be in the non-classical region. A reasonable definition of “large part” is that the probability f_h for being in the non-classical region should be at least 0.5, but where one draws the line between a halo system and a non-halo system is to some extent a matter of taste. The important point is that halo-specific features will be more pronounced the larger f_h is. It is of course of interest to know where the best examples of halos are found, but it will clearly also be interesting to see less developed halos and understand for a given nucleus what might prevent the appearance of halo features. Section 3 will look more into this question.

Note that there is no specific quantum number connected with halos. To get a sizeable tunnelling out into the non-classical region the system needs to be loosely bound, so the halo structure is in some sense a threshold effect, the important parameter being the separation energy S_h needed to remove the halo particle(s) from the core. The structure depends so sensitively on S_h that the most useful theoretical approach is to use few-body structure models that are tuned to fit S_h . More complex models that could also describe the internal structure of the core will typically not be able to reproduce S_h precisely and will then give wrong predictions for important quantities such as the total spatial size. A brief discussion of which theoretical approaches have been attempted was given in [14]. (When reading accounts of these more complex models one must be prepared to digest phrases such as “coupling to the continuum”, they sound impressive and tend to be connected to specific assumptions made in, or limitations inherent in, the model.)

2.1 Two-body Halos

To see in detail how halos appear, we shall look first at the wave function for the simplest two-body halos, consisting of a core and one neutral halo particle. With r being the distance between the two components, μ the reduced mass and l the angular momentum quantum number, the radial wave function at distances large enough that the potential can be neglected is simply proportional to $k_l(\kappa r)$, where the “inverse decay length” κ is given by $S_h = \hbar^2 \kappa^2 / (2\mu)$ and k_l is the spherical modified Bessel function. The latter is simple to express analytically,

$$\begin{aligned} k_0(x) &= \frac{\exp(-x)}{x}, \quad k_1(x) = \frac{\exp(-x)}{x} \left(1 + \frac{1}{x}\right), \\ k_{l+1}(x) &= \frac{2l+1}{x} k_l(x) + k_{l-1}(x). \end{aligned} \quad (1)$$

The outer wave function must be matched to the internal wave function to solve the problem completely. This is done numerically and it is a good exercise to try this for oneself.

For completeness we can also write down the momentum wave function. The Fourier transform of a wave function $\psi(\mathbf{r}) = R(r)Y_{lm}(\hat{r})$ is

$$g(\mathbf{k}) = \sqrt{\frac{2}{\pi}} i^l Y_{lm}(\hat{k}) \int_0^\infty R(r) j_l(kr) r^2 dr, \quad (2)$$

j_l being the usual spherical Bessel function. When, for $r > r_0$, we have $R(r) = B k_l(\kappa r)$ one gets the following contribution from the outer part

$$\begin{aligned} &\sqrt{\frac{2}{\pi}} \int_{r_0}^\infty B k_l(\kappa r) j_l(kr) r^2 dr \\ &= B \sqrt{\frac{2}{\pi}} \frac{r_0}{k^2 + \kappa^2} [k r_0 j_{l+1}(k r_0) k_l(\kappa r_0) - \kappa r_0 j_l(k r_0) k_{l+1}(\kappa r_0)]. \end{aligned} \quad (3)$$

For $r_0 \rightarrow 0$ (when the inner part is neglected completely) this becomes

$$B \sqrt{\frac{2}{\pi}} \frac{1}{k^2 + \kappa^2} \frac{k^l}{\kappa^{l+1}}. \quad (4)$$

However, in this limit the spatial wave function is only normalizable for $l = 0$ in which case it is simply the Yukawa wave function and $B = \sqrt{2} \kappa^{3/2}$.

The spatial extent of a system can be characterized by the radial moments $\langle r^n \rangle$. It is actually not necessary to know the internal wavefunction in order to derive how the radial moments scale as the binding energy goes to zero. The results [27] are:

$$\langle r^n \rangle \propto \begin{cases} S_h^{(2l-1-n)/2} & n > 2l-1 \\ -\ln(S_h) & n = 2l-1 \\ \text{const.} & n < 2l-1 \end{cases}. \quad (5)$$

Exceptionally large values can occur, in principle arbitrarily large values, if the binding energy is small and the angular momentum not too high. The probability distribution itself ($n = 0$) diverges for $l \leq 1/2$ (this is strictly speaking meaningless, it means that the normalization constant – B in the above notation – must go to zero) and the mean square radius ($n = 2$) for $l \leq 3/2$, so good halos will only occur for s and p waves.

If one is looking at an extreme halo system where most of the wavefunction lies outside of the potential it doesn't really matter what potential shape one uses (this was observed early on for the deuteron, where low-energy scattering can be described using only the scattering length and an effective range). For convenience many papers in the literature therefore use square well potentials. For consideration of less extreme systems one can of course employ more realistic potentials, but this is often not necessary: It turns out that for most of the interesting halo systems scaling laws apply, so we shall turn now to an explanation of what the scaling is and how to make use of it.

2.2 Scaling for Halo Systems

The key ingredient in deriving the scaling is to choose the appropriate scale R on which to measure a given system. We take it as the classical turning point, the distance where the potential energy is equal to the binding energy. If one uses a square well potential R is simply the well radius. The classical turning point is an obvious theoretical choice since it separates the “tunnelling region” from the region where the potential is large and the wavefunction could depend on details of the potential shape. We shall discuss below how one can find R experimentally.

The quantity most often used to measure the size of a system is its mean square radius (or the square root of it, the rms radius). It is also used here, but one should not use it uncritically for halos. To see why, consider an s -wave with negligible core radius where the wavefunction has the simple Yukawa form. The total mean square radius here is $\hbar^2/(4\mu S)$ (see e.g., [14]), but almost half the contribution to $\langle r^2 \rangle$ comes from the outer one third of the wavefunction (measured in terms of the integrated probability distribution). The outer tail contributes even more heavily, distances beyond κ^{-1} contain only 13.5% of the total wavefunction but contribute 27% to $\langle r^2 \rangle$. We shall come back to this point several times.

The dimensionless measure for the size is $\langle r^2 \rangle/R^2$, and the dimensionless measure for the binding energy S can be chosen as $\mu SR^2/\hbar^2$. In these units the pure Yukawa form gives a size equal to the inverse binding energy divided by 4. A potential of finite size will deviate from this behaviour at larger binding energies where the core size becomes noticeable, but calculations show [28] that different potential shapes gives almost the same deviation expressed in scaled units. This universal behaviour is one reason for using the scaling units, the other main reason is that it allows systems on widely different absolute scales (such as nuclei and molecules) to be compared. Figure 1 shows both

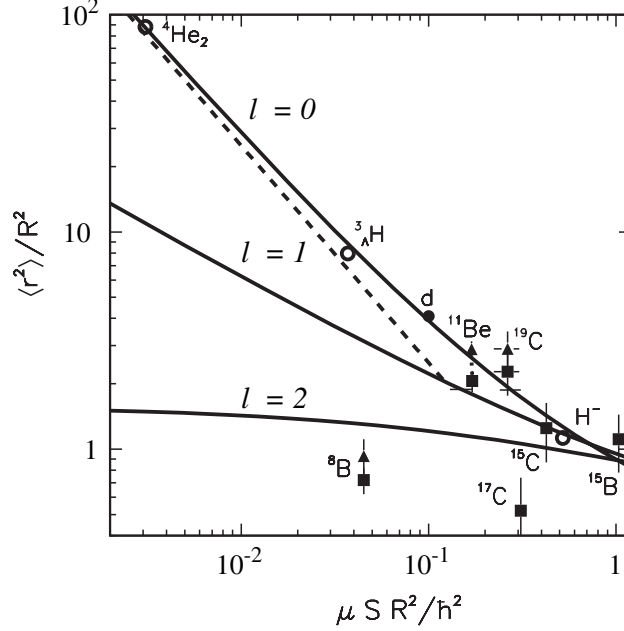


Fig. 1. Scaling plot for two-body halos. The mean *square radius* of the halo is plotted versus the binding energy in scaled, dimensionless units. *Open points* are theoretical values. *Filled points* are derived from experimental data as explained in the text, *squares* and *triangles* are from optical limit and few-body approaches, respectively. *Solid lines* are for pure *s*-, *p*- and *d*-waves, the *dashed line* corresponds to a pure Yukawa wavefunction

the Yukawa behaviour and the scaling curves for angular momenta $l = 0, 1, 2$. The thin horizontal lines on the $l = 0, 1$ curves just below $\langle r^2 \rangle / R^2 = 2$ indicate where there is 50% probability of being in the non-classical region. Good halo systems will lie above this value.

The theoretical values on the plot (reproduced from [29]) are for three systems outside the field of exotic beams: the H^- atom (an electron bound to the neutral hydrogen atom), the hypertriton (a Λ -particle bound to a deuteron) and the He dimer (the molecule consisting of two neutral He atoms). As the last two show some of the best halo systems are actually found outside of our field.

The experimental data are placed in the plot in the following way. All masses are taken from the 2003 Atomic Mass Evaluation [30]. Apart from this one needs to know the mean square radius $\langle r^2 \rangle$ and the scale R . Let us look first at the deuteron, the classical example of a loosely bound system. The size of the deuteron [31] is known very accurately from atomic spectroscopy. However, the value of R will depend on what potential shape one assumes. The point in the figure corresponds to a value of 1.93 fm as used in [32], but

slightly different values would be as justified and would correspond to moving the point parallel to the dashed line (the Yukawa curve). Two conclusions can be drawn from this. First, the position of a nucleus in the plot may be subject to some (small) uncertainty. Secondly, it is nevertheless important to use as trustworthy a value of R as possible since varying R will move the point on and off the scaling curve, even for s waves (unless one is dealing with as extreme systems as the He dimer).

For all other points the radii are deduced from measurements of total reaction cross sections performed at high energy, typically around 1 GeV/u. We shall look a bit more at these experiments in Sect. 5, but an earlier lecture [14] discussed these measurements in much detail explaining the Glauber model and the two versions used of it, the optical limit and the few-body model. The distinction is important for halo nuclei, where the few-body model will be more reliable. A compilation of cross sections and of radii deduced from them (in both approaches) can be found in [33]. Figure 1 employs these data except for the few-body model radii for ^8B and ^{11}Be that are taken from [34]. Analyses of other types of measurements have also led to deduced radii; these are in my opinion not as securely founded as the ones from the reaction cross sections and are therefore not included, but these other data of course can contribute in a decisive way in determining how much halo character a given nucleus has. We shall return to this later.

The mean square radius of the halo is given by:

$$\langle r^2 \rangle = \frac{m_{tot}^2}{m_c m_h} \langle r^2 \rangle_{tot} - \frac{m_{tot}}{m_h} \langle r^2 \rangle_{core} , \quad (6)$$

where the total mean square radii of the halo nucleus and the core nucleus enter, m_{tot} is the mass of the total system and m_c and m_h the masses of the core and the halo particle(s). The mean square core radius also determines the value of R if one corrects for the difference between the size of the density distribution and that of the potential. The relation used [29] is $R^2 = \frac{5}{3}(\langle r^2 \rangle_{core} + 4 \text{ fm}^2)$.

2.3 Three-body Halos

Many two-neutron halos are genuine three-body systems and are most efficiently described using relative coordinates between the three constituents. They are often Borromean [1], meaning that all two-body subsystems are unbound. The binding of the three-body system is then due partly to all three interactions being present at once, partly to the changed kinematics, see [35] for a thorough analysis. Jacobi coordinates (see e.g., [14]) are often introduced, but it is convenient here to go directly to the hyperspherical coordinates that are presented in detail in [1, 8] – the transformation is to one radial and five angle coordinates. The so-called hyperradius ρ and the corresponding scaling length ρ_0 are defined as

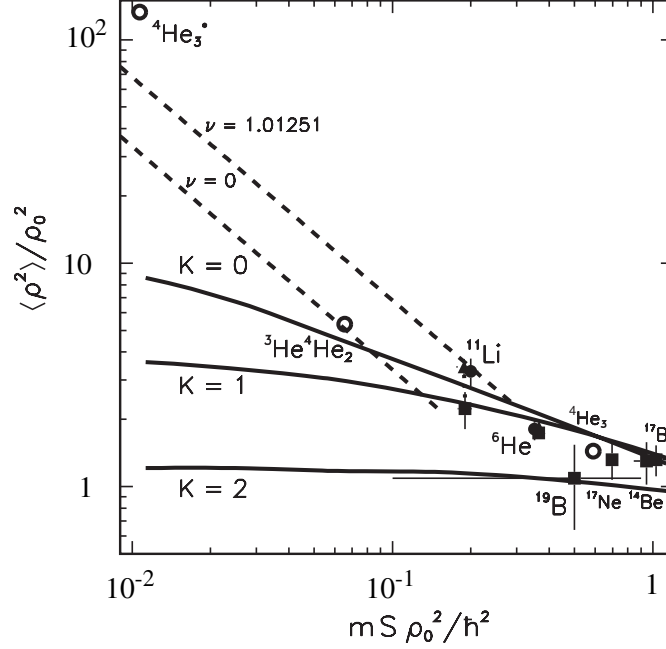


Fig. 2. Scaling plot similar to Fig. 1, but for three-body halos. *Solid lines* correspond to states with hypermomentum $K = 0, 1, 2$. The *dashed lines* indicate where the extreme Efimov states may appear, see the text for details

$$m\rho^2 = \sum_{i < k} \frac{m_i m_k}{m_{tot}} (\mathbf{r}_i - \mathbf{r}_k)^2, \quad m\rho_0^2 = \sum_{i < k} \frac{m_i m_k}{m_{tot}} R_{ik}^2, \quad (7)$$

where R_{ik} is the scaling length in the two-body subsystem, m_i ($i = 1, 2, 3$) are the masses of the three components and m_{tot} the total mass. The mass unit m is in principle arbitrary, it cancels in the ratio ρ/ρ_0 and with a scaled energy $mS\rho_0^2/\hbar^2$ (S now being the three-body binding energy, e.g., the two-neutron separation energy) it also does not enter there. Using these coordinates it is straightforward to construct a scaling plot for three-body halos, as done in Fig. 2. Scaling can not be expected to be as good for three-body systems as for two-body systems but is in practice working fairly well. The borderline for good halo systems is also here just below $\langle \rho^2 \rangle / \rho_0^2 = 2$.

The radial equation in hyperspherical coordinates (shown in [14]) is rather similar to a two-body radial equation. It contains a centrifugal barrier $(K + 3/2)(K + 5/2)/\rho^2$ where the hypermomentum K takes the place of the two-body angular momentum l . In contrast to the two-body system, K will not be a good quantum number, but it is often used for expansions of the full solutions and the behaviour of states with good K values is therefore indicated in Fig. 2. The structure of the radial equation implies that even for the lowest K value of 0 systems will be much less divergent than the most

extreme two-body systems ($K = 0$ has an “effective l ” of $3/2$). However, three-body systems have more degrees of freedom and correlations between the particles can in several cases give systems of larger sizes. A complete discussion of this has only been given recently [12, 36]; I will here only present the main conclusions and refer to the original papers for details.

First of all it turns out that the $K = 0$ curve underestimates the mean square radius a system can obtain. There is still a universal curve, but it lies above $K = 0$ by more than a factor 6 at very low binding energies and still about 20% at a scaled binding energy of 0.1. The “problem” is that enforcing $K = 0$ leaves too little freedom for how the relative distances in the system are distributed. Secondly, a two-body subsystem may be bound in which case the behaviour will resemble more a total two-body system, one body being composite. A good example of this is the hypertriton that was placed in Fig. 1 (examples of the transition from an effective two-body to a three-body system can be found in [36]). Thirdly, a peculiar effect will occur if two subsystems at the same time happen to have a large scattering length, both when they are just bound and just unbound. Many – in the extreme case infinitely many – three-body bound states of very large size will then exist, they are called Efimov states [8, 37]. Roughly speaking, the large spatial extent of the two-body systems allow bound three-body states to be built. If Efimov states turn up in two-neutron halo nuclei they would appear between the two dashed lines in Fig. 2, the ν parameter enters in the effective three-body potential.

The theoretical points in Fig. 2 are for various Helium trimers, molecules composed of three He atoms. In the system with three ^4He atoms the ground state is fairly bound whereas the first excited state is a Efimov state, it is actually a halo built upon a Efimov state [38]. The molecule with two ^4He atoms and one ^3He has a large ground state and no bound excited states. The experimental points for nuclei are based on high-energy reaction experiments as for two-body systems, the triangles again using radii derived from a few-body analysis [34] and the squares from an optical limit analysis [33]. The filled circles use radii derived from elastic scattering experiments on protons at 0.7 GeV/u [39, 40], a well-tested probe of density profiles. It is important to note that although the scattering data were of very good quality they actually did not allow to determine the density profile sufficiently far out that the $\langle r^2 \rangle$ was determined accurately, cf. the discussion at the beginning of Sect. 2.2. One has to assume a specific functional behaviour of the tail of the density distribution to extract a final value, the value in the figure (a total ^6He radius of 2.45(10) fm) uses an exponential tail as is appropriate for a halo system. Using a Gaussian tail the derived radius is smaller (2.30(7) fm). Incidentally, the deduced core radius fits very nicely with the significantly more precise ^6He charge radius measured recently [41] using laser spectroscopy.

2.4 Other Aspects of Halos

What about four-body halos? Or N -body? One can do a transformation similar to the one leading to the hyperspherical coordinates and obtain a generalized centrifugal barrier. We know that correlations between particles will probably be important and give rise to the formation of clusters inside the system, but in the absence of such clustering the system gets an effective l of at least $3(N - 2)/2$ and halos will therefore only appear for $N \leq 3$. Therefore, it suffices to look for systems composed of effectively two or three clusters.

In nuclear physics, we are mainly concerned with attractive potentials that have a short range. Several papers in the literature have looked at more general potentials, in particular potentials with a power-law behaviour at large distances. The division between short-range and long-range potentials is at the r^{-2} potential that has a somewhat pathological behaviour. For more information on these matters the interested reader is referred to the overviews in [12, 29].

Repulsive long-range potentials are relevant for nuclear halos, namely when the halo particles are charged. The larger the Coulomb barrier is, the more the tunneling into the nonclassical region will be suppressed, as an extreme example we do not expect α -decaying nuclei to have a halo structure even though they are actually unbound. One-proton and two-proton halos are thereby disfavoured and will always reside in the scaling plots close to (or below) the limit. The structure will therefore be more sensitive to details in the potential and one needs to look carefully at each case. General rules are harder to make, but a simple numerical example as given in [2] indicates that claims for proton halos much above a core charge of 10 should be scrutinized carefully.

Finally, there do exist rather general (and consequently not always too competitive) theoretical results for when two- or more-body systems are bound [42, 43] as well as for the sizes of halo systems [44, 45]. Those with an interest in such general analyses should also take a look on the monograph by Grosse and Martin [46].

3 Where Can One Find Halos?

Turning now to the question of where halos can be found, it is obvious from Figs. 1 and 2 that some molecular systems have very pronounced halo structure. There are not so many other physics subfields where halos are found, they seem rare in atoms although they have a theoretical chance of being found in negative atomic ions. When the structure of molecular halos can be studied experimentally in detail (this is yet to come) it will be interesting for detailed tests of our ideas on how halos behave, but in a way finding halo systems in molecules is “too easy” since the clustering, the first of the two

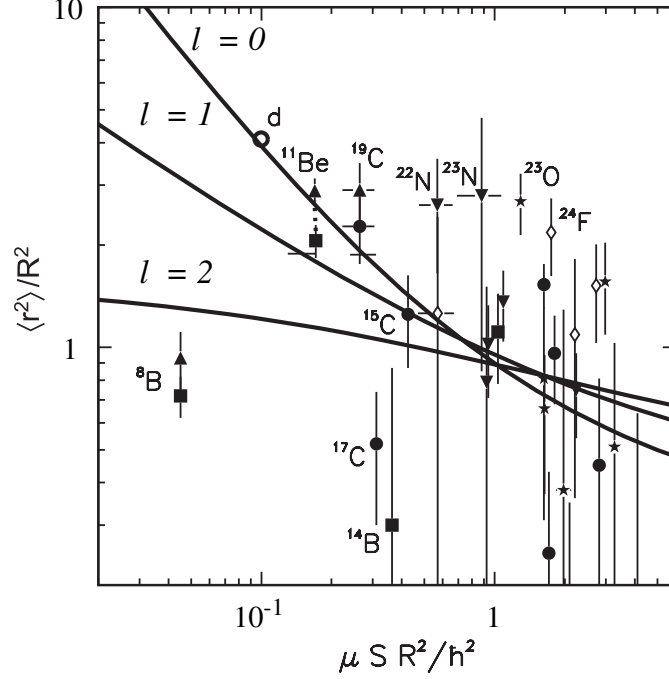


Fig. 3. Two-body scaling plot, as Fig. 1, but for heavier neutron-rich nuclei. The symbols denote Be and B isotopes (*full square*), C isotopes (*full circle*), N isotopes (*full triangle*), O isotopes (*star*) and F isotopes (*diamond*)

halo criteria mentioned in the beginning of Sect. 2, is almost automatic. In contrast, clustering is far from being trivial for nuclear systems.

To illustrate how things can go wrong for nuclei, Fig. 3 shows the two-body scaling plot for many neutron-rich nuclei based on systematic measurements of their radii [47]. In each case a “core plus one neutron” model is assumed to hold. Although the scaling curves ought to give upper limits for the size of a system, several nuclei with scaled binding energy above 0.4 violate this, among them nuclei very close to the neutron dripline such as $^{22,23}\text{N}$ and ^{23}O . The assumption that these nuclei can be described as simple two-body systems must fail.

As we approach ‘normal nuclei’ the inert core approximation is increasingly likely to be wrong. Configuration mixing can become important and would invalidate the simple model, the final nucleus would rather be described as a superposition of different core excited states with a surrounding neutron, and the spatial size of the excited states must be known to calculate halo radii. In the specific case of the “too large” nuclei shown in Fig. 3 it has been argued [48] that the core is enlarged. Even in the cases where a core plus neutron approximation still holds the core might become deformed.

Deformation by itself does not prevent halo formation [49], but changes in deformation will of course spoil the simple relation (6). Finally, the nucleus might be clustered, but with a different cluster division than core plus neutron.

The effects of configuration mixing can be estimated qualitatively once the level density is sufficiently high [50]. The point is that unless conserved quantum numbers prevent it (and there are no halo specific quantum numbers) different states in a nucleus will mix. For a given excitation energy in a given nucleus one can estimate the level density and the size of the coupling matrix element between states, taking into account the dilute wavefunction of the halo state, and it is then a standard exercise to evaluate how much the original halo wavefunction is spread out and how much remains in one state. It turns out that mixing most likely is strong except for the lowest few states in a given nucleus, so halos will occur mainly close to the groundstate, and since the binding energy should be low this restricts us to nuclei close to the driplines except possibly in the very light nuclei.

3.1 Where are the Driplines?

The most remarkable feature of the nuclear chart (when one compares to other physical systems) is that there are only a limited number of combinations of neutron and proton numbers that give a particle stable nucleus. Figure 4 shows two theoretical estimates of where nuclei are bound,

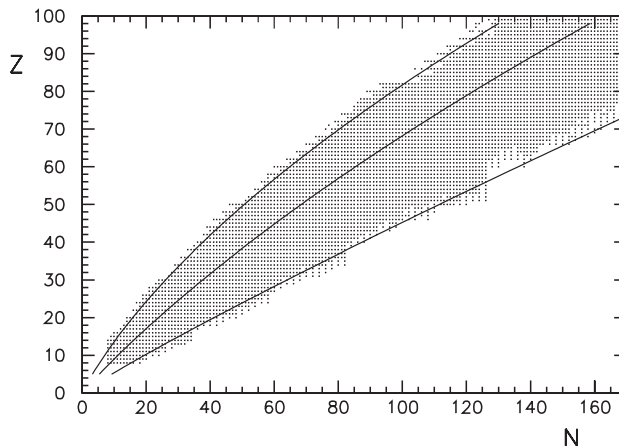


Fig. 4. The extent of the chart of nuclei according to two theoretical estimates. The points mark nuclei that are stable towards prompt particle emission according to the Hartree-Fock calculations in [51]. The lines show the valley of beta-stability and the proton and neutron driplines according to a simple version of the liquid drop model [32] (excluding pairing to get smooth curves). Some of the heaviest nuclei shown will be unstable with respect to prompt fission

experimentally we are only about half-way in having seen (not to mention studied) these nuclei. Fortunately, there is a recent good review [52] on the limits of particle stability. As shown there in detail we have by now reached the proton dripline for all odd Z elements up to $Z = 91$, whereas for even Z there are still missing isotopes for Z in the range 32–64 and above 82 (actually, as explained in more detail in [52], the exact position of the proton dripline is not always known since nuclei that are just unbound to proton emission will still be long-lived and therefore observable in experiments). On the neutron-rich side the dripline is known up to about Fluorine ($Z = 9$), but it is clearly more appropriate to discuss its position in terms of N : we know it up to $N = 20$ and might have reached it up to $N = 27$, above this value the expected position rapidly go beyond what experiments can do presently.

The position of the neutron dripline must therefore be taken from theory for all but the lightest nuclei. The two estimates in Fig. 4 fit reasonably well to each other, but the question is how reliable the extrapolation from known nuclei to the dripline is. Some comparisons of different models give the impression that the position is fairly well established (at least when looking for a given value of N rather than a given value of Z [52]), but extrapolations of masses tend to diverge as one leaves the region where they have been fitted, examples of this can be found in many references (e.g., [53]). The basic reason is that the physics effects that become important for determining the exact position of the neutron dripline give very small variations for normal nuclei and therefore are hard to determine with sufficient accuracy. Two examples of physics effects that could be important are the isospin-dependence of terms that are constant in the simple liquid drop model (a good discussion is given in [54] of the droplet model where the surface term is modified) and the neutron skin that we shall return to in the next subsection. A recent overview of the theoretical models used today is given in [53].

The two things to remember from this is firstly that one should still be sceptical about extrapolations all the way to the neutron dripline and, secondly, that if we are only interested in getting a picture of how very neutron rich nuclei could behave it is sufficient to use the very simple liquid drop model.

3.2 Established and Future Nuclear Halos

We can now return to the question of where halos are positioned in the chart of nuclei. In the spirit of the conclusion just made the first step is to look at where simple models would place halo candidate states, i.e. where we would have s - and p -states at zero binding energy. Such an analysis was presented in [50] using a simple nuclear potential (even disregarding the spin-orbit splitting of p -states) and the results are shown in Fig. 5. Both one- and two-neutron halos need to include substantial amounts of s - or p -wave in order to be large, so the favoured regions for finding halos are where the dotted and dash-dotted lines in the figure cross the dripline. Note that the

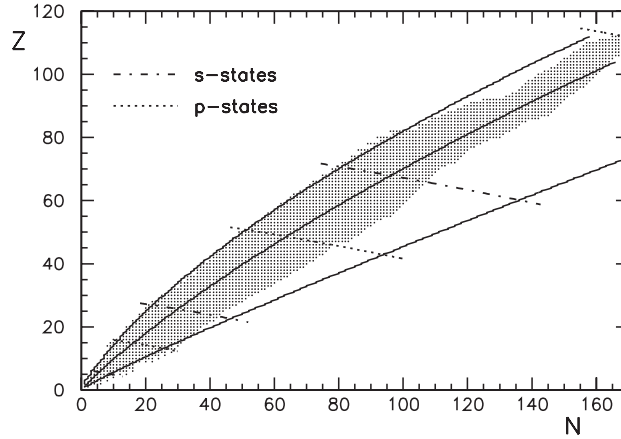


Fig. 5. The chart of nuclei as in Fig. 4 except that points here mark nuclei where the masses are experimentally known or extrapolated from known masses [30]. The lines show where s - and p -wave neutron states would be situated at zero binding energy in a square well estimate [50]

first s - and p -states (the orbits up to $N = 8$) are not shown. The detailed results of this simple estimate cannot be expected to be very reliable, but the main features will be robust. We should therefore expect halos to occur in specific mass regions rather than being spread out uniformly along the dripline.

A selected list of some halo states is given in Table 1. More extensive tables that also include halo candidates can be found in [12]. The relative nucleon-core angular momentum l is given for each state; for two-nucleon halos, l denotes the dominating component. Values in parenthesis are not firmly established. Not all separation energies are well known; the uncertainty is given for the nuclei where it is larger than 1 keV. The deuteron is a special case, both due to its small reduced mass and to it being a “true” two-body system (as long as we neglect the quark substructure). Among the other one-nucleon halos ^8B , ^{11}Be and ^{19}C were shown in Fig. 1. The last two seem good halos whereas the proton halo ^8B is somewhat smaller, partly due to it being a p -wave. The excited states in ^{11}Be and ^{17}F probably have values of $\langle r^2 \rangle / R^2$ just above and below 2, respectively. All listed two-nucleon halos were shown in Fig. 2, here ^6He and ^{11}Li are good halos, whereas ^{14}Be and ^{17}Ne are somewhat smaller.

There does not seem to be any really large proton halo, but the listed systems (and others: ^{26}P was much promoted at a certain stage [55]) are of course interesting to compare to neutron halos of various sizes. Halos in excited states have only been considered occasionally and are much harder to access experimentally, but could give a valuable insight both by testing

Table 1. Halo states. For each state the excitation and separation energies and the angular momentum of the halo particle(s) are listed

Nucleus	E_x (MeV)	$S - E^a$ (keV)	Configuration	l
d	g.s.	2225	n+p	0
^8B	g.s.	138	p+ ^7Be	1
^{11}Be	g.s.	504 ± 6	n+ ^{10}Be	0
^{11}Be	0.32	184 ± 6	n+ ^{10}Be	1
^{17}F	0.50	105	p+ ^{16}O	0
^{19}C	g.s.	580 ± 90	n+ ^{18}C	0
^6He	g.s.	972	n+n+ ^4He	1
^{11}Li	g.s.	300 ± 19	n+n+ ^9Li	0,1
^{14}Be	g.s.	1260 ± 130	n+n+ ^{12}Be	(0,1,2)
^{17}Ne	g.s.	973 ± 27	p+p+ ^{15}O	(0)

^a From [30].

our understanding of where halo formation can occur and by giving more examples of halo states.

There might not be that many more good halo states to be found in the near future. Proton halos above ^{26}P are unlikely to exist. The nuclei at the neutron dripline with $N = 15, 16$ (where the second s -orbit normally would appear) might show halo structure, but the only one with a potentially really low binding energy is ^{22}C whose mass is uncertain by 1 MeV. One has to recall that since the scaled energy contains the factor SR^2 the limit for where good halos occur is not constant in MeV, but rather decreases as $A^{-2/3}$. Further up the dripline the next candidates will have a p -orbit that in the standard filling of the shell model appears above $N = 28$, a region where we are still not certain about the dripline position. However, as many textbooks will tell you (see e.g., Fig. 9.13 in [32]) the ordering of nuclear orbits is not fixed as we move around in the nuclear chart, there is a tendency for low angular momentum orbits to go down in energy as one gets closer to the threshold (due to tunneling out from the nucleus, the same effect that gives the halo structure). On top, the residual nuclear forces can move orbits around to the extent that even magic numbers are changing, see e.g., the discussions in [10, 56]. It is already clear from the examples in Table 1 that the second s -orbit can play a role in the whole region $N = 8$ –16 so one should look out for the second p -orbit before $N = 28$. An optimist would also search for new excited state halos.

The way I have been discussing halo nuclei there is a clear distinction between them and nuclei with a neutron skin, for halos it is the outer tail that expands, for neutron skins it is the bulk of the neutron matter that extends beyond the proton matter [2, 3, 4]. When going through the literature one must be aware that not all authors agree to this distinction, partly for historical reasons. Some just look at neutron radii (e.g., [57] and references therein) or seem satisfied with the neutron density being larger than the proton density in the tail of the nucleus no matter how far out this happens. This is a questionable approach, even stable nuclei will eventually have a neutron density larger than the proton density simply because of the extra Coulomb barrier for the protons. It can be measured with antiprotons [58] that annihilate on the neutrons or the protons in the outer surface, roughly 2.5 fm beyond the half density charge radius. The composition of the nuclear density in this extreme situation is of course interesting by itself, but when we discuss halos a more restricted view is required: the facts that neutrons outnumber protons in the surface or that the system is spatially large are facets of halos, but they are not exhaustive characteristics. The clustering implies that there is a large amount of single-particle (or few-particle) behaviour. We now turn to these other consequences of the halo structure; they will also illustrate the ways halos can be identified and probed.

4 Probes of Halo Structure

The first step in experimental investigations of a nuclear halo system is to create it, the deuteron (that occurs naturally) being the only exception. Production of radioactive nuclei has been a recurrent lecture subject at the Euroschools. There are two basic routes, in-flight separation of projectile fragments and Isotope Separation On Line (ISOL) [59, 60]. An earlier lecture described in-flight separation in detail [61], here the separation happens faster and one immediately has a secondary beam that can be used in nuclear reaction studies. Most halo experiments have taken place at in-flight facilities, but ISOL facilities [62] are advantageous for precision experiments and for reaction studies at very low energy.

A list of probes of halo structure would to a large extent be the list of probes of nuclear structure. We can use all three interactions, the strong, the electro-magnetic and the weak. The last two are well understood on the nuclear scale and can be applied immediately for investigations of halos. In contrast much theoretical work has had to be put into understanding the reaction mechanism in nuclear reactions involving halos. The nuclear reactions and β -decay probes are discussed in more detail in Sects. 5 and 6, respectively. Here some general comments will be given that naturally leads to a discussion of electro-magnetic transitions.

Whether we are creating the halo state during the experiment or probing (destructively or non-destructively) a halo coming in a secondary beam at a

radioactive beam facility, we can model many of the probes by an operator sandwiched between an initial and final state wavefunction. To give concrete examples, for a two-body halo with charges $Z_h e$ and $Z_c e$ for the halo and the core the electric multipole operators are [27]

$$\mathcal{M}(E\lambda, \mu) = \left[Z_c e \left(\frac{m_h}{m_h + m_c} \right)^\lambda + (-1)^\lambda Z_h e \left(\frac{m_c}{m_h + m_c} \right)^\lambda \right] r^\lambda Y_{\lambda\mu}, \quad (8)$$

i.e. a weighted sum of the charges times the usual one-particle multipole operator for the relative coordinate. The factor r^λ enhances the large distances. For proton halos all multipoles are enhanced, whereas for neutron halos with $Z_h = 0$ the mass factors in (8) reduce the effect drastically when $\lambda \geq 2$. The magnetic transition operators contain factors of $r^{\lambda-1}$ and enhance less. The explicit expression for the M1 operator for a two-body halo is

$$\mathcal{M}(M1, \mu) = \mu_N \sqrt{\frac{3}{4\pi}} \left[g_s s_\mu + l_\mu \left(\frac{Z_h}{A_h} \frac{m_c}{m_c + m_h} + \frac{Z_c}{A_c} \frac{m_h}{m_c + m_h} \right) \right], \quad (9)$$

where μ_N is the nuclear magneton, s_μ and l_μ are the μ^{th} spherical components of the halo spin and angular momentum, A_h and A_c are the two mass numbers and g_s is the spin g-factor for the halo particle. In both the above cases the core is treated as inert and structureless so the contributions from the intrinsic core degrees of freedom are omitted.

When the operator directly enhances the tail of the wavefunction the matrix element for halos can become much larger than for normal states. The enhancement is illustrated in the left panel of Fig. 6 using the simple square well model for a two-body halo. The expectation values of $\langle r^n \rangle^{1/n}/R$ (the root mean square radius for $n = 2$) are shown versus the probability that the halo particle is outside of the potential well of radius R . This probability will go to 1 for very loosely bound s -waves, but remains less than 2/3 for p -waves. Nevertheless the expectation values all diverge, the more so the larger the power n . (For the root mean square radius we have already discussed this tail enhancement for the pure Yukawa wavefunction in Sect. 2.2.)

The same conclusion is of course reached when there is a halo both in the initial and the final state. The standard example of this is the E1 transition between the two bound states in ^{11}Be [23] that is unusually large.

If we consider now matrix elements M_n of r^n between halo states and a normal well-bound nuclear state there will be two conflicting effects: the fact that the halo extends further out would increase the contribution from the large distances, but since at the same time the probability of being inside the nucleus diminishes, the contribution from this part decreases. The final outcome is shown in the right panel of Fig. 6, where the bound state is chosen to have $\mu S R^2 / \hbar^2 = 1$ corresponding to about 5 MeV binding for nuclei with $A \simeq 10$. The main thing to note is that it actually takes quite a while even for an s -wave halo state to “disappear” from a nucleus (a part of the p -wave

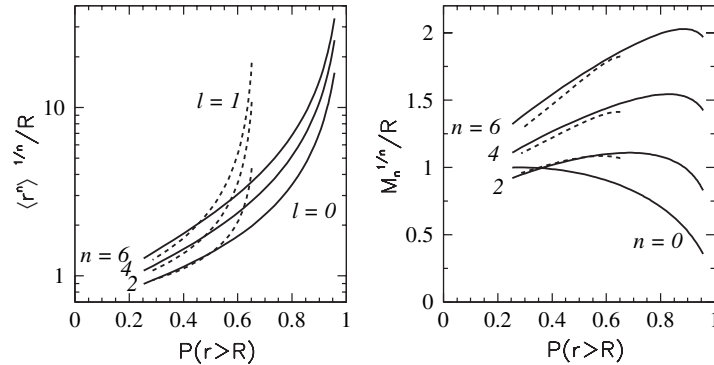


Fig. 6. Radial expectation values plotted versus the probability of being outside the potential for a two-body halo in s - (full lines) or p -wave (dashed lines), estimated in a square well model. Left panel: The n^{th} radial moment of the halo state. Right panel: The matrix element of r^n between the halo state and a well bound state with $\mu SR^2/\hbar^2 = 1$. In both cases the n^{th} root scaled by the well radius R is shown

would always remain). Even the case $n = 0$ with no enhancement of the outer part, where the matrix element is simply the overlap between the two wavefunctions, gives a rather large overlap until we get to extreme halo sizes. Up to this point also the other matrix elements increase, although of course not as drastically as the radial moments of the halo itself.

The calculation just presented is very schematic in that it assumes a very simple structure of the bound state wavefunction, but the conclusions are general since they show the importance of the tail behaviour when discussing halos. As stressed earlier it is dangerous to use a “classical interpretation” of the rms radius when we are dealing with halos. Due to the very extended tail a halo system can have a large rms radius and at the same time an appreciable overlap with core states. This point is elaborated a bit more in [27].

Finally, let us consider the case where a halo state is connected to a continuum state. If the halo state is the final state in a reaction this could correspond to radiative capture of protons [22] or neutrons [27, 63]. Here the enhancement of the halo is felt fully, in particular at low energies of the incoming particle when the extent of the continuum wavefunction is much larger than the nuclear size. The inverse process, the halo in the initial state, would be Coulomb break-up. It shows the same halo enhancements and is the experimentally doable one for neutron halos. There is a large literature on Coulomb processes and loosely bound systems, and there is here only space to refer to some reviews [12, 64, 65, 66, 67]. A key concept here is the distribution of electromagnetic strength in the nuclei, a recent paper has shown how one from an effective-range approach can derive an almost universal expression valid for one-neutron halo nuclei [68]. These distributions will have a

strong peak just above threshold (see e.g., [66]) from the halo single-particle contribution – in the early days there were speculations as to whether collective resonances played a role (a so-called “soft dipole resonance”), but we know now [69] that it is a single-particle effect. That the halo contribution has to appear at low energy was clear from the beginning [26] due to sum rules. As an example the total E1 strength for the halo scales as $\langle r^2 \rangle$ whereas the sum of the energy-weighted strength is a constant, so the larger the radius the lower the strength has to reside (for more details, see e.g., [70]).

Before turning to nuclear reactions, I should point out that we also need precision measurements of masses of halo nuclei and that data on their ground state properties often are very valuable for tests of models of their structure. Some specific examples can be found in [7, 41, 53], but readers interested in these types of experiments are encouraged to browse through the other lectures in this series, see in particular [71].

5 Reaction Studies of Halos

Reactions in inverse kinematics (i.e. the halo nucleus is the beam) at intermediate to high energy (above 20 MeV/u) have so far been the main tool for exploring nuclear halos. The earlier lecture by Al-Khalili [14] treated reaction models for these processes in detail and his lecture should be read along with the present section if one wants a complete picture. I shall focus more on experimental questions and on the (simple as well as sophisticated) methods of analysis and data presentation that are used. For a final interpretation we will often be in the hands of the theoreticians, but I shall present some simple models and point out possible pitfalls in their use. For more details one can consult several reviews that among them cover both the experimental and the theoretical situation [3, 4, 5, 6, 10, 11, 12, 13, 14, 15, 16, 17, 18, 19, 20, 65].

We shall first discuss the “simple” experiments one can do with beams of about one ion/sec, then go to more advanced experiments that need higher count rates and end with a brief look at possibilities with low energy beams and future machines.

5.1 Some Experimental Considerations

Radioactive beams have low intensities and reaction experiments have to be designed accordingly. A compromise has often to be found between resolution and counting rate so the targets employed become quite thick. This gives higher counting rates, but will, due to energy loss in the target and straggling, worsen the resolution in angle and (unless the target itself is a detector with position resolution: an active target) energy too.

Unless the target is very thick we can neglect the decrease of beam intensity as it passes through the target. The total reaction probability P in a given target is then given by $P = 0.6022\sigma d/A_t$ (the factor in front is from the

Avogadro number) where d is the thickness of the target measured in g/cm^2 , σ is the total cross section in barns and A_t is the atomic weight of the target in units of u. The formulae for specific channels and for differential cross sections are similar. The total reaction probability is of order 1% for cross sections of order one barn and targets about one g/cm^2 , but the probability will, due to the factor A_t^{-1} , decrease for heavier targets unless the target thickness (or the cross section) increase drastically. This will normally not be the case and one therefore must compensate with longer running times for heavy targets.

The practical limit to the target thickness is given partly by the energy loss in it, and partly by the angular straggling. The ranges and stopping-powers for heavy ions in various materials are listed in the tables [72] or can be found using the SRIM [73] software. In many applications one can approximate the range by $R = a(E/A)^\alpha + b$ (the determining factor in stopping – apart from the charges – is the ion velocity, which is why it is the energy per nucleon that enters). The power α depends (for light ions) very little on the ionic charge and only slightly on the target material and is roughly 1.75. The factor a is directly proportional to the ion mass so that light isotopes are stopped faster than heavy ones. If the energy loss ΔE and the total energy E of a particle is measured, one can construct a “particle identifier” $PI = (\Delta E + E)^\alpha - E^\alpha$ that can be used to separate different isotopes. One normally does not detect where in the target the reaction took place. Since a beam particle and the reaction products will lose energy differently, this uncertainty leads also to a smearing (on top of the one coming from beam spread and from the reaction itself) of the energy of the outgoing charged fragments.

Multiple scattering in the target will also deteriorate the resolution. The deflection of the beam due to multiple scattering can at small angles often be described by a Gaussian, but the outer tails (of relative intensity a few percent) will be larger than given by a Gaussian shape. A brief account can be found in [74]. There is still active research going on in the field of stopping of heavy ions at energies up to several GeV/u. A recent experimental overview is given in [75]. Practitioners and users of energy loss data should note that there will soon be a new evaluation out, a book on the theory behind has already appeared [76] and a large report is in print [77].

We are dealing with radioactive beam experiments and the intensities are sometimes very low. If the total number of events at the end of an experiment is low one needs to be very careful about the statistical analysis, since the assumption behind standard analysis methods that data are distributed as a Gaussian (implicit e.g., in the χ^2 method) breaks down. This is clearly an important subject, but it is not possible here to give all the necessary details. A brief useful overview of statistics is given by the Particle Data Group [74], more details on analysis methods to be used in nuclear physics for low count numbers are given in [78]. One specific danger of having few counts is that it can be harder to separate signal and background. In statistics

many techniques have been developed to deal with this situation. Such robust methods have been applied, at least in one case [79], for the analysis of shapes of the momentum distributions, which we shall discuss in a short while.

5.2 First Generation Experiments

The first reaction observable to measure was the total interaction cross section, a review on such experiments is given in [33]. From this nuclear radii can be extracted, as explained in [14]. In brief, at high energies the collision happens so fast that the basic process is that individual nucleons in the projectile interact with individual nucleons in the target nucleus. Since the nuclear interaction is strong, nucleons will “shadow” each other and the total reaction probability therefore depends on how much the nucleons are “spread out” and can, through model calculations, be converted to a value for the radius. This is in contrast to weak interactions – a neutrino would not see the difference between a halo nucleus and a normal nucleus, only the total number of nucleons.

For heavier targets the Coulomb interaction will also be important and must be included in the calculations. The resulting target dependence gives another way of checking the structure of the projectile. The low-lying strength mentioned at the end of Sect. 4 that is the most obvious halo contribution would actually be felt stronger at lower beam energies [67]. We return in Sect. 5.4 to low reaction energies where a description in terms of nuclear potentials is to be used.

The total interaction cross section can be measured by simply detecting the disappearance of the beam, the reasoning being that since no (or at least very few) bound excited states exist in halos any excitation leads to break-up. By looking at what emerges from the collision one obtains the cross sections for individual channels. One would expect that the major part of the effects from the halo appear in the “halo removal” channel. In accordance with this it has been shown [80] that good halo nuclei satisfy an approximate additivity relation

$$\sigma_{tot}(\text{halo state}) = \sigma_{-halo}(\text{halo state}) + \sigma_{tot}(\text{core}) , \quad (10)$$

where σ_{-halo} is the halo removal cross section. Independent confirmation comes from looking at channels that involves reactions of the core, e.g., charge changing reactions [81], when going through an isotopic series, such cross sections tend to vary slowly whereas the total cross sections show the enhancement from the halo.

As an illustration, Fig. 7 shows state-of-the-art calculations compared with data (taken from [82]) for ^{11}Li reacting on C and Pb targets. Apart from the total interaction cross section the separate contributions from two-neutron removal reactions and core break-up reactions are shown. Note that for the Pb target the large increase at low beam energies that comes from the contribution of Coulomb break-up to the two-neutron removal channel.

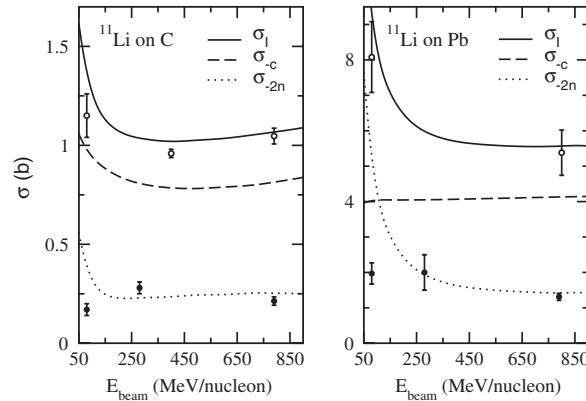


Fig. 7. Total interaction cross sections for ^{11}Li reactions on C and Pb targets shown as a function of beam energy. The contributions from two-neutron removal and core break-up are shown separately. The figure is taken from [82] where the theoretical calculations are described in more detail

Considerable effort, starting with one-neutron and one-proton halo nuclei, has gone into understanding one-nucleon knockout processes [83]. For the detailed interpretation of the data it is very important to include the effect of particle bound excited states in the remaining nucleus (the core for a halo nucleus). If the configuration mixing mentioned in Sect. 3 is large one would expect core excited states to be prominent in the one-nucleon knockout, but such final states can also be caused by knockout in the core itself and corrections must be applied to isolate the halo effects. A similar situation can of course be expected for two-nucleon halos.

Turning now to the second “classical” observable for halos, momentum distributions of particles emerging from a reaction, we have gone through a corresponding gradual refinement of understanding. The main emphasis has been on the neutrons and the core in the “halo removal” channel, but neutrons emerging from reactions where the core break up have also been considered [84] (in this case there can also clearly be a contribution of neutrons from the core, but their typical momentum is likely to be significantly higher). At first, the influence of the reaction was neglected, in which case the three-dimensional momentum distribution for an extreme halo is simply the square of the momentum wavefunction (4). For an s -wave this is a Lorentzian function that, including the Y_{00} part and expressed in terms of $k = p/\hbar$, is

$$\rho_3(\mathbf{k}) = \frac{1}{\pi^2} \frac{\kappa}{(k^2 + \kappa^2)^2} . \quad (11)$$

Normally one does not display this three-dimensional distribution, but rather the projection on one coordinate. Integrating in Cartesian coordinates over two components gives the following results, relevant e.g., for the longitudinal

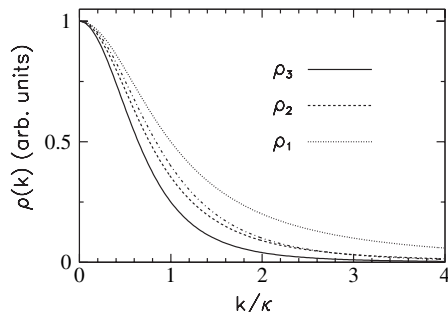


Fig. 8. The ρ_i distributions for a Lorentzian are shown as a function of k/κ , the extra curve (*dash-dotted line*) is the ρ_3 distribution modified for finite (here 50%) transmission. The distributions have been normalized to coincide at zero momentum. The difference in shape and width is seen clearly

component of the momentum:

$$\rho_1(k_{\parallel}) = \frac{1}{\pi} \frac{\kappa}{k_{\parallel}^2 + \kappa^2} . \quad (12)$$

The two distributions, ρ_3 and ρ_1 , have a different shape, which implies that experiments with different acceptance of transverse momentum in principle will measure different longitudinal momentum distributions! This effect is illustrated in Fig. 8, which also shows the ρ_2 distribution (integration over one Cartesian component) and the longitudinal distribution obtained if only the central 50% are transmitted in the transverse dimension. That this effect survives also when effects of the reaction mechanism are included has been shown both theoretically and experimentally [79, 85, 86]. That does not mean that theoreticians now need to become experts on experimental acceptance effects, but they at least should be aware of this possible experimental “distortion”.

In the approximation where reaction mechanisms are disregarded we see that halos give a narrow momentum distribution with width parameter $\hbar\kappa$, i.e. the less bound the system, the larger it is in space and the narrower in momentum space. However, the reaction does influence the measured momentum distribution and (at least) the following three effects must be considered. Firstly, momentum is transferred between target and projectile, so some of the outgoing fragments will have changed momentum. One often neglects this, since the unperturbed fragments are so localized in momentum space (a distribution ρ_3 around the beam momentum) that at least perturbed neutrons are unlikely to contribute much here. The effect is of course felt stronger for neutrons than for the core. A good example of this effect is the diffractive dissociation illustrated in [4]. Secondly, since the core reacts strongly with any part of the projectile that passes it, only the part of the wavefunction that goes clear of the target will contribute to the observed spectrum (this

is sometimes called “shadowing”, and differs from the effect mentioned at the beginning of this section). Finally, if more than one fragment emerges from the collision with small relative momentum, final state interactions are likely to play an important role. This is important in particular for detailed description of two-neutron halos.

If the halo state is large the first two effects will not be prominent, but for less pronounced halos the “shadowing” can become a problem [87, 88]. For instance, in the halo removal channel it would remove the part of the wavefunction where the halo and the core overlaps. Thus the resulting distribution tends to become narrower, and one cannot simply convert momentum width to a value of κ . Inclusion of other observables in the analysis, such as the cross section for the halo removal channel, is needed to identify good halo states. A good example of how to use this procedure for a systematic survey of many neutron-rich nuclei is the experiment by Sauvan et al. [89, 90], these data were shown already as Fig. 6 in [14]. An example of a one-neutron halo where these effects are small is ^{11}Be , a recent complete experiment on this nucleus was performed at GSI [91]. One needs here instead to consider contributions from core excited states as discussed above, see also Sect. 4 in [15].

An example for a two-neutron halo is shown in Fig. 9. The core momentum distribution is shown for ^{11}Li breaking up upon C and Pb targets. The transverse distributions are widened with respect to the longitudinal ones since contributions from the in- and outgoing parts of the reaction tend to add in the transverse direction and cancel in the longitudinal. As mentioned above Coulomb break-up is very important on the heavy targets. The main

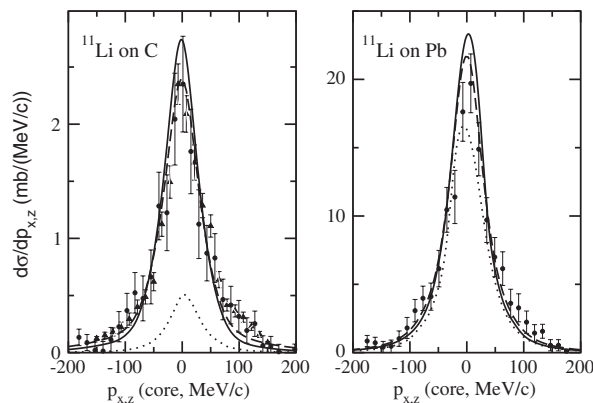


Fig. 9. Momentum distributions of ^9Li emerging from reactions of ^{11}Li on C and Pb targets. The theoretical distributions are from [82], both longitudinal (*solid line*) and transverse (*dashed line*) distributions are shown, as is the contribution (*dotted line*) where the core is reacting with the target. The experimental distributions [92], taken at 280 MeV/u, have been scaled to the calculations

Coulomb contribution actually comes at rather large impact parameters, the small kick the core receives is, even at large distances, sufficient to dissociate ^{11}Li . This is in contrast to nuclear break-up that takes place at the nuclear surface.

5.3 More Advanced Methods

Luckily, technical advances have during the last two decades given a steady increase in the available intensities of radioactive beams and simultaneously detection systems have become more complete in coverage. The combined effect of this progress is that it is now possible to do coincident detection of several outgoing fragments and get sufficient statistics for detailed studies.

A particularly important extracted quantity is the invariant mass m^* of a system. An example could be break-up of ^{11}Li into ^9Li and two neutrons, where the interesting systems to study would be $^9\text{Li}+n$, $n+n$ and $^9\text{Li}+n+n$ depending on which particles actually are detected in the reaction. Given a number of particles with total rest-mass (i.e. with threshold) $M = \sum m_i$ the invariant mass is to a good approximation (when the excitation energy is small compared to M) given by

$$m^* - M = \sum_{i < j} \frac{E_i E_j - m_i m_j c^4 - \mathbf{p}_i \mathbf{p}_j c^2}{M c^4}, \quad (13)$$

where E_i and $\mathbf{p}_i$ are the measured energies and momenta of the particles. Defined in this way, the invariant mass is always positive. *If* (and this is a big “if”) the system on its way to detection goes through a resonance, this resonance will be seen in the invariant mass spectrum. In our example the relevant systems would be ^{10}Li , the di-neutron and excited (continuum) states in ^{11}Li . Some people interpret even high-energy reactions in terms of intermediate resonances [93], but this seems hard to defend in view of the very short timescales involved [94] and runs against the standard reaction models used [12, 14, 65]. (It is interesting to note that the “intermediate resonance” approximation nevertheless for some nuclei might give a fair description. Quantum mechanically one can always insert a full set of intermediate states, and an approximation based on only one or two state can be good if these have a large overlap with the final state. As explained in [94] this will indeed tend to be the case for two-neutron halos, since it is the same interactions that are at play in the full halo system and in the subsystems.) A technical remark on resonances: if an s -wave state is positioned so low in energy that it overlaps with the threshold to bound states, its “resonance shape” will be essentially featureless with a monotonic decrease from the threshold, and it is called a virtual state rather than a resonance (see [95] for a deeper explanation). Whatever the reaction mechanism, the invariant mass spectrum will show how much energy is going into the breaking-up of a nucleus or into a specific subsystem, so it is a useful quantity to consider.

Just as a three-body state has more degrees of freedom than a two-body state, there are also plenty of observables to look at in an experiment with three particles in the final state. One standard tool that can be used to get an overview of the data is Dalitz plots as used in particle physics [74]. (Essentially the same information comes out in a two-dimensional plot of the energy of one particle versus the energy of one of the others, but the specific representation in a Dalitz plot has “nicer” overall properties.) Again, one should be careful about interpreting structures seen in a Dalitz plot as resonances, but the effects of final-state interactions can be seen as applied in an experiment at GANIL [18, 96]. Other new methods applied by the same group is two-neutron interferometry [18, 97] that with care can be used to extract the neutron-neutron distance in the halo.

Another possible observable is shown in Fig. 10. One way of interpreting this plot is as the angular distribution between the two halo neutrons, one direction being given by the core-n relative momentum and the other by the direction of the other neutron (in nuclear break-up where typically only one neutron emerges, the last direction is assumed to be aligned with, but opposite to, that of the remaining system). Such a distribution will in a naive picture simply be given by the relative angular momentum. The striking thing about the distribution shown is that it is asymmetric. This is only possible by having two simultaneous components of opposite parity and thus proves directly that the ^{11}Li halo contains both s - and p -wave components [98]. The more complete calculations (actually performed before the measurement) shown confirm this interpretation.

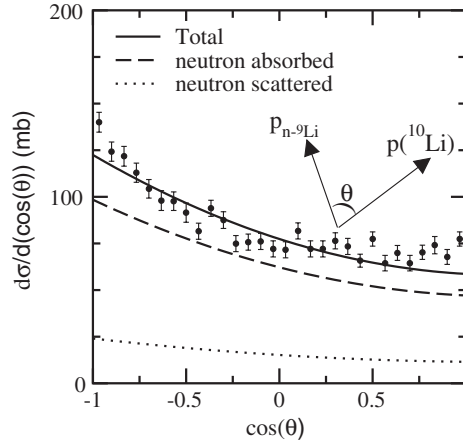


Fig. 10. Angular distribution extracted from break-up of ^{11}Li on a C target, as shown in the inset θ is the angle between the ^9Li -n relative momentum and the centre of mass momentum of ^{10}Li . The data are from [98]. For the theoretical prediction from [82], components from different reaction mechanisms are shown

The level of sophistication reached by experiments today is such that one can even combine different analysis methods, so that the angular distributions just mentioned can be extracted for different energy bins in the invariant mass spectrum [99]. It can be shown in this way that for both ^{11}Li and ^{14}Be the dominant contribution just above threshold in the subsystems ^{10}Li and ^{13}Be will be s -waves, whereas p -waves (and d -waves for Be) enter at higher energies. A detailed reaction modelling of such data has not been done yet. Much more details on the recent experimental developments and what the status is today are given in [11].

5.4 Other Types of Reaction Experiments

The reaction experiments discussed so far have taken place at high beam energy. It is somewhat striking that reactions at lower energy have only recently begun to contribute to halo physics. This is not due to theoretical or interpretational problems – there is a large experience with elastic scattering and transfer experiments from work with stable nuclei – but rather due to a “gap” in the range of beam energies where we can deliver radioactive beams. This gap is being closed now through postacceleration of ISOL beams and physics results are starting to appear.

To some extent, experiments at slightly higher beam energies have already started to look into what might change when one replaces a stable nucleus with a halo nucleus. Much of this work has been discussed already in the lecture by Alamanos and Gillibert [15], so I will mainly refer to them. Results from elastic scattering on protons at high energy were mentioned already in Sect. 2.3.

There are plans at several future radioactive beam facilities (FAIR and RIKEN) for doing electron scattering on the radioactive ions. This is a technically challenging project, but theoreticians have already looked into [100, 101] how one could gain new information on halo nuclei in this way.

6 Beta-decay and Studies on Halos

An introduction to β -decay was already given in two earlier lectures [102, 103]. I will refer to these and to recent reviews of β -decay of exotic nuclei [104] and halo nuclei [7] for more details and here only focus on two aspects, namely β -delayed particle emission and specific halo features in β -decay. The reason for discussing the β -delayed processes is that they will be prominent in the decays of nuclei situated close to the driplines, including halo nuclei.

The only place where weak interactions otherwise play a role in halo physics is in the nuclear muon-capture on ^{11}B [105] to states in ^{11}Be , but effects are small.

6.1 Beta-delayed Particle Emission

As we move towards the driplines particle separation energies become smaller and Q_β -values larger and emission of one or more particles directly from an excited state fed in β -decay becomes possible. Many different processes have been seen already, Fig. 11 shows where the processes occur on the nuclear chart. We shall look at estimates of the Q -values for some of the processes in a moment, but first I should point out some simple relations [104] for them.

N-5	N-4	N-3	N-2	N-1	N	
$\beta 4n$	$\beta 3n$	$\beta 2n$	βn	β		Z+1
		βt	βd	βp		Z
		$\beta \alpha$				Z-1

Fig. 11. The position in the nuclear chart of nuclei that can be reached in β - and β -delayed decays of neutron rich nuclei

For βxn processes (where x is an integer indicating the number of emitted neutrons) the Q -value is the difference between Q_β and S_{xn} for the daughter nucleus so it is obvious how these processes become allowed as we move to more neutron-rich nuclei. For βp , βd and βt the Q -values depend directly on the neutron separation energies of the mother nucleus:

$$Q_X = c - S, \quad (14)$$

where the parameters are given in Table 2. At least for βp and βd this implies that the processes will take place close to the driplines.

To see the systematic behaviour of the Q -values we can use the same simple model as in Sect. 3.1: the simple liquid drop formula [32] neglecting pairing so that the odd-even effect is smeared out. We shall mainly consider neutron-rich nuclei. Figure 12 shows, as an example, the separation energies for the Kr isotope chain. The last β -stable isotope is ^{86}Kr and the last one for

Table 2. Parameters in (14)

X	c (keV)	S	X	c (keV)	S
$\beta^- p$	782	S_n	ECn	-782	S_p
$\beta^- d$	3007	S_{2n}	ECd	1442	S_{2p}
$\beta^- t$	9264	S_{3n}	EC ^3He	6936	S_{3p}

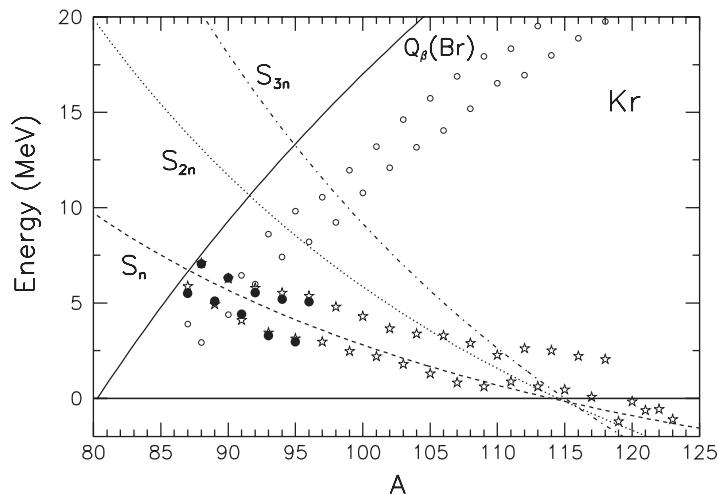


Fig. 12. Neutron separation energies for neutron-rich Kr isotopes. The *filled circles* are experimental single neutron separation energies [30], the *stars* are theoretical values from Hartree-Fock calculations [51]. The *dashed* and *dotted lines* are results from a simple liquid drop mass formula where pairing has been neglected, the *solid line* gives for reference the corresponding Q_β value for Br isotopes. The *small circles* show more realistic Q_β -values taken from [30] up to $A = 96$ and from [51] above this

which the half-life is known is ^{99}Kr [106]. For comparison, the experimental S_n values and the ones from a more realistic model are shown as well, apart from the odd-even effect one also sees the $N = 82$ shell at $A = 118$, but the overall trend is the same. The conclusion that the single-, double- and triple-neutron driplines lie very close is therefore robust. The Q_β of the Br isotopes that would feed Kr is also shown both in the simple model and for the more reliable approach. Even though these two sets of values diverge by several MeV one would still expect to see e.g., $\beta 3n$ processes for at least the last ten nuclei before the dripline.

The simple liquid drop estimate of the regions where β -delayed multi-neutron emission is energetically allowed is shown in Fig. 13, again neglecting pairing. In particular when discussing emission of many neutrons one should keep in mind that these processes could well be sequential – actually, we have at present very little experimental information on this even for $\beta 2n$, but the corresponding $\beta 2p$ processes that are easier to investigate in detail are predominantly sequential [104]. For a sequential process we need appropriate levels to be situated inside the decay window. On top of this the branching ratios to different excitation energies in the daughter nucleus will be biased towards lower-lying levels due to the phase space factor (the matrix elements will tend to favour higher-lying levels, but this is typically a smaller effect). All in all, βn and $\beta 2n$ should be prominent for many neutron rich nuclei, $\beta 3n$

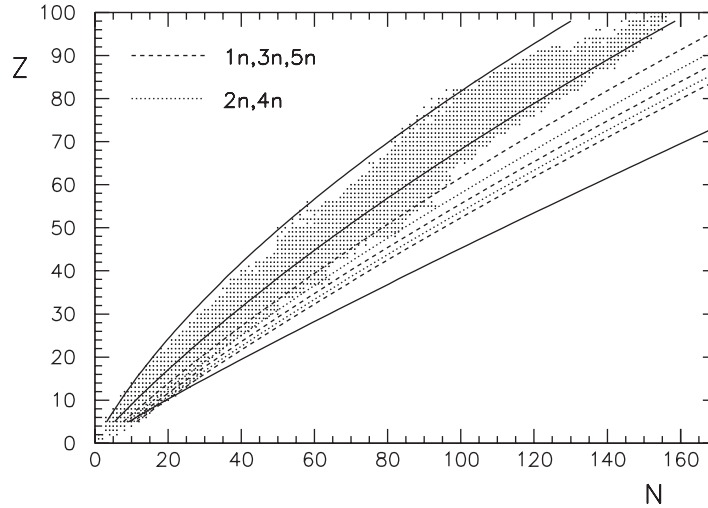


Fig. 13. The nuclear chart with liquid drop estimates for dripline positions as in Fig. 4. The *dashed* and *dotted* lines shown where β -delayed multi-neutron emission becomes energetically allowed according to a simple estimate from the liquid drop model. The points mark nuclei stable towards particle emission for which the mass is experimentally known [30]

will certainly be seen closer to the neutron dripline, $\beta 4n$ and more extreme processes might be more marginal. For completeness, the limits for allowed β -delayed multi-proton emission are almost symmetrically placed with respect to the line of beta-stability, but since the proton dripline is significantly closer than the neutron dripline the multi-proton process will be much less likely.

The corresponding estimates of where the βp , βd and βt processes (and the ECd and EC^3He) are energetically allowed are shown in Fig. 14. As mentioned above, these decay modes become allowed closer to the driplines, but as they are “single step” particle decays they can in principle take place as soon as the Q -value is positive, although the phase space factor of course will give a suppressed branching ratio. Only βd and βt have been observed so far, the most interesting process in the present context being βd that has only been seen in the two-neutron halo nuclei ^6He and ^{11}Li .

6.2 Halo Signals in Beta-decay

Apart from using β -decay to probe the presence of specific orbits, as is also done for non-halo nuclei, there are two halo specific signals [7]. First, since the halo extends beyond normal nuclear radii the spatial overlap with daughter states will be reduced. This is a rather small effect as seen from Fig. 6 in Sect. 4 and would give a general reduction of β -strength. Secondly, since halo particles are removed from the core the two clusters might β -decay

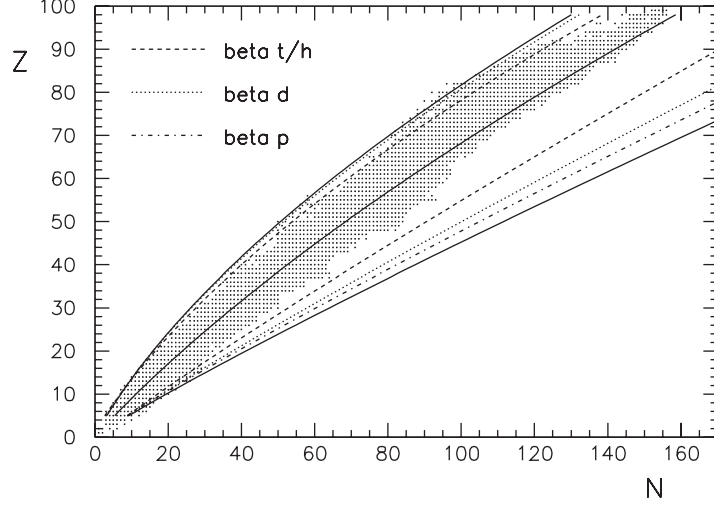


Fig. 14. As Fig. 13, but showing instead where β -delayed triton, deuteron and proton emission becomes energetically allowed according for neutron-rich nuclei. For proton-rich nuclei the lines are for EC-delayed ${}^3\text{He}$ and deuteron emission

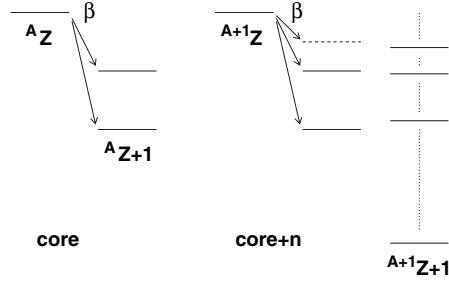


Fig. 15. Schematic model of beta decay of a one-neutron halo. Ideally the decay of the core (*left*) and the halo contribute separately (*solid and dashed lines at right*). The true eigenstates in the daughter nucleus will probably differ somewhat (*extreme right*)

independently from each other. If $\mathcal{O}_\beta$ is the beta-decay operator one formally would have:

$$\mathcal{O}_\beta|\text{halo state}\rangle = \mathcal{O}_\beta(|\text{core}\rangle|\text{halo}\rangle) = (\mathcal{O}_\beta|\text{core}\rangle)|\text{halo}\rangle + |\text{core}\rangle(\mathcal{O}_\beta|\text{halo}\rangle), \quad (15)$$

as illustrated in Fig. 15. Note that there would be a rather small energy release in transitions corresponding to “halo decay”, so these components would be mainly at high excitation energy. One should not expect this splitting into two parts to be strictly true, in most cases larger or smaller mixing with other states in the final nucleus is likely to occur, but one might find transitions

where (15) is a good approximation. A particular interesting one is Gamow-Teller decay of a two-neutron halo, where the last term can describe a core and a deuteron. For the part of the halo where the two neutrons are outside of the nuclear potential it then makes more sense to describe the decay as going directly to a free deuteron state rather than through an excited state in the daughter nucleus. Examples of this will be shown below. It should be noted that such a decay directly to a continuum state differs qualitatively from what has been seen in other β -decays.

One case where both terms in (15) will clearly contribute at the same time is when the core has $N \neq Z$, i.e. non-zero isospin. Components where a halo neutron is changed to a proton and where a core neutron is changed to a proton are both needed to get good isospin in final state. (See Fig. 5 in [6] for an experimental example, the isospin multiplet including ^{11}Be states.) Isospin conservation can of course be broken, but calculations [107, 108] show that isospin mixing will be quite small even for halos. However, the “overlap effect” can still be present and move strength around between states with the same isospin.

6.3 Experimental Data

Nuclei close to the dripline will typically have Q_β -values in the range 15–20 MeV and half-lives that are typically below 10 ms. As mentioned in Sect. 4 the half-life would favour production at in-flight facilities, and many first studies of β -decay properties are also done there. However, once one wants to move beyond measurements of emitted neutrons and γ -rays one needs very thin sources in order that energy losses of emitted charged particles are minimized. This requirement is much easier to fulfill at ISOL-facilities, where beams are well-defined in space and in energy. All measurements of emitted charged particles (including recoil nuclei following neutron emission, a technique that for light nuclei can be competitive to direct detection of neutrons) has therefore taken place at ISOL-facilities.

Experiments on ^6He have taken place at ISOLDE [109, 110], TISOL [111] and Louvain-la-Neuve [112]. The main interest is in the βd branch, observed here for the first time. Even though different theoretical approaches have been used to describe the decay [113, 114, 115] they all agree that the decay takes place directly to the continuum and that cancellation effects make it hard to predict the intensity. At the same time the experiments do not agree on the absolute branching ratio, but the situation will hopefully be clarified soon.

There have also been quite a few experiments on ^{11}Li (references can be found in [7] and through the latest papers [116, 117]). The experimental difficulty is that the decay includes $\beta\gamma$, βn , $\beta n\gamma$, $\beta^2 n$, $\beta^3 n$, βd and βt branches (and of course the corresponding recoil nuclei). In particular the multi-neutron branches are hard to get a handle on, some experimental attempts have been made but nothing is published yet. The deuteron branch

has been seen, but the energy spectrum is not known. Still, we are progressing towards a reliable decay scheme and β -decay has already helped in determining the amount of $(p_{1/2})^2$ in the neutron halo [7, 118] as at most 50%.

The only other nucleus where several experiments have been performed is ^{14}Be . Again many channels are energetically allowed, but the present decay scheme is actually surprisingly simple ([119] and references therein). The main decay branch is to a low-lying 1^+ state, exactly as in the core ^{12}Be and in good accordance with (15), but apart from this no halo features have been seen – in particular there is too little strength seen at high excitation energies where a “halo decay” component might appear. For completeness it should be mentioned that the decays of ^8B , ^{11}Be and ^{17}Ne all are well known. Here the only place where halo effects might enter is a first-forbidden transition from ^{17}Ne , see [7] for details.

7 Summary and Outlook

Halo states in nuclei will mainly occur in ground states or low-lying excited states of some nuclei close to the neutron dripline. Both one-neutron and two-neutron halos are known, if described in terms of neutron orbitals the halo neutron will mainly be in s - and p -waves. Only a limited number of proton halos will exist, they are smaller in size compared to neutron halos.

Apart from the large spatial extent the key feature of halos is their pronounced single-particle behaviour. Many experimental probes take advantage of this.

Most halo experiments so far have employed break-up reactions with beam energies above 20 MeV/u at radioactive beam facilities. In particular the halo-removal channel, where the core nucleus and perhaps some outgoing neutrons are detected, have given much information. The absolute size of cross-sections as well as the width of momentum distributions for the detected fragments are important indicators for halo structure, but more sophisticated analysis methods have been developed for high quality data. For two-neutron halos the structure of the unbound neutron-core subsystem is interesting both experimentally and theoretically. Actually “physics beyond the dripline” has become an important field of research as seen in the recent review [11].

Concerning β -decay, the halo structure will also give rise to changes, as already observed in several nuclei. The major problem here is that β -delayed processes become dominant in near-dripline nuclei and makes experiments harder to perform. A major advantage of this probe is that its results are completely independent from the ones from reaction experiments. Also other types of experiments, e.g., measurements of the mass or of static properties of the halo state, are important, but they have not been treated in detail in this lecture.

The nuclei with the most pronounced halo structure presently known are ^{11}Li and ^{11}Be . The mass-region $A = 6\text{--}19$ contains nearly all known halo

states, but not all halos have been sufficiently well studied and there are still several states whose status have not been clarified yet (this holds even more when excited state halos are considered as well). There is therefore still much work to be done before the next generation of radioactive beam facilities start up and we can have a closer look at dripline nuclei with N around 28 where the next halos, built upon the second p -orbital, are likely to be found.

Acknowledgements

I am grateful to Aksel Jensen and to my experimental colleagues, in particular Björn Jonson and the students in our groups, for many discussions through the years. I would like to thank Eduardo Garrido for providing several figures.

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Isotope Separation On Line and Post Acceleration

P. Van Duppen

Instituut voor Kern- en Stralingsfysica, K. U. Leuven, 3001 Leuven, Belgium

Abstract. In this Lecture the production of radioactive isotopes using the isotope separator on line (ISOL) method is discussed. General properties of the method, the different production and ionization mechanisms are presented and the way these are implemented in the target-ion source systems are highlighted. Mass separation and post acceleration together with an overview of the different facilities are included.

1 Introduction

As emphasized in [1], the radioactive nuclei building up the nuclear chart have, since the discovery of radioactivity, been corner stones of studies in fundamental nuclear physics research and of many applications in various fields of science. The availability of mature techniques for the production and study of these nuclei makes it possible to explore their properties in an unprecedented way and as a consequence, these isotopes will continue to play this crucial role. If one moves away from the valley of stability the production of these so-called exotic nuclei, however, is confronted with difficulties that stem from the

- extremely low production cross sections
- overwhelming production of unwanted species in the same target
- very short half lives of the nuclei of interest.

In general there exist two complementary ways to make good quality beams of exotic nuclei: the in-flight separation technique and the isotope separation on line (ISOL) technique. The latter can eventually be followed by post-acceleration. Both methods transport the nuclei of interest away from their place of production, where a large background from nuclear reactions is present, to a well-shielded experimental set-up, where the nuclear properties can be explored. Apart from creating low-background conditions for the experiment, the transport serves at the same time to purify the beam and to prepare it in the necessary conditions with respect to energy, time and ion optical properties for the experiments. The in-flight method makes use of the reaction kinematics, some combination of magnetic and electrical fields, and atomic processes to identify the nuclei and to separate the isotopes of interest from the primary beam or from other isotopes produced in

the reaction. The basic principles of in-flight separation are discussed by D.J. Morrissey and B.M. Sherrill [2]. The ISOL techniques rely on the availability of the radioactive species produced in a target and thermalized in a catcher consisting of solid, liquid or gas material. Often the target and catcher are one and the same system. The isotopes are subsequently extracted from the catcher material and ionized in an ion source. After extraction from the ion source the species are mass analyzed using a dipole magnet and subsequently accelerated to the required energy. A general comparison of these two techniques including some historical note is given in the Lecture of M. Huyse [1], especially Fig. 3 from [1] highlights the interesting differences. Progress in the field of in-flight separation and ISOL techniques is regularly reported in the conference series on “Electromagnetic Isotope Separators and Techniques related to their Applications (EMIS)” [3] and a detailed overview of on line mass separators including many technical aspects were reported by H.L. Ravn and B.W. Allardyce in [4].

The ultimate aim for ISOL systems is the production of beams of exotic nuclei that are abundant, pure, of good ion optical quality and variable in energy essentially from rest (meV/u) to intermediate energy (a few 100 MeV/u). Therefore the whole production sequence must possess the following properties:

- **High Production Rate**

The production cross-section of a particular reaction is energy dependent. But it is a nature given number, to which one has to optimize the beam-target combination. Accelerators have to be used that can deliver the highest beam intensities and target systems have to be developed that can cope with the power deposition of the primary beam and of the secondary reaction products.

- **Efficient**

The production rate of the very exotic nuclei will always be marginal. Therefore any manipulation of the reaction products – e.g., ionization, purification, acceleration, transport to the detection system – has to be very efficient, otherwise one loses the few precious nuclei.

- **Fast**

As one is dealing with short-lived exotic nuclei, the losses due to radioactive decay between the moment of production and the arrival at the experimental set-up should be kept to a minimum.

- **Selective**

In the nuclear reaction process the unwanted – in general more stable – nuclei are produced much more abundantly. Furthermore ISOL systems often produce beams of isotopes from the target material itself or from other components of the target-ion source system. Thus the separation process should distinguish between the wanted and unwanted species in an effective way. One defines the unwanted species as **contaminants** which,

if they have the same mass-over-charge ratio as the ions of interest, are called isobaric contaminants.

This lecture gives a basic description of the ingredients of the ISOL technique and discusses its properties. Chapter 2 starts with a schematic presentation of the technique and defines relevant parameters like production rates, efficiencies and delay times. It also introduces the ionization mechanisms used at ISOL systems. More details on the different reaction mechanisms used are given in Chap. 3 which is followed by a more technical chapter discussing the targets, catchers and ion sources used. The mass separation and post-acceleration principles are briefly discussed in Chap. 5 while an overview of the different facilities currently operational is presented in Chap. 6.

2 Isotope Separation On Line: Schematic Presentation

The basic ingredients of an ISOL system, shown in Fig. 1, represent different steps: production, thermalization, ionization, extraction, mass separation, cooling, charge-state breeding and acceleration. These processes are governed by physical and chemical processes, thus both physical (e.g., production cross-section, decay-half life, ionization potential) and chemical (e.g., molecular formation probability, volatility) properties of the nuclei of interest and of the target material are important. It will be discussed later how these properties are used to obtain the necessary efficiency and selectivity, and to reduce the delay time. The latter is the average time the radioactive atoms spent from the moment of production to the moment of arrival at the experimental position. The fact that all items in Fig. 1 are shown separately does not imply that every item represent one mechanical unit. Some of the ingredients are embedded in a single unit but for the clarity of the discussion this approach is chosen. To be noted as well is that in many cases only a selected number of steps are to be taken.

- **Production Target**

A beam of light (e.g., protons) or heavy (e.g., ^{58}Ni) ions, neutrons or electrons hits a target to produce the radioactive ion beams of interest through different reaction mechanisms (see Sect. 3). In order to distinguish between the different beams one denotes the former beam as **primary beam** and the latter as **secondary beam**. In some cases a converter is used to transfer e.g., a primary beam of protons or electrons into a flux of neutrons or γ rays, respectively. This flux is subsequently sent onto the target and induces the nuclear reaction. In choosing the target it is crucial to optimize the beam-target combination with respect to the highest production cross-section and the lowest amount of contaminants, and to use target material that can stand the highest possible beam currents.

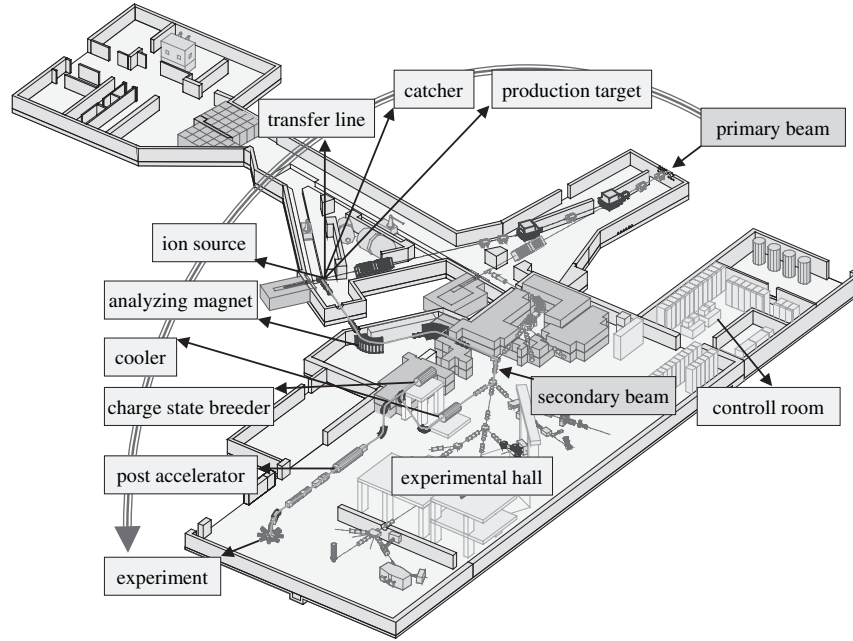


Fig. 1. The different ingredients for the production of radioactive ion beams using the ISOL method are shown on a background of the layout of the ISOLDE-CERN hall. ISOLDE has two target stations that can operate in beam sharing mode. The different stages along the path of a radioactive ion from the target-ion source system to the final experimental set-up after post acceleration are indicated. Depending on the requirements, the radioactive ion beams can be sent to an experimental set-up after the analyzing magnet (e.g., decay studies [5]), the cooler (e.g., mass measurements [6] or laser spectroscopy [7]), the charge-state breeder (e.g., deeper implantation of radioactive ions into material) and the post-accelerator (e.g., for Coulomb excitation measurements [8] or reaction studies [9, 10, 11]); cf. Plate 3 in the Colour Supplement

- **Catcher**

Once the exotic nuclei are produced in a nuclear reaction, they are stopped in a solid or gaseous catcher. Often the word “thermalized” instead of “stopped” is used to indicate that the radioactive atoms are cooled down from e.g., about 200 MeV kinetic energy in case of a fission reaction to ~ 300 meV average kinetic energy corresponding to a catcher temperature of 2000°C . In some cases the target and the catcher are the same. After thermalization the radioactive nuclei of interest (partly) escape from the catcher and are transported towards the ion source. In case of a solid catcher, diffusion from the target/catcher material is followed by effusion towards the ion source. The former process is governed by the diffusion properties of the atoms of interest and of the target/catcher matrix. The effusion is controlled by a combination of factors: pumping through the

pores of the target/catcher material and through the connecting tube, sticking times when the atom hits a wall and diffusion properties of the atoms with respect to the wall material. In case the sticking time is long compared to the diffusion time of the atom inside the wall material, it will diffuse and eventually disappear in the walls. To reduce the loss of nuclei the speed of these processes is increased by heating the target/catcher system, the transfer line and the ion source to a suitable – in most cases high – temperature. In this way the diffusion process to the surface of the solid matrix is accelerated and the sticking time at the walls of the target or catcher material and container is reduced. In case of a gaseous catcher the radioactive atoms of interest stay in the gas phase and there is no need for high temperature systems.

- **Transfer Line**

Through the transfer line the nuclei are transported from the target/catcher system to the ion source. This line is kept generally at a high temperature to avoid lingering of the atoms to the walls of the line.

- **Ion Source**

Depending on the requirements, several different ionization mechanisms are used. In general singly-charged positive or negative ions are produced. In a few cases multiply charged ions are produced as well. All above mentioned parts are commonly named as the target-ion source system. It is put on a positive (in case of beams of positively charged ions) or negative (negatively charged ions) high voltage of typical values between 40 and 60 kV.

- **Analyzing Magnet**

After extraction from the ion source the low-energy ion beam is mass separated by an analyzing magnet and transported to the focal plane. The focal plane is where in ideal conditions the ion beams are focussed after the analyzing magnet. An important property that expresses the quality of the system is the mass resolving power, defined as $R = M/\Delta M$. Here ΔM is the full width at half maximum of a beam of ions with mass M in the focal plane of the separator. The resolving power depends on the properties of the magnet and of the ion optical properties of the ion beam. Typical values for ISOL systems vary from a few hundred to a few thousand. Other properties of the mass analyzing system are discussed in [4].

- **Cooler**

A “cooler” device is used to literally cool down the ion beam in order to improve its ion optical properties and in some cases to bunch the ion beam. Cooling should be understood in terms of reducing its axial and radial temperature or more precisely its axial and radial momentum or energy spread. Bunching of the ion beam is often required to increase the peak to background ratio of certain experiments like laser spectroscopy experiments or to inject the beam into a charge-state breeder. There are

two distinctive devices used for cooling and bunching : Penning traps [12] and radio frequency (RF) coolers [13]. Both systems are based on the storage of the ions using electrical DC and RF fields (RF-coolers) or a combination of magnetic and electrical fields (Penning traps). Buffer gas like helium or argon is introduced and through collisions between the buffer gas atoms and the radioactive ions energy is lost and cooling occurs. This cooling technique has been developed over the last years yielding excellent performance. Efficiencies, defined as the beam intensity of the cooled beam versus that of the injected beam, over 50% have been reached. Coolers are installed at different ISOL systems [14, 15, 16, 17] for efficient injection into traps as well as for solid-state studies. A more general description of the principles of radioactive ion beam manipulation can be found in the Lecture of G. Bollen to the first volume of these lecture notes [6].

- **Charge-State Breeder**

As the simplicity, efficiency and cost of the post-accelerator are directly related to the charge-state of the ions, it is usefull to produce a multiple charge state ion beam before injection into an accelerator. A charge-state breeder performs a transformation from a singly charged ion beam into a multiple charged one. Two types of charge-state breeding ion sources are used: the Electron Beam Ion Source (EBIS) and the Electron Cyclotron Resonance (ECR) ion source. Both rely on the principle of intense bombardment of the ions with energetic electrons, with electron impact ionization yielding ions in higher charge states. The plasma of ions and electrons is confined through electrical and strong magnetic fields.

- **Post Accelerator**

The highly charged ion beam from the charge-state breeder or the beam of singly charged ions is subsequently injected into the accelerator. This process is referred to as post-acceleration. Finally the energetic beam of radioactive ions is sent to experimental set-ups where a multitude of different properties can be studied.

2.1 Production Rates and Intensities

The intensity ($I_{reaction\ products}$) of the reaction products synthesized in the target is calculated using the following equation

$$I_{reaction\ products} = \sigma N_{target} \Phi \quad (1)$$

with σ being the reaction cross-section (cm^2), N_{target} the number of target atoms per surface area (cm^{-2}) and Φ the primary beam intensity. The cross section is often expressed in units of barns, mbarns or μbarns , with $1\text{ barn} = 10^{-24}\text{ cm}^2$. Since the cross section is energy dependent (e.g., in heavy-ion fusion evaporation reactions) and the primary beam loses energy while passing through the target, the intensity has to be calculated by integrating over the target thickness taking into account the energy loss of

the beam. In order to make a general but quantitative comparison between different beam-target combinations the luminosity (L), defined as

$$L = N_{target} \Phi \quad (2)$$

is often used and expressed in units of $\text{s}^{-1} \text{cm}^{-2}$. Because of all the steps described in the previous paragraph the final intensity of the secondary beam (I) will be reduced due to the different loss processes involved. One can express that final intensity of the secondary beam as

$$I = \sigma N_{target} \Phi \epsilon \quad (3)$$

with ϵ being the efficiency of the whole process (some authors used η as symbol for efficiency). This efficiency is defined as the ratio of the final secondary beam intensity that arrives at the experimental set-up (I) versus the intensity of the reaction products ($I_{reaction\ products}$). It is a product of a series of partial efficiencies, as given in Table 1,

$$\epsilon = \epsilon_{delay} \epsilon_{ion} \epsilon_{trans} \epsilon_{cool-bunch} \epsilon_{breeding} \epsilon_{accelerator} \quad (4)$$

Although ϵ_{delay} is defined for the extraction of the ions from the target-ion source system, the cooling, bunching and charge-state breeding induces a delay as well. This loss factor will be discussed separately below when discussing the relevant process.

Table 1. Partial efficiencies

ϵ_{delay}	probability of survival against radioactive decay during the time needed to extract the ion from the target-ion source system
ϵ_{ion}	ionization efficiency
ϵ_{trans}	efficiency of mass analysis and transport to the experimental set-up
$\epsilon_{cool-bunch}$	cooling and bunching efficiency
$\epsilon_{breeding}$	charge-state breeding efficiency
$\epsilon_{accelerator}$	efficiency of the post-accelerator
ϵ	total efficiency: the product of the above mentioned terms

2.2 Delay Times

The parameter ϵ_{delay} expresses the relative amount of ions that survive, in spite of their radioactive decay, the time elapsed from the moment of their production to the moment of their extraction from the ion source. This time is often referred as **release time** as in case of solid and liquid catcher systems it is determined by the diffusion from the target/catcher material, the

desorption from the material surface and the effusion to the ion source exit hole (or any other hole in the target ion source system). In case of a gaseous catcher the term “release time” is replaced by “**delay time**”. It is determined by the gas flow and, in case electrical fields are used to obtain a faster evacuation of the gas cell, the mobility of the ions. The radioactive isotopes leave the ion source as neutral, singly- or multiply charged atomic or molecular ions. Once an ion beam is formed the mass separation and transport goes fast compared to the delay time and half life of the nuclei and no significant losses are expected. If cooling and bunching is involved (see Fig. 1) an extra loss factor due to radioactive decay might occur. It is thus clear that ϵ_{delay} will strongly depend on the half life of the exotic nucleus of interest.

To quantify these loss factors one defines the delay-time distribution ($P(t)$) whereby $P(t)dt$ is the probability for an atom of a given element created at $t = 0$ to be released from the ion source between time t and $t + dt$. It is important to note that this function depends e.g., in case of high temperature target ion source systems on the chemical and physical properties of the element but **not** on the half life of the isotope. One can subsequently calculate ϵ_{delay} by folding $P(t)dt$ with $\exp(-\lambda t)$

$$\epsilon_{delay}(T_{1/2}) = \int_0^\infty P(t) \exp(-\lambda t) dt \quad (5)$$

with $\lambda = 1/\tau = \ln 2/T_{1/2}$. Note that $P(t)$ is normalized thus that $\epsilon_{delay}(T_{1/2} = \infty) = 1$, which means that all produced atoms will in the end have escaped from the target-ion source system. One obtains the average delay time as

$$\tau_{delay} = \int_0^\infty t P(t) dt \quad (6)$$

Finally one defines the fractional release function $F(t)$ as the fraction of the atoms produced at $t = 0$ that has escaped from the target-ion source system after a time t

$$F(t) = \int_0^t P(t') \exp(-\lambda t') dt' \quad (7)$$

Often in release studies of target-ion source systems the fractional release function is determined experimentally, which allows one to deduce diffusion and effusion coefficients.

It is very illustrative to consider two peculiar release processes. For the first one a delay-time distribution equal to a properly normalized single exponential function is taken, $P(t) = \lambda_{delay} \exp(-\lambda_{delay} t)$, with a release rate constant equal to $\lambda_{delay} = 1/\tau_{delay} = \ln 2/T_{delay}$. In this case it is easy to show that:

$$\epsilon_{delay}(T_{1/2}) = T_{1/2}/(T_{1/2} + T_{delay}) \quad (8)$$

In this case half of the radioactive atoms will escape from the ion source if the delay time (τ_{delay}) equals the mean life time ($\tau = 1/\lambda$) but still a fraction of 10% escapes even if the delay time is nine times longer than the half life. The latter is due to the exponential nature of the release process in this case. A second example is the case where the atoms are released all at once after a fixed time t_0 . In this case $P(t)$ is a delta function and

$$\epsilon_{delay}(T_{1/2}) = \int_0^\infty \delta(t - t_0) \exp(-\lambda t) dt = \exp(-\lambda t_0) \quad (9)$$

and

$$\tau_{delay} = \int_0^\infty \delta(t - t_0) t dt = t_0 \quad (10)$$

In Sect. 4 some examples of ϵ_{delay} measurements will be discussed.

2.3 Ionization Mechanisms

The **ionization efficiency** (ϵ_{ion}) is defined for a specific isotopes as the ratio of the number of ions extracted from the ion source to the number of atoms injected into the ion source. Losses due to leaks in the transfer line between target and ion source and in the ion source itself are included in this number. In order to distinguish them from losses due to the radioactive decay, the ionization efficiency is always defined for stable isotopes of a specific element. The loss factor due to radioactive decay is taking into account by multiplying ϵ_{ion} with ϵ_{delay} .

Different ionization mechanisms are implemented in the ion sources of ISOL systems. The use of them depends essentially on the ionization potential (W_i) of the element of interest, the required charge state and selectivity. Below a summary of the three main ionization mechanisms is given and ion survival in gas catchers is discussed. The way these mechanisms are implemented in the target-ion source systems is discussed in Chap. 4. Reviews on the target-ion source systems used at on line isotope separators can be found in [4, 18, 19, 20, 21].

Electron Impact Ionization

For isotopes of elements with $W_i > 7\text{ eV}$ and for the creation of multiply charged ions, electron impact ionization is mostly used. The atoms or ions are bombarded by energetic electrons, thereby loosing one or more of their outer electrons. Figure 2 presents the cross section for electron impact ionization of argon atoms as a function of electron energy. The bombardment of atoms by mono-energetic electrons and the related evolution of the atom and ion density as a function of time can be described by a differential equation

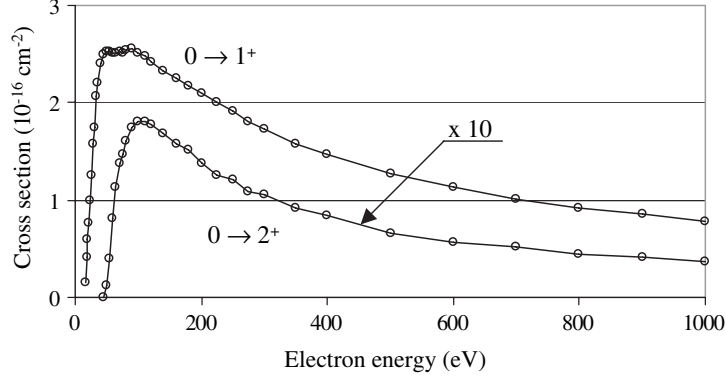


Fig. 2. Cross section for electron impact ionization of argon atoms as a function of electron energy. Data for neutral to singly charged ($0 \rightarrow 1^+$) and to doubly charged ($0 \rightarrow 2^+$) ions are shown [22]. Note the different threshold values for the $0 \rightarrow 1^+$ and $1^+ \rightarrow 2^+$ ionization. The ionization potential of 1^+ ions is 15.75 eV while it is 27.62 eV for 2^+ ions

$$dn_i/dt = (n_{i-1} \sigma_{i-1 \rightarrow i} - n_i \sigma_{i \rightarrow i+1}) j_e \quad (11)$$

with n_i being the ion density with charge state $Q=i$, j_e the electron current density and $\sigma_{i-1 \rightarrow i}$ the cross section for impact ionization from charge state $i-1$ to i . This equation is valid for i from 1 to $i_{max} - 1$. For the neutral atom density (n_0) the first term of the right part of (11) vanishes, while for the fully stripped ion ($n_{i_{max}}$) the last term vanishes. Note that (11) does not take into account charge exchange, multiple charge ionization or recombination effects. Figure 3 shows the results of the atom and ion density as a function of the product of $n_e v t$, with t being the time the neutral atoms have been bombarded with mono-energetic electrons of density n_e and velocity v . From this figure one can determine the conditions ($n_e v t$) that are needed to achieve a certain efficiency for a specific charge state. For example to obtain $\epsilon_{ion}(5^+) = 30\%$ the neutral atoms have to be bombarded for 6.7 ms with a mono-energetic electron beam of 17 keV and 120 A/cm². Electron impact ionization is used in high-temperature gaseous discharge ion sources, ECR ion sources and EBIS. These sources will be discussed in some detail in Chap. 4. In the former two types of ion sources the electron energy is distributed over a wide energy range and therefore (11) should be adapted. Because of the very unselective nature of this ionization process, ion sources based on this principle do in general not offer much chemical selectivity.

Surface Ionization

A second ionization mechanism used in ISOL facilities is based on surface ionization. When an atom interacts with a heated surface it can lose or gain

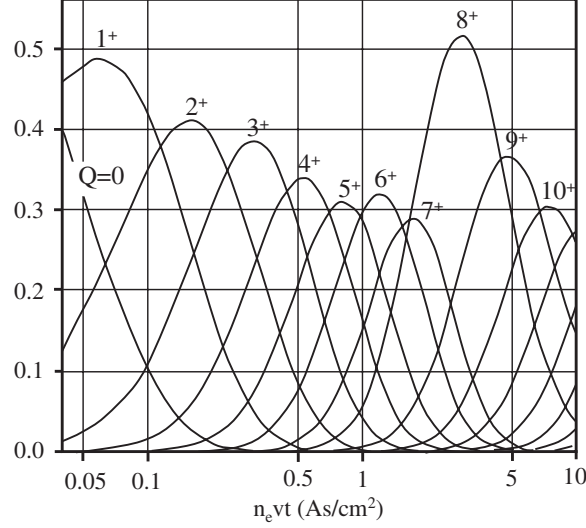


Fig. 3. Calculated relative atom ($Q = 0$) and ion density of argon under bombardment of a mono-energetic 17 keV, 120 A/cm² electron beam as a function $n_e v t$ in As/cm². The electron density and velocity is denoted by n_e and v , respectively. At $t = 0$ the relative atom density was 1 while the ion density was 0. The calculation assumes only one-step electron impact ionization. Multi-step electron impact ionization is not taken into account. The cross sections were calculated using [23]. The neon-like properties of Ar^{+8} reduces the cross section for electron impact ionization from the 8^+ to 9^+ compared to 7^+ to 8^+ ionization and create a peak for the 8^+ charge state. This peak is reached after about 17 ms of electron bombardment

an electron before leaving the surface as a positive or negative singly charged ion. This technique can be used efficiently for elements with $W_i < 7$ eV for the creation of positive ions (positive surface ionization) and with electron affinity $E_A > 1.5$ eV for the creation of negative ions (negative surface ionization). The ratio between the ion density (n_i) and the neutral density (n_0) of a certain element, with W_i or E_A , at a heated surface, with temperature (T) and work function (φ), is given by the Langmuir equation:

$$n_i^+ / n_0 = (g_i^+ / g_0) \exp[(\varphi - W_i) / kT] \quad (12)$$

$$n_i^- / n_0 = (g_i^- / g_0) \exp[(E_A - \varphi) / kT] \quad (13)$$

with g_i and g_0 being the statistical weight of the ionic and atomic ground state. The positive and negative sign indicate positive and negative ions, respectively. The efficiency is then described as

$$\epsilon_{ion} = n_i / (n_i + n_0) \quad (14)$$

depending strongly on the temperature and the difference between the work function of the surface material and the ionization potential or the electron

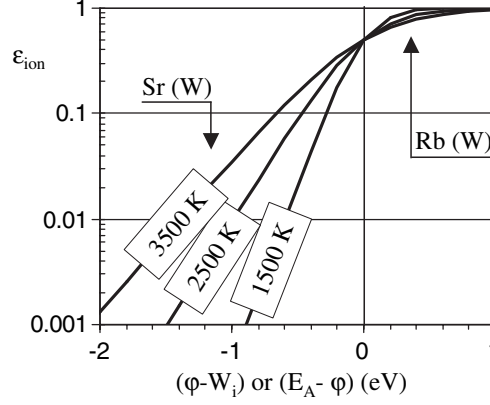


Fig. 4. Theoretical surface ionization efficiency according to (14) is shown for different temperatures as a function of the difference of the ionization potential (W_i) or the electron affinity (E_A) and the work function (φ) of a fictitious element (based on [18]). The $(\varphi - W_i)$ value for strontium and rubidium with a tungsten ionizer is indicated. The curves are calculated for $g_i^\pm/g_0 = 1$. When the ionization potential or the electron affinity equals the work function, the surface ionization efficiency is 50%

affinity of the element of interest. Figure 4 shows the theoretical surface ionization efficiency as a function of the difference of the ionization potential or the electron affinity and the work function. Materials with high work functions and that can be heated to high temperatures like tantalum ($\varphi = 4.19$ eV), tungsten (4.53 eV) and rhenium (5.1 eV) are used to construct positive surface ionization sources. While compounds with low work function like LaB_6 ($\varphi = 2.6$ eV) or BaO (1.7 eV) are used for negative surface ionization.

Using this ionization technique extreme selectivity can be obtained if the elements of which isotopes are produced in the same nuclear reaction have very different ionization potentials. For example the neighboring elements krypton, rubidium and strontium are often produced in the same nuclear reaction. The ionization potentials of krypton (14.0 eV), rubidium (4.18 eV) and strontium (5.70 eV) are very different. Using tungsten as surface ionizing material at a temperature of 3000 K results in $\epsilon_{ion} = 0.77$ for rubidium while ϵ_{ion} is about 80 times smaller for strontium and almost sixteen orders of magnitude smaller for krypton ($g_i^+/g_0 = 1$ was assumed for this calculation). In the former case a very pure beam of rubidium ions would result. In general the alkali elements (group IA of Mendeljev's Table: lithium, sodium, potassium,...) are very efficiently ionized using positive surface ionization while the halogen elements (group VIIA of Mendeljev's Table: fluorine, chlorine, bromine,...) are good candidates for negative ionization.

When the surface ionization mechanism is applied in a hot cavity a quasi-neutral plasma in thermal equilibrium can be created under certain conditions of positive surface ionization and electron emission from the walls. This means that for the reaction of the type $[atom \leftrightarrow ion + e^-]$ the equilibrium constant can be calculated using the Eggert-Saha equation that was originally developed to interpret the spectra from stellar atmospheres:

$$n_i^+ n_e / n_0 = (2 g_i^+ / g_0) (2\pi m_e kT h^2)^{3/2} \exp[-W_i / kT] \quad (15)$$

with n_e being the electron density, m_e the electron mass and h Planck's constant. From this equilibrium constant, ϵ_{ion} can be calculated resulting in a strong dependence on the ionization potential, the temperature and the pressure of the plasma ($= kT (n_i^+ + n_e + n_0)$). In case of low pressures much higher efficiencies are obtained compared to pure surface ionization. Taking the above example of rubidium and strontium production and assuming a realistic number for the plasma pressure of 9.10^{-5} mbar in a tungsten cavity the efficiency for rubidium is close to 100 % while for strontium a value of about 70 % is reached. The ϵ_{ion} for strontium increases by several orders of magnitude compared to pure surface ionization, but the ratio of the ϵ_{ion} values rubidium and strontium decreases by roughly the same factor. More details on the surface ionization and this so-called thermoionization sources can be found in a paper by R. Kirchner [18].

Laser Ionization

A method that has recently been successfully implemented at ISOL systems is resonant laser ionization [24]. During this process the atoms are stepwise excited by laser photons, leading finally to the continuum, to auto-ionizing states or to highly excited states close to the continuum. In the latter case the ionization step is achieved through infra-red irradiation, an electrical field or atomic collisions. The ionization process consists of typically two or three steps and, because of the resonant nature of most of these steps, resonant laser ionization is very efficient and chemically selective, resulting in isobarically and, if the laser bandwidth is narrow enough, isomerically pure beams. Details on the principles and applications of resonant laser ionization can be found in [25, 26].

Figure 5 shows an idealized two-step ionization scheme that illustrates the laser conditions needed to reach ϵ_{ion} close to 100%. As the cross section for the first step excitation is large the atomic ground and excited state are in equilibrium during laser irradiation. Under this condition the evolution of the number of atoms ($N(t)$) can be obtained from

$$-dN(t)/dt = (\sigma_I F + \beta) N(t) \quad (16)$$

with σ_I (cm^2) being the cross section for ionization from the excited state to the continuum, F ($\text{cm}^{-2}\text{s}^{-1}$) the flux of photons, and β the depopulation rate

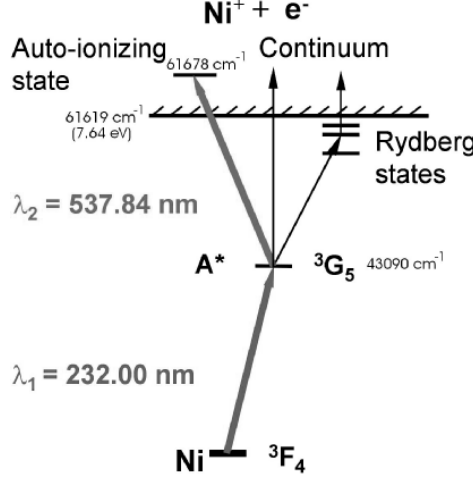


Fig. 5. Schematic atomic level scheme showing the principles of resonant photoionization of nickel. The atoms are stepwise excited to the continuum, auto-ionizing states or highly excited, so-called Rydberg states from which ionization is achieved by infrared irradiation, electrical fields or gas collisions; cf. Plate 4 in the Colour Supplement

of the excited state due to natural decay or other loss mechanisms. Moreover, it is assumed that the laser pulse fires at $t = 0$. From (16) conditions on the flux (the number of photons per second and unit area) and on the fluence (the number of photons during the laser pulse per unit area) of the laser photons are obtained. The flux-condition requires that the depopulation rate of the excited state by ionization into the continuum is much larger than that of the natural decay towards, e.g., a metastable state. Thus

$$\sigma_I F \gg \beta \quad (17)$$

The fluence-condition requires that in total (i.e. in one laser pulse) enough photons irradiate the atom to induce the ionization step:

$$\sigma_I \phi g_2 / (g_1 + g_2) \gg 1 \quad (18)$$

with ϕ (cm^{-2}) being the fluence, and g_i the statistical weight factor of the ground state ($i = 1$) and excited state ($i = 2$), respectively. When both (17) and (18) are fulfilled, complete saturation is reached, i.e. essentially all the atoms irradiated by the laser light are also ionized. What do these conditions mean with respect to the laser power needed? The ionization cross-section (σ_I) is of the order of 10^{-17} cm^2 when the second step is a non-resonant one that leads directly into the continuum, and β has a typical value of 10^6 s^{-1} . This means that for a continuous wave (CW) laser of 1 mm^2 cross section the flux needed to fulfil (17) must be much larger than 10^{21} photons/s which for

a typical photon energy of 3 eV corresponds to about 500 W. Pulsed lasers (with typical pulse width of 10 ns) need to deliver 5 mJ/pulse to fulfil the flux condition. For the fluence condition (18), ϕ should be much larger than 0.5 mJ/pulse. Thus pulsed lasers, whose photon wavelength is tuneable, are commonly used for efficient laser ionization. CW laser systems with high enough power are continuously improved and might thus be used for certain elements. In case the ionization step is also a resonant excitation (e.g., to auto-ionizing states or high lying Rydberg states) the conditions are more relaxed as σ_I is typically a factor of 100 larger. The above mentioned conditions can be reached with commercially available pulsed lasers like Nd-Yag lasers (typical repetition rate 10 Hz), Excimer laser (100 Hz), copper vapour laser (10 kHz) and solid-state lasers (10 kHz). A review of recent developments using the latter for radioactive beam production can be found in [27]. However, as the radioactive nuclei are essentially produced in a continuous mode, the low-duty cycle of the pulsed lasers limits ϵ_{ion} to about 10^{-4} . Therefore, apart from the laser development as such, much effort has been devoted to construct an effective storage device for radioactive atoms in between laser pulses. These devices, however, reduce the inherent selectivity of laser ionization when implemented in an ion source. Laser ionization can result in a selectivity of 10^4 per resonant step whereas the selectivity of currently used on line laser ion sources varies between 10 to 1000. Typical applications of these devices will be discussed in Chap. 4.

Finally, it should be mentioned that recently non-resonant high-power laser ionization has been used to create highly charged ion beams and even energetic ion beams. This technique has only been used for surface and plasma analysis as well as for the production of stable isotope beams. However, applications towards the production of radioactive ion beams might be envisaged.

Ion Survival in a Noble Gas Environment

An important new class of target ion source systems is based on the gas catcher. The reaction residues are thermalized in a gas cell filled with noble gas (helium or argon). Together with the gas, part of the atoms and ions are evacuated through the exit hole of the gas cell and injected into the first acceleration stage of the isotope separator or into a radio-frequency ion guide [6]. Ionization is in first instance not really created by external means as in the above mentioned mechanisms (like using electron bombardment or surface ionization). Still this mechanism has important applications in ISOL systems. Because of the large recoil velocities, the reaction products are in a high charge state when entering the gas cell. When the reaction products slow down, their charge state is lowered by charge-exchange reactions with the buffer gas atoms. Eventually, the reaction products end up in a 1^+ or 2^+ charge state because of the high ionization potential of the noble gases. The surviving ions are then extracted from the gas cell and mass analyzed. The way the 1^+ ions are formed and survive the transport through the gas cell

depends strongly on the experimental conditions, namely evacuation time, purity of the buffer gas, electron and ion density created in the cell and, in some cases, the chemical properties of the isotopes. In the so-called ion guide isotope separator on line (IGISOL) systems fast evacuation is the key issue as it prevents the ions from recombining and neutralizing [28, 29]. In other gas catcher systems electrical RF and DC fields are used to decrease the evacuation time of the gas cell [30, 31, 32] thereby reducing the diffusion losses. The latter systems are discussed in the framework of large gas catcher systems for the conversion of high-energy radioactive ion beams from fragmentation or heavy-ion fusion evaporation reactions into low-energy ones. Recently the efficiency and selectivity of gas cell has been improved using resonant laser ionization [24].

3 Production of Exotic Nuclei

A wide variety of reactions and beam-target combinations are used to produce the radioactive isotopes of interest. The physics behind the different reaction mechanisms will not be discussed here but a comprehensive list of the reactions used at ISOL facilities will be given together with some examples of typical intensities available at today's accelerators or reactors. The key parameter is the reaction cross-section, see (1). While often this quantity is unknown experimentally, it can be reliably calculated for many different reactions. Calculation codes can be found for heavy-ion fusion evaporation in [34, 35] and for fission, spallation and fragmentation reactions induced by beams of relativistic energies, in [36]. Different projectiles with a large variety of energies are used to induce the nuclear reaction in the target, namely low-energy protons and neutrons of 30 to 100 MeV, high-energy protons of 500 to 1500 MeV, heavy ions of 4 to 100 MeV/u, thermal neutrons and electron beams of 50 MeV.

The choice of the nuclear reaction, used to produce an ion beam of a particular isotope, depends in the first place on its position on the chart of nuclei, but also on its physical and chemical properties. For example neutron-deficient rhodium isotopes (e.g. ^{91}Rh , $T_{1/2} = 1.47$ s) can be produced in copious amount by high-energy proton spallation on a cadmium target. However because of the refractory nature of rhodium and the low-melting point of cadmium, the target can not be heated to high enough temperature to induce fast diffusion and thus substantial losses are encountered. On the other hand a heavy-ion induced fusion evaporation reaction on a thin target (e.g. $^{58}\text{Ni}(1\text{ mg/cm}^2)(^{36}\text{Ar}, 1p2n)^{91}\text{Rh}$) leads to only 300 atoms per second using a 1 particle- μA ^{36}Ar beam. However in this case the rhodium isotopes recoil out of the target and can be stopped in a noble gas and laser ionized. This yields a ^{91}Rh beam of ~ 2 atoms/s which is orders of magnitude more than what can be reached by using a high-energy proton induced reaction and a thick high-temperature target.

Light- and Heavy-ion Fusion Evaporation Reactions or Direct Reactions

With light-ion induced fusion reactions nuclei close to the line of stability on the neutron-deficient side are produced. The advantages are the high cross section of the reaction and the high intensity of the primary beam available. An example is the $^{13}\text{C}(p, n)^{13}\text{N}$ reaction that has a peak cross section of 224 mbarn at an energy of 6.6 MeV. At the Louvain-la-Neuve cyclotron center [37, 39, 38], a proton beam with intensity up to 500 μA and an energy of 30 MeV was directed onto an enriched ^{13}C target. The cross section drops from its maximum to about 10 mbarn at 30 MeV. With a typical target density of 1.8 g/cm² the proton beam is stopped after 0.55 cm. This gives a conversion rate of primary to secondary ions, defined as the number of ^{13}N produced per incoming proton, of $1.6 \cdot 10^{-3}$. This means that for every 625 incoming protons one ^{13}N atom is formed. For a primary proton beam of 500 μA , $5 \times 10^{12} {}^{13}\text{N}$ atoms/s are obtained.

Heavy-ion fusion evaporation reactions have typically a much lower cross section but produce neutron-deficient nuclei very far from the line of stability. The primary beam energy is of the order of 5 MeV/u while its intensity is limited to about one particle- μA due to target heating. The energy loss in the target material is much higher for heavy ions compared to protons resulting in a lower number of target atoms, see (1). For example the $^{52}\text{Cr}(225 \text{ MeV/u})(^{142}\text{Nd}, 4n)^{190}\text{Po}$ reaction was used to produce the very neutron-deficient isotope ^{190}Po with a cross section of about 100 nbarn. A target of 1 mg/cm² thickness results then in a conversion rate of only 4.2×10^{-13} , which means that 400 particle-nA of primary ^{54}Cr beam results in about 1 ^{190}Po atom/s.

It should be noted that these reactions produce on average a limited number of different isotopes (a typical value for the heavy-ion reaction mentioned above is 10) while the reaction mechanisms discussed below produce hundreds of different isotopes. This puts much more constraints on the selectivity of the target-ion source systems used with the reactions discussed below.

Fission

To produce neutron-rich nuclei in a wide mass range fission of $^{235,238}\text{U}$, ^{232}Th and long lived actinides is used. The different particles to induce fission are low- and high-energy protons, heavy ions, fast or thermal neutrons and electrons, with photo-fission being used in the latter case. Thermal neutrons are produced in a reactor while fast neutrons are obtained using the “so-called” convertor method [41]. In the original concept a flux of fast neutrons was generated by stripping an intense deuteron beam at 200 MeV. However fast neutrons can also be produced by, e.g., high-energy proton beams. The high power of the primary beam is dumped in a well-cooled neutron production target while the neutrons whose power deposition is much smaller than that

of protons or deuterons create the radioactive isotopes. In this way problems related to high power deposition can be decoupled from the target used to produce the radioactive isotopes. Figure 6 shows the cross section for the production of rubidium isotopes from fission of ^{238}U using different projectiles. For comparison the cross section from thermal neutron induced fission of ^{235}U is shown as well. Interesting to note is that the cross section of a fission reaction is not very much dependent on the energy of the incoming particle, and that most primary particles have a large range in the target material. Therefore, in contrast to heavy-ion fusion evaporation reactions, thick targets can be used, resulting in a much higher intensity of the reaction products. At ISOLDE a uranium target of 100 g/cm^2 is bombarded with a 1 GeV , $2\mu\text{A}$ proton beam resulting in an intensity after mass separation of 3×10^{10} atoms/s for $^{85-90}\text{Rb}$ isotopes (see Fig. 6).

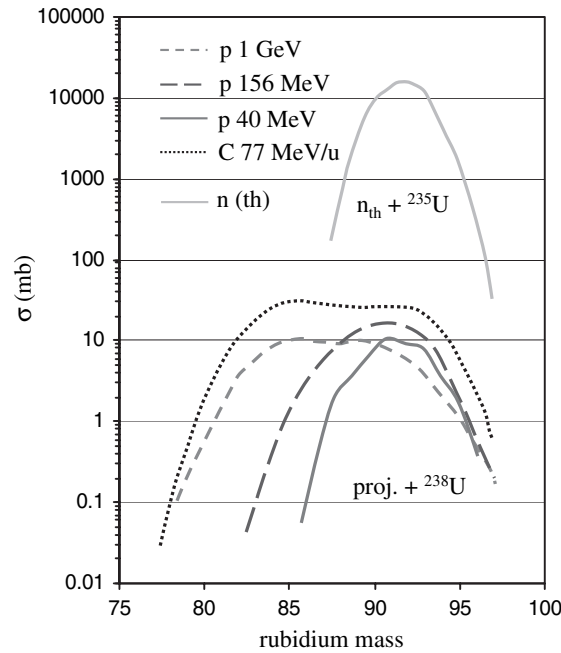


Fig. 6. Cross sections for the production of rubidium isotopes from fission of ^{238}U using various probes [42]. For comparison cross sections from thermal neutron induced fission of ^{235}U are shown. Note that the latter reaction has a large total fission cross section of 586 barns; cf. Plate 5 in the Colour Supplement

Spallation

When a high energy proton beam hits a target, three kind of reactions can occur: fission, spallation and fragmentation. In a spallation reaction a large number of protons, neutrons and α particles are ablated from the target nucleus. For example, a 1 GeV proton beam on a ^{238}U target produces ^{200}Fr , thereby ablating 5 protons and 23 neutrons. Spallation products can span a large part of the nuclear chart on its neutron-deficient side. The production cross-sections for nuclei very far from stability are modest, but the high proton energy allows one to use thick targets (of the order of 100 g/cm^2). Figure 6 shows the cross section for producing neutron-deficient rubidium isotopes via spallation. While the dominant production process is fission for neutron-rich isotopes and spallation for neutron-deficient, the transition from one type of reaction to the other is not clearly distinguishable. Thus for the high-energy proton induced reactions a smooth distribution of production cross-sections results.

Fragmentation

Projectile fragmentation reactions induced by a high-energy heavy ion beam with energies above 50 MeV/u are used at fragment separators [2]. For ISOL systems both target fragmentation, whereby a heavy target is bombarded with a high-energy proton beam [43, 44], and projectile fragmentation, whereby a high-energy (energy $> 50\text{ MeV/u}$) heavy-ion beam is sent onto e.g., a ^{12}C target [45], are used. These reactions produce again a wide variety of isotopes, close to the initial target or projectile nucleus as well as very light nuclei. The very short-lived ^{11}Li ($T_{1/2} = 8.6\text{ ms}$) is produced in this way. An overview of the different reactions on ^{238}U are given in Fig. 1 of the Lecture of J. Benlliure [46] showing the large capabilities using fission, spallation and fragmentation reactions [47].

4 Target and Ion Source Systems

The main aim of the target systems is to produce as much isotopes as possible and to get them released in gaseous form from the target matrix as fast and as efficient as possible. This implies that the target systems are generally speaking not very selective with respect to the chemical properties of the reaction products. The diffusion in the target material, however, depends strongly on these properties. Apart from bulk diffusion, release from the surface (part of the effusion process) is another parameter that depends strongly on the properties of the atoms or molecules. For short-lived isotopes of elements like titanium, vanadium, zirconium, niobium, molybdenum, hafnium, tantalum and tungsten, for example, the efficiencies for surface release are small. This means on the one hand that the radioactive ion beams will not

be contaminated with isotopes of these elements, but on the other hand that beams of short-lived isotopes of these elements must be produced by other means. A gas catcher system is here the most appropriate solution.

In a target-ion source system, the target/catcher can be connected to the ion source with a transfer tube or can be part of the ion source itself. The way this integration is done allows for implementing different selection criteria, as will be discussed together with the relevant ion source in Sect. 4.2. For clarity in the following section the target and catcher systems will be discussed separately from the ion sources even though they are sometimes intimately linked. Detailed reviews on target-ion source systems can be found in several papers, the most recent ones being [20, 21, 48].

4.1 Targets and Catchers

High-temperature Thick Targets

Thick targets are mainly used in combination with high-energy proton beams but for light-ion fusion reactions [39, 40] and fragmentation reactions [45] thick carbon catchers are exploited as well. Typical targets have thicknesses of a few 100 g/cm² and consist of foils (e.g., tantalum), fibres, liquids (e.g., mercury) or powders (e.g., uranium-carbide). The target container is heated, e.g., by sending an electrical current through the container, reaching temperatures around 2000°C. Figure 7 shows a picture of the ISOLDE target system.

The thick target-ion source systems used at ISOLDE have been studied in great detail to obtain the ϵ_{delay} of the system. Because of the pulsed structure of the proton beam (one 2.4 μ s long proton pulse of about 3×10^{13} protons every 1.2 s) the production of the radioactive ions can be measured as a function of time after the proton beam impact. Figure 8 shows a typical release curve for ^8Li ($T_{1/2} = 840$ ms) produced by target fragmentation of tantalum foils. These data can be described by analytical functions using diffusion and effusion time constants [50]. This allows one to predict the ϵ_{delay} for other isotopes of the same element and to optimize the target configuration by systematically varying it. More recently Monte Carlo methods have been implemented to calculate the loss in thick target-ion source systems. Processes like diffusion and effusion including sticking to the target material were included in the simulation and different type of target materials were studied (foils, fibres and powders) [42, 52, 53]. Figure 9 shows the simulated path of one beryllium atom produced in a target consisting of tantalum foils. On average the atoms make up to a million collisions before leaving the ion source. These simulations have reached a high level of perfection and are now used to optimize the target systems for different isotopes.

For the new generation of radioactive ion beam projects, multi MW target systems are under discussion [54]. These target developments are also needed for the muon-neutrino factory [55]. For the radioactive beams two options are

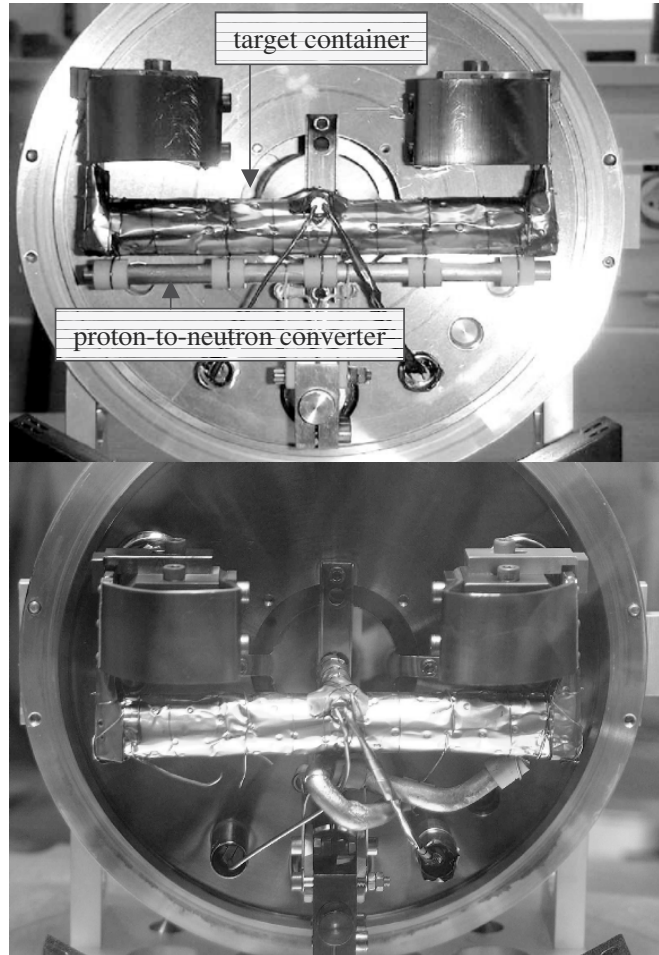


Fig. 7. Two pictures of an ISOLDE uranium-carbide target [49]. The *top* photo shows the target before proton irradiation while the *bottom* one displays the same target unit after having received 2.5×10^{18} protons from a focused beam of 3×10^{13} protons per pulse. The tantalum target container, that is wrapped in heat shields, is shown together with the leads to supply the high electrical current that heat the target. The proton beam comes from the left while the reaction products escape from the target material and container through a narrow tube installed on top of the container. The ion source, in this case a hot cavity thermal source, followed by the extraction system and beam line towards the analyzing magnet is situated behind the target container. Below the target container a tantalum rod is shown that serves as proton to neutron converter. When the proton beam is directed onto this rod, neutrons are produced which hit the target inducing fission. Because of the high instantaneous power deposition in the converter, the latter was twisted after the irradiation (see bottom photo); cf. Plate 6 in the Colour Supplement

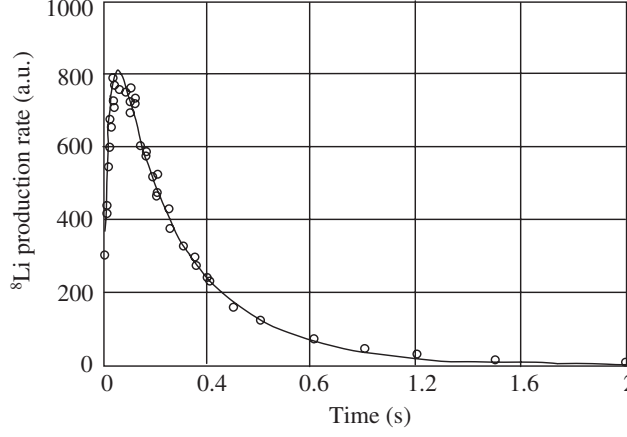


Fig. 8. Release curve of ^8Li ions from a tantalum foil target. The data points represent a measure of the production rate of ^8Li ions as function of time after proton impact. A proton pulse (typical width $2.4\mu\text{s}$) impinges on the target at $t = 0$ s. The line represents a fitted analytical expression taking into account a delay time due to diffusion ($\tau_{\text{diffusion}} = 1$ s), effusion through the target container ($\tau_{\text{effusion}} = 83$ ms) and through the ion source ($\tau_{\text{effusion}} = 13$ ms) [50]. These data have also been simulated using Monte-Carlo simulations [53]

foreseen, namely direct target irradiation with about $100\mu\text{A}$, 1 GeV protons or a proton beam of a few mA directed on a proton to neutron converter. For the latter purpose and for the muon-neutrino factory liquid mercury-jet targets are under consideration.

Thin Targets and High-temperature Catcher Systems

In heavy-ion fusion evaporation reactions the recoil products are stopped in a thin catcher foil heated to high temperatures. For several applications carbon is used as catcher material as for many elements the diffusion process is fast and graphite can stand very high temperatures. Still other materials like niobium, tantalum, tungsten and rhenium are used as well. This configuration was used extensively at the GSI ISOL facility [56]. An very interesting method to perform systematic release studies and to determine the total efficiency of the target-ion source system ($= \epsilon_{\text{delay}}\epsilon_{\text{ion}}\epsilon_{\text{trans}}$) was developed at GSI [57]. A heavy-ion beam of stable isotopes is implanted in the target-ion source system with a constant intensity. During the implantation the time evolution of the release of the isotopes implanted is registered after mass separation and from the ratio of the mass-separated to the primary beam intensity the release profile and the efficiency can be determined. The release profile is subsequently fitted to determine diffusion and effusion parameters that are used to determine the efficiency as a function of the half-life for different

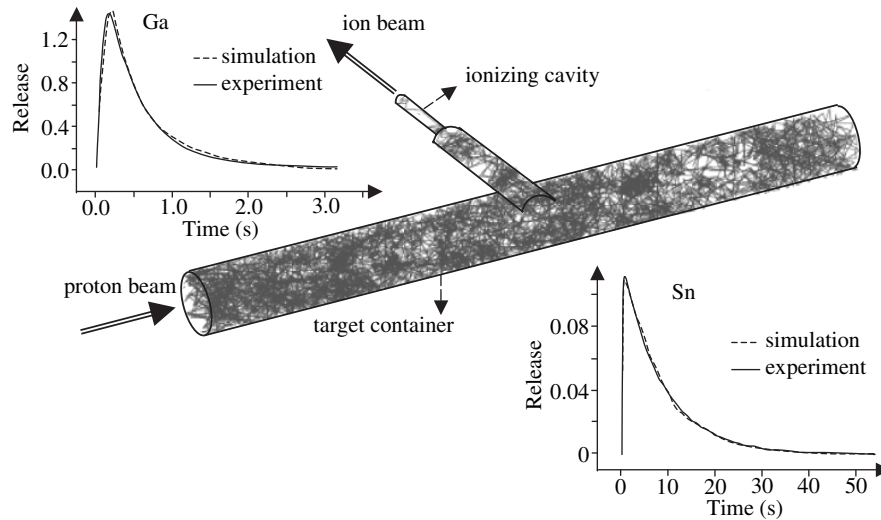


Fig. 9. Results obtained by simulating the path of one atom (not an ion!) produced in a tantalum foil target connected with a surface ionizing cavity. In this particular case, the number of collision was 78597 which correspond to a fast evacuation path. Note that the number of wall collisions in the tube where thermo-ionization takes place for this particular atoms was only 10. The influence of the plasma sheet that prevent the ions from reaching the wall of the ionizer tube is not taken into account in these simulations. The inset shows the experimental and simulated release curve of gallium and tin atoms from a uranium-carbide powder target. Note the different time scales for gallium and tin atoms. The agreement between simulation and experiment is excellent and allows reliable calculations for other target systems and other isotopes [42, 51, 52]; cf. Plate 7 in the Colour Supplement

elements. Figure 10 gives an example of the results obtained with graphite and niobium catchers at a temperature of 2300 K while the rest of the ion source was kept at 1850 K [57].

Gas Catcher Systems

Gas catchers, in which the reaction products from fission or heavy-ion fusion evaporation reactions are stopped in a noble gas, are widely implemented at ISOL systems. Mostly helium, but also argon gas at a typical pressure of 0.5 to 1 bar are used to stop the products with energies ≤ 1 MeV/u. The first gas catcher systems used were He-jet systems. In these systems the reaction products were transported with the gas flow – sometimes over distances of a few meters – towards a high-temperature ion source [58]. Because of their rather poor performances, these combinations are nowadays rarely exploited. In contrast, the IGISOL systems [28, 29] continue to be used successfully. They make use of the fact that a fraction of the thermalizing atoms is still in

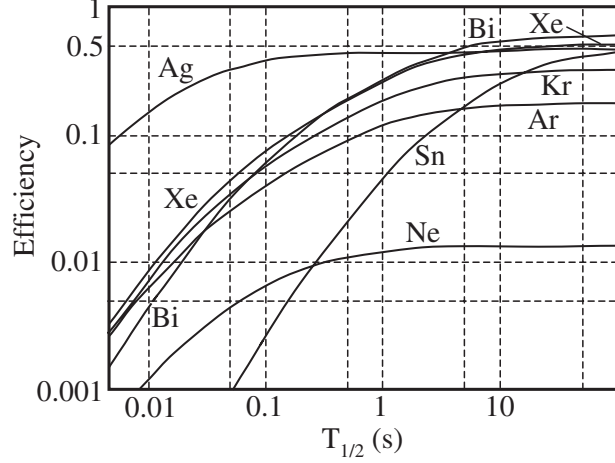


Fig. 10. Half-life dependence of the overall efficiency of a series of elements and catcher materials [57]. For silver a graphite catcher and for the other elements niobium foils were used. In all cases a discharge ion source of the FEBIAD type was used. For large values of the half-life the loss due to delay is negligible leading to a constant overall efficiency $\epsilon_{ion}\epsilon_{trans}$ that reaches 50% in some cases. One notices a very fast diffusion/effusion for the silver isotopes: even for isotopes with half-lives as short as 100 ms, ϵ_{delay} is close to 100%. This is in contrast to isotopes of tin that exhibit a very long sticking time resulting in an ϵ_{delay} of only 0.5% for isotopes with a 100 ms half-life. Note the rather low efficiency of neon which is an inherent property of the ion source used for these measurements (see Sect. 4.2)

ionic form so that an ion source in the technical meaning of the word is not necessary. Some of the second generation radioactive beam projects [31, 32] are partly based on this technique. In this way the beam quality of fragment separator beams can be improved, thus making beams of isotopes from elements available that are difficult to produce with other target-catcher systems.

4.2 Ion Sources

A comprehensive and complete overview of ion sources can be found in the book edited by B. Wolf [59]. Here the discussion is restricted to the most commonly used ion source systems at ISOL facilities, including a comment on emittance.

High-Temperature Plasma Ion Sources

Electron impact ionization is used to ionize the atoms that are present in the gas phase inside the ion source. The electron flux is created by a discharge in a

low-pressure environment. In this way a plasma is produced in which the ions are confined, preventing them from wall collisions and neutralization. Different type of so-called arc discharge ion sources have been developed (for an overview see [4, 21]). One of the most successfully and widely used sources is the forced-electron beam induced arc-discharge (FEBIAD) ion source. Here the electrons are extracted from a heated cathode and accelerated into a low-pressure plasma (see Fig. 1 of [21]). Very high efficiencies up to 50% are obtained with this source (see Fig. 10). The lightest atoms (e.g., helium, neon) have a lower efficiency as their residence time inside the plasma is short and thus the product $n_e vt$ is too low for efficient ionization (see Fig. 3). In general, arc-discharge ion sources are not selective. The energy spectrum of the electrons is broad and allows ionization of virtually every element. This property is nicely demonstrated in Fig. 10 which clearly shows that the ionization efficiency for the heavy elements (bismuth, lead, xenon, tin, silver) is close to 50% while it is lower for krypton, argon and certainly for neon. The latter element has an ionization efficiency just over 1%. At ISOLDE a FEBIAD-type ion source is coupled with the target cylinder through a transfer tube (see Fig. 7 where the case of a hot cavity thermal source is shown). By changing the temperature of the transfer tube, isotopes of less volatile elements adsorb on its walls while isotopes from gaseous elements or gaseous molecules reach the ion source. In this way a very high degree of selectivity can be obtained and pure beams of noble gas isotopes are produced. Another elegant way to obtain selectivity was the development of a FEBIAD ion source with bunched release at GSI [60]. The method is based on the fact that different elements have different enthalpies of adsorption (ΔH_a); the higher the enthalphy of adsorption, the lower is the release efficiency of the respective element from the surface. The principle of this method is as follows. During a certain time all isotopes produced in the reaction are condensed on a cold spot. Subsequently the spot undergoes a fast thermal cycle and is heated. As a consequence the isotopes from elements with the lowest enthalphy of adsorption will be released first, followed by the isotopes from the element with the second lowest value and so on. The power of this technique is illustrated in Fig. 3 of [60] where time profiles of thallium ($\Delta H_a = 3.1$ eV), silver ($\Delta H_a = 3.4$ eV) and bismuth ($\Delta H_a = 3.6$ eV) beams are shown. During an on line experiment, such cooling-heating cycle are repeated many times, with the time periods for measuring the nuclear properties of interest being optimized with respect to the corresponding release profile.

Another technique used in combination with various ion sources is the so-called decay release method. It was applied at ISOLDE making use of the pulsed structure of the proton beam when investigating neutron-rich nuclei around ^{208}Pb . These nuclei are produced in spallation reactions of uranium which, however, yield contaminants such as francium and radium isotopes, with orders of magnitude higher rates. Fortunately, these isotopes have half-lives in the μs region. Thus, introducing a delay of a few milliseconds after

the proton beam impact allowed a suppression of these unwanted isotopes by many orders of magnitude and enabled spectroscopic studies in this region of the nuclear chart [61].

Finally, it should be noted that the creation of molecular sidebands can yield very pure beams. This delicate technique has been applied for a very long time. Recent success has been obtained by adding sulphur to the ion source producing very pure beams of tin isotopes from fission [62]. More details on the molecular-ion techniques can be found in [21].

Hot Cavity Ion Sources

When a tungsten cavity or tube is heated to temperature well above 2300 K high ionization efficiencies for elements with modest ionization potentials can be reached, see (15) and Fig. 4. As pointed out before these ion sources produce a high degree of selectivity for selected elements. Attempts to create negative ion beams using surface ionization are also made. Results from an early example can be found in [63].

Electron Cyclotron Resonance Ion Sources

In ECR ion sources a plasma is confined in a bottle-type magnetic structure consisting of a solenoid and a radial multipole field. Electrons are confined by the magnetic field and ions through the charge neutrality condition of the plasma. The electrons of the plasma are heated by injecting RF power at a frequency that is suitable for obtaining a resonant transfer of energy between the RF and the electrons. Because of the excellent confinement and availability of powerful RF generators, high plasma densities and electron energies can be reached in this way. High ionization efficiencies are obtained also for the lightest elements [64]. ECR sources have been successfully used for the production of 1^+ or low-charge state radioactive ion beams of gaseous elements at Louvain-la-Neuve and Triumf [65, 66]. Because of its low-temperature this type of ion sources is very robust but beams of non-gaseous elements or molecules are difficult to produce (see below). Moreover, the source does not exhibit selectivity and produces strongly contaminating beams of stable isotopes. This is a limiting feature when performing, e.g., experiments with post-accelerated exotic beams of low intensity, say 1000 atoms/s. The high plasma densities and electron energies make these sources very efficient for the production of high charge states. At GANIL, high efficiencies have been obtained for multiply charged ions of radioactive isotopes of gaseous elements, using a projectile fragmentation reaction on a carbon target that was coupled through a room temperature transfer tube with the source [67].

Laser Ion Sources

Laser ionization has been implemented in two ways at ISOL systems, namely in a hot cavity and in a gas cell. A third method based on first condensing the

atoms on a cooled surface, subsequently ablating them by a high-intensity laser pulse and finally laser ionizing them has been used for laser spectroscopy purposes but has not yet been installed in an on line target-ion source system [68]. As pulsed lasers with a very low duty factor are employed, the aim of these systems is to store the atoms in between laser pulses in a gaseous phase and in their atomic ground state or metastable state.

In the hot cavity approach the system is in essence identical to a hot cavity ion source (see Sect. 4.2). After diffusion in the target and effusion towards the cavity, the atoms are kept in the gaseous phase in the cavity and irradiated several times with the laser pulses. As the residence time in the cavity is short (a few hundred μs), lasers with high repetition (≥ 10 kHz) are required. Typically copper-vapour lasers are used, but new developments involve solid-state lasers as well [27]. Although the selectivity of the photo ionization process is inherently very high due to the high temperature of the cavity, isotopes of elements with a low ionization potential are ionized as well, leading to isobaric contamination. Efficiencies close over 10% are routinely obtained. New cavity materials with lower work function to reduce surface ionization are under investigation. More details can be found in [69].

In the gas cell approach the fast and universal thermalization of reaction products in a buffer gas, the high selectivity and efficiency of photo ionization and the ion/atom-storage capacity of noble gases are combined. After thermalization a major fraction of the reaction products are neutralized after a few milliseconds due to the large ion-electron pair production of the primary and secondary beam. These atoms are then laser ionized and extracted from the gas cell [72]. The laser ionization happens in a place with a lower electron density or in between beam pulses. Because of the delicate interplay between the necessary electron density for fast neutralization and the unwanted production of electrons after the atoms of interest are photo ionized, the efficiency depends on the primary beam intensity. The total selectivity of gas cell laser ion source is determined by the production of ions in the reaction, their thermalization in the gas cell and their survival during the evacuation from the gas cell. Extensive tests have been performed to map out and to understand the working area of these laser ion sources. Detailed descriptions on their performances can be found in [70, 71, 72].

The purity of the resulting radioactive ion beam is an important parameter and laser ion sources play a crucial role in reducing these contaminations. Figure 11 shows γ -ray spectra with and without lasers from a Coulomb excitation measurement of ^{74}Zn .

Gas Catcher Sources

In order to overcome the limitation of high-temperature target-ion source systems concerning the production of radioactive beams of isotopes of certain type of elements, in particular the very short-lived ones, the IGISOL method

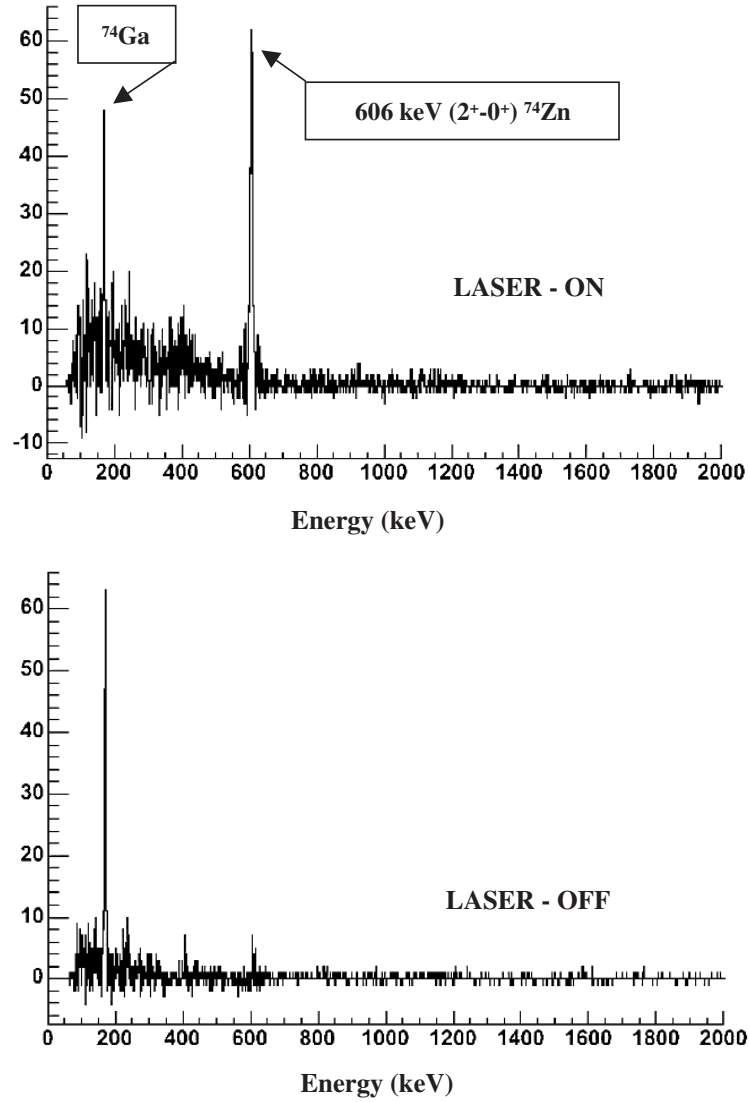


Fig. 11. Gamma-ray spectra obtained by the germanium MINIBALL array from a Coulomb excitation measurement using a post accelerated (2.8 MeV/u) ^{74}Zn beam from ISOLDE. In the laser-on spectrum the 2^+-0^+ transition at 606 keV is clearly present, while it is absent in the laser-off spectrum. The line at 171 keV from Coulomb excitation or a transfer reaction of the contaminant ^{74}Ga is present in both spectra. Having the laser frequency as a supplementary parameter allows to assign contributions from beam contamination in a reliable way [73]

has been developed [28, 29]. This system has been very successful in producing short-lived isotopes and isomers of elements, like tantalum, tungsten, molybdenum and rhodium that are not accessible to high-temperature ion sources. As discussed before the key issue in ion guide systems is fast evacuation, avoiding neutralization. When neutralization is allowed by inducing a longer evacuation time, laser ionization can be applied in gas cells [70]. Attempts to place a gas catcher after an fragment separator and thus removing the primary beam as well as the major part of the secondary beams are currently made [30, 31, 32]. Because of the large energy spread of the fragmentation beams (about 5 MeV/u [2]) large gas cells of about 1 m length and 30 cm diameter filled with 500 mbar helium are considered. As the evacuation time using the gas flow only is much too long, DC- and RF-electrical fields are used to speed it up. Furthermore the electrical fields collect the electrons created during the slowing down process and thus reduce the neutralization rate of the ions. While these very promising schemes are tested, other solutions like adding laser ionization to this system are considered as well.

Emittance

An important property of the ion source is the ion optical quality of the ion beams, expressed as emittance. In essence one considers the velocity of the beam in a plane perpendicular to the ideal beam direction. Instead of the velocity vector profile, one plots the divergence angle in x direction as a function of the x position ($x' = v_x/v_z$), and similarly for the y direction. This results in a typical plot as presented in Fig. 12. The 95% radial emittance of an ion beam is then defined as the area of an ellipse that contains 95% of the beam intensity (see Fig. 12)

$$\epsilon_x = \int \int dx dx' / \pi \quad (19)$$

and this for the x and y direction. The unit of emittance is π mm mrad. The quantity π resulting from the surface of an ellipse is included in the units: the area of an ellipse with its axis along the main axis would have as surface $x_{max} x'_{max} \pi$. The measured emittance diagram (see Fig. 12) can have very peculiar shape. Still it is the surface of the ellipse that includes a specific percentage (in this particular case 95%) of the beam intensity. The latter quantity is generally considered for specifying the emittance as the standard ion-optical elements (lenses, beam pipes, accelerators) have in most cases ellipsoidal acceptances. Ion optical elements like lenses or deflection plates can manipulate an ion beam, thereby turning the ellipse or reshaping it, but the total surface stays constant. The task of focussing an ion beam through a small collimator needs ideally an ellipse with its long axis along the y -axis thereby reducing the spatial spread of the beam. It is easy to understand

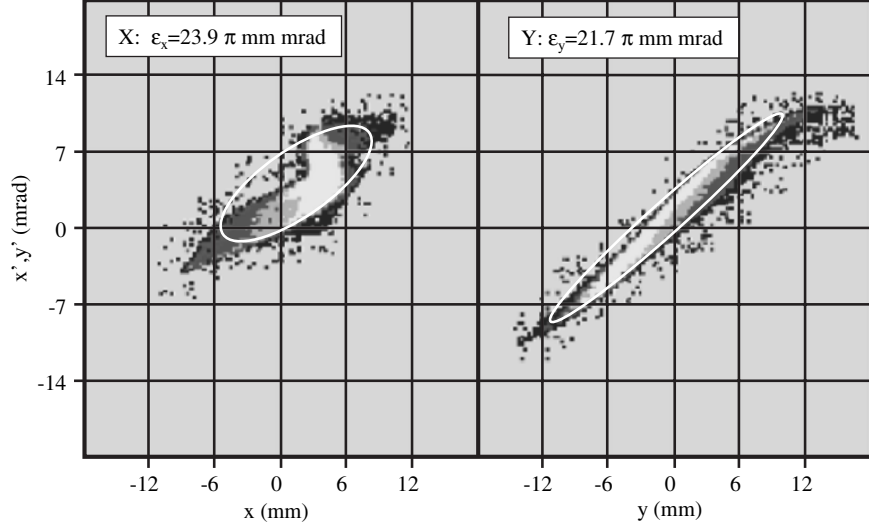


Fig. 12. Example of the results of an emittance measurement for a ^{120}Sn beam from the ISOLDE laser ion source [33]. The ellipsoidal contour covers 95% of the beam intensity, the surface of the ellipse representing the final value for the emittance as indicated in the inset. Although the relative difference between ϵ_x and ϵ_y is small, the shape of the two emittance curves is distinctly different; cf. Plate 8 in the Colour Supplement

that after acceleration of the ion beam the absolute emittance decreases and the contour plot in Fig. 12 becomes smaller. The normalized emittance is a conserved quantity of the ion beam, defined as

$$\epsilon_{nx} = \beta \epsilon_x \quad (20)$$

with $\beta = v/c$. Along similar lines, to slow down an ion beam needs specific precautions as the emittance will increase, resulting in a larger beam spot. Apart from the radial emittance, the longitudinal emittance, ϵ_{long} , is important when e.g., considering the injection of beam pulses into a trap or accelerator structure. It is defined as the product of the time width of the pulse, Δt , and its energy spread, ΔE ,

$$\epsilon_{long} = \Delta E \Delta t \quad (21)$$

After extraction from the source the ion beam has to be transported, mass analyzed and eventually post-accelerated. The transport system, analysing magnets and post-accelerator have limited acceptances. In order to reduce the beam losses to a minimum the radial and longitudinal acceptance of these devices must be equal to or larger than the emittance of the beam. Recently RF-coolers and Penning traps have been developed to cool the ion beam, thus reducing the emittance of an ion beam (see Sect. 2). These very

successful devices, which are described in [6], allow, among other applications, to slow down ion beams and capture them in ion traps, to softly deposit radioactive ions on surface for solid-state studies or to adapt the emittance to the acceptance of spectrometers, beam transport systems and accelerators.

5 Mass Separation and Post Acceleration

In this section the methods for mass separation and post-acceleration are presented. It aims at describing some developments in the field of radioactive ion beam manipulation and post acceleration, defining a few important parameters and showing their impact on the conditions for radioactive beam experiments.

5.1 Mass Separation

After the ions are created in the ion source, they are extracted and accelerated in a DC electrical field created by a high-voltage potential difference between the ion source and an extraction electrode or the beam line. The ion source, and in most cases the target system as well, are polarized with a positive voltage between 20 and 60 kV resulting in energies of 20 and 60 keV, respectively, for 1^+ ions. In case of negative ions the ion source is put on negative potential. With an ion optical system of electrostatic or magnetic lenses the beam is subsequently transported to an analyzing magnet. In ISOL systems a dipole magnet is used. Details on the transport system and the properties of the analyzing magnet can be found in [4]. As mentioned above the quality of the mass analyzing system is expressed by its mass resolving power that typically varies from a few hundred to a few thousand. High-resolution mass separators have been constructed which reach resolving powers beyond 10^4 but their transmission often suffer from the large emittance of the ion sources used. By limiting the acceptance of the mass analyzing system the mass resolving power can be increased, but the transmission is correspondingly reduced, which is in most cases not acceptable. Another pitfall is that the mass resolving power does not give the complete picture as shown in Fig. 13. RF-coolers can deliver pencil like ion beams with small radial and longitudinal emittance, which should allow one to obtain the ultimate mass resolving power of the analyzing system [6]. These systems are currently under development, with interesting results being expected in the near future. The mass analyzing dipole magnets used at ISOL systems allow also for beam sharing as different masses are available at the same time in the focal plane of the mass separator. These can be sent to different experimental set-ups. Especially when using spallation, fission or fragmentation reactions, high cross sections for the production of a series of isotopes can be obtained, resulting in intense beams at several different masses.

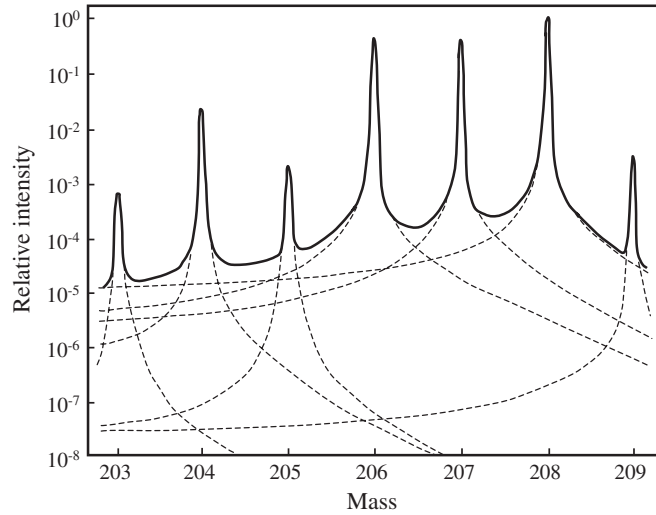


Fig. 13. Spectrum of an ion beam obtained after mass separation. The mass peaks have long and asymmetric tails. These are caused by collisions of the ions with gas molecules when the ions are extracted and transported through the beam lines [4]. The collisions induce a spread in energy and change the direction of the ions resulting in different positions at the focal plane. As a consequence at every mass contributions of nearby (intense) mass peaks can be observed, even when the mass resolution is very good. It is illustrative to see that the contamination due to the low-mass tail of the $A = 208$ peak at the position $A = 207$ is a few times 10^{-5} of the intensity of the mass $A = 208$ peak and stays above 10^{-5} even at $A = 203$. Narrow slits around $A = 205$ result in a beam with contaminations of $A = 208$ and 206 at a level of 1% of the $A = 205$ intensity

5.2 Post Acceleration

The beam of singly charged radioactive ions is accelerated by using a variety of different post accelerators. As mentioned in Sect. 2.1, highly charged ions are needed in most cases. At Triumf and Oak-Ridge the singly charged ions are accelerated directly and stripping is performed during the acceleration process. For all other projects, charge state breeders are required. Three type of accelerators are used for post acceleration: cyclotrons, linear accelerators and tandems. It is outside the scope of this lecture to give an overview of accelerators and their properties. Thus only specific aspects relevant for radioactive ion beam acceleration are addressed.

Charge State Breeders

Highly charged ions are needed for acceleration in cyclotrons and their use reduces the cost of a linear accelerator. Charge state breeders have been developed to boost the 1^+ ion beam to a higher charge state. Again this process

has to be efficient and fast without inducing extra beam contamination from stable isotopes.

ECR sources and EBIS are used to obtain higher charge states through electron impact ionization. The ECR source has already been described in Sect. 4.2. Normally the ECR source is coupled via a cold transfer line to the target. But when used as a charge-state breeder, a 1^+ ion beam is directly injected into the ECR source and subsequently slowed down and captured by the dense plasma of the ECR source. The 1^+ to n^+ ($n \geq 2$) efficiencies of ECR sources can reach 12% with a breeding time of 120 ms for indium and lead ions [74]. The efficiency summed over all charge states larger than 1 is 65%. The beam contamination due to stable isotopes can, however, be the limiting factor for the post-acceleration of very weak radioactive ion beams.

An EBIS consists of a strong magnetic field that squeezes a high intensity electron beam along the field axis to a very small radius. This results in electron beam densities over 100 A/cm^2 . After retardation the ion beam is injected into the EBIS and starts counterpropagating the electron beam. After injection the ions are captured in an axial potential well created by an electrode structure and are radially confined by the negative space charge of the electrons. Under continuous electron bombardment, the ions get highly ionized. As can be seen in Fig. 3 argon ions reach an 8^+ charge state after 25 ms breeding time with an 17 keV electron beam of 120 A/cm^2 . When the desired charge state is reached the potential is lowered on the ejection side allowing the highly charged ions to escape from the EBIS. Charge breeding efficiencies of 50% are reached for sodium and magnesium after 20 ms breeding time, the breeding efficiency for one particular charge state being 15% [75]. In contrast to an ECR source, that can capture ions in continuous mode and delivers continuous beams, an EBIS needs a pulsed ion beam for injection and yields a pulsed beam (width about $100 \mu\text{s}$). At REX-ISOLDE the 1^+ ion beam is injected in a Penning trap to bunch the ions prior to injection into the EBIS [6, 43]. This pulsed ion beam structure is needed for linear accelerators as they often work with a specific duty factor (see below).

Cyclotrons

The radioactive ion beam facility that produced post accelerated radioactive beams for the first time is situated at Louvain-la-Neuve [37, 39, 40]. There as well as at the SPIRAL facility at GANIL [45], cyclotrons are used for the primary beam production and for post-acceleration. Because of their limited A/Q acceptance cyclotrons need highly charged ions. At Louvain-la-Neuve and GANIL an ECR source is directly coupled to the target through a cold transfer line. Because of the magnetic field used in the accelerating process of a cyclotron, the mass analyzing power of these accelerators is very good. This results in a strong suppression of isobaric contaminants without loss in transmission. However, as the ECR source is connected to the cold transfer line, essentially only beams of isotopes of gaseous elements or very long-lived

isotopes of non-gaseous elements can be produced. Recent developments aim at using ECR sources as 1^+ to n^+ charge-state breeders which will allow to accelerate also radioactive isotopes of non-gaseous elements.

Tandem Accelerators

At the Holifield post accelerator facility HRIBF [62] a cyclotron is used for the primary production and a tandem post accelerator is employed. As a tandem accelerates negative ions, one needs negative ion sources or a charge exchange system that has to be installed in the beam line between ion source and accelerator. The latter option is chosen in most cases at HRIBF.

Linear Accelerators

Several radioactive ion beam facilities use linear accelerators. At ISOLDE the beams from the EBIS are accelerated to 3.1 MeV/u while at TRIUMF 1^+ beams are accelerated via a stripping stage to 1.5 MeV/u. Both facility are currently upgraded to reach energies over 5 MeV/u.

Time Structure

Tandem accelerators deliver DC beams. Apart from the RF time structure, cyclotrons do as well. Typical RF frequencies used at cyclotrons are a few 10's of MHz resulting in a beam pulse every ~ 100 ns with a width of a nanosecond or less. On the other hand, most room temperature linear accelerators are operated with a certain duty factor. Its value depends on the amount of RF power (heating) that the accelerating structure can stand. Super-conducting accelerators do not have this problem and can run in continuous wave (CW) mode. At Triumf the room temperature linear accelerator (ISAC-I) delivers a beam with 100% duty cycle while at ISOLDE the accelerator is on for 1 ms followed by a 19 ms off period. As this sequence is synchronized with the cooling and charge state breeding process no losses due to the duty factor are encountered.

6 Overview of the ISOL Facilities

At different places ISOL facilities are operational and some have extended their energy domain towards higher energy to allow reaction studies with radioactive beams. Table 2 gives an overview of the on line isotope separators producing low-energy (~ 50 keV) radioactive ion beams. First generation ISOL based facilities for the production of energetic beams are briefly described below including some technical details. Some projects are preparing for an upgrade of their facility to increase the intensity and energy of the final beam or to widen the spectrum of beams available.

Table 2. Overview of the ISOL facilities producing low-energy (~ 50 keV) radioactive ion beams. At IGISOL, Triumf and PARNNE-ALTO the development of a laser ion source is underway. At the latter facility an electron accelerator for photo fission is currently under construction. Information on these facilities can be found in [3]. Apart from the facilities mentioned below, developments are pursued at Beijing, China, and Warsaw, Poland. The GSI ISOL facility has been closed in 2004

Facility	Beam Targets	Reactions Ion Source
ISOLDE-CERN Switzerland	p(1.4 GeV, 2 μ A)	fragmentation, spallation, fission
	high-temp., thick targets	plasma, hot cavity, laser
GSI Germany	heavy ions	fusion evaporation, transfer
	thin targets	plasma, hot cavity
LISOL Belgium	p(30 MeV, 10 μ A), heavy ions	fusion evaporation, fission
	thin targets	gas catcher, laser
IGISOL Finland	p(30 MeV, 30 μ A), heavy ions	fusion evaporation, fission
	thin targets	gas catcher
ISAC-TRIUMF Canada	p(500 MeV, 100 μ A)	fragmentation, spallation
	high-temp., thick targets	plasma, hot cavity
OSIRIS Sweden	n(3×10^{11} cm $^{-2}$ s)	fission
	high-temp., thick ^{235}U target	plasma
PARNNE France	d(26 MeV, 0.25 μ A) $\rightarrow$ n convertor	fission
	high-temp., thick ^{238}U targets	plasma
IRIS Russia	p(1 GeV, 0.2 μ A)	fragmentation, spallation, fission
	high-temp., thick targets	plasma, hot cavity, laser
JAERI Japan	p(18–36 MeV, 3–10 μ A)	fission
	gas jet and high-temp., thick ^{238}U targets	plasma, hot cavity

RIB Project at Louvain-la-Neuve, Belgium, Operational Since 1989

The first ISOL-based energetic radioactive beams have been produced by means of the radioactive ion beam (RIB) project at Louvain-la-Neuve in 1989. It combines two existing cyclotrons with a target-ion source system. The radioactive isotopes are produced with a 30 MeV proton beam from a compact cyclotron that can deliver beam currents up to 500 μA . Due to the power limit of the targets, proton beam intensities are restricted to about 200 μA . Direct reactions on carbon, lithium and fluorine targets are used to produce intense beams of light radioactive isotopes close to the line of stability (e.g. ^{13}N , $^{18,19}\text{Ne}$, ^6He , ^{11}C). The radioactive atoms are transported to an ECR ion source via a cold transfer line and, after ionization to a modest charge state, accelerated in a $K = 110$ or $K = 44$ cyclotron. The latter has been built for low-energy acceleration (0.2-0.8 MeV/u) while the former delivers energies up to 10 MeV/u (www.cyc.ucl.ac.be).

HRIBF at Oak-Ridge, U.S.A., Operational Since 1998

The HRIBF uses a cyclotron as primary accelerator delivering beam intensities up to 13 μA (44 MeV protons and 44 MeV deuterons) and 3 μA (85 MeV α particles) to various target materials. Direct reactions as well as fission reactions are used. As post acceleration is performed by a 25 MV tandem accelerator negative ions are needed. High-temperature negative ion sources as well as plasma ion sources are used. In the latter case, the beam is passed through a cell with cesium vapor, in which the 1^+ ions are changed into 1^- ones prior to injection into the tandem. Laser ion source developments are underway. For $A \leq 100$, energies up to 5 MeV/u are reached, for heavier masses supplementary stripping inside the tandem is required (www.phy.ornl.gov/hribf).

ISAC Facility at Triumf, Canada, Operational Since 1998

At the ISAC facility, a primary proton synchrotron accelerator delivers 500 MeV beams up to 100 μA . The radioactive isotopes are produced in spallation and fragmentation reactions. Thick targets and hot-cavity ion sources are used, with the 1^+ ions being accelerated in a linear accelerator up to 1.5 MeV/u. The use of ECR and laser ion sources as well as actinide (e.g., uranium) targets for fission reactions is under development. With the ISAC-II upgrade the beam energy will be boosted to about 6.5 MeV/u (for masses up to 150) by extending the linear accelerator. A new experimental hall has been constructed to host the new post-accelerator section and related experimental equipment (www.triumf.info/facility).

REX-ISOLDE Facility at CERN, Switzerland, Operational Since 2001

At CERN, proton beams up to 1.4 GeV are delivered by the proton-synchrotron booster accelerator with an average intensity of $2\mu\text{A}$. Target spallation, fragmentation and fission produce a wide spectrum of radioactive ion beams. Plasma, hot-cavity and laser ion sources are used to produce 1^+ ion beams. These beams are cooled and bunched in a buffer gas filled Penning trap, charge state bred in an EBIS and accelerated with a linear accelerator to 3.1 MeV/u. With the REX-ISOLDE post-accelerator system in principle all available low-energy beams of ISOLDE can be accelerated. The HIE-ISOLDE upgrade aims at increasing the intensity primary proton beam to $10\mu\text{A}$ and the energy of the radioactive beam to values above 5 MeV/u (isolde.web.cern.ch/isolde).

SPIRAL Facility at GANIL, France, Operational Since 2001

Within the framework of the SPIRAL facility at GANIL, heavy-ion beams at intermediate energy (carbon up to 95 MeV/u, uranium up to 24 MeV/u) are sent onto a hot carbon target to produce radioactive isotopes through fragmentation reactions. Via a cold transfer line gaseous atoms are injected into an ECR source producing multiply charged ion beams. Acceleration is performed with a cyclotron up to 25 MeV/u. GANIL prepares a major upgrade called SPIRAL2. A new injector will be constructed to produce intense beams of deuterons (up to 5 mA) and heavy ions (up to 1 particle-mA) at 14.5 MeV/u. Via fusion evaporation and fission reactions a wide range of isotopes will be produced. Sources for singly charged ion will be combined with an ECR source for charge state breeding. Acceleration will be performed with the existing cyclotron CIME (www.ganil.fr).

Other ISOL Facilities

At the 88" cyclotron at Berkeley [76] and at the ATLAS linear accelerator at Argonne National Laboratory [77], a selective number of energetic radioactive ion beams are produced in a "batch" mode operation. Several other projects are currently under discussion or preparation. These include DRIBS at Dubna, Russia, EXCYT at Catania, Italy, MAFF at Munich, Germany, SPES at Legnaro, Italy, and VEC-RIB at Calcutta, India. The MAFF project will use thermal neutrons with very high intensity ($8 \times 10^{14} \text{ s}^{-1} \text{ cm}^{-2}$) to produce beams of fission products from ^{235}U [78].

On the longer term the construction of second generation facilities are considered. In North America the Rare Isotope Accelerator (RIA) is proposed (www.ria.com). Here a primary heavy-ion accelerator capable of delivering intense beams of protons up to uranium at 400 MeV/u is used for fragmentation and fission reactions. Both an in-flight separator and an ISOL system

will be constructed. Apart from a conventional target-ion source system, a gas catcher program is proposed whereby the beams delivered by a fragment separator will be stopped in a large gas cell for thermalization. The post accelerator will bring the beam up to an energy of 12 MeV/u. In Europe the European Isotope Separator On Line (EURISOL) is in a design study phase [54]. A multi MW proton accelerator will be used as primary accelerator to induce spallation, fragmentation and fission reactions. Two target options are considered, i.e. direct irradiation of the target with maximum intensities of 100 μ A and irradiation of a proton to neutron converter with beams up to 1 mA (see Sect. 3). A linear post accelerator will deliver a beam energy of up to 100 MeV/u (the latter design aim is under discussion). In Japan the J-PARC project is in preparation, planning to use the 3 GeV proton beam for the production of radioactive isotopes and to accelerate them with a linear accelerator up to 9 MeV/u (jkj.tokai.jaeri.go.jp).

7 Outlook

The field of ISOL systems is now more than 50 years old. Still it is very much “alive and kicking” as witnessed by the many new projects to produce energetic radioactive ion beams. Especially the new techniques to produce purer and more intense beams as well as to post-accelerate them in combination with the renewed quest for understanding atomic nuclei with very unusual proton to neutron ratios, have made this field to blossom again. The recent development of laser ionization, the possibility to obtain primary beams with much higher intensities and the development of efficient charge state breeders are only a few examples of the recent technical developments. Also the progress in instrumentation (segmented germanium and silicon detectors, traps and coolers, digital electronics to name a few) has led to a vast extension of the experimental possibilities. Furthermore the use of pure radioactive ion beams in other fields of science, such as the cost effective production of radioactive isotopes for medical purposes and the use of radioactive tracers in solid state physics, push this field further. The first ISOL beam became available in 1951 [79], the first post-accelerated radioactive ion beam was produced in 1989 [37, 38] and since 1998 four new radioactive beam post accelerators have been commissioned (HRIBF, ISAC, REX-ISOLDE, SPIRAL) and are taking data. Intermediate plans to upgrade the currently running facilities and the long term plans to construct second generation facilities all point to the importance of this research which will yield a better insight into the structure of exotic nuclei, the manifestation of the strong and weak interaction in atomic nuclei and the way the elements are formed in astrophysical nucleosynthesis processes.

Acknowledgements

I would like to thank several of my colleagues with whom I discussed the subject and in particular H.L. Ravn and R. Kirchner for the excellent training they gave me in the field of ISOL systems.

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Gamma-Ray Arrays: Past, Present and Future

W. Gelletly¹ and J. Eberth²

¹ Physics Department, School of Electronic and Physical Sciences, University of Surrey, Surrey GU2 7XH, U.K.

² Institut für Kernphysik, Universität zu Köln, 50937 Köln, Germany

Abstract. In this lecture the development of arrays of γ -ray detectors is described. The starting point is the reason why we want to study γ -rays emitted in nuclear reactions. An account of how gamma rays interact with matter follows. The characteristics of the γ -ray detectors needed for the study of both exotic nuclei and high spin states are then outlined. The development of γ -rays arrays is followed up to the present introduction of detectors based on gamma-ray tracking. Examples taken from experiment are used throughout.

1 Introduction

In these lectures we shall be concerned with the development of arrays of γ -ray detectors for studies of nuclear spectroscopy. Although the topic is thus limited in scope we shall find that it touches on many other aspects of nuclear spectroscopy, since it cannot be isolated from them. Where, then, should one begin? This is always a difficult question: the choice is often just a matter of taste. Since the need and desire for improved multi-detector germanium (Ge) arrays has its roots in spectroscopy we shall begin there.

The study of any physical system is likely to follow a general pattern: nuclear structure studies are no exception. We have tried to identify key parameters in the nuclear system and then study how nuclear properties vary with these parameters.

Figure 1 shows schematically a few of the quantities of importance in the nuclear system. One of the key parameters is angular momentum, which dictates how rapidly the nucleus is rotating. This is important because atomic nuclei are not really influenced by electric and magnetic fields of the strengths we can generate in the laboratory. Applying such fields to other physical systems is often one of the principal tools for studying them. In nuclei the Coriolis force plays this role. It takes the same form mathematically as the force exerted by a magnetic field and so it mimics the effects. As a result these lectures are very much concerned with the study of so-called “High Spin” states; so known because they are rotating rapidly. Most of the drive to build and improve γ -ray arrays has come from such studies. However, as often happens in science, they have found many other uses. In particular they have turned out to be important for the study of exotic nuclei, the subject of the Euroschool Lectures.

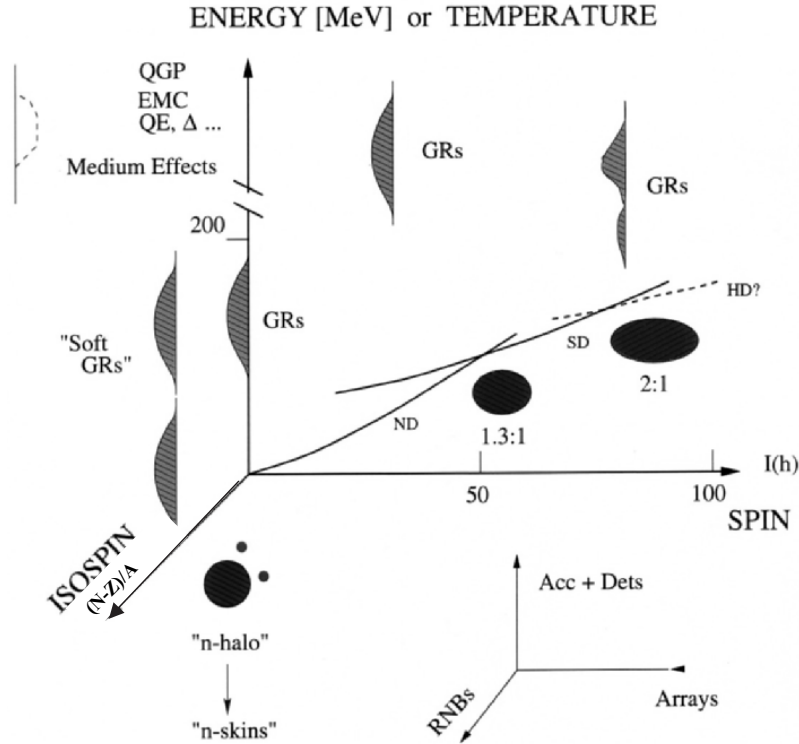


Fig. 1. Diagram showing the space defined by three key parameters describing the atomic nucleus. Studies of nuclear reactions allow one to vary these parameters, namely the excitation energy, angular momentum and the ratio of neutrons to protons. The various cartoons are intended to give an impression of some of the phenomena observed as the key parameters are varied. To indicate just a few of the many phenomena observed as we vary the parameters, a series of abbreviations are used. They are as follows: GR-Giant resonance, QGP-Quark gluon plasma, EMC-The EMC effect, Δ -Delta resonance, ND-Normal deformed states, SD-Superdeformed states, HD-Hyperdeformed states, n-halo-Neutron halo, n-skins-Neutron skins, RNBS-radioactive nuclear beams, Acc + Dets-accelerators and detectors, SPIN-total angular momentum. The reader should further note that the *scales* on the axes are notional and the break in the energy/temperature axis is simply meant to indicate that the various effects indicated occur at much higher excitation energies. The small set of labelled axes in the *bottom right hand corner* is meant to indicate the key element that needs to be developed to progress further with examining how each parameter affects nuclear properties and phenomena. As will be seen from the main text these technical developments are also important for exploring these and other parameters; cf. Plate 9 in the Colour Supplement

Our starting point is the studies [1] of Morinaga and his co-workers in the 1960s. In these experiments α -particle beams of 27 MeV energy were used to bombard targets of rare-earth elements and the prompt γ -rays released were detected with NaI scintillators clustered around the target. An example of the γ -ray spectra recorded in such studies is shown in Fig. 2. It turns out that we see a series of peaks sitting on an underlying, continuous background of γ -rays: they are the successive transitions de-exciting levels in the ground-state rotational band of ^{162}Dy . From many such studies, then and later, we have built up a picture of so-called fusion-evaporation reactions of which the $^{160}\text{Gd}(\alpha, 2n)^{162}\text{Dy}$ reaction of Fig. 2 is an example.

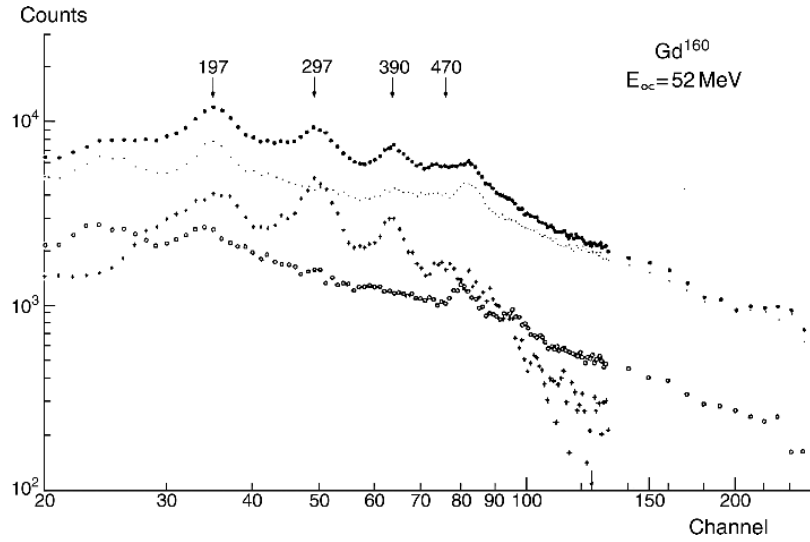


Fig. 2. Example of the γ -ray spectra recorded in early studies of fusion-evaporation reactions. In this case the γ -rays from the $^{160}\text{Gd}(\alpha, 2n)^{162}\text{Dy}$ reaction recorded with a NaI detector in singles. The γ -ray photo-peak energies are labelled in keV. The target was composed of an oxide of Gd and the various spectra correspond to measurements with various thicknesses of Lead between the target and the counter (see [1] for details). This spectrum is shown courtesy of Nuclear Physics

A large fraction of the studies which have been carried out with large γ -ray arrays involve fusion-evaporation reactions, although Coulomb excitation and other reactions have been used as well. The picture we now have of such reactions is summed up in Fig. 3. In the initial stage of the reaction the two nuclei fuse to form a compound system that exists for a long time on a “nuclear” time scale; usually taken to be the time for the projectile nucleus to cross the target nucleus, i.e. $\approx 10^{-22}$ s. This compound system is “hot”, i.e. it is at high excitation energy and rotating rapidly. The angular momentum (J) can be as large as 50–80 $\hbar$, i.e. a rotational frequency of $\sim 10^{20}$ Hz. In

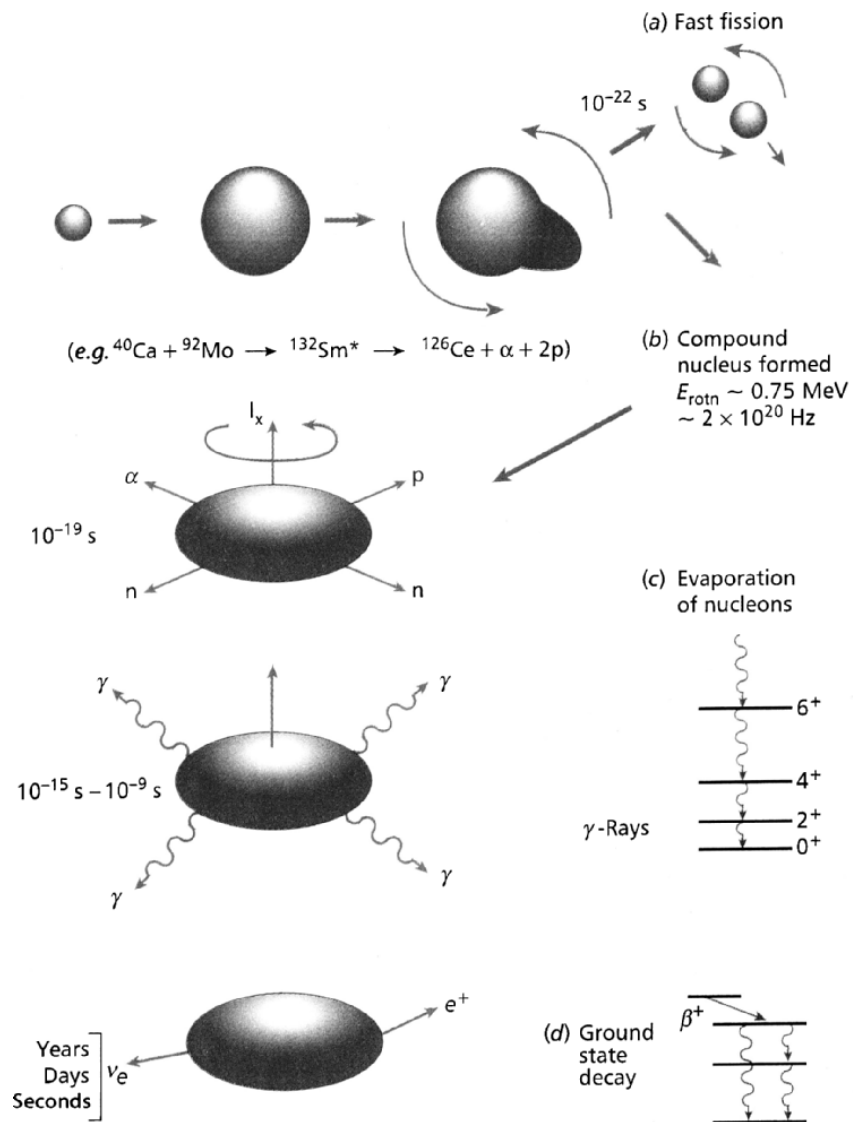


Fig. 3. Schematic view of what happens in a fusion-evaporation reaction (see text)

essence we now have a hot, electrically charged, rapidly rotating liquid drop. Not unexpectedly, it does what any liquid drop would do to reduce its temperature: it evaporates particles. These neutrons, protons and alpha particles carry away energy but little angular momentum since the centrifugal barrier they face means that the emission of particles with high angular momentum is strongly inhibited. The evaporation comes to an end when the nucleus is left in an excited state with an energy below the lowest particle separation energy. The nucleus is now cooler but still rotating very rapidly. The only way it can lose more energy is by emitting γ -rays, which also carry away angular momentum. Thus a long sequence of 30–40 γ -rays follows, leading, through a series of excited states, to the ground state. In even-even nuclei the lowest states normally belong to the rotational band built on the 0^+ ground state and have $J^\pi = 2^+, 4^+, 6^+$ etc. The many possible γ -ray cascades from the initially populated states funnel down into this ground state band or its equivalent in odd-A nuclei. It is these intense γ -rays which we see in Fig. 2, albeit in poor resolution and with poor signal-to-noise ratio. It is studies of the γ -rays from such reactions, which are the backbone of our understanding of the properties of high spin states.

It is not our purpose here to pursue our understanding of these states: rather we want to show how advances in methods of detecting them have gradually given us more and more sensitive tools to study γ -rays from many different types of reaction. This will lead us to an appreciation of their importance for the study of exotic nuclei.

Julin [2] has summarised how Ge detectors, with or without ancillary detection systems, have been used to study the properties of exotic nuclei: the interested reader is referred to his lecture. We can, however, briefly summarise the essence of what we have discovered using even-even nuclei as the vehicle. In the ground state the nucleons are paired off in time-reversed orbits. In excited states with non-zero values of J , the Coriolis force acts to try to align the angular momenta associated with these pairs with the axis of rotation.

Figure 4 shows some of the effects observed and interprets them in cartoon form. Firstly we should recall that the Coriolis force takes the form $-\omega \times j$ where ω is the angular frequency of rotation and j is the angular momentum of an individual nucleon. Any effect of the force clearly increases with both rotational frequency and the angular momentum of the nucleon. Figure 4 shows the spectrum of γ -rays de-exciting levels in the ground state band of ^{158}Er produced in the $^{146}\text{Nd} + ^{16}\text{O}$ reaction at 84 MeV bombarding energy. In such experiments we do not observe the levels directly: we infer their positions and properties from the de-exciting transitions between them. It is necessary to use a medium- or heavy- mass projectile such as ^{16}O in order to create a compound nucleus of high angular momentum, as the reader will appreciate from examination of Fig. 3. For nuclei such as ^{158}Er with $A \approx 160$ and $Z \approx 68$ the $i_{13/2}$ neutrons and $h_{11/2}$ protons form the intruder orbitals in the Nilsson scheme.

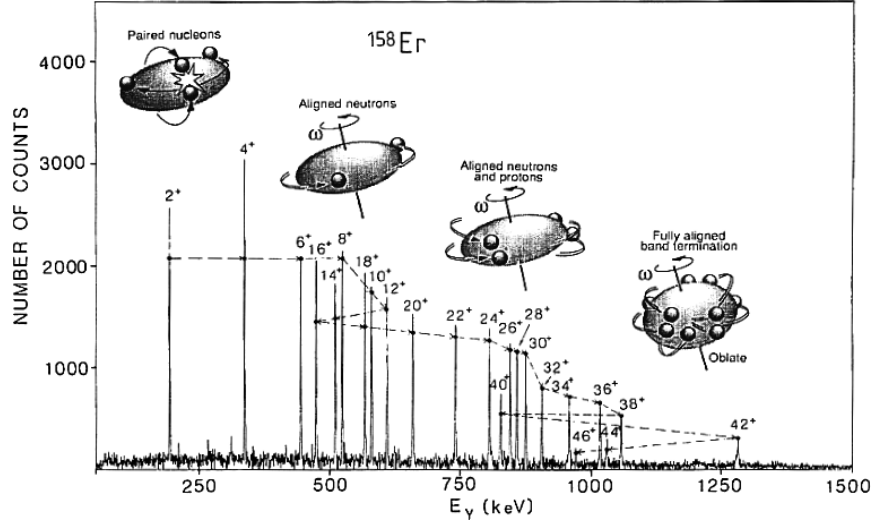


Fig. 4. Spectrum [3] of γ -rays from the $^{146}\text{Nd}(^{16}\text{O}, 4n) ^{158}\text{Er}$ reaction recorded with the Tessa array [4] (see text)

Starting on the left of Fig. 4 we see regularly spaced γ -transitions which correspond to the de-excitation of levels in the band built on the ground state. In these states the nucleons are paired off in time-reversed orbits as shown on the top left. As ω increases the Coriolis force increases and is, of course, largest for particles in the high- j orbitals. Eventually the force becomes too strong and a pair of particles is “broken”, i.e. the angular momenta of the two particles become aligned with the collective spin vector. The abrupt change in level energy is reflected in the spectrum: the $14^+ - 12^+$ transition is lower in energy than the $8^+ - 6^+$ transition. This is a so-called backbend because of the behaviour of the nuclear moment-of-inertia if it is plotted against ω . In ^{158}Er the $i_{13/2}$ neutrons are the first to break and the corresponding cartoon in Fig. 4 shows the two neutrons in an orbit with their angular momenta aligned. Similarly at $J = 32$ we see the effect of the alignment of an $h_{11/2}$ proton pair with both $i_{13/2}$ neutron and $h_{11/2}$ proton pairs now shown as being aligned. Above $J = 38$ we see several more changes in structure, which are due to the successive alignment of the remaining available pairs of particles in ^{158}Er . The full story behind this spectrum is explained in [3] and the interested reader is referred to this article.

This example shows what one might call the basic phenomenon of the decoupling of particle pairs under the stress of rotation. Many other phenomena are also observed such as superdeformation, identical bands, chiral symmetry, shears bands, Coulomb energy displacements, etc. The reader is referred to [5] for a full description. As stated earlier it was the desire to study these phenomena that has driven the development of γ -ray arrays. If

this were their only use they might still be of interest to those studying exotic nuclei since we are anxious to know about high spin phenomena there too. However one of the main design aims for γ -ray arrays for high spin studies is high efficiency. It turns out that detecting γ -rays from exotic nuclei is one of the best ways of learning about their structure. However, exotic nuclei produced in reactions have small production cross-sections and so the detection methods used to study them must have high efficiency. It turns out, therefore, that γ -ray arrays are proving to be of great importance in the study of exotic nuclei [2, 6].

2 The Interaction of γ -Rays with Matter

The detection of γ -rays is of considerable importance, not just in studies of nuclear physics but in many applications as well. In general two main characteristics are of importance in the design of systems to detect them, namely energy resolution and efficiency. More recently it has become important to have detectors, which allow us to determine the initial direction of the interacting γ -rays. In application to nuclear physics this allows us to correct the measured energy according to how fast the emitting source is moving. In other applications it is important in terms of pinpointing the source of radiation.

To understand the main ideas behind the construction of γ -ray detection systems it is essential that the reader knows how γ -rays interact with matter. Many textbooks deal with this and we will concern ourselves only with the essentials here. One of the best sources of information is Evans [7]. This is an old textbook but it covers everything the student of nuclear physics needs to know in this area. In broad, general terms it tells us that if γ -rays are incident on an atom there are some twelve ways that it might interact.

To see this is relatively straightforward. There are four types of interaction between a photon and an atom. It can interact with (a) atomic electrons, (b) nucleons, (c) the electric field surrounding nuclei or electrons and (d) the mesonic field surrounding nucleons. We can further classify the interactions by the result of the interaction. Thus we can have (a) complete absorption of the photon, (b) elastic scattering and (c) inelastic scattering. Combining these we have twelve possible modes of interaction altogether.

Amongst these twelve only three are of real importance for our story, namely the photoelectric effect, Compton scattering and pair production.

The three interactions are illustrated in Fig. 5. In the photoelectric effect the incident γ -ray is totally absorbed by an atomic electron with the resulting photo-electron having energy $E_e = E_\gamma - E_{BE}$, where E_γ is the γ -ray energy and E_{BE} is the binding energy of the atomic shell where the interaction occurs. The process leaves behind an ionised atom. It should be noted that it is not a process that occurs with free electrons as simple consideration of conservation laws will confirm. It is, therefore, not a surprise that

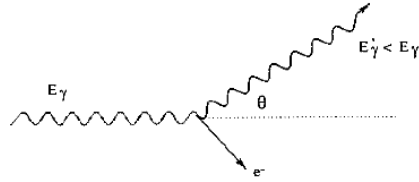
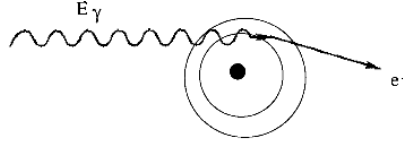
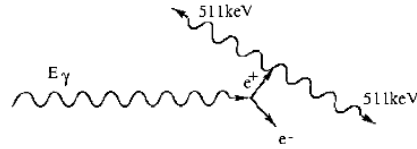
a) Photoelectric effect**b) Compton scattering****c) Pair production and annihilation**

Fig. 5. Illustration of the three main interactions between gamma rays and atoms (see text)

it mainly involves the K shell and hence the most tightly bound electrons. Compton scattering is different: it involves the inelastic scattering of γ -rays from electrons. After the collision we have both a γ -ray of reduced energy and an electron ejected from the atom. Such a process can occur with free electrons and so outer electrons are mainly involved in such interactions. In this process the initial energy and momentum is shared between the electron and scattered γ -ray. For our purposes it is useful to know that the energy of the scattered photon ($h\nu$) is given by

$$h\nu = h\nu_0 / [1 + h\nu_0/m_0c^2(1 - \cos\theta)] , \quad (1)$$

where $h\nu_0$ is the initial energy of the photon, m_0c^2 the rest mass of the electron and θ the angle of scattering. The angular distribution of the scattered γ -rays, i.e. the distribution in intensity as a function of scattering angle, is shown in Fig. 6 for scattering from Pb ($Z = 82$) for γ -rays of energies from 0–10 MeV. It should be noted that at non-zero energies the γ -rays are scattered predominantly in the forward direction, as one would expect from consideration of momentum conservation. The third process of importance here is pair production. In this process the incident γ -ray photon is completely absorbed, with the creation of a positron-electron pair whose total energy is equal to $h\nu$, the energy of the photon. Thus we can write

$$h\nu = (T_+ + m_0c^2) + (T_- + m_0c^2) , \quad (2)$$

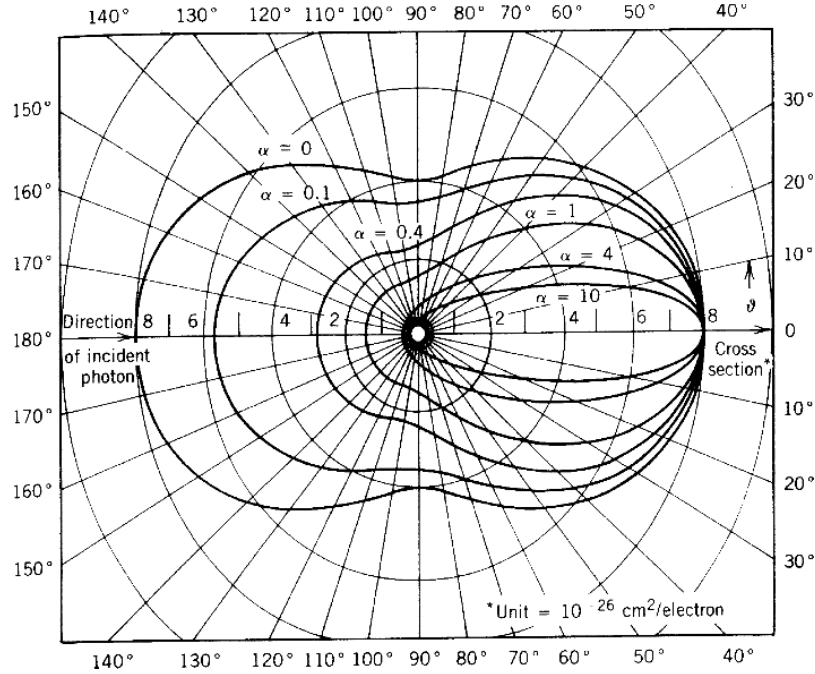


Fig. 6. Angular distributions of Compton scattered γ -rays of various energies, where the energy of the incident photon is given in terms of $\alpha = h\nu_0/m_0c^2$. Note that for high energy photons the scattered γ rays lie mainly in the forward direction. This figure is reproduced from [7]

where m_0c^2 is again the rest mass of the electron and T_+ and T_- represent the kinetic energies of the positron and electron, respectively. This process can only occur in the field generated by charged particles. In general this means that it occurs in the field associated with the nucleus but it can also be in the field of an electron. In any case, the presence of the other particle is necessary, otherwise the conservation laws are not satisfied.

Figure 7 shows the cross-sections for all three processes as a function of the incident photon energy for the specific case of germanium ($Z = 32$). It should be noted that it is plotted on a log-log scale. A number of features are worth noting.

The first lies below the energy range shown here. At such low energies there are several abrupt increases in the cross-section at specific energies; they are the result of the photon reaching sufficient energy to be absorbed in the photoelectric effect on more tightly bound electrons. These “jumps” are called absorption edges. They are important in terms of designing the shielding of both detectors and other things from low energy γ -rays. Naturally, absorption in, say, the K shell by the photo-electric effect is followed by the emission of a K X-ray or an Auger electron, a process in which the K shell vacancy is filled

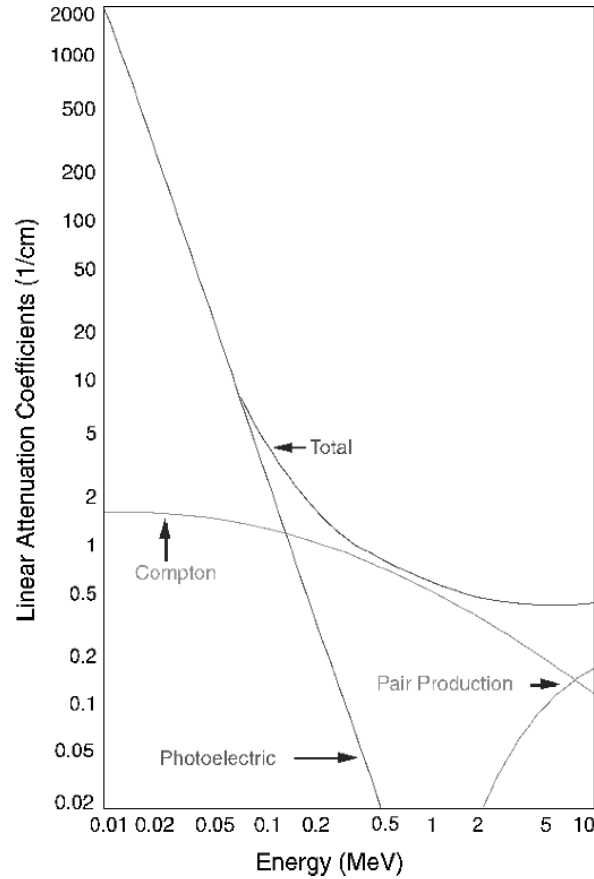


Fig. 7. Linear attenuation coefficient for Germanium ($Z = 32$) in inverse centimetres versus γ -ray energy in MeV on a logarithmic scale

by an electron from a higher shell. The possible atomic processes involved are discussed in detail by Burhop [8]. The photoelectric cross-section falls very rapidly with increasing energy. This is not surprising. For high-energy γ -rays we are moving towards the situation where the binding energy is almost negligible and we approximate the situation of absorption on a free electron, which is forbidden.

Secondly the Compton effect takes over as the dominant process in a region of particular importance for studies of nuclear structure from 0.5–5 MeV. The main things to note are that the scattered γ -rays in the Compton process are predominantly emitted in the forward direction. Such scattered γ -rays may escape from the scatterer if they do not interact for a second time. This is of great importance in designing γ -ray arrays as we shall see.

Thirdly once we exceed an energy of 1.022 MeV, equal to twice the rest mass of an electron ($2m_0c^2$), the γ -ray can create an electron-positron pair in interacting with the electromagnetic field generated by the nucleus. These particles will lose energy in slowing down in the medium and if the positron slows down sufficiently it will annihilate with an electron in the material to create a pair of γ -rays, each having an energy of 511 keV ($=m_0c^2$), emitted at 180° to one another (see Fig. 5c). The reader is referred to the textbooks [7] for details. Again, if the annihilation occurs close to the edge of the material then one or more of the two 511 keV γ -rays may escape from the material carrying its energy with it. This is also of importance for the design of γ -ray arrays.

3 The Development of Arrays

3.1 The Beginnings

The early experiments [1] on high spin states involved just a few NaI detectors as we saw in Sect. 1. One obvious line of development was simply to surround the target with an array of such scintillation detectors: this would clearly lead to a much higher efficiency of detection. If a large number of detectors was used one would also benefit from the granularity of the system in terms of recording coincidences between the large number of γ -rays emitted in a single fusion-evaporation reaction. Various such arrays were developed. With NaI detectors one finds amongst others the Spin Spectrometer [9] at Oak Ridge with 72 detectors and the Crystal Ball [10] at Heidelberg. BaF_2 was also used, with the primary example [11] being the Chateau de Crystal. In Fig. 8 we see spectra, taken with this spectrometer, of the Coulomb excitation of ^{76}Ge

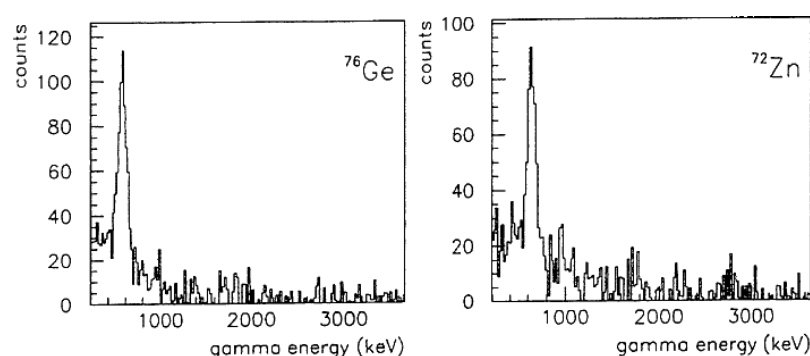


Fig. 8. Gamma-ray spectra from the Coulomb excitation of ^{72}Zn and ^{76}Ge recorded with the Chateau de Crystal [11], an array of BaF_2 scintillation detectors. The energy of the 2^+-0^+ transition is 653 keV in ^{72}Zn and 563(45) keV in ^{76}Ge (Courtesy of F. Asaiez)

and ^{72}Zn at GANIL. The former is a stable nucleus with a well known $B(E2; 2^+-0^+)$, and the experiment allows a comparison with the $B(E2; 2^+-0^+)$ of the radioactive species ^{72}Zn .

All of these spectrometers have the advantage of high efficiency. They operate well as calorimeters since they pick up a large fraction of the total γ -ray energy emitted in the γ -ray cascades in fusion-evaporation reactions. The BaF_2 has the further advantage in that it allows fast timing and it is possible to discriminate between neutrons and γ -rays using variations in the pulse shape for the two types of radiation (pulse shape discrimination). However both NaI(Tl) and BaF_2 suffer from poor energy resolution, typically 8% FWHM for 1 MeV γ -rays in the case of NaI and even worse for BaF_2 . Thus their use is limited for discrete line spectroscopy.

3.2 The Introduction of Ge Detectors

This field was transformed by the development [12] in the early 1960s of the reverse-biased Ge detector. These semiconductor diodes have been steadily refined over the years and have formed the basis of most spectrometers and arrays devoted to discrete-line, γ -ray spectroscopy ever since. A description of the basic device, either in the form of lithium (Li)-drifted devices or hyperpure-Ge (HpGe) detectors can be found in textbooks [13]. Initially the small Ge diodes available had a resolution of about 6 keV for the γ -rays emitted by ^{60}Co at 1332 keV energy. Now, in either form, Ge detectors have good energy resolution, typically ≤ 1.0 keV at 122 keV and ≈ 2 keV at 1332 keV. This is not comparable with the energy resolution one can achieve with diffraction spectrometers (see e.g., [14]) but they have efficiencies that are larger by many orders-of-magnitude. At first these detector efficiencies were also very small (only about 1% of that of a $3'' \times 3''$ NaI detector at 25 cm distance). But over the next decade or so there was a steady improvement in the sizes of Ge(Li) detectors. The situation was then transformed by the introduction of high-purity Ge (HpGe) detectors. Now it was no longer necessary to use the Li drift process and one could make much bigger detectors. The resulting increase in efficiency is of course particularly important in coincidence measurements. The timing characteristics of Ge detectors are modest but good enough for their use in the arrays described here.

At a very early stage Ge detectors were used in coincidence measurements with considerable success. The discovery of backbending, the alignment of a pair of particles in time-reversed orbits that we described earlier, was first observed by Johnson et al. [15] using two such detectors. A whole series of related measurements soon followed. In essence the better energy resolution of Ge detectors meant that we could see weak γ -rays from states with spins up to about $16-20 \hbar$ (see e.g., [16]).

All of these early experiments were limited, however, by a fundamental feature of γ -ray interactions, namely that Compton scattering and pair production can lead to the incomplete collection of the energy of a γ -ray that

interacts with the crystal. As a result the peak-to-background ratio in the spectrum is poor. One possible answer to this problem, the answer that was widely adopted, is described in the next section.

3.3 Compton Suppression

Figure 9, which shows γ -ray spectra from ^{60}Co decay, may serve to demonstrate the method of reducing the background from Compton scattering in such spectra. The two full-energy peaks, due to events where all the γ -ray energy is deposited in the crystal, are clearly seen in both spectra. However, in the spectrum taken with a single Ge crystal they sit on a background of incomplete collection events, where a Compton-scattered γ -ray escapes from the crystal. This does not look to be too serious a matter in this spectrum, but it becomes important in studies of reaction γ -rays when we are hunting for weak γ -rays sitting on the accumulated Compton backgrounds from many higher energy γ -rays. How can we improve this situation? The answer lies in *Compton suppression*. In the incomplete events the Compton-scattered γ -ray escapes from the crystal. If we detect it in another detector surrounding the Ge crystal then we can veto the event electronically. This will eliminate from the spectrum their contribution to the background. The idea is illustrated in Fig. 10, where we see the Ge crystal surrounded by a scintillator or scintillators. The scattered γ -ray entering the scintillation detector and interacting

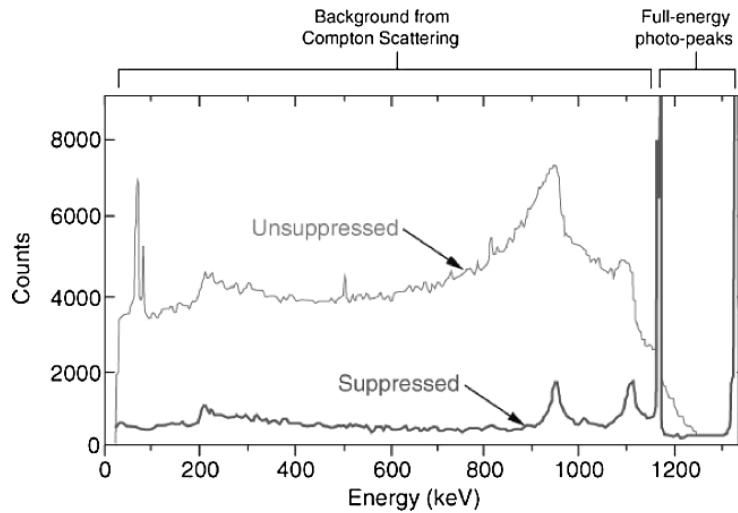


Fig. 9. Spectrum of γ -rays from ^{60}Co decay recorded with and without Compton suppression. One sees clearly the reduction in background due to the suppression of events in which γ -rays have scattered from the central Ge detector into the surrounding active shield. Note that the photopeaks are off-scale in order to show the suppression effect clearly (Courtesy of P. Sellin)

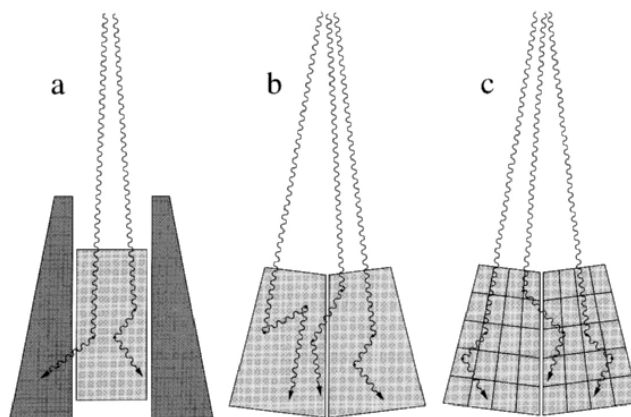


Fig. 10. Figure taken from Julin [2], showing (a) the basic idea of how the Compton-suppression shield surrounding an individual Ge detector works, (b) how signals from Compton scattered events detected in various parts of a composite detector can be added together (see Sect. 3.5) and (c) the idea of γ -ray tracking (see Sect. 4.2)

with it provides a signal which we can use to veto the associated Ge detector signal. The second spectrum in Fig. 9 shows the effect of applying this Compton suppression technique to measure the ^{60}Co spectrum. Roughly speaking, doing this increases the ratio of counts in the full-energy peak to all the counts in the spectrum from 20% to 50–60%, a considerable improvement.

This technique had been used with scintillation detectors. It was first applied with Ge detectors around 1980 in Copenhagen [4] when an array of five Ge detectors, each surrounded by a NaI scintillation detector shield, was used to study γ -rays from ^{84}Zr , $^{129,130}\text{Ce}$, $^{157,158}\text{Er}$ and ^{168}Hf produced in fusion-evaporation reactions. Quite quickly thereafter a more efficient scintillator was found to replace the NaI. This scintillating material, bismuth germanate (BGO), has both a higher density and a higher average Z value than NaI. Crudely speaking, this means that it is about three times more efficient per unit length for detecting the scattered γ -rays than NaI. It was now possible to assemble a collection of Compton-suppressed Ge detectors into an array with a much more compact geometry.

Both in North America and Europe such compact arrays of BGO-suppressed spectrometers were assembled and used to great effect. At Daresbury Laboratory in the U.K. a series of such arrays, called the TESSA arrays, were developed with different configurations. The TESSA arrays and their U.S. equivalent, HERA [17], were much more sensitive than the detection systems they superseded.

Applied to the study of fusion reactions induced by beams of heavy ions they allowed us to observe high-spin states up to about spin $50\hbar$.

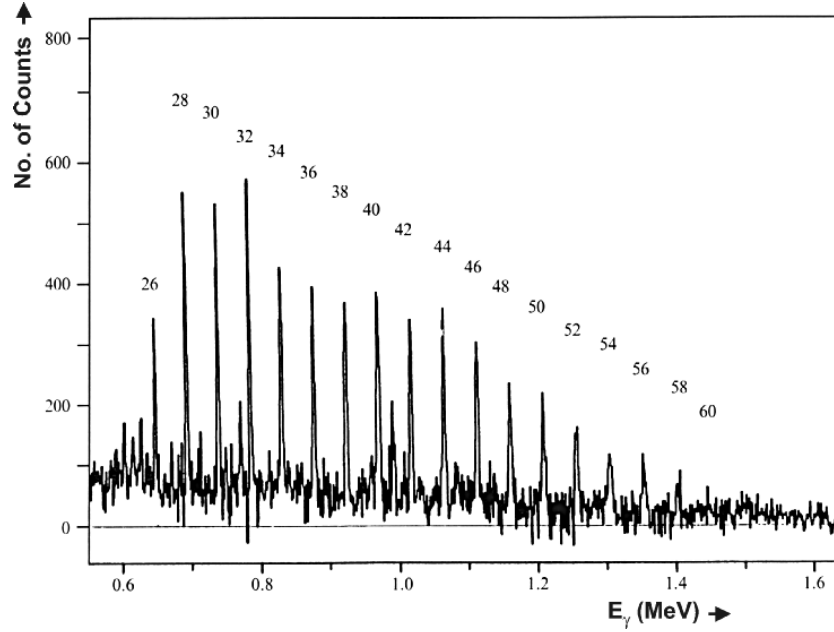


Fig. 11. Part of the γ -ray spectrum showing the decay of levels in the superdeformed band in ^{152}Dy recorded [18] with the TESSA2 spectrometer

This led to many new discoveries. The most striking of these was the observation of “superdeformed states” at high spins [18]. Figure 11 shows the spectrum of γ -rays de-exciting the levels in a superdeformed band in ^{152}Dy . In effect what we see in the spectrum is the decay of one of the best nuclear rotors that we know. The close spacing of the γ -rays reveals that they have a much larger moment-of-inertia than the corresponding ground state bands. They are thought to result from the rapid rotation of an elongated, axially deformed nucleus and are due to the rotation of this “superdeformed” shape based on a second minimum in the nuclear potential energy surface at large deformation. In the case shown they result from a nucleus, which is axially deformed but with a ratio of the major to minor axes which is 2:1. Following the initial observation of the γ -rays de-exciting this band, the deformation was confirmed [19] by measurements of the transition probabilities and hence quadrupole moments for these states. Soon, other examples were found in the rare-earth nuclei (see e.g., [20]) and later, elsewhere in the Nuclear Chart. Measurements on the light rare-earth nuclei revealed “superdeformed” bands but with an axis ratio of 3:2 rather than 2:1. In all cases we see the characteristic picket fence structure, evident in the spectrum of Fig. 11, due to the decay of the levels in the band. This is a key feature in their observation and is used to help pick them out from complex spectra, since in most cases they are populated in only

about 1–2% of the reactions leading to the nucleus of interest. It is clear that to observe and study them in detail requires spectrometers with both high efficiency and high resolving power, i.e. the ability to resolve closely-spaced transitions (see below). A review of all the cases found to date can be found in the “Table of Superdeformed Nuclear Bands and Fission Isomers” [21].

The picture we have painted so far is a simplified one, relying as it apparently does on larger Ge detectors and improved Compton-suppression alone. However dedicated ancillary detectors also led to improved sensitivity. BGO and other scintillators can be used, not just to veto Compton-scattered γ -rays, but also to act as a calorimeter. In the TESSA arrays an inner ball of BGO scintillation detectors was developed which picked up a large fraction of the cascade γ -rays. This allowed a measure of the total γ -ray sum energy and the multiplicity of the γ -rays. By gating on those events detected with high fold (the number of γ -rays detected from a single fusion-evaporation reaction) in the BGO inner ball one could enhance the selection of γ -rays in the Ge detectors associated with high spin states. In other words, one uses the inner ball to select events associated with long cascades of γ rays, which are, in general, the cascades involving high spin states. In Fig. 12 we see fold distributions on the left and the corresponding sum energies on the right for γ -rays from the ^{16}O plus ^{94}Zr reaction at 92 MeV. The upper, middle and lower spectra are, respectively, for a) the evaporation of five neutrons from the compound nucleus leading to ^{105}Cd , b) the evaporation of four neutrons leading to ^{104}Cd and c) the total projections. It is clear from these spectra that one can discriminate between the various reaction channels by selecting events with particular ranges of fold and sum energy.

It is not the purpose of the present article to explore what has been learned about superdeformed bands or, more generally, about high spin states. However [21] tells us that it is a relatively common feature of nuclei. Bohr and Mottelson [23] showed, using an axially-symmetric, anisotropic harmonic oscillator potential, that they could predict the occurrence of the superdeformed “magic numbers” quite well. However, to understand the structure associated with these bands well requires more sophisticated calculations. To first order it is possible to interpret the configurations associated with most superdeformed bands as being based on single quasi-particle excitations in an axially-symmetric, quadrupolar mean field. However this clearly neglects other degrees-of-freedom such as a stable, tri-axial, superdeformed minimum [24] or an octupole deformation.

3.4 The Key Characteristics of Arrays

We are now in a position to consider what characterises the ideal spectrometer for the efficient measurement of the energies and intensities of γ -rays in cascades de-exciting sequences of high spin states and determine the correlations between the γ -rays. It seems obvious that the best sensitivity will be obtained if we maximise the number of events we can study by maximising

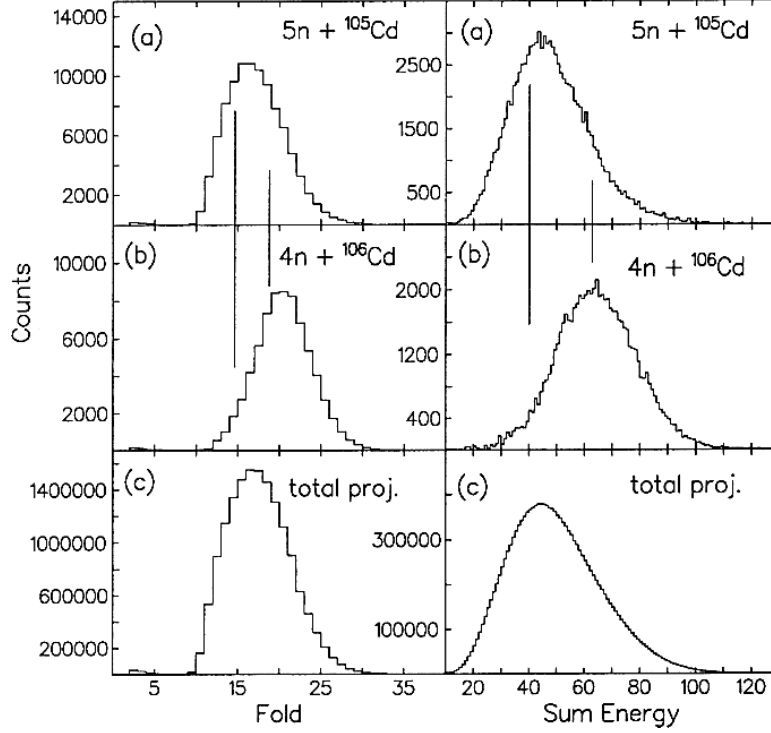


Fig. 12. Example of the γ -ray folds and sum energies recorded [22] in a study of the $^{16}\text{O} + ^{94}\text{Zr}$ reaction at 92 MeV. The *upper*, *middle* and *lower* spectra are for the evaporation of five neutrons from the compound nucleus leading to ^{105}Cd , the evaporation of four neutrons leading to ^{104}Cd and the total projections, respectively (Courtesy of P.H. Regan)

the total photopeak efficiency of the Ge detectors. This we do by using a large number of detectors while maintaining their good energy resolution. This will lead to high statistics in the recorded spectra. We also want to reduce the background of incomplete events, i.e. events in which some energy is lost; this we achieve using Compton suppression.

At the same time we want to be able to isolate sequences of γ -rays in very complex spectra such as those seen in Fig. 11. Accordingly we would like to have some measure of the resolving power (R) of a detection system. It turns out that there is no universal measure of R but we can determine it for the specific case of a spectrum consisting of γ -ray sequences de-exciting rotational bands.

Following Beausang and Simpson [25] let us assume that SE_γ is the average energy separation of the γ -rays in the cascade of interest, PT is the peak-to-total ratio for the detector system and ΔE_γ is the energy resolution of the detector, defined as the full width at half maximum for the γ -ray peaks

in the spectrum. Then R is defined by

$$R = 0.76(SE_\gamma/\Delta E_\gamma)PT . \quad (3)$$

Here we have taken 0.76 as the fraction of a Gaussian peak included in setting a coincidence gate on the peak. Each of the factors in (3) clearly influences the quality of the final spectrum but it does not really answer the crucial question: “What is the limit of observation of a peak in the spectrum?” The limit of observation (α_0) of a detection system is the minimum intensity of a γ -ray transition that can be detected in a spectrum. The peak-to-background ratio depends on R according to

$$(N_p/N_b)_F = \alpha_0(R)^F, \quad (4)$$

where $F - 1$ is the number of selection criteria placed on the spectrum. It should be noted that the higher the number of γ -ray coincidences demanded, the γ -ray fold, the lower the limit of observation calculated from (4).

We have already mentioned the use of a BGO inner ball to improve the selection and hence the sensitivity. There are many other possible ancillary devices we can use to help select the transitions of interest. For example an inner array of charged particle detectors, a set of neutron detectors, a recoil mass separator etc. can all be used to select γ -rays from a given reaction channel with suitable coincidence arrangements. The use of any of these devices improves the peak-to-background ratio and we take account of it in (4) by introducing a further factor R_0 , thus modifying (4) to give

$$(N_p/N_b)_F = \alpha_0 R_0(R)^F. \quad (5)$$

This gives us the peak-to-total ratio but the peak must also have sufficient counts to be seen. If we are dealing with a cascade where we have a multiplicity of emitted γ -rays of M_γ then the number of counts in the peak after gating on F other γ -rays and some other ancillary device is given by

$$N_p = (0.76PT)^F \alpha_0 R_0 (N_0)^F / M_\gamma (M_\gamma - 1) \dots \dots (M_\gamma - F + 1) . \quad (6)$$

It should be noted that this depends on $(N_0)^F$, the total number of F -fold events; that is the number of events obtained when we gate on F different γ -rays. As a result it also relies on the total photopeak efficiency of the array.

It should be emphasised that although this kind of expression is commonly used to characterise the sensitivity of γ -ray arrays it is only applicable to their use to detect γ -rays emitted in cascades de-exciting the levels in rotational bands. We shall return to an example of the use of this formula in comparing different arrays after we have discussed those arrays currently in use. Other experimental situations really demand some other measure of sensitivity.

In studying prompt γ -rays from nuclear reactions the energy resolution and hence the sensitivity are also affected by the Doppler effect. Since the

emitting nuclei are moving with velocity v at an angle of θ to the detector we measure not E_0 , the rest energy, but

$$E_\gamma = E_0(1 + v/c \cos \theta) . \quad (7)$$

At the same time the finite opening angle of the detector means that we measure a spread in energy of

$$\Delta E_\gamma = E_0 v/c \sin \theta \Delta \theta . \quad (8)$$

For a nucleus moving with a velocity of 3% of the velocity of light, emitting a 1 MeV γ -ray, detected at 90° in a detector with an opening angle $\Delta\theta$, the observed spread in energy is $0.03\Delta\theta$. Under the same conditions with a velocity of $0.5 v/c$ the spread in energy is $0.5\Delta\theta$. In this latter case the average measured energy E_γ is 1.5 and 0.5 MeV in detectors at 0° and 180° .

In practice this is not the only factor affecting the resolution. Other effects also lead to an increase in the final, observed energy resolution. In addition to the spread in energy (ΔE_{op}) due to the finite opening angle given in (8), there is also a further Doppler broadening due to a) the angular spread of the recoiling nuclei in the target (ΔE_{rec}) and b) the variation in velocity of the recoils depending on where they are produced in the target (ΔE_{vel}). The latter effect arises because the beam particles lose energy as they pass through the target. If we designate the intrinsic energy resolution of the system as ΔE_{IN} then the effective energy resolution (ΔE_{eff}) is obtained by combining these various factors in quadrature. Thus

$$\Delta E_{eff} = ((\Delta E_{op})^2 + (\Delta E_{rec})^2 + (\Delta E_{vel})^2 + (\Delta E_{IN})^2)^{1/2} . \quad (9)$$

Figure 13, taken from Beausang and Simpson [25], shows the case of the reaction $^{30}\text{Si} + ^{124}\text{Sn}$ at a beam energy of 158 MeV. Here $v/c = 0.021$. Part of the spectrum, recorded with the EURO GAM I and EURO GAM II arrays, centred on a peak at 1450 keV, is shown. Although $\Delta E_{IN} = 2.4$ keV the resulting resolution is 6.2 keV for the former and 4.8 keV for the latter array. The details are given in [25].

This is not the end of the story as far as spectrum quality is concerned. Beausang and Simpson consider a variety of other effects that reduce R by reducing the peak-to-total ratio and/or reducing the total photopeak efficiency. These effects result mainly from the high multiplicity of the events and the interaction of neutrons with the detectors. These effects reduce both the resolving power and the photopeak efficiency. We will not expand on this topic here and refer the reader to refs [26, 27]. Suffice to say that in many reactions, neutrons are produced copiously and it is obvious that γ -rays and neutrons will scatter from collimators and other passive shielding into the BGO suppression shields.

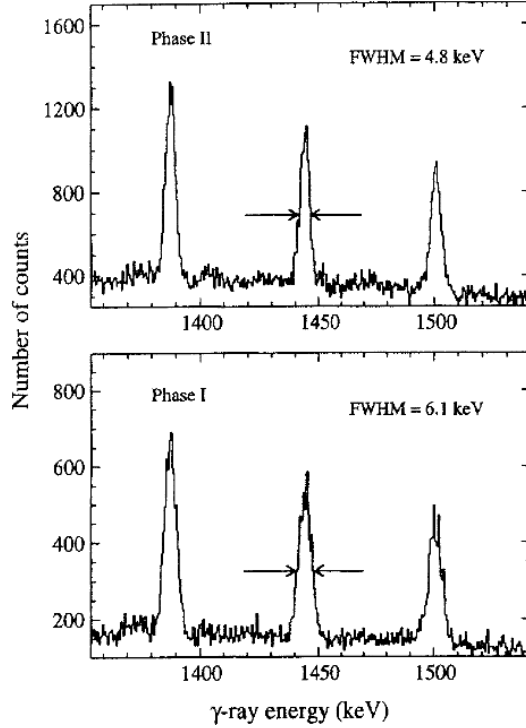


Fig. 13. Spectra showing part of the decay of the yrast superdeformed band in ^{149}Gd following the reaction $^{30}\text{Si} + ^{124}\text{Sn}$ with beam energy 158 MeV and relative recoil velocity $v/c = 2.1\%$, in EUROGAM phase I and II. The resolution of the 1450 keV transition reduces from 6.1 keV in EUROGAM I to 4.8 keV in EUROGAM II [25]

3.5 Modern γ -ray Spectrometers

In the last section we saw that the features that determine the performance of an array are the energy resolution, the peak-to-total ratio and the full-energy peak efficiency. The initial development of Ge detector arrays led to an improvement in all three of these quantities, hence their success. In the late 1980s and early 1990s a series of spectrometers were developed which took these ideas further and made use of several other technical developments in Ge detector manufacture. At the same time the new arrays concentrated on filling as much of the 4π solid angle as possible. This was greatly aided by the progressive manufacture of bigger individual detectors. Later in the 1990s it also became possible to make composite detectors, notably the clover [28] and cluster [29] detectors. This led to another leap forward in sensitivity and efficiency.

In what follows we shall use two different definitions of the photopeak efficiency, namely the *absolute efficiency* for a source positioned at the centre of an array and the *relative efficiency* defined with reference to a $3'' \times 3''$ NaI detector at a distance of 25 cm from the source. Both quantities as well as the energy resolution will be given for the 1332 keV γ -ray of ^{60}Co unless otherwise specified.

The new arrays GaSP (Italy) [30], EUROGAM (UK-France) [31] and GAMMASPHERE (U.S.A.) [32] were focussed on filling the full solid angle with the larger detectors that had become available. Typically the individual detectors used had a relative efficiency of 75% and also an improved peak-to-total ratio. These arrays retained the Compton-suppression principle and used shared suppression. In this extension of the Compton suppression technique a scattered γ -ray detected in the shield of an adjacent Ge detector can also be used to veto the incomplete energy event. The most powerful of these is GAMMASPHERE with 110 detectors covering the full solid angle. The absolute efficiency is 9.4%, a big advance on the previous generation of spectrometers. Two notable features of this highly successful spectrometer are its symmetry about the beam axis, a feature it shares with GaSP, which makes certain types of analysis easier, and the fact that about 70 of the 110 detectors are segmented. Electrical segmentation of the outer electrode separates it into D-shaped halves in these detectors. For studies of reaction γ -rays, where the effective energy resolution is determined by Doppler broadening due to the motion of the source, this segmentation leads to better resolution by reducing the finite opening angle of the detector (see below). The GaSP spectrometer installed at Legnaro is similar in design to GAMMASPHERE. It has 40 detectors, symmetrically placed with respect to the beam line with an absolute efficiency of $\approx 3\%$ when there is an inner ball of BGO detectors in place. A second geometry, in which the inner ball is removed, has an absolute efficiency of 6%. In GaSP none of the detectors are segmented. Both GaSP and GAMMASPHERE can operate with a range of ancillary detectors including neutron detectors, charged particle arrays, plungers for measuring lifetimes by the recoil-distance technique etc. As we learned from (5) this improves their sensitivity even further.

In parallel with these developments a number of arrays were conceived based on composite detectors. Two main types of composite detector were created, the clover and the cluster. In the arrays based on these detectors the idea of Compton suppression with BGO scintillators is also retained. Again the main design aim, as we would anticipate from (4), is to maximise the full-energy efficiency and the peak-to-total ratio. The designation, clover detector, reveals its form. It consists of four large, co-axial Ge detectors with their sides cut so that they can fit together like a four-leaf clover. This form of detector was developed by the CRN-Strasbourg laboratory and the company EURYSIS-MESURES [28]. The EUROGAM II array was composed of this form of detector with the crystals in the clover being 50 mm in diameter and

70 mm long, prior to shaping. These detectors have a number of advantages. They have a large relative efficiency of 140%, but since the signals are, at worst, identified as being from one of the four detectors the effective resolution is that of one component detector. Later developments have meant that one can locate the first interaction rather better than that and so the Doppler broadening can be reduced even further. The clover detector, like all composite detectors, has an efficiency given not just by the sum of the efficiencies of the individual detectors but a greater efficiency, since one can add the energy signals from events where there is scattering between the crystals (see Fig. 10b). For the EUROGAM clovers this “add-back” factor adds 50% to the efficiency. These detectors are also sensitive to the linear polarisation of the γ -rays since one can use the scattering between the various elements of the clover. The sensitivity has been shown [33] to be comparable to that of other polarimeters. Further, these detectors are less sensitive to neutron damage because the individual crystals are smaller in size. They also have relatively good timing properties. The details of the various versions of the EUROGAM array are not relevant here. In EUROGAM II there were 126 detectors yielding an absolute efficiency of 8.1% and an observation limit of $\approx 1 \times 10^{-4}$. The latter quantity gives the intensity of the weakest γ -ray observed, normalised to the strongest one occurring in the de-excitation scheme of interest.

The second significant type of composite detector that has been developed is the cluster detector [29]. These detectors were developed by a collaboration involving the University of Köln, KFA-Jülich and EURYSIS-MESURES. They consist of seven closely-packed, tapered, hexagonal crystals housed in a common cryostat. The individual crystals are 70 mm in diameter and 78 mm long before they are shaped. Each crystal has a relative efficiency of 60%. In order to achieve the close packing of the detectors each crystal is permanently, hermetically sealed in its own aluminium can of wall thickness 0.7 mm with the crystal sitting 0.5 mm from the inner surface of the can. The can provides electrical shielding. The crystals sit at a distance of 2.7 mm from each other so that the overall packing is optimised but the system is flexible since each can may be removed separately for repair. Maintenance of these detectors is relatively easy. They can be handled safely and are easily annealed to repair neutron damage. The energy resolution is better than 2 keV. As with the clovers they benefit from “add-back”, the Doppler spreading is determined by the opening angle of the individual detectors, one can use them to determine linear polarisations and again the timing characteristics are relatively good.

Figure 14 shows the spectrum of γ -rays from the β decay of the ^{150}Ho 2^- state to ^{150}Dy , recorded with the so-called “Cluster Cube”, a collection of six cluster detectors in close geometry with an absolute efficiency of 13%. The observation of weak β -delayed γ -rays, in particular those of high energy, is important when trying to deduce reliable information on the entire β -intensity distribution.

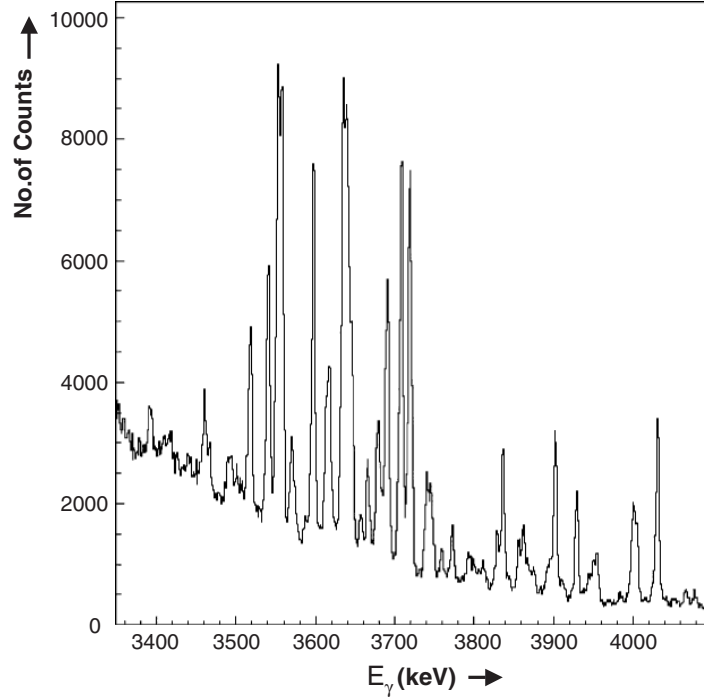


Fig. 14. Portion of the spectrum of γ -rays from the radioactive decay of the ^{150}Ho 2^- isomer recorded with the Cluster Cube, an array of cluster detectors. The spectrum is recorded in coincidence with the 803 keV, 2^+-0^+ transition in ^{150}Dy . (Courtesy of B. Rubio)

The encapsulation technology developed for these detectors has many advantages which are now being exploited for other purposes such as the space missions Mars Odyssey [34] and INTEGRAL [35], where they are used for studying γ -ray emissions from astronomical objects.

Cluster detectors form a major part of Europe's equivalent of GAMMA-SPHERE, the EUROBALL array. This array consists of 30 large volume Ge detectors, 26 clovers and 15 clusters. Each of these detectors has its own suppression shield. The mono-crystals were placed in the forward direction, to allow easy exchange with neutron detectors when they were needed, and the clusters in the backward direction with the clovers arranged in two rings at 90° . In total this is equivalent to an array of 239 individual detectors. A photograph of the array is shown in Fig. 15. In its final manifestation at Strasbourg the array had an inner ball and a number of other important ancillary detectors, to detect neutrons [36] and charged particles [37]. This particular inner ball is a triumph of engineering design since it is highly complex. This is because of the differing shapes of the Ge detectors in the three sections of the array. EUROBALL is a very powerful spectrometer; it has an

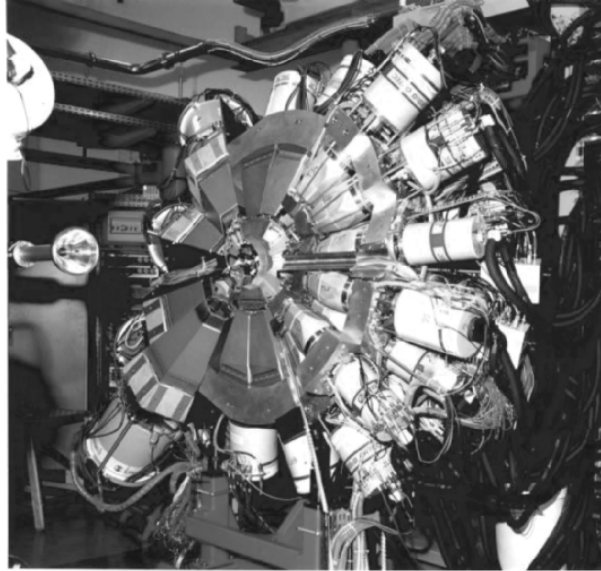


Fig. 15. Photograph of one half of the EUROBALL array pulled back from the target chamber(see Sect. 3.5); cf. Plate 10 in the Colour Supplement

absolute efficiency of greater than 9% and excellent resolution. Together with GAMMASPHERE it represents the culmination of the line of development involving Compton suppression. The designs of these two spectrometers differ significantly and they complement each other in terms of their strengths. To go further, however, we need a new approach. Before we look at the most recent developments we need to consider the significance of these arrays for the study of exotic nuclei. As we have indicated earlier this was not the main motivation for building them but it turns out that they have been vital to progress in this area too.

4 Gamma Rays from Exotic Nuclei

Much of our knowledge of the properties of exotic nuclei comes from γ -ray spectroscopy. It is garnered from studies of a variety of reactions and also from β decay. The latter is often the first source of information about the most exotic nuclei.

As we saw earlier, this was not the main reason to develop γ -ray arrays but their main characteristics, namely good resolution, high photopeak efficiency and good peak-to-total ratio turn out to be ideal for many studies of nuclei far from stability. As we move away from stability the cross-sections for producing individual species in almost every reaction process decrease

rapidly. For example the CERN-ISOLDE facility produces short-lived nuclear species far from stability via spallation or fission. Far from stability the intensities of the separated 60 keV beams of radioactive nuclei it produces fall by approximately one order-of-magnitude for each neutron we remove. This is not a universal rule since greater losses occur as we go to very short-lived nuclei, where the half life is so short that many of the nuclei produced initially do not survive the diffusion and effusion processes in the target. It does, however, give some idea of the difficulties of studying such nuclei.

We have already seen in Fig. 14 how the power of the modern arrays can be used to good effect in β decay studies. In this spectrum the effects of good resolution, high efficiency and good peak-to-total ratio are all clearly manifest. They are also a powerful tool in studying prompt γ -rays from reactions. In this application channel selection is vital and the γ -ray detectors have to be combined with ancillary detectors that allow us to select a particular channel cleanly with suitable coincidence criteria. This can be done in a variety of ways. Here we will briefly discuss just two.

One powerful method of channel selection is the use of a recoil mass separator to identify the recoiling nuclei produced in a fusion-evaporation reaction. Figure 16 shows schematically an early example of this type of experimental set-up. The γ -ray array is set up around the target. In this technique it is a thin target and the nuclei recoil out of the target in the forward direction. The details of the recoil mass separator vary but they have common features. There is a set of lenses to collect the maximum number

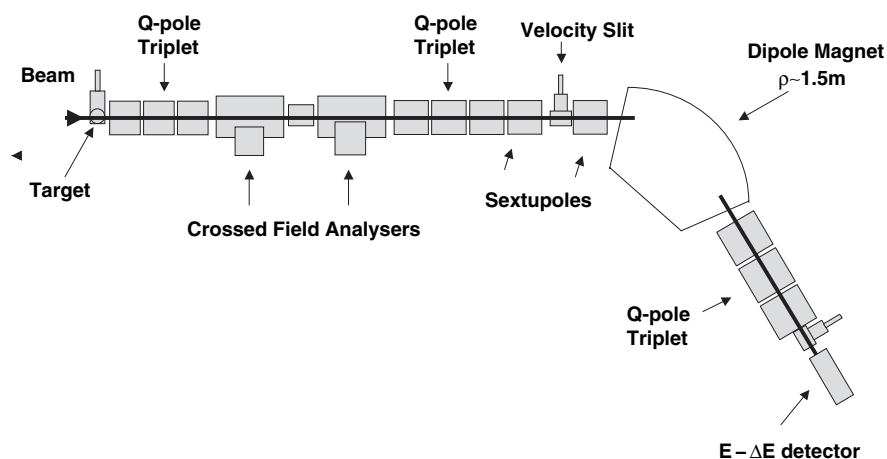


Fig. 16. Schematic picture showing the arrangement of a typical recoil mass separator. It consists of a series of elements. It is designed to collect ions emitted over as wide a solid angle as possible, followed by crossed electric and magnetic fields to separate beam and recoil particles and a magnetic sector to separate the recoils by mass

of ions leaving the target. A system of crossed electric and magnetic fields is then used to separate the recoils from the primary beam. The latter are then separated by A/q in the field of a dipole magnet and finally detected in an ion chamber. By measuring the energy loss and total energy of the ion in the last detector and knowing A/q we can identify the recoil by A and Z . We record the detection of the γ -rays in coincidence with the identified recoils. Some other elements are often included to deal with optical aberrations.

Good separation in Z is only possible, however, if the ion is moving fast enough. Even if this is achieved, the Z separation is only partial. We can see all of these features in a series of experiments [38] carried out at Daresbury laboratory with detectors from the TESSA arrays and the Daresbury recoil mass separator. Figure 17 shows the final spectrum of γ -rays from ^{80}Zr extracted from a study of the $^{24}\text{Mg} + ^{58}\text{Ni}$ reaction at 190 MeV. Here the Z separation is incomplete as we can see from the inset to the figure, which shows a plot of ΔE for the recoils vs. γ -ray intensity for several γ -rays seen in the reaction in coincidence with the recoils. The final spectrum is only obtained in this case by careful subtraction of spectra gated by different ranges of ΔE . Using this technique it was possible to observe prompt γ -rays from all the $N = Z$, even-even nuclei from ^{64}Ge to ^{84}Mo [38].

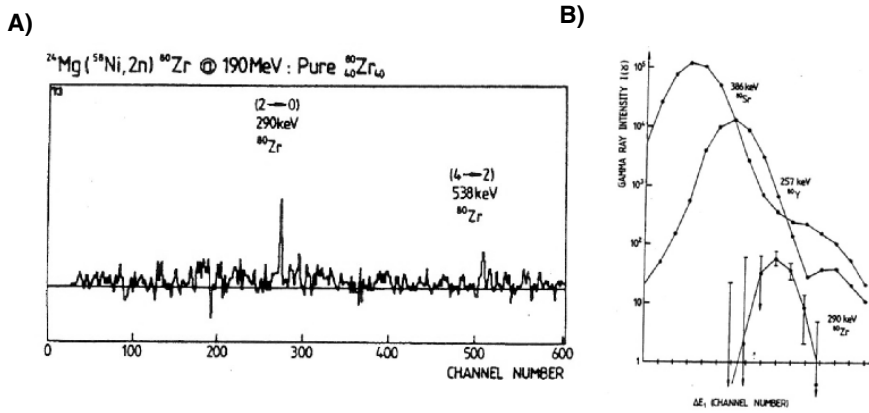


Fig. 17. (a) Part of the γ -ray spectrum from ^{80}Zr [38] recorded in coincidence with recoiling ions identified by mass and Z with the Daresbury recoil mass separator; (b) intensities of γ -rays from mass-80 nuclei plotted as a function of the differential energy loss in an ionisation chamber at the end of the Daresbury Recoil Mass Separator

An alternative approach is to select the reaction channel of interest by some other means. One possibility is to detect the charged particles and neutrons from the reaction in coincidence with γ -rays. From the multiplicities of the particles we can define the reaction channel. Excellent examples of this

approach can be found in studies of the same $N = Z$ nuclei with the GaSP spectrometer at Legnaro. This array has an inner, charged-particle detector based on Si detectors. By adding a set of neutron detectors it was possible to confirm the earlier results for ^{84}Mo [38] and extend the level scheme and also see the γ -rays from ^{88}Ru [39].

More recently GAMMASPHERE has been used together with the Fragment Mass Analyser (FMA) at Argonne to look at some of these nuclei. The increased solid angle of the FMA compared with the Daresbury separator and the ten-fold increase in efficiency provided by GAMMASPHERE is reflected in the spectrum of Fig. 18. Here we see the ^{76}Sr prompt spectrum taken with GAMMASPHERE in coincidence with recoils identified by the FMA. Comparing this with Fig. 17 we see the greatly improved performance of the modern γ -ray arrays.

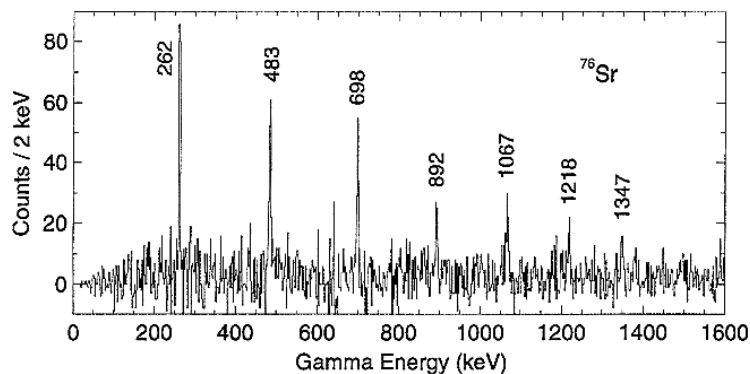


Fig. 18. Spectrum of γ -rays from ^{76}Sr recorded with GAMMASPHERE [31] in coincidence with recoiling ions identified with the Fragment Mass Analyser at Argonne National Laboratory

4.1 Radioactive Ion Beams

There are many examples of the type discussed in the last section. The article by Julin in volume 1 of this series [2] shows some excellent examples of transactinide nuclei studied in this way. In particular it shows examples of how the recoils can be cleanly identified by first separating them from beam particles in a gas-filled separator and then “tagging” them electronically with the particles subsequently emitted in radioactive decay. Much more can be done using these techniques but the information is limited because fusion-evaporation reactions select the so-called yrast states, the states of lowest energy for a given spin. We can also study β decay but again the selection rules in β decay limit the levels we see to a small range of spin. In brief, to study these nuclei properly requires *radioactive ion beams*.

As one can read elsewhere [40, 41] we have found two main ways of creating such beams. The ISOL method produces beams of similar quality to the stable-isotope beams already available. ISOL beams are readily accelerated to the Coulomb barrier and above. However not all nuclear species can be produced in this way since the techniques used depend on the chemistry of the element involved. The in-flight method [40] of producing radioactive ions has no chemical sensitivity but the beam quality is poor. One generally gets a beam that is a cocktail of neighbouring species and the ions are identified by A and Z on an individual basis. The methods are, in essence, complementary and one uses that best suited to answer a particular question. Here we will look briefly at examples of the kind of experiment opened up by the availability of radioactive ion beams produced by both methods. Our concern in this lecture is the study of γ -rays and both examples involve γ -ray detection.

In the first example [42] beams of ions were produced at GANIL by fragmentation of a ^{92}Mo beam with an energy of 60 MeV/u on a natural Ni target. The ions were analysed by the LISE3 spectrometer and implanted and brought to rest in a Si detector telescope. The ions were identified in A and Z on a one-by-one basis, from measurements of their time-of-flight and energy loss ΔE , and recorded in coincidence with γ -rays detected in Ge detectors set in close geometry around the telescope. In the left part of Fig. 19 we see the time-of-flight versus energy loss for the reaction products

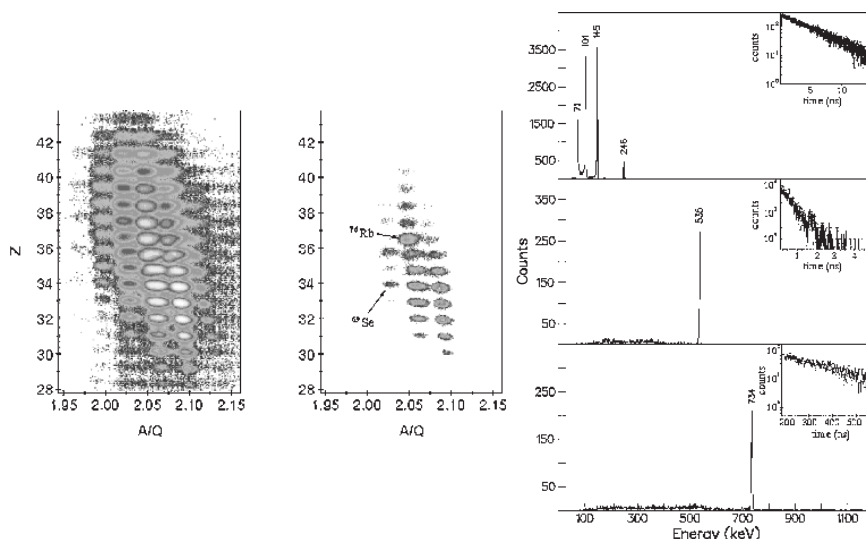


Fig. 19. Results of studies at GANIL of delayed γ -rays from the ions produced in the fragmentation of ^{92}Mo ions at 60 MeV/u on a natural Ni target [42]. On the left we see two-dimensional spectra with and without coincidences with delayed γ -rays. On the right we see examples of delayed γ -ray spectra from selected ions; cf. Plate 11 in the Colour Supplement

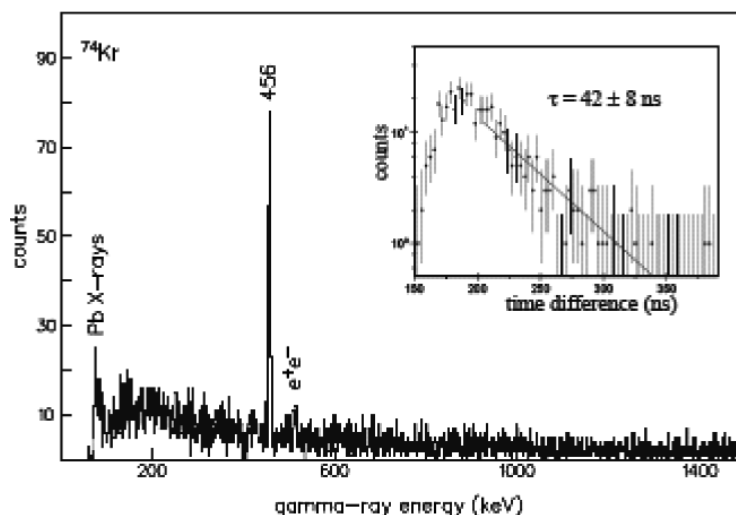


Fig. 20. Delayed γ -rays from ^{74}Kr [42]

emerging from LISE3 in this experiment. In the middle part we see the same plot but now we also demand that a time-delayed γ -ray is detected in one of the Compton-suppressed Ge detectors placed around the stopping point. In the right part of the figure we see the delayed γ -ray spectra in coincidence with three of the groups of identified ions. These plots reveal the presence of γ -decaying isomers associated with some of the fragments. The well-known isomer in ^{76}Rb clearly enhances the peak from this nucleus in the central plot. The time distributions associated with these γ -rays are shown in the insets to the spectra on the right. Figure 20 shows the γ -ray spectrum associated with the ^{74}Kr ions. Here we see the 456 keV γ -ray de-exciting the first 2^+ state. However, this state has a lifetime of only 25 ps [43] and it appears to have a lifetime of 42(8) ns in the figure. With the normal lifetime it would not survive passage through the LISE3 spectrometer. The only explanation is that the isomeric level is not the 2^+ state but a level with spin and parity 0^+ , located less than 60 keV above it, which decays mainly by an E0 transition to the ground state. Since the ions are fully stripped during passage through the spectrometer and E0 transitions only decay by internal conversion, the level cannot decay until it stops in the Si detector and the ion is neutralised. This experiment was followed up by Becker et al. [44] who observed the isomeric decay and confirmed this explanation. The tentative interpretation of the observations is that there is considerable prolate-oblate mixing of these states.

How can one do better? The answer is to use an ISOL-produced beam of ^{74}Kr and Coulomb excite it. Doing this will allow a measurement of all the relevant matrix elements. This has been done with a beam of ^{74}Kr ions

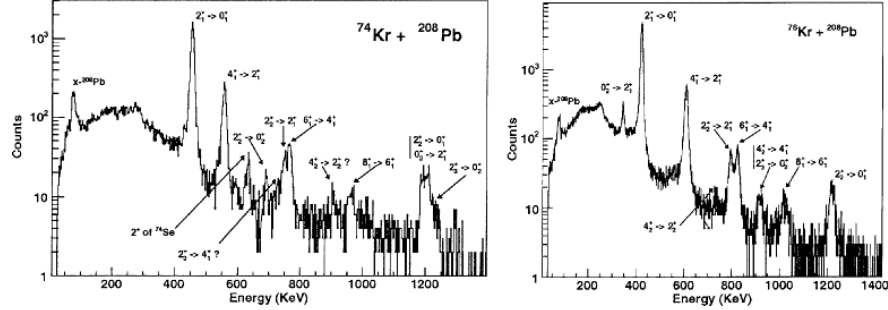


Fig. 21. Gamma-ray spectrum [45] observed after Coulomb excitation of ^{74}Kr (*left*) and ^{76}Kr (*right*). The spectra were recorded in coincidence with scattered ions and represent the integrated yields over the entire range of scattering angles covered by the Si detector used to detect the scattered ions. The radioactive ions were delivered by the SPIRAL facility. (Courtesy of E. Bouchez and W. Korten)

from the SPIRAL facility [45] at the GANIL. The spectrum of γ -rays from Coulomb excitation of a beam of ^{74}Kr ions from SPIRAL incident on a Pb target at 4.7 MeV/u is shown in Fig. 21. This provides the matrix elements needed and confirms that the 0^+ ground state of this nucleus is a 50:50 oblate-prolate mixture. More recently, a β decay study [46] using Total Absorption spectroscopy has confirmed this result.

These examples show only the tip of the iceberg: Radioactive ion beams will revolutionise our attempts to study nuclei far from stability. They bring new demands on our γ -ray detection systems. In the case of fragmentation beams the most obvious is that, if we are to study prompt γ -rays from reactions induced by them, the nuclear source will be moving at very high velocity and the resulting Doppler effects will be large. Secondly the beams will be very weak compared with stable-isotope beams and this puts a premium on high efficiency. We shall return to this in the next section.

5 Present Arrays and γ -Ray Tracking

5.1 MINIBALL and EXOGAM

Various arrays have been devised to meet the problems outlined at the end of the last section. Two prominent examples are the MINIBALL [47] and EXOGAM [48] arrays deployed at REX-ISOLDE and SPIRAL respectively. The ideas and techniques they have pioneered will be used in the next generation of arrays based on the idea of γ -ray tracking. Both MINIBALL and EXOGAM make use, for the first time, of position-sensitive Ge detectors. In both cases the detectors are segmented longitudinally. This allows the localisation of the first interaction of the γ -ray in two dimensions for Doppler

correction. However they are not capable of full γ -ray tracking (see below) because they lack segmentation in depth.

Electrically segmenting the Ge crystals is one route to improving the granularity of an array. Now instead of the opening angle of the whole crystal dictating the Doppler broadening it is the opening angle of a single segment that matters. A segmented detector is a standard n-type Ge detector with the main, high-resolution signal derived from a central contact and the position information obtained from signals from isolated outer contacts. This development was incorporated into the central EUROBALL clovers to improve their performance. Here each leaf of the clover was segmented into four parts so that overall, each clover detector had 16 active elements.

Larger clovers segmented in the same way have been used in the construction of the EXOGAM array. This array can be used in various geometries. With only four detectors set at 68 mm it has an absolute efficiency of 10%, a peak-to-total ratio of 60% and is ideal for decay studies. With the full 16 detectors set at 114 mm it has an absolute efficiency of 20% and consists, in essence, of 256 elements giving it high granularity and good resolution and at the same time the Doppler effects are limited. The Coulomb excitation spectra of Fig. 21 were recorded with 7 and 11 EXOGAM detectors, respectively, in coincidence with recoils detected in a Si detector.

The MINIBALL array is an excellent example of a modern array designed specifically for studies of reactions of low γ -ray multiplicity with radioactive ion beams. It consists of 40, six-fold segmented, encapsulated Ge detectors made using the encapsulation method developed for the EUROBALL clusters. The 40 detectors are clustered in eight cryostats with three detectors each and four with four detectors each. This arrangement gives an optimum coverage of Ge in a 4π detector arrangement. It uses the encapsulation technology developed for the EUROBALL cluster detectors. This encapsulation has proved itself thoroughly in terms of reliability with a very low failure rate since 1993 with all of the detectors having been annealed several times to repair the effects of neutron damage. Segmented detectors are particularly sensitive to contamination of the passivated, intrinsic surfaces and the encapsulation helps to preserve their properties. The encapsulation also offers another bonus. It separates the detector vacuum from the vacuum containing the cold parts of the pre-amplifiers. It turns out that the positioning of the cold components, shielding of components and wiring is crucial to the performance; one has to eliminate cross-talk and oscillation of the pre-amplifiers. Encapsulation means that one can readily optimise these arrangements without endangering the Ge detectors themselves. The absolute efficiency of MINIBALL is 8% for eight triple clusters at 12 cm from the target and 14% for the final configuration with 40 detectors.

The position sensitivity of the MINIBALL detectors is based on their segmentation and the analysis of the pulse shape of the core signal for the radial position and the mirror charges induced in the segments for the

azimuthal position. We will not enter into detail here but, in essence this involved the development of pre-amplifiers with a large bandwidth to transfer all the information associated with the signal and the digitising of the pulses for processing.

MINIBALL is in use at REX-ISOLDE and Fig. 22 shows the spectrum of γ rays from the Coulomb excitation of ^{31}Mg [49]. The MINIBALL and EXOGAM arrays represent the state-of-the-art in γ -ray detection for use with radioactive beams based on the ISOL technique.

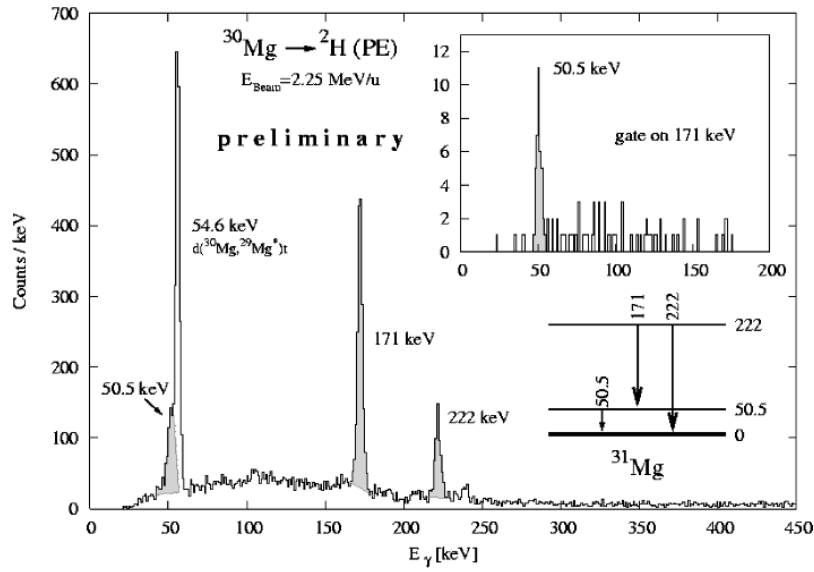


Fig. 22. Coulomb excitation of ^{31}Mg in inverse kinematics with beams from REX-ISOLDE [49]

5.2 Tracking Arrays

The ultimate γ -ray array could be seen as a ball of Ge detectors covering 4π in solid angle. This is not possible at present. If we are to improve arrays further and approach this ideal, we need a new idea, a new technique.

The clover and cluster detectors improve the resolving power and efficiency of arrays but they still suffer from some of the limitations of conventional, individual detectors. Compton suppression improves the quality of the spectrum but the suppressor uses up some of the useful solid angle. The arrays composed of these composite detectors have improved efficiency but their sensitivity is limited by the various effects discussed by Beausang and Simpson [25]. For example if two γ -rays interact simultaneously in the same

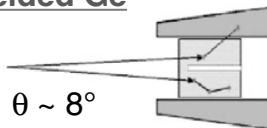
segment of the detector they cannot be distinguished and we see the summed effect. This means that the peak-to-total ratio is not as good as it should be.

In both North America and Europe the way forward has been identified as γ -ray energy tracking [50] in an array of electrically segmented detectors. In an array based on tracking detectors the complete solid angle will be filled with Ge detectors and there will be no suppression of scattered γ -rays. Instead the events will be recovered by reconstructing the scattering path. Tracking has been widely used in particle detection for a long time. It is readily possible to follow the trajectory of a charged particle or the time sequence of its position since it continuously loses energy and ionises the medium it is moving through. As we saw in Sect. II this is not possible for γ -rays. For example if the γ -ray Compton scatters we get the simultaneous appearance of signals in several places at once. The idea of γ -ray tracking is illustrated in Fig. 23. Here we use the energy deposited at each point of interaction and the positions of interaction in three dimensions. The energy deposited is given by the difference in energy of the incident and scattered γ -rays (see (1)). This gives the angle of scattering and we also get this from the positions of the interaction points. These two numbers must be consistent. Thus γ -ray tracking relies on the Compton scattering formula to allow us to develop an algorithm to reconstruct the event and determine the position of its first interaction and its total energy. This will enable us to deal with the results of the Doppler effect and provide high efficiency. It will require highly segmented Ge detectors and digital electronics. The detectors must enable the determination of the interaction points with good spatial resolution.

Idea of γ -ray tracking

Compton Shielded Ge

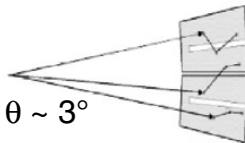
$\epsilon_{ph} \sim 10\%$
 $N_{det} \sim 100$
 $\Omega \sim 40\%$



large opening angle
means poor energy
resolution at high
recoil velocity

Ge Sphere

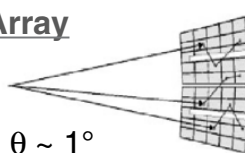
$\epsilon_{ph} \sim 50\%$
 $N_{det} \sim 1000$



too many detectors
are needed to avoid
summing effects

Ge Tracking Array

$\epsilon_{ph} \sim 50\%$
 $N_{det} \sim 100$
 $\Omega \sim 80\%$



Combination of:

- segmented detectors
- digital electronics
- pulse processing
- tracking the γ -rays

Fig. 23. Illustration of the idea of γ -ray tracking; cf. Plate 12 in the Colour Supplement

Given the advantages of such a system it is not surprising that several efforts to create such systems are underway. In the U.S.A there is the GRETA project [51], in Italy we have MARS [52] and there is the main European effort AGATA [53]. The last two are closely linked. All three have much in common.

The technique of γ -ray tracking has two main steps. As indicated above, the first step involves the identification of all interaction positions and the energy deposited there. This information is required as input for the reconstruction of the scattering path using the tracking algorithm. To do this successfully the points of interaction have to be determined to within a few mm.

It turns out that the interaction points can be determined by analysing the pulse shapes of the signals induced on the electrodes during charge collection. As in any detector the free charge carriers created by the electrons produced in the interactions drift towards the electrodes under the influence of the electric field. The shape of the signals is determined by the form of the electric potential. Once the potential is known one can deduce the initial interaction point from the pulse shape. This can be used to make the Doppler correction, improve the energy resolution and enhance the peak-to-total ratio [54].

One can do much better than this, however, by analysing the transient signals induced in adjacent segments. It will not surprise the student of physics that the charge drifting towards an electrode induces a signal in other neighbouring electrodes and that the nature of this signal depends on the position of the interaction relative to the boundaries of the various electrodes. Figure 24 shows an example [55]. Here we have a detector with six-fold segmentation along the central axis as well as on the front face. The front-face segments are labelled as shown on the right of the figure. The interaction positions of two hypothetical events are marked. On the left we see the digitised signals induced on the various electrodes in a 6×6 segmented prototype detector from EURYSIS MESURES. Each signal has 256 samples, one every

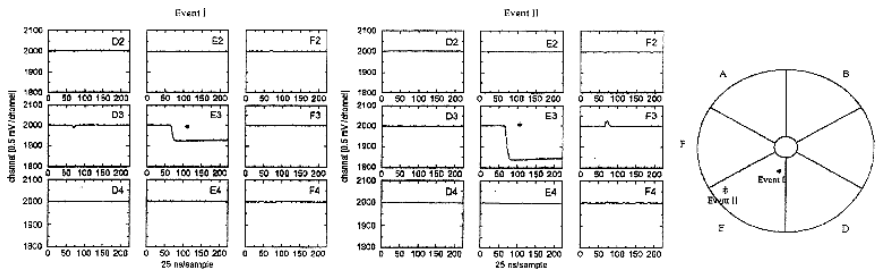


Fig. 24. An example of digitised signals from a 36-fold segmented Ge detector[54]. The *right hand side* shows the segmented front face of the detector and two events occurring in Sect. E. In both cases the main signal comes from Sect. E3 of the detector but the mirror signals are from Sect. D3 and F3 for events 1 and 2, respectively, indicating where the interaction occurred relative to the segment boundaries in the two cases

25 ns, covering around 5 μ s. It is clear that both events show a net charge in segment E3. Event I shows a γ -ray interaction, which deposited approximately 200 keV. It manifests a negative mirror charge in D3. In other words the main interaction occurred in segment E3 close to the inner contact near to the boundary with D3. Event II deposits approximately 430 keV and there is a positive mirror charge in F3. Here the main interaction occurred in segment E3 close to the outer contact, near the boundary with F3. These examples show crudely that it is possible to say where the interaction occurred by looking at these signals. Accurate definition of the interaction position can be achieved by the use of methods of artificial intelligence, genetic algorithms, artificial neural networks and the Discrete Wavelet Transform method. Much work remains to define the best method but it is clear that tracking can work. We can anticipate that we will meet the required position resolution of a few mm.

Based on these ideas it is possible to build a full array with greatly enhanced performance. In the case of AGATA [53] and GRETA [51] the aim is to achieve the best possible spectrum quality with the maximum efficiency for ions moving with velocities up to $v/c = 0.7$. This is of great importance in the case of AGATA since the array is intended for use with relativistic RIBs from the FAIR project [56] at GSI as well as elsewhere. Much of the improvement relative to EUROBALL and GAMMASPHERE will come from the much improved angular definition of the events in AGATA. This array is based on hexagonal, 36-fold segmented, encapsulated Ge crystals. Each detector has a length of 9 cm and a diameter, before shaping, of 8 cm. They are then tapered to form the hexagonal geometry. Three of these with digital front-end electronics and a common Dewar, form an AGATA module. The optimum design requires a great deal of simulation to decide on the best geometry. The geometry chosen for AGATA involves tiling the sphere with 180 hexagonal clusters, grouped into 60 identical triple modules, and 12 pentagons. There are, of course, holes for the beam to enter and leave. Simulations suggest an excellent performance for AGATA. With a γ -ray multiplicity of one the absolute efficiency is expected to be 39% with a peak-to-total ratio of 53%. For high multiplicity events (multiplicity 30) this falls to 25% and 46% respectively. The spectrum quality will still be excellent in the latter case. The geometry chosen has a very high granularity with an angular resolution of 1.25° . This will be of considerable importance when we are carrying out experiments involving high recoil velocity. Figure 25 shows the resolving power of Ge arrays starting with the TESSA arrays and ending with the projected values for AGATA and GRETA. Enormous strides have been made in sensitivity and there are more to come.

Both AGATA and GRETA will have a huge impact on γ -ray spectroscopy. However if we are to study exotic nuclei formed in reactions with tiny cross-sections it is vital that they are built. Their excellent angular resolution will allow spectroscopy with relativistic fragmentation beams and thus allow

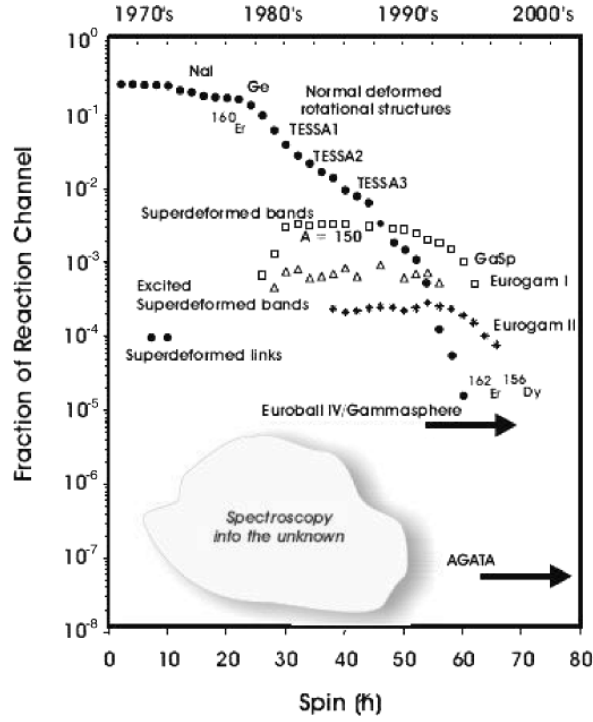


Fig. 25. History of the improvement in sensitivity of γ -ray arrays

studies of the most exotic nuclei. Their capability of operating at high event rates will allow us to cope with the high backgrounds likely to be encountered in many such experiments. Highly segmented coaxial detectors are not the only possible route to a tracking array. Our brief description of the main developments in tracking detectors neglects other developments [57] involving planar Ge detectors. In principle one can achieve good position resolution without pulse shape analysis by using planar detectors with sufficient pixellation. At Argonne National Laboratory [58] the idea of stacked planar detectors is being developed with the aim of obtaining good spatial resolution. The idea is to use stacks of thin planar detectors, each with one highly pixellated electrode or a double-sided Ge strip detector where opposite electrodes are in orthogonal strips. In essence the position resolution is achieved if we can ensure that the first interaction occurs in only one pixel. This means that one must use small strips (pixels), as well as thin detector layers. Alternatively one can use thin strips and a thicker detector and determine the interaction depth from the measured drift time. The second detector must also have good position resolution in order to determine the Compton scattering angle precisely. For the second detector layer one can use either a segmented Ge detector or a stack of planar de-

tectors. In terms of creating a tracking array composed of planar detectors they have certain limitations at present. In any array one wants to maximise the efficiency by maximising the solid angle. However planar detectors have dead material, such as the guard rings at the edges of the crystals and the mechanical structure needed to provide cooling. Future efforts will be devoted to minimising or eliminating them. Accordingly, for the present, segmented coaxial detectors will be the main development route for tracking arrays. However there is already a range of applications for position/direction sensitive detector systems. Gamma-ray astronomy is amongst the most important of these and therefore such systems have been actively developed for at least a decade. In nuclear physics they may play a strong role in studying the decay of short-lived activities from the new radioactive beam facilities and this means that they are also of interest for detecting weak activity for security or environmental purposes.

6 Summary

Arrays of Ge detectors have proved to be very powerful tools for nuclear spectroscopy. Steady improvements in resolving power and efficiency have produced many results on the properties of atomic nuclei at high spin. It turns out that they are also well suited for use in the study of γ rays from exotic nuclei. They will be amongst the most powerful tools available to exploit the beams from the next generation of radioactive beam facilities.

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Colour Supplement



Plate 1. Photo taken at the 2004 Surrey School



Plate 2. Photo taken at the 2005 Darmstadt-Mainz School

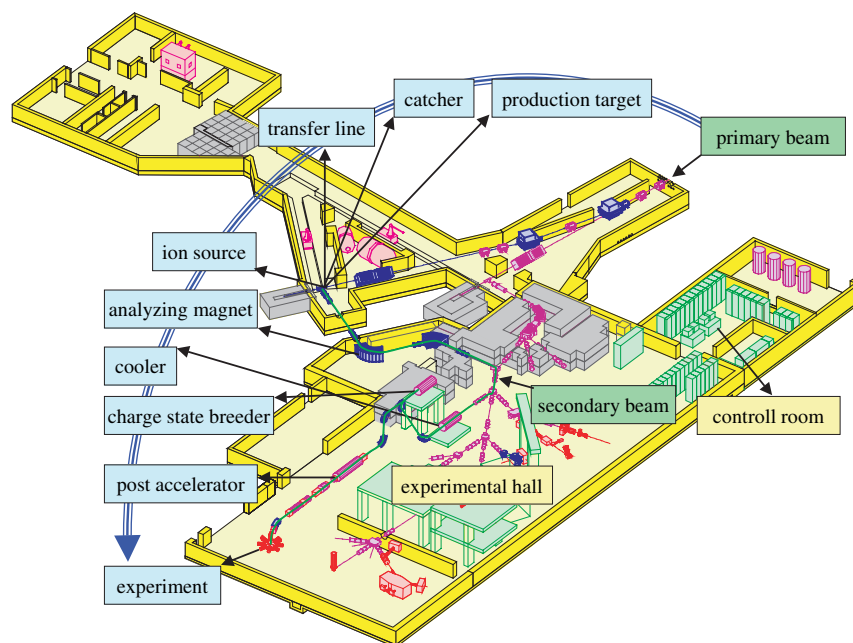


Plate 3. Layout of the CERN-ISOLDE hall; see Fig. 1 of the chapter by P. Van Duppen

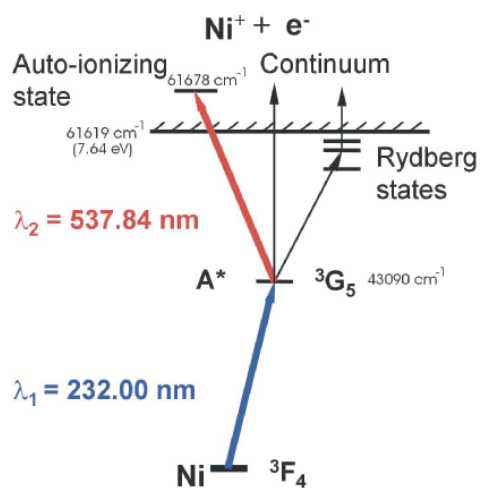


Plate 4. Principle of resonant photo-ionization of nickel; see Fig. 5 of the chapter by P. Van Duppen

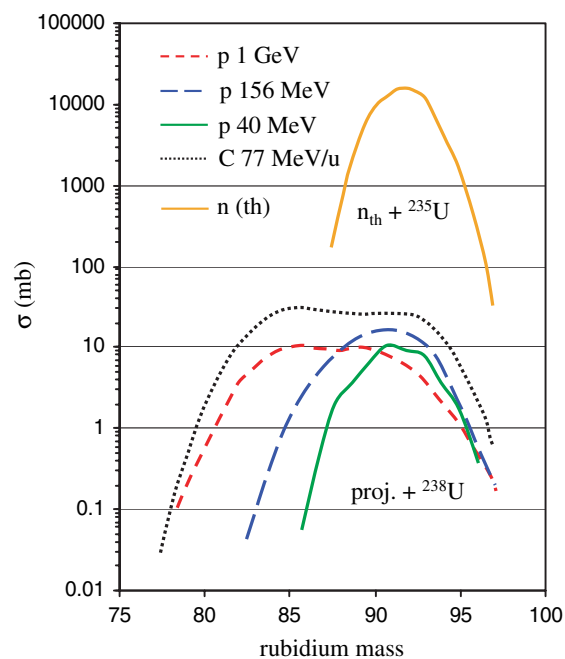


Plate 5. Cross sections for the production of rubidium isotopes from fission of ${}^{238}\text{U}$ using various probes; see Fig. 6 of the chapter by P. Van Duppen

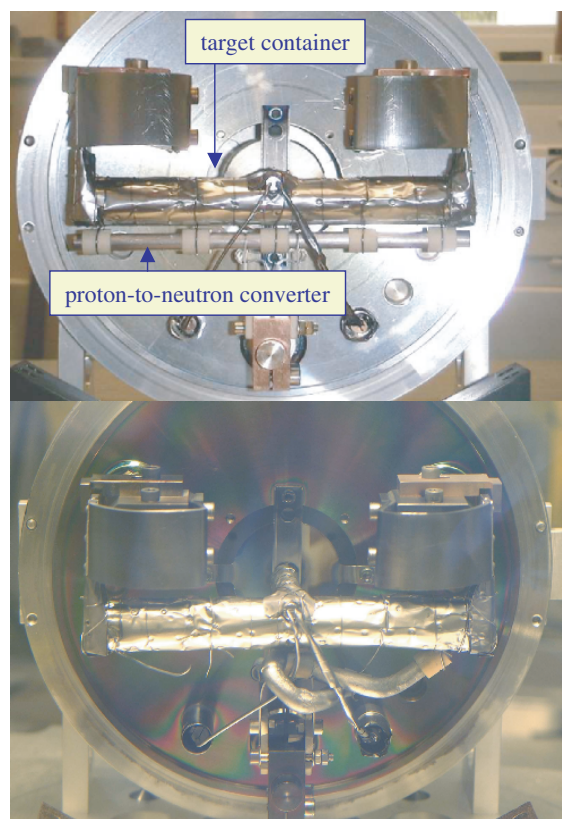


Plate 6. ISOLDE uranium-carbide target, see Fig. 7 of the chapter by P. Van Duppen

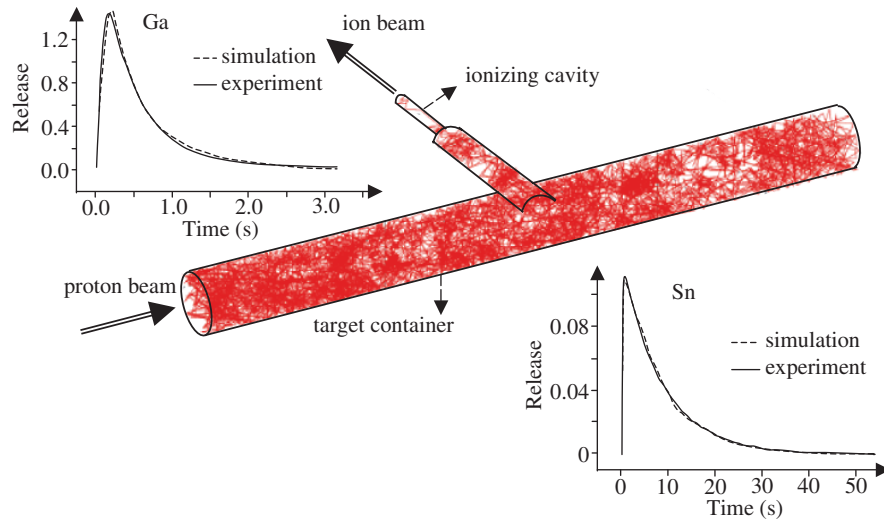


Plate 7. Results obtained by simulated the path (red line) of one atom produced in a tantalum foil target connected with a surface ionizing cavity; see Fig. 9 of the chapter by P. Van Duppen

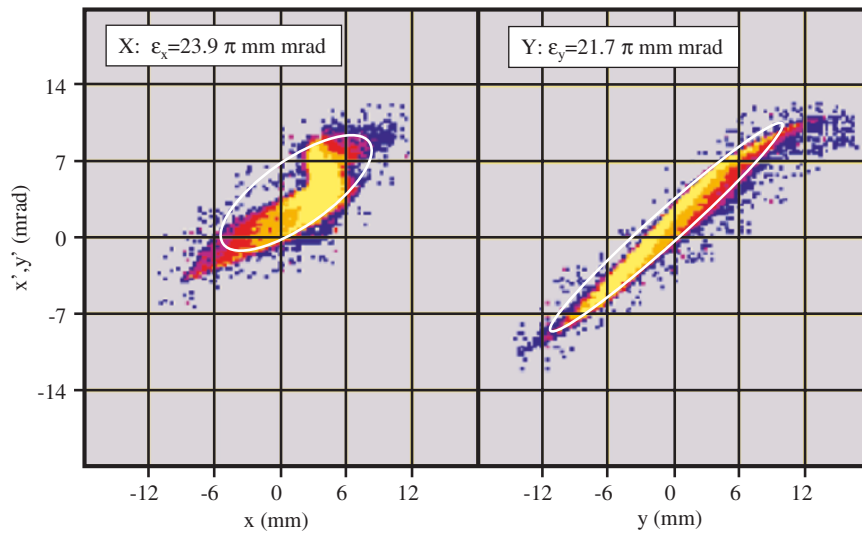


Plate 8. Example of the results of an emittance measurement for a ^{120}Sn beam; see Fig. 12 of the chapter by P. Van Duppen

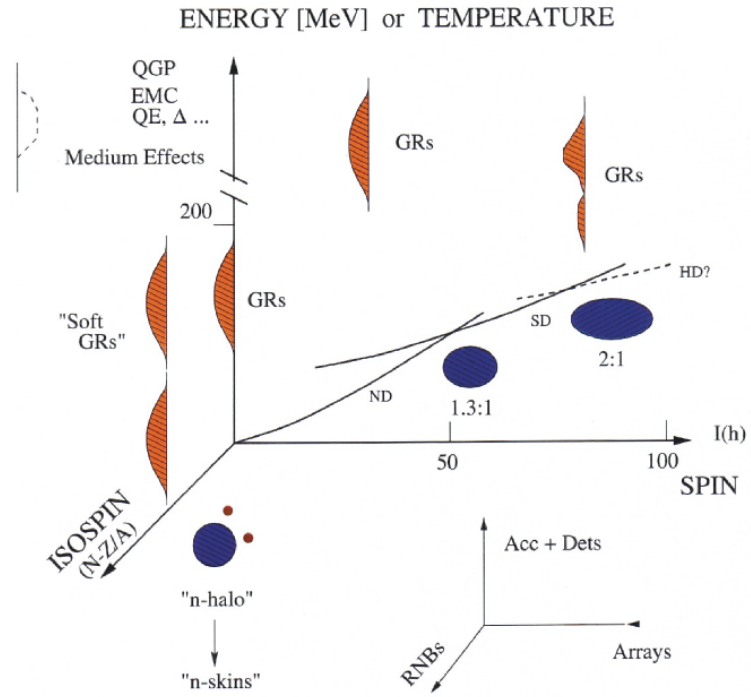


Plate 9. Key parameters describing the atomic nucleus; see Fig. 1 of the chapter by W.B. Gelletly and J. Eberth

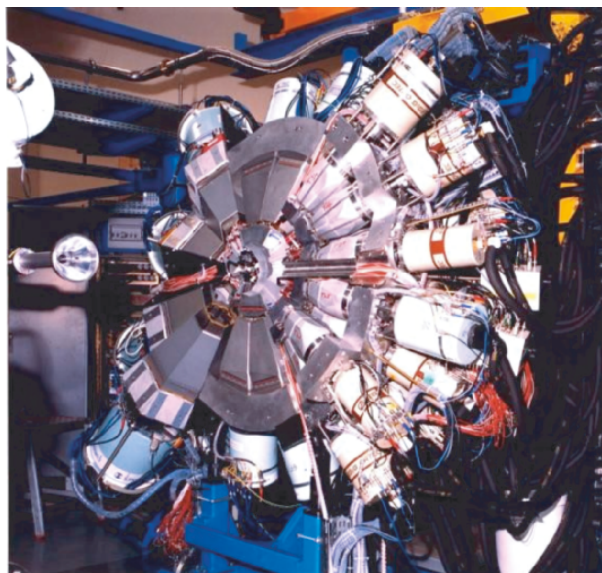


Plate 10. EUROBALL array; see Fig. 15 of the chapter by W.B. Gelletly and J. Eberth

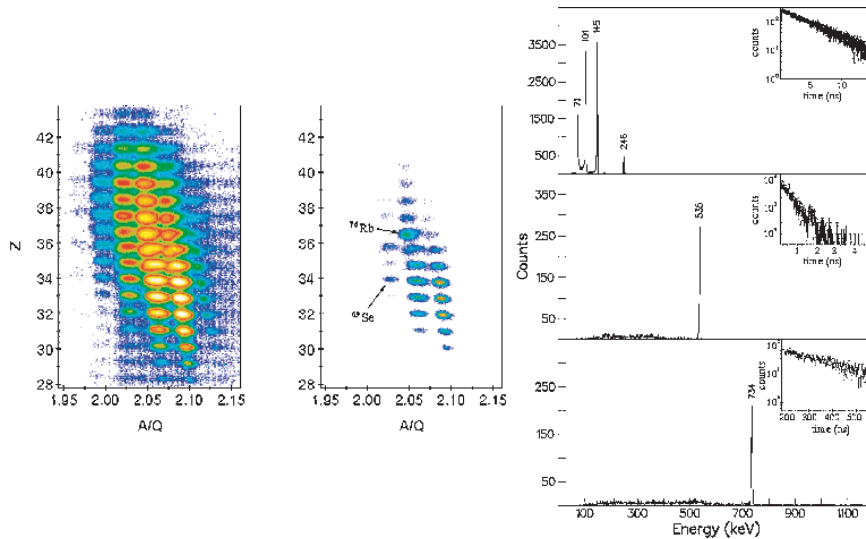
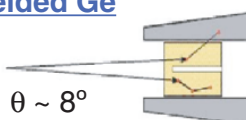


Plate 11. Isotope identification and γ -ray spectroscopy of neutron-deficient isotopes produced in the fragmentation of ^{92}Mo ions; see Fig. 19 of the chapter by W.B. Gelletly and J. Eberth

Idea of γ -ray tracking

Compton Shielded Ge

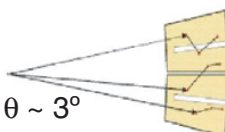
$\epsilon_{ph} \sim 10\%$
 $N_{det} \sim 100$
 $\Omega \sim 40\%$



large opening angle
means poor energy
resolution at high
recoil velocity

Ge Sphere

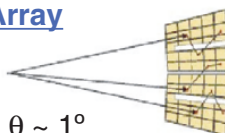
$\epsilon_{ph} \sim 50\%$
 $N_{det} \sim 1000$



too many detectors
are needed to avoid
summing effects

Ge Tracking Array

$\epsilon_{ph} \sim 50\%$
 $N_{det} \sim 100$
 $\Omega \sim 80\%$



Combination of:

- segmented detectors
- digital electronics
- pulseprocessing
- tracking the γ -rays

Plate 12. Illustration of the idea of γ -ray tracking; see Fig. 23 of the chapter by W.B. Gelletly and J. Eberth

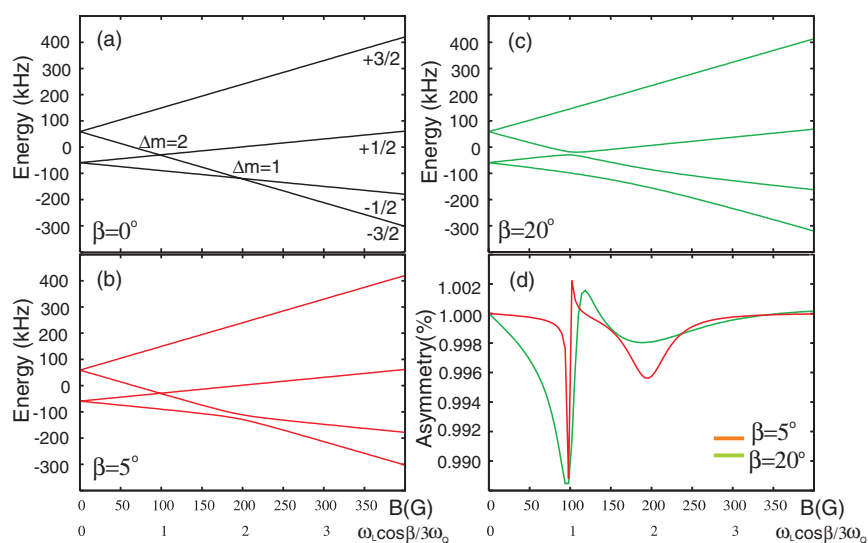


Plate 13. Crossing or mixing of hyperfine levels; see Fig. 6 of the chapter by R. Neugart and G. Neyens

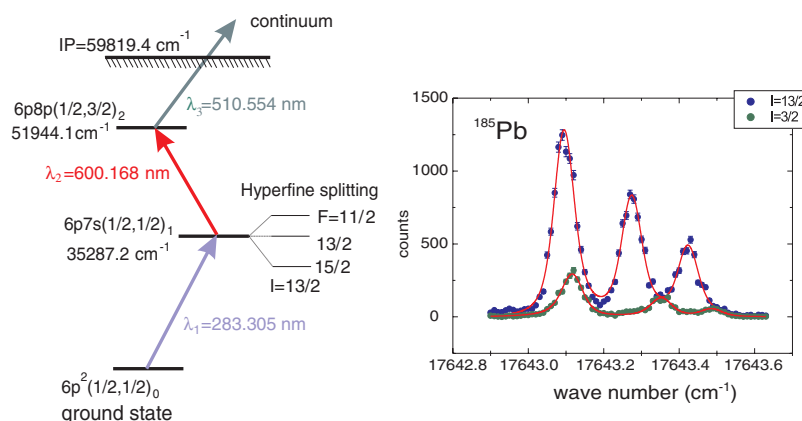


Plate 14. Laser ionization scheme for Pb isotopes and hyperfine structure observed for two isomeric states in ^{185}Pb ; see Fig. 13 of the chapter by R. Neugart and G. Neyens

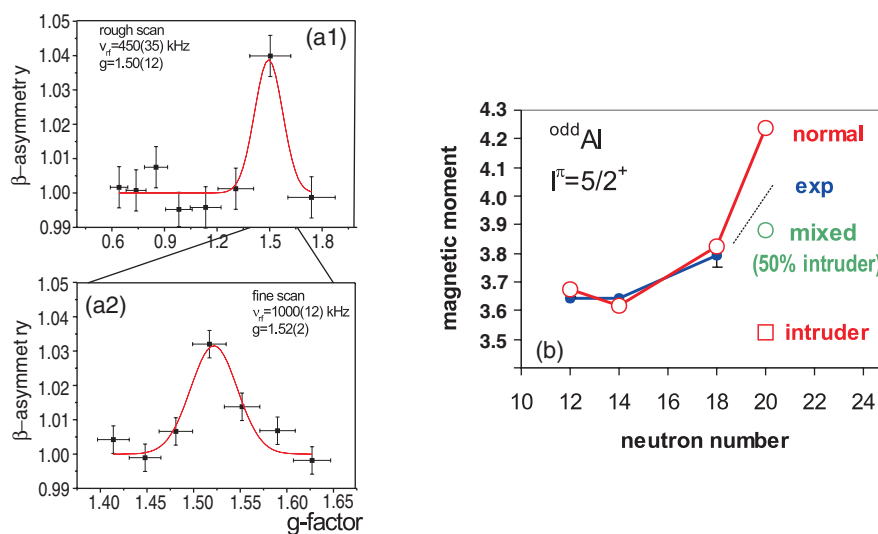


Plate 15. Nuclear magnetic resonances observed for ^{31}Al (a1 and a2) and comparison of experimental magnetic-moment data for ^{33}Al with model predictions; see Fig. 17 of the chapter by R. Neugart and G. Neyens

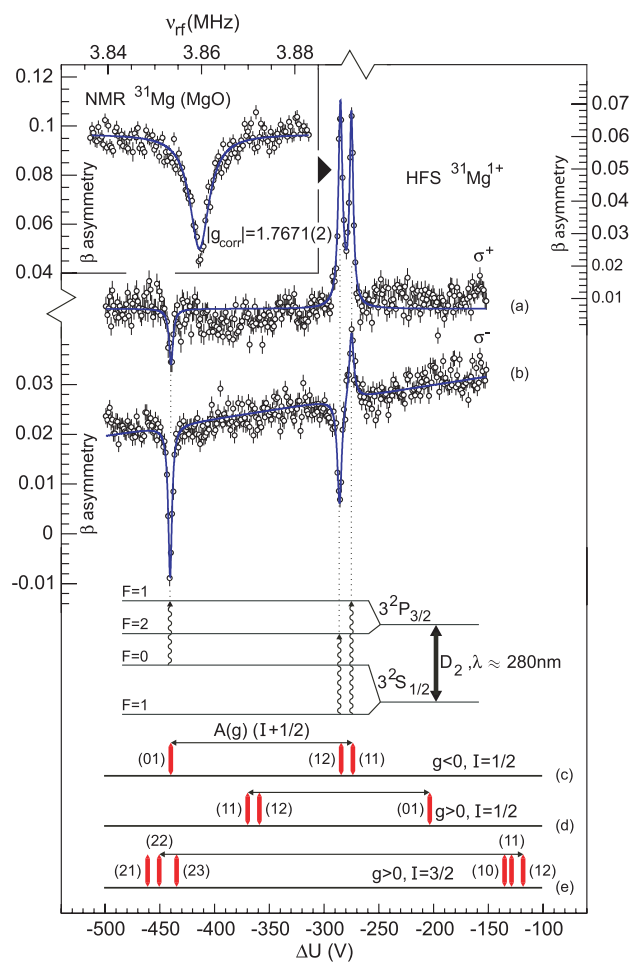


Plate 16. Hyperfine structure measured for the D_2 line of $^{31}\text{Mg}^+$; see Fig. 18 of the chapter by R. Neugart and G. Neyens

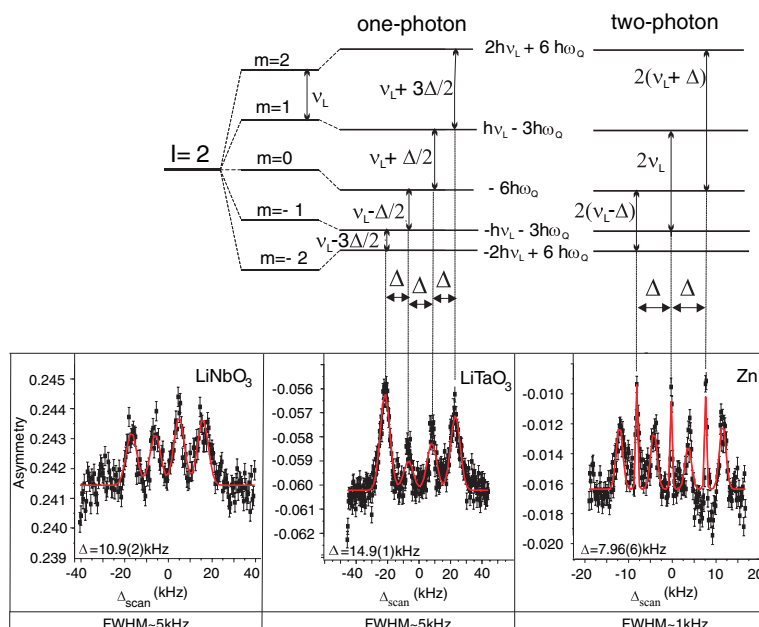


Plate 17. Nuclear magnetic resonances observed for ^8Li ; see Fig. 19 of the chapter by R. Neugart and G. Neyens



Plate 18. Illustrative representation of a spallation reaction; see Fig. 1 of the chapter by J. Benlliure

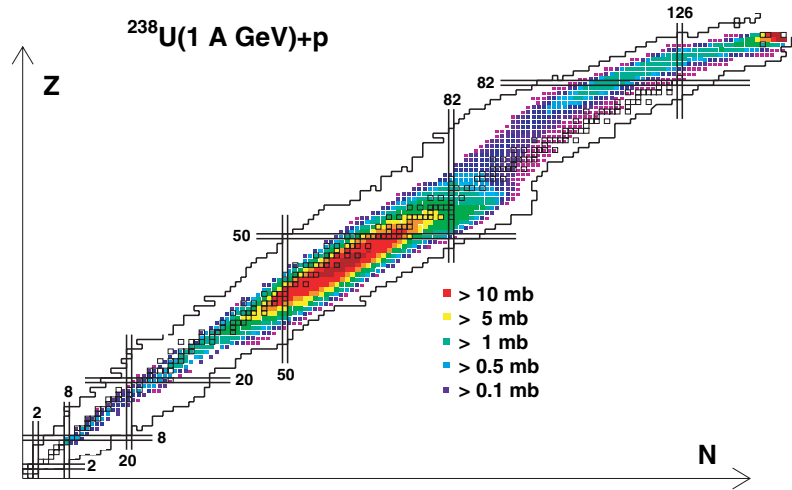


Plate 19. Residual nuclei produced in the interaction of 1 GeV protons with ^{238}U ; see Fig. 3 of the chapter by J. Benlliure

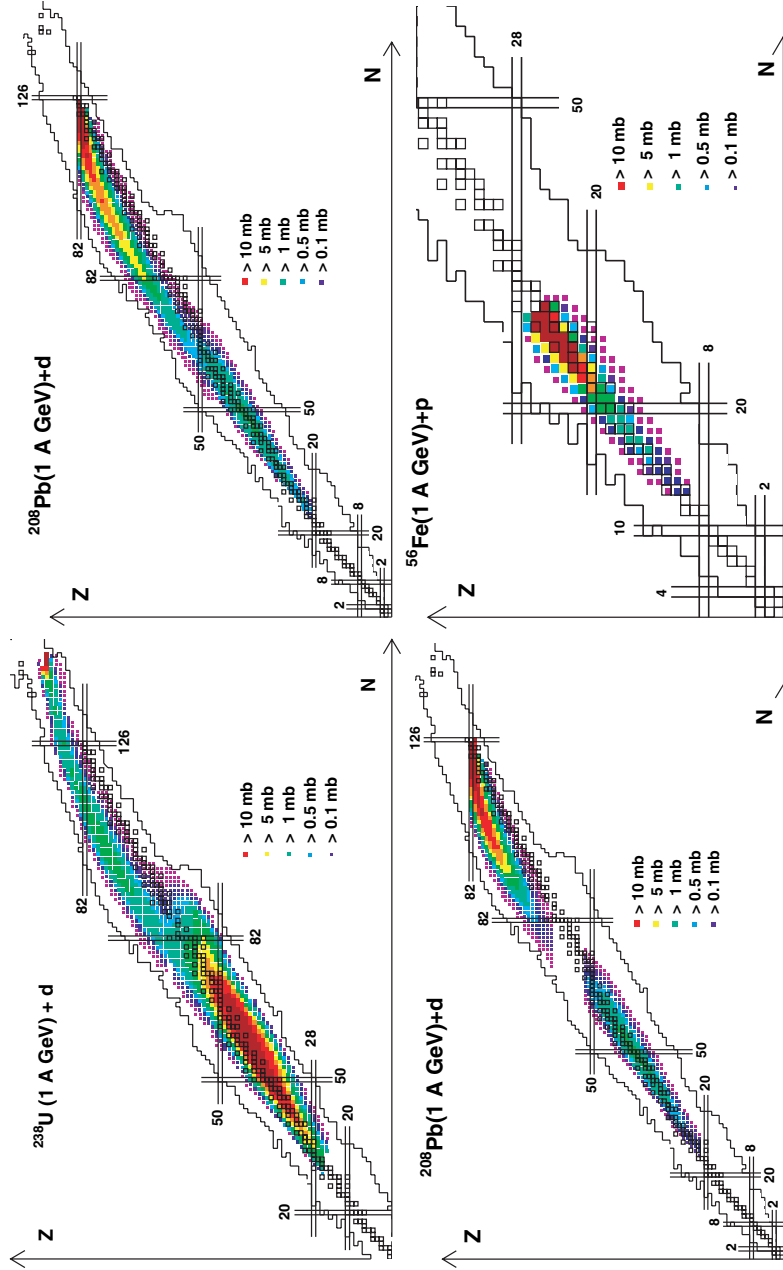


Plate 20a–d. Residual nuclei produced in the reactions $^{238}\text{U}(1\text{ A GeV}) + \text{d}$, $^{208}\text{Pb}(1\text{ A GeV}) + \text{p,d}$ and $^{56}\text{Fe}(1\text{ A GeV}) + \text{p}$; see Fig. 8 of the chapter by J. Benlliure

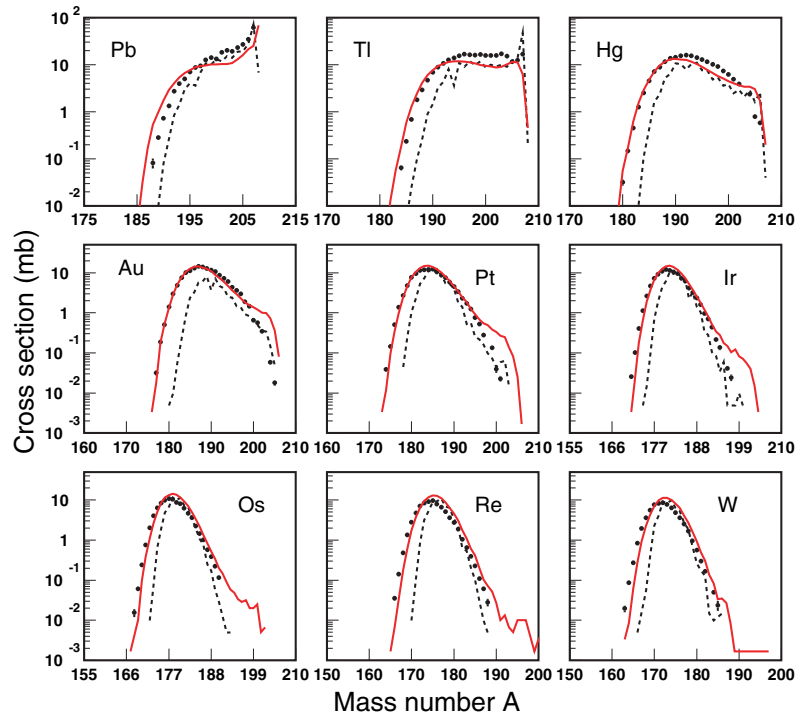


Plate 21. Isotopic distributions of production cross-sections for tungsten to lead isotopes from the ^{208}Pb (1 A GeV) + p reaction; experimental data are compared to predictions obtained from the Lahet code (dashed lines) and an intra-nuclear cascade model; see Fig. 9 of the chapter by J. Benlliure

Nuclear Moments

R. Neugart¹ and G. Neyens²

¹ Institut für Physik, Universität Mainz, 55099 Mainz, Germany

² Instituut voor Kern- en Stralingsfysica, K.U. Leuven, 3001 Leuven, Belgium

Abstract. Spins, electromagnetic moments and radii are basic observables characterizing a particular nucleus. We review a variety of experimental techniques currently used to measure these quantities for unstable nuclei far from stability and for isomeric nuclear states. The experiments can be traced back to two different principles: Optical high-resolution spectroscopy reveals the interaction between nuclear moments and the electromagnetic fields produced by the shell electrons. On the other hand, β - or γ -ray detection reveals the movement of nuclear spin systems in externally applied fields or hyperfine fields produced by a crystal lattice surrounding. The relevance of such measurements for nuclear structure research is discussed in connection with selected examples of results.

1 Introduction

The structure of nuclei is a fingerprint of how protons and neutrons in these many-particle systems interact in order to form a bound nucleus. Measuring nuclear properties is of utmost importance to understand the interactions which bind the protons and neutrons together in an isotope with a mass $A = Z + N$ (= bound nucleus with Z protons and N neutrons). Next to the few hundred stable and long-lived nuclei that occur in our universe, several thousands of unstable nuclei are predicted to be bound. In our laboratories, more than 2000 of them have been produced to date, and for several hundreds of them basic properties such as mass, lifetime, decay and excitation scheme, spins and moments have been investigated thoroughly.

These basic properties allow us to derive directly or indirectly information on the nuclear structure, as well as on the strong nuclear force. Comparison of the experimental properties of very exotic nuclei to calculations performed with different nuclear models allow testing the predictive power of these models when going to the extremes, or give a hint on how to further improve the nuclear models and their parameters. Some properties are immediately sensitive to the pairing interaction between nucleons, while others reveal more information about the proton-neutron interaction or about nuclear deformation and other phenomena, as discussed e.g., by Grawe in [1].

To get a clear understanding of the single-particle structure or the collective nature of nuclear states, static nuclear moments are crucial ingredients.

The magnetic moment is sensitive to the single-particle nature of the valence nucleon, while the nuclear quadrupole moment is also sensitive to the deformation. Since more than 40 years, with a peak in the seventies and eighties, nuclear moments of radioactive nuclei have been measured, both in the ground state and in some isomeric states [2]. Experiments were mainly done for nuclei close to the valley of stability and on its neutron-deficient side, because the nuclei were produced mostly in fusion-evaporation or spallation reactions.

Since the nineties, there has been a renewed interest to measure nuclear moments. This is because it has become possible to produce and select very exotic nuclei in sufficient quantities ($> 10^2\text{--}10^3\text{ s}^{-1}$), as discussed by Huyse [3] and by Morrissey and Sherrill [4]. Fascinating new features have been discovered in nuclei far from stability, the most striking one being the existence of “halo” structures in nuclei with a very asymmetric proton to neutron ratio (see Al-Khalili [5] for an introduction to this subject). Also the appearance and disappearance of magic numbers has been a topic of several recent investigations, both experimentally and theoretically (see [1] and [6] for an introduction and some examples).

In this lecture, we shall present some techniques which allow studying the static nuclear moments of exotic nuclei in their ground state or isomeric states, and give some examples of recent experiments on exotic nuclei. Comprehensive review papers have been published by Otten [7] on the use of atomic spectroscopy methods for nuclear moments research (see also more recent reports [8, 9, 10]), and by Neyens on the use of nuclear spin-oriented radioactive beams [11].

2 Monopole, Dipole and Quadrupole Moments

2.1 The Nuclear Mean Square Charge Radius

In atomic physics the nucleus is usually considered to be a positive point charge Ze forming the center of gravity of a system of electrons. The nuclear mass and size affect very little the atomic energy levels, and this is reflected in the isotope shifts which are accessible only to very high resolution atomic spectroscopy methods. The finite nuclear mass gives rise to a small nuclear recoil energy which is observed as a “mass shift” between the atomic transition frequencies of different isotopes. This effect is not used to deduce information on nuclear properties or nuclear structure since nuclear masses can be measured directly with high accuracy, as discussed by Bollen in [12].

A second part of the isotope shift is due to the nucleus not being a point charge, but having an extended charge distribution which changes from one isotope to another. Traditionally this is called “nuclear volume shift” or “field shift”. Roughly speaking, the extension of the nuclear charge can be expressed by a charge radius R_p . Due to this nonzero radius the binding energy of a

real atom is smaller than that of a hypothetical atom containing a point-charge nucleus. Of course, this effect cannot be measured directly, because a point-charge nucleus does not exist, but it is visible in the shift between the transition frequencies of different isotopes for which R_p is different. Vice versa, information about differences of radii can be obtained from the measurement of isotope shifts. As nuclear radii are sensitive to details of the nuclear structure such as shell effects and deformation, this has become an important tool for exploring the behavior of nuclear systems.

Considering the nucleus as a liquid drop, with the protons homogeneously distributed over the sphere of the nucleus, one can assume the radius of the proton distribution R_p to be equal to that of the nuclear mass distribution,

$$R = R_0 A^{1/3}, \quad (1)$$

where $R_0 = 1.2$ fm and A is the atomic mass number, i.e. the number of nucleons in the nucleus. A refined model would distinguish between protons and neutrons, because the proton radius R_p and the neutron radius R_n develop differently as a function of the proton number Z or neutron number N . This distinction is made in the “droplet” model [13, 14] which is often used to interpret isotope shift data on nuclear radii. Still, the liquid-drop model can serve as a guide, while it is clear that the interesting physics is found in local deviations from the global formula (1).

To a good approximation, the nuclear quantity governing the isotope shift is the second radial moment of the nuclear charge distribution [15],

$$\langle r^2 \rangle = \frac{\int_0^R \rho(\mathbf{r}) r^2 d\mathbf{r}}{\int_0^R \rho(\mathbf{r}) d\mathbf{r}}, \quad (2)$$

where the integral in the denominator is just the nuclear charge Ze . This quantity, called “nuclear mean square charge radius”, can be considered as a monopole moment of the nucleus. Isotope shifts contain information on the change of this quantity as a function of the neutron number. For the liquid-drop nucleus with radius R , one has

$$\langle r^2 \rangle_{\text{LD}} = (3/5)R^2 = (3/5)R_0^2 A^{2/3}, \quad (3)$$

which leads to the differential effect for small changes of A :

$$\delta \langle r^2 \rangle_{\text{LD}} = (2/5)R_0^2 A^{-1/3} \delta A. \quad (4)$$

It is mainly the deviation from this liquid-drop behavior of radii that reveals interesting features of nuclear structure such as the development of neutron (or proton) skins, shell closures, pairing and deformation. As a collective phenomenon the effect of deformation can already be incorporated into the liquid-drop description.

Nuclei are not necessarily spherical. The shell structure causes deformed equilibrium shapes mainly in the regions between the shell closures (magic numbers) of protons and neutrons. Conventionally this nuclear deformation is described by a quadrupole deformation parameter β defined by an angular dependence of the length of the radius vector to the nuclear surface expressed in spherical harmonics. For the important case of rotational symmetry one can express this by using an expansion of the nuclear shape in spherical harmonics Y_{kq} of which Y_{20} represents the quadrupole deformation term:

$$R(\theta) = R_1 [1 + \beta Y_{20}(\theta)] . \quad (5)$$

R_1 is chosen such that the nuclear volume is constant, i.e. independent of β . With these definitions the mean square radius of a deformed nucleus [15] becomes

$$\langle r^2 \rangle = \frac{3}{5} R^2 + \frac{3}{4\pi} R^2 \beta^2 . \quad (6)$$

More generally – without the restriction to a sharp nuclear surface – the right-hand side of (6) can be expressed by the mean square radius $\langle r^2 \rangle_{\text{sph}}$ of a spherical nucleus which has the same volume,

$$\langle r^2 \rangle = \langle r^2 \rangle_{\text{sph}} + \frac{5}{4\pi} \langle r^2 \rangle_{\text{sph}} \beta^2 , \quad (7)$$

and the differential effect becomes

$$\delta \langle r^2 \rangle = \delta \langle r^2 \rangle_{\text{sph}} + \frac{5}{4\pi} \langle r^2 \rangle_{\text{sph}} \delta \beta^2 . \quad (8)$$

This expression shows how a change in nuclear deformation affects the change in the nuclear mean square radius. Note again that all quantities refer to the charge (proton) distribution in the nucleus and that in particular β is the charge deformation related to this charge distribution. This should be kept in mind in particular when comparing the experimental results to nuclear models.

2.2 The Nuclear Magnetic Dipole Moment

The magnetic moment of a nucleus is induced by the orbiting charged particles (the protons) giving rise to an orbital magnetic field (characterized by g_l) and by the intrinsic spin $s = 1/2$ of the nucleons, inducing their own intrinsic magnetic field (characterized by g_s). The dipole operator, expressed in terms of these two contributions, is given by

$$\boldsymbol{\mu} = \sum_{i=1}^A g_l^i \mathbf{l}^i + \sum_{i=1}^A g_s^i \mathbf{s}^i . \quad (9)$$

The free-nucleon gyromagnetic factors for protons and neutrons are $g_l^{(p)} = 1$, $g_l^{(n)} = 0$, $g_s^{(p)} = +5.587$, $g_s^{(n)} = -3.826$. The magnetic dipole moment μ_I is the expectation value of the z -component of the dipole operator $\boldsymbol{\mu}$:

$$\mu(I) = \langle I, m = I | \boldsymbol{\mu}_z | I, m = I \rangle. \quad (10)$$

It is related to the nuclear spin $\mathbf{I}$ via the gyromagnetic ratio g_I : $\boldsymbol{\mu} = g_I \mathbf{I} \mu_N$, with μ_N being the nuclear magneton.

Within the shell model picture (see [1]) the properties of odd- A nuclei near closed shells are described by the characteristics of the unpaired valence nucleon. The magnetic moment of such a nuclear state with its valence nucleon in an orbit with total angular momentum $\mathbf{j}$ and orbital momentum $\mathbf{l}$, can be calculated as a function of the free-nucleon g -factors, and they are called the Schmidt moments:

$$\mu(l + 1/2) = \left[\left(j - \frac{1}{2} \right) g_l + \frac{1}{2} g_s \right] \mu_N, \quad (11a)$$

$$\mu(l - 1/2) = \frac{j}{j + 1} \left[\left(j + \frac{3}{2} \right) g_l - \frac{1}{2} g_s \right] \mu_N. \quad (11b)$$

In a real nucleus the magnetic moment of a nucleon is influenced by the presence of the other nucleons. This can be taken into account by using “effective” proton and neutron g -factors to calculate the *effective single-particle magnetic moment* $\mu(lj)^{\text{eff}}$ for a nucleon in a particular orbit. The single-nucleon g_s -factors are reduced to typically about 70% of their free-nucleon value in heavy nuclei, while in light nuclei the experimental numbers are rather well reproduced with free-nucleon g -factors. A detailed discussion on the different corrections required to account for the nuclear medium can be found in [16] or [17].

The dipole operator is a one-body operator. Its expectation value for a nuclear state with spin I is given by

$$\mu(I) = \langle I(j_1, j_2, \dots, j_n), m = I | \sum_{i=1}^n \boldsymbol{\mu}_z(i) | I(j_1, j_2, \dots, j_n), m = I \rangle. \quad (12)$$

For a nuclear state composed of valence nucleons in orbits $\mathbf{j}_i$ coupled to a state with spin $\mathbf{I}$, some general additivity rules for the magnetic moment as a function of the single-particle magnetic moments can be deduced from (12) by decomposing the wave function $|I(j_1, \dots, j_n), m\rangle$ into its single-particle components $|j_i, m_i\rangle$. Here one uses Clebsch-Gordon coefficients in case of two particles [18] and Coefficients of Fractional Parentage for more particles [19] in an orbit. For a nuclear state described by a weak coupling between protons and neutrons, the magnetic moment can be calculated [20] as

$$\mu(I) = \frac{I}{2} \left[\frac{\mu_p}{j_p} + \frac{\mu_n}{j_n} + \left(\frac{\mu_p}{j_p} - \frac{\mu_n}{j_n} \right) \frac{j_p(j_p + 1) - j_n(j_n + 1)}{I(I + 1)} \right]. \quad (13)$$

Examples of the additivity of magnetic moments in odd-odd nuclei are given e.g., in [11] or [20].

2.3 The Nuclear Electric Quadrupole Moment

The non-spherical distribution of the charges in a nucleus gives rise to an electric quadrupole moment. The classical definition of the charge quadrupole moment in a Cartesian axis system is given [20] by

$$Q_z = \sum_{i=1}^A Q_z(i) = \sum_{i=1}^A e_i (3z_i^2 - r_i^2), \quad (14)$$

with e_i being the charge of the respective nucleon and (x_i, y_i, z_i) its coordinates. In a spherical tensor basis the z -component of the quadrupole operator is more easily expressed as the zero order tensor component of a rank 2 tensor:

$$Q_2^0 = Q_z = \sqrt{\frac{16\pi}{5}} \sum_{i=1}^A e_i r_i^2 Y_2^0(\theta_i, \phi_i). \quad (15)$$

The nucleus is a quantum mechanical system that is described by a nuclear wave function, characterized by a nuclear spin I . In experiments we observe the spectroscopic quadrupole moment, which is the expectation value of the quadrupole moment operator, defined as

$$Q_s(I) = \langle I, m = I | Q_2^0 | I, m = I \rangle = \sqrt{\frac{I(2I-1)}{(2I+1)(2I+3)(I+1)}} (I || Q || I). \quad (16)$$

This expression shows that the spectroscopic quadrupole moment Q_s of a nuclear state with spin $I < 1$ is zero. Thus, although a nucleus with spin $I = 0$ or $1/2$ can possess an intrinsic deformation, one can not measure this via the quadrupole moment.

The spectroscopic quadrupole moment can be related to an intrinsic quadrupole moment Q_0 reflecting the nuclear deformation β , only if certain assumptions about the nuclear structure are made. An assumption that is often made (but is not always valid!), is that the nuclear deformation is axially symmetric with the nuclear spin having a well-defined direction with respect to the symmetry axis of the deformation (strong coupling). In this case, the intrinsic and the spectroscopic quadrupole moment are related as follows:

$$Q_s = \frac{3K^2 - I(I+1)}{(I+1)(2I+3)} Q_0, \quad (17)$$

with K being the projection of the nuclear spin on the deformation axis.

This intrinsic quadrupole moment Q_0 , induced by the non-spherical charge distribution of the protons, can then be related to the nuclear charge deformation β as follows [21]:

$$Q_0 = \frac{3}{\sqrt{5\pi}} Z R^2 \beta (1 + 0.36\beta) , \quad (18)$$

with R as defined in (1) and β related to the mean square charge radius as in (6).

In the shell model, the nuclei are described by a nuclear mean field (core) in which some individual valence nucleons move and interact with each other through a residual interaction. For calculating the spectroscopic quadrupole moment, the sum over all nucleons in expressions (15) and (16) is then reduced to the sum over the valence particles. To take into account the interaction with the mean field, an effective charge is attributed to the nucleons, neutrons as well as protons. The effective charges are model dependent: if a smaller model space is taken for the valence nucleons, the effective charge needed to reproduce the experimental quadrupole moments deviate more from the nucleon charges. However, for a large enough model space, the effective charges are found to be constant in a broad region of nuclei and deviations between the calculated and experimental moments can then be a signature for unaccounted nuclear structure effects. Effective charges have been determined in several regions of the nuclear chart by comparing experimental quadrupole moments of nuclei, whose proton or neutron number deviates from doubly-magic by ± 1 , to their values calculated in the shell model. Typical values vary from $e_\pi^{\text{eff}} \approx 1.3 e$, $e_\nu^{\text{eff}} \approx 0.35 e$ in light nuclei [22] to $e_\pi^{\text{eff}} \approx 1.6 e$, $e_\nu^{\text{eff}} \approx 0.95 e$ in the lead region [11].

Similar as the magnetic moment operator, the quadrupole moment operator is a one-body operator. Thus, for a multi-nucleon configuration of weakly interacting nucleons one can deduce some additivity rules for quadrupole moments, starting from the general definition in (16) and using angular momentum coupling rules and tensor algebra [19, 20]. With these additivity rules one can predict rather reliably the quadrupole moments of configurations of which the moments of the proton and neutron part have been measured [11].

In the same way, by decomposing the single-particle wave function into its radial, spin and orbital part, the single-particle quadrupole moment for an unpaired nucleon in an orbit with angular momentum $\mathbf{j}$ can be deduced from (16)

$$Q_{\text{s.p.}} = -e_j \frac{2j-1}{2j+2} \langle r_j^2 \rangle . \quad (19)$$

Here e_j is the effective charge of the nucleon in orbital j and $\langle r_j^2 \rangle$ is the mean square radius of the nucleon in that orbital.

Note that free neutrons have no charge, $e_\nu = 0$, and therefore do not induce a single-particle quadrupole moment. However, neutrons in a nucleus

interact with the nucleons of the core and can polarize the core, which is reflected by giving the neutrons an effective charge. Because the nuclear energy is minimized if the overlap of the core nucleons with the valence particle (or hole) is maximal, a particle (respectively hole) will polarize the core towards an oblate (respectively prolate) deformation, as demonstrated in Fig. 1.

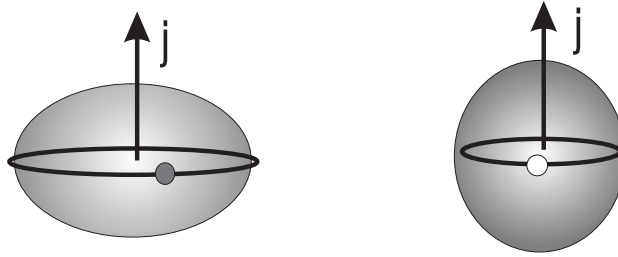


Plate 1. Graphical representation of a particle in a orbital j , polarizing the core towards oblate deformation with a negative spectroscopic quadrupole moment (*left*), and a hole in an orbital giving rise to a prolate core polarization (*right*)

A change of the quadrupole moment as a function of N or Z can be either a signature for a change in the core polarization or, if the change is drastic, an indication for an onset of a static nuclear deformation. For example, the systematic increase of quadrupole moments with decreasing neutron number in neutron-deficient Po isotopes (see Fig. 2) has been explained by an increase of the quadrupole-quadrupole interaction between the proton particles and the increasing amount of neutron holes [23, 24]. Figure 2 also displays the measured quadrupole moments of isomeric states in neutron-deficient Pb isotopes, having the same spin and valence configuration as those in Po. They are much larger than those of their Po isotones which is consistent with the assumption of a statically deformed nuclear potential [25]. This was the first prove that the intruder isomers in the Pb isotopes are indeed deformed due to the strong deformation-driving effect of the two proton holes in the $Z = 82$ shell. The intruder isomers coexist with the normal near-spherical structures at low energy.

3 Measuring Static Nuclear Moments: Basic Principles

Static moments of nuclei are measured via the interaction of the nuclear charge distribution and magnetism with electromagnetic fields in its immediate surroundings. This can be the electromagnetic fields induced by the atomic electrons (described in Sect. 3.1) or the fields induced by the bulk electrons and first neighboring nuclei for nuclei implanted in a crystal, usually in combination with an external magnetic field (described in Sect. 3.2).

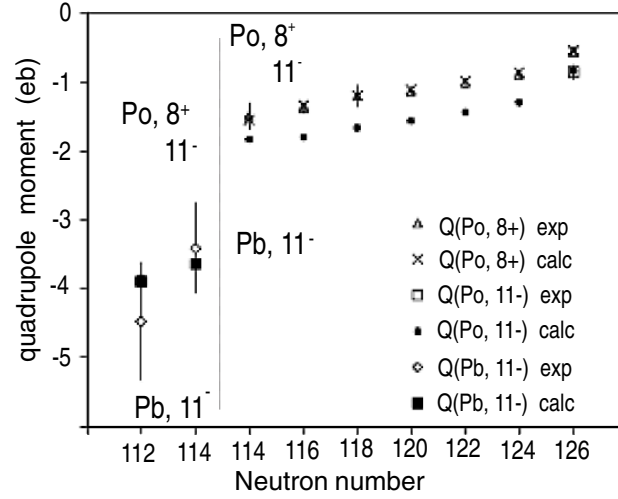


Plate 2. The increase in the absolute value of the quadrupole moments of isomers in the Pb region has been understood as due to a coupling of the valence particles with quadrupole excitations of the underlying core. The observed large quadrupole moments of the intruder isomers in the Pb isotopes are the first experimental evidence for the deformation of these intruder states

The interaction of the three lowest multipole-order nuclear moments with the surrounding fields, is an interaction between two scalars, two vectors and two tensors, respectively:

- The radial nuclear charge distribution influences the interaction with the charge of the shell electrons, inducing an overall shift of the electronic fine structure levels. This energy shift is different for each isotope (as well as for each isomer) and is called the “isotope (isomer) shift”.
- The magnetic dipole vector μ interacts with the dipole magnetic field $\mathbf{B}$. The energy of this interaction is defined by the scalar product of the two vectors.
- The electric quadrupole moment Q_2^n , a tensor of rank 2, with the spectroscopic moment Q_s being the zero order component of the tensor [20], interacts with the second order derivative of the electric field, being the tensor of the electric field gradient.

3.1 Electromagnetic Fields in an Atom: The Atomic Hyperfine Structure

Optical Isotope Shift

As pointed out in Sect. 2.1, information on isotopic differences between nuclear mean square charge radii is contained in the isotope shifts of optical

spectral lines [26]. Let A , A' and m_A , $m_{A'}$ be the mass numbers and atomic masses of the isotopes involved. Then for an atomic transition i the isotope shift, i.e. the difference between the optical transition frequencies of both isotopes, is given by

$$\delta\nu_i^{A,A'} = \nu_i^{A'} - \nu_i^A = F_i \delta\langle r^2 \rangle^{A,A'} + M_i \frac{m_A' - m_A}{m_A' m_A}. \quad (20)$$

This means that both the field shift (first term) and the mass shift (second term) are factorized into an electronic and a nuclear part. The knowledge of the electronic factors F_i (field shift constant) and M_i (mass shift constant) allows one to extract the quantity $\delta\langle r^2 \rangle$ of the nuclear charge distribution. These atomic parameters have to be calculated theoretically or semi-empirically.

For unstable isotopes high-resolution optical spectroscopy is a unique approach to get precise information on the nuclear charge radii, because it is sensitive enough to be performed on the minute quantities of (short-lived) radioactive atoms produced at accelerator facilities. Other techniques are suitable only for stable isotopes of which massive targets are available. Elastic electron scattering [27] even gives details of the charge distribution, and X-ray spectroscopy on muonic atoms [28] is dealing with systems for which the absolute shifts with respect to a point nucleus can be calculated. Thus both methods give absolute values of $\langle r^2 \rangle$ and not only differences. Eventually, the combination of absolute radii for stable isotopes and differences of radii for radioactive isotopes provides absolute radii for nuclei all over the range that is accessible to optical spectroscopy.

Typical orders of magnitude for the mass shift are between GHz (light elements, $Z \approx 10$) and 10 MHz (heavy elements, $Z \approx 80$) and for the field shift between 10 MHz (light elements, $Z \approx 10$) and 10 GHz (heavy elements, $Z \approx 80$). Laser spectroscopy methods providing a resolution better than 10^{-8} on the optical frequency scale of order 5×10^{14} Hz can indeed cope with such small effects. This is the reason why optical isotope shift data for $\delta\langle r^2 \rangle$ are usually very accurate. For a chain of isotopes nuclear structure effects are often very clearly observable. They are reflected by irregularities in the (relative) changes of radii, even though the (absolute) values of $\delta\langle r^2 \rangle$ may suffer from a poor knowledge of the atomic constants F_i and M_i .

An example of the isotopic changes of mean square charge radii deduced from optical isotope shifts is given in Fig. 3. This is a compilation of data for the elements close to the magic proton number $Z = 82$ with many isotopes below the magic neutron number $N = 126$. These were obtained from several experiments using different methods of laser spectroscopy, including the ones discussed in Sects. 4.1 and 4.2. According to (7) the curves not only reflect a regular increase of $\langle r^2 \rangle_{\text{sph}}$ with the neutron number, but in addition they show pronounced effects of nuclear deformation. This figure is extensively discussed in the review by Otten [7]. It nicely illustrates the staggering between near

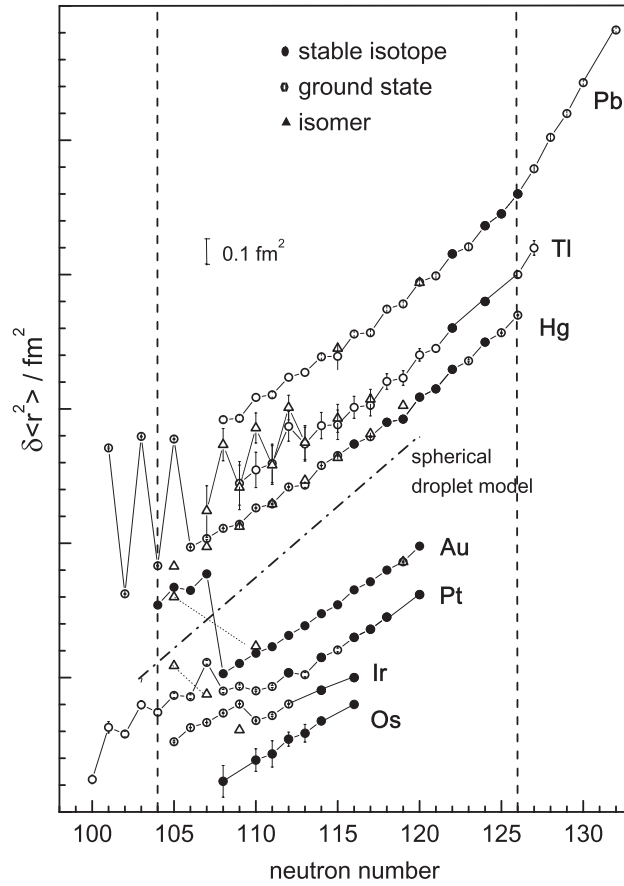


Plate 3. (a) Isotopic changes in mean square charge radii for elements in the Pb region ($Z \leq 82$). The large irregularities observed for the most neutron-deficient isotopes correspond to an onset of strong deformation

spherical oblate and strongly deformed prolate shapes of the neutron-deficient Au and Hg isotopes. Some of the phenomena, such as the different slopes of the Pb charge radii below and above $N = 126$, were explained only several years after the experiments had been performed [29].

Hyperfine Structure

Not only the radial distribution of the nuclear charge (monopole moment), but also the higher multipole electromagnetic moments of nuclei with a spin $I \neq 0$ influence the atomic energy levels. By interacting with the multipole fields of the shell electrons they cause an additional splitting called hyperfine structure. For all practical purposes it is sufficient to consider only the

magnetic dipole and the electric quadrupole interaction of the nucleus with the shell electrons. We have seen that nuclei with $I \geq 1/2$ possess a magnetic moment. On the other hand, the shell electrons in states with a total angular momentum $J \neq 0$ produce a magnetic field at the site of the nucleus. This gives a dipole interaction energy $W_D = -\boldsymbol{\mu} \cdot \mathbf{B}$. The spectroscopic quadrupole moment of a nucleus with $I \geq 1$ interacts with an electric field gradient produced by the shell electrons in a state with $J \geq 1$ according to $W_Q = eQ_s (\partial^2 V / \partial z^2)$.

For a particular atomic level characterized by the angular momentum quantum number J , the coupling with the nuclear spin I gives a new total angular momentum F according to the vector operator formula

$$\mathbf{F} = \mathbf{I} + \mathbf{J}, \quad (21)$$

meaning that $|I - J| \leq F \leq I + J$. The hyperfine interaction removes the degeneracy of the different F levels and produces a splitting into $2J + 1$ or $2I + 1$ hyperfine structure levels for $J < I$ and $J > I$, respectively (see Fig. 4).

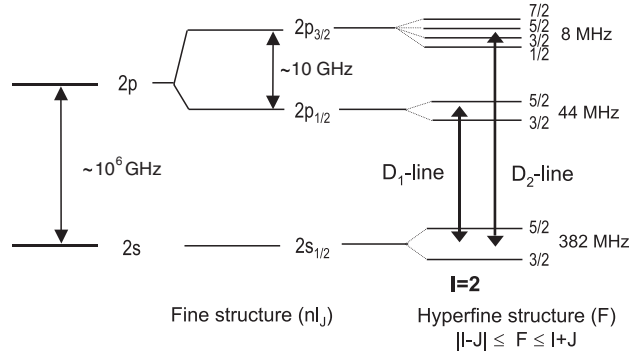


Plate 4. Example of the atomic fine and hyperfine structure of ^8Li . For free atoms the electron angular momentum J couples to the nuclear spin I , giving rise to the hyperfine structure levels F . The atomic transitions between the $^2S_{1/2}$ ground state to the first excited ^2P states of the Li atom are called the D_1 and D_2 lines (in analogy to Na)

Using quantum mechanical vector coupling rules one obtains an expression for the hyperfine structure energies of all F levels of a hyperfine structure multiplet with respect to the atomic fine structure level J (see, e.g., [30] and Fig. 4):

$$W_F = \frac{1}{2}AC + B \frac{\frac{3}{4}C(C+1) - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}, \quad (22a)$$

$$C = F(F+1) - I(I+1) - J(J+1). \quad (22b)$$

Within the multiplet, these energies depend on only two parameters, the magnetic dipole interaction constant A (not to be mistaken for the atomic mass number) and the electric quadrupole interaction constant B . Both these parameters contain as a nuclear part the respective nuclear moment and as an electronic part a quantity which is independent of the isotope, given a chemical element in a particular atomic state:

$$A = \mu_I B_e(0)/(IJ), \quad (23a)$$

$$B = eQ_s V_{zz}(0). \quad (23b)$$

$B_e(0)$ is the magnetic field and $V_{zz}(0)$ is the electric field gradient of the shell electron at the site of the nucleus.

The determination of nuclear moments from hyperfine structure is particularly appropriate for radioactive isotopes, because the electronic parts of (23a) and (23b), namely $B_e(0)$ and $V_{zz}(0)$, are usually known from independent measurements of moments and hyperfine structure on the stable isotope(s) of the same element.

We have implicitly assumed that A and B are the only unknown parameters in (22a). However, for isotopes far from stability very often not even the nuclear spin is known. It is easily seen that this basic quantity characterizing a nucleus is directly determined by a hyperfine structure measurement, either from the number of components (for $I < J$) or from their relative distances.

The size of a hyperfine structure depends very much on the coupling of spins and orbital angular momenta of the shell electrons in a particular atomic state. The large magnetic splitting produced by a single s electron varies from the order 100 MHz in light elements to about 50 GHz for the heaviest elements. The measurement of the quadrupole term requires $J \geq 1$ states with non-vanishing orbital angular momentum of at least one electron in an open shell. For light elements ($Z \simeq 20$), where not only the electric field gradients, but also the quadrupole moments are small, this effect can usually not be resolved. The possibility to measure quadrupole moments of light nuclei is one of the strengths of the β -ray detected nuclear magnetic resonance (β -NMR) technique (see Sect. 5.3).

3.2 Externally Applied Electromagnetic Fields

When a nucleus with spin I is implanted into a solid (or liquid) material, the interaction between the nuclear spin and its environment is no longer governed by the atomic electrons. For an atom imbedded in a dense medium, the interaction of the atomic nucleus with the electromagnetic fields induced by the medium is much stronger than the interaction with its atomic electrons. The lattice structure of the medium now plays a determining role [31]. This “hyperfine interaction” is observed in the response of the nuclear spin system to the internal electromagnetic fields of the medium, often in combination with externally applied (static or radio-frequency) magnetic fields.

This is why “nuclear hyperfine levels” are defined as the energy eigenstates of the nucleus in this combination of fields. The simplest case can just be an external static magnetic field.¹

The Magnetic Dipole Interaction

If atoms are implanted into a crystal with a cubic lattice structure (for example a body-centered-cubic BCC or face-centered-cubic FCC), no electric field gradient is induced if the nucleus is on a regular lattice position (substitutional site). In this case, the magnetic substates m_I of a nucleus with spin I remain degenerate.

The degeneracy can be lifted by applying a static magnetic field. This can be an externally applied field, typically of the order of a few hundred Gauss up to several Tesla (1 Tesla = 10^4 Gauss) or the internal hyperfine magnetic field of a host material, typically of the order of 10–100 Tesla.

The interaction of a nucleus with spin I immersed into a static magnetic field $\mathbf{B}_0$ is described by the Zeemann Hamiltonian [32]

$$H_B = -\frac{g_I \mu_N}{\hbar} \mathbf{I} \cdot \mathbf{B}_0 = -\omega_L I_z, \quad (24)$$

with $\omega_L = g_I \mu_N B_0 / \hbar$ being the Larmor frequency and g_I the nuclear g -factor. In an axis system with the z -axis parallel to the magnetic field (often called the laboratory (LAB) system), the magnetic substates m_I of the spin operator are eigenstates of the Zeeman Hamiltonian. The energies of the Zeeman levels are proportional to m_I (see Fig. 5(a)):

$$E_m = -\hbar \omega_L m_I. \quad (25)$$

This equidistant splitting of Zeeman levels is typically of the order of a few 100 kHz to several MHz, depending on the strength of the applied magnetic field and on the nuclear dipole strength, characterized by the g -factor. Several experimental methods have been developed to measure this Larmor frequency for radioactive nuclei and for their isomeric states. A very precise value of the nuclear g -factor can be deduced (see Sect. 5), provided the magnetic field at the site of the implanted nucleus is known accurately enough.

The Electric Quadrupole Interaction

In a material with a non-cubic lattice structure, the implanted nuclei interact with an electric field gradient (EFG) induced by the atoms and electrons in

¹In atomic systems one usually distinguishes between “hyperfine structure” arising from the interaction of the nucleus with the shell electrons and “Zeeman effect” reflecting the much weaker interaction with external fields.

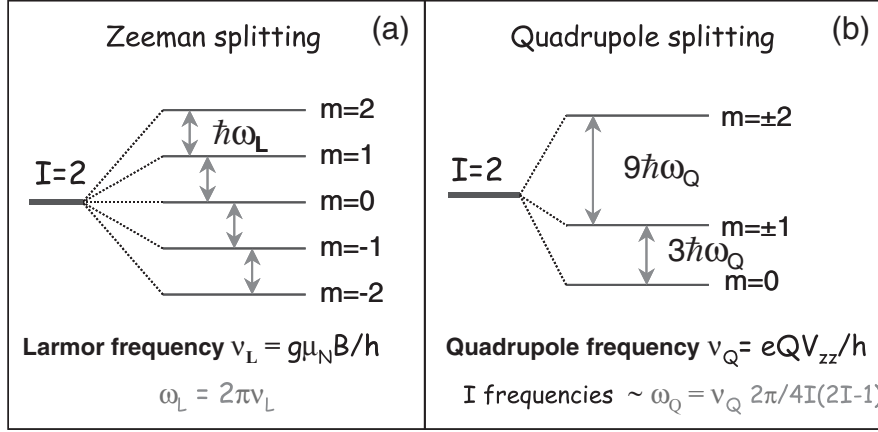


Plate 5. (a) Hyperfine levels of a nucleus immersed into a static magnetic field. The Zeeman splitting is equidistant and proportional to the Larmor frequency ν_L . (b) Hyperfine levels of a nucleus implanted in a crystal with an electric field gradient. The nuclear level splitting is not equidistant, and proportional to the quadrupole frequency ν_Q .

the first atomic layers around the implanted impurity. For an impurity on a substitutional (regular) lattice position, the induced EFG has the symmetry of the lattice structure (e.g., in a material with a hexagonal-close-packed (HCP) lattice structure, the nucleus interacts with an axially symmetric EFG with its symmetry axis along the crystal c-axis, the 0001 axis). If the implanted impurity ends up at an interstitial or defect-associated lattice position in the crystal, the symmetry of the induced EFG can be different from the lattice symmetry.

The strength of the EFG is the same for all isotopes of a particular element implanted at similar lattice positions in a crystal. Thus, if the quadrupole frequency $\nu_Q = eQ_s V_{zz}/h$ is measured for several isotopes of one element in a crystal, the ratio of these frequencies gives the ratio of their quadrupole moments. To deduce an absolute value for the spectroscopic quadrupole moment of each isotope, the EFG strength V_{zz} needs to be determined. This can be done e.g., by using a known EFG (measured by another method or calculated). Nowadays, it is possible to perform ab-initio calculations of the EFG induced on whatever impurity in any type of crystal. It has been shown that most of such calculations are accurate to about 10%, provided the lattice parameters are known well enough [33]. This is an important step forward in quadrupole moment measurements. For example, such calculations have revealed that some earlier deduced quadrupole moments in Fe isotopes are wrong, due to a wrong determination of the EFG in some crystals [34].

If a single crystal is used as implantation host, the orientation of the crystal symmetry axis (and thus of the EFG principal axis) is the same over the whole crystal volume, hence for each implanted nucleus the EFG strength and orientation are the same. If a polycrystalline material is used as implantation host, each implanted nucleus will interact with an EFG that has the same strength, but with a random orientation of the principal axis. In this case an integration over all possible principal axis directions with respect to a chosen reference frame needs to be made.

The Hamiltonian for a nucleus with spin I submitted to an EFG is easiest expressed with respect to the Principal Axis System (PAS) which has its z -axis along the c -axis of the crystal. We restrict the discussion here to crystals with an axially symmetric EFG (asymmetry parameter $\eta = 0$) along the z -axis of the PAS [32]:

$$H_Q = \frac{\omega_Q}{\hbar} (3I_z^2 - I^2). \quad (26)$$

The magnetic states m_I of the spin operator are good eigenstates of the quadrupole interaction Hamiltonian. Their energy is degenerate in m_I^2 (see Fig. 5(b)):

$$E_m = \hbar\omega_Q [3m_I^2 - I(I+1)], \quad (27)$$

and proportional to the quadrupole coupling constant $\omega_Q = \frac{eQ_s V_{zz}}{4I(2I-1)\hbar}$ which is related to the quadrupole frequency $\nu_Q = 4I(2I-1)\omega_Q/2\pi$.

The m_I levels of a nucleus interacting with an EFG are not equidistant. Several transition frequencies occur, which vary again in the range of 10 kHz up to several MHz, depending on the strength of the EFG and on the nuclear spectroscopic quadrupole moment. For nuclei with a spin $I < 1$ no quadrupole splitting occurs, corresponding to the fact that the expectation value (16) defining the spectroscopic quadrupole moment vanishes. Several experimental methods have been developed to measure the quadrupole frequency for radioactive nuclei and for their isomeric states. The typical accuracy on quadrupole moments deduced from such measured frequencies depends on how well the EFG at the site of the nuclei in the crystal lattice is known. It can vary from 1% to 15%.

Combined Static Interactions

When both a magnetic field and an electric field gradient interact with the implanted nuclei, the energy of the nuclear hyperfine states can be calculated analytically only if these interactions are axially symmetric and aligned with each other. In this case the energy of the magnetic substates m_I is simply the sum of the dipole and quadrupole energies given in (25) and (27). When plotting these energies as a function of the magnetic field strength (see Fig. 6(a)), one obtains at equidistant positions a crossing of nuclear quantum levels [35].

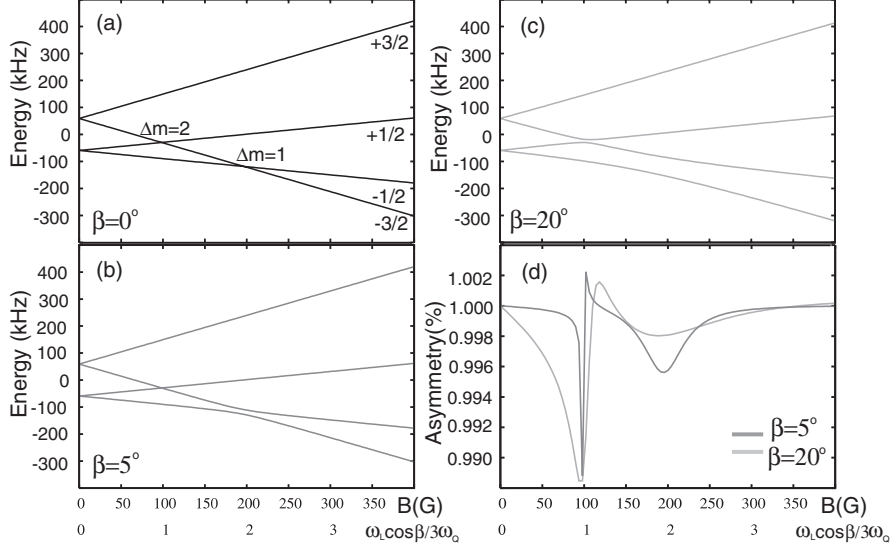


Plate 6. Nuclear hyperfine levels of a nucleus with spin $I = 3/2$ submitted to a combined static magnetic interaction and an axially symmetric quadrupole interaction: (a) for collinear interactions, $\beta = 0^\circ$; (b) and (c) for non-collinear interactions with $\beta = 5^\circ$ and $\beta = 20^\circ$, respectively. Crossing or mixing of hyperfine levels occurs at well-defined values for the ratio of the involved interactions frequencies, if $\frac{\omega_L}{2\pi\nu_Q} = \frac{3(m+m')\cos\beta}{4I(2I-1)}$. (d) At these positions, resonances are observed in the decay angular distribution of oriented radioactive nuclei, from which the nuclear spin and moments can be deduced; cf. Plate 13 in the Colour Supplement

If the two interactions are non-collinear, one needs to diagonalize the Hamiltonian of the combined interaction in order to determine its eigenvalues. The combined interaction Hamiltonian can be expressed with respect to the LAB reference frame which is chosen with its z -axis parallel to the static field. Then the quadrupole interaction Hamiltonian needs to be described with respect to this frame as well (as in [36]). However, to get a better insight into the physical behavior of a combined interaction system, it is easiest to describe the quadrupole interaction with respect to the PAS reference frame which forms an angle (β, γ) with respect to the LAB frame (as in [37]). In the latter case the Hamiltonian is given by

$$H_{combined} = -\omega_L I_z \cos \beta + \omega_L I_z \sin \beta + \frac{\omega_Q}{\hbar} (3I_z^2 - I^2). \quad (28)$$

Diagonalization of this matrix gives the “mixed” eigenstates with respect to the chosen reference frame and the energy levels of these “mixed” quantum states. In Fig. 6(b) the energy levels of the non-collinear combined interaction are calculated for a misalignment angle $\beta = 15^\circ$ of the crystal c -axis with respect to the LAB frame. One can see that near the crossing points, the nuclear

levels now repel each other (a phenomenon known from atomic physics as anti-crossing). For small misalignment angles, an analytical expression for these anti-crossing energy levels can be deduced. This is done in a two-level approximation using first order perturbation theory [37, 38]. The calculated eigenstates are a mixture of the “unperturbed” m_I substates. Therefore they are called “level-mixed states”. It can be shown that in the level mixing region the populations of the two levels are equalized. This gives rise to a resonant change in the spin orientation of a spin-oriented ensemble, which can be observed as resonances in the angular distribution of γ - or β -decaying oriented nuclei (see Sect. 5). A simulation of such resonances for spin-polarized beams from a projectile-fragmentation reaction is demonstrated in Fig. 6(a). Since the eighties, these level-mixed states have been used for developing new methods to measure moments of exotic nuclei [39, 40, 41, 42, 43].

4 Methods Based on the Measurement of Atomic Hyperfine Structure

The beginning of research on atomic hyperfine structure and isotope shift dates back to the thirties of the past century, as reviewed recently in [44] and [45]. It was discovered that nuclei (with an odd number of protons or neutrons) possess intrinsic angular momentum, called nuclear spin, and that this is associated with electromagnetic properties known as magnetic dipole moment and electric quadrupole moment. The tool for such studies was a combination of an electric discharge lamp and a high resolution spectrometer arrangement consisting of a conventional prism or grating spectroscopy and a Fabry-Pérot interferometer. Of course, the objects of investigation were stable isotopes of those elements that happened to show large effects.

This type of optical spectroscopy had a serious limitation. Even with the highest-resolution spectrometer one was unable to overcome the resolution limit given by the light source. The spectral lines are broadened by the Doppler effect from the thermal movement of the atoms. Given the Maxwell-Boltzmann distribution of velocities for a gas temperature T , the optical frequency ν_0 and the atomic mass m , one obtains a Gaussian-shaped spectral line with the Doppler width (FWHM)

$$\delta\nu_D = \frac{\nu_0}{c} \sqrt{8kT \ln 2 / m}. \quad (29)$$

This width is typically of the order of 1 GHz, i.e. in many cases larger than the hyperfine structure and isotope shift effects described in Sect. 3.1. Therefore, accurate hyperfine structure measurements became a domain of radio-frequency and microwave spectroscopy. Instead of measuring extremely small effects on the optical frequency scale, one used the possibility to induce (magnetic dipole) transitions between the hyperfine structure levels. There are a

number of methods how these transitions can be detected. Detailed expositions of the principles and applications of these early (pre-laser) methods of high-resolution spectroscopy are given in the literature, e.g., [30, 46, 47, 48]. Brief descriptions of the different techniques with references for further reading can be found in earlier compilations of nuclear moments [49].

When it became possible in the seventies to produce appreciable quantities of radioactive isotopes in the form of mass-separated beams, lasers were about to revolutionize the classical field of atomic spectroscopy. Laser spectroscopy methods were invented to overcome the resolution limit given by the Doppler broadening of spectral lines. Since then, many experiments were performed with an optical resolution close to the natural linewidth. This width is given by the mean lifetime τ of the optically excited state, via the Heisenberg uncertainty principle:

$$\delta\nu_n = 1/(2\pi\tau), \quad (30)$$

yielding the order of magnitude of 10 MHz for strong optical transitions ($\tau \approx 10^{-8}$ s).

Again, it is impossible in this lecture to discuss all the methods of “sub-Doppler” or “Doppler-free” spectroscopy, the most important of them being characterized by the expressions “saturation spectroscopy”, “two-photon spectroscopy” and “atomic beam spectroscopy” (see [50, 51]). For sensitivity reasons only the latter has played a notable role in investigations of unstable isotopes.

Here, we restrict our discussion to the outstanding combination of an isotope production scheme and a method of laser spectroscopy that has been most successful in giving access to the properties of unstable isotopes in many regions of the nuclear chart. This is the combination of isotope separation on-line (ISOL) concept [52]) and collinear laser spectroscopy [53], which has been demonstrated for the first time by measurements on a few neutron-rich Cs [54] and Rb [55] isotopes.

4.1 Collinear Laser Spectroscopy

On-line Isotope Separation

Over three decades ISOLDE at CERN [56] has been a unique facility, delivering intense beams of radioactive isotopes, particularly of short-lived ones far from stability. The idea of ISOLDE is to get a multitude of products from reactions of high-energy (typically 1 GeV) protons in a thick (~ 20 cm) target. The products diffuse out of the hot target material and are guided to the ion source of an electromagnetic mass separator. Such a separation according to the mass gives beams of isobars, i.e. mixtures of isotopes of neighboring elements with different Z , but with the same mass number A . Different diffusion times of atoms in the target material may give rise to an element-selectivity of the system. A further tool for controlling the element composition of the

mass-separated beam is the ion source. Depending on what element should be ionized, one has the choice between thermal surface ionization or a plasma ion source. Moreover, in recent years also a highly element-selective laser ion source [52, 57] has been employed very successfully.

Thus from an appropriate combination of a target and an ion source, one can obtain rather pure beams of the isotopes of a single chemical element. The beam energy is typically 60 keV, and electrostatic lenses and deflectors are used to guide the beam to a number of experimental stations.

Principles of Collinear Laser Spectroscopy

It is a gift of nature that beams such as those delivered by ISOLDE are suitable for a conceptually very simple method of high-resolution laser spectroscopy: Take the given ion beam and superimpose it on a narrow-band laser beam by an electric deflector. If the laser is in resonance with an atomic transition from the ground state (resonance line), excitation takes place, which can be detected by observing the fluorescence light from the decay back to the ground state or to a third state. Spectroscopic information can be obtained by sweeping the laser frequency across the resonance(s).

The line width in this particular situation [58, 59] can be calculated from the distribution of velocity components in the direction of the beam, as shown in (29) for the Doppler width of a thermal ensemble of atoms. Irrespective of the details of this distribution (which depend very much on the conditions in the ion source) one starts with a kinetic energy spread δE of the ions of mass m leaving the source (see [53]). This ensemble of ions is exposed to an electrostatic acceleration field, the key point being that all ions get the same increase eU in kinetic energy, while their initial energy spread remains unchanged. For the velocities along the beam direction we can calculate the velocity spread using

$$\delta E = \delta \left(\frac{1}{2} m v^2 \right) = m v \delta v. \quad (31)$$

As $\delta E = \text{const}$ this means that the velocity spread decreases while the velocity increases. Now the Doppler shift for co-propagating or counter-propagating beams is $\Delta \nu_D = \nu_0 \beta = \nu_0 (v/c)$ and the Doppler width is $\delta \nu_D = \nu_0 \delta \beta = \nu_0 (\delta v/c)$. This gives

$$\delta E = \frac{m c^2}{\nu_0^2} \Delta \nu_D \delta \nu_D, \quad (32)$$

i.e. the product of the Doppler shift and the Doppler width is a constant of motion. With the relation $v = \sqrt{2eU/m}$ for the velocity v as a function of the acceleration voltage U , one obtains the Doppler width

$$\delta \nu_D = \nu_0 \frac{\delta E}{\sqrt{2eU m c^2}}. \quad (33)$$

Realistic numbers inserted in (33) yield a Doppler width of about 10 MHz for a beam-energy spread of about 1 eV. Here the surprising fact is that a resolution close to the natural linewidth is inherently present in collinear laser spectroscopy with an ion beam formed by electrostatic acceleration.

For spectroscopy on a limited number of accelerator-produced atoms this has another important consequence. At resonance all atoms in the beam participate in the interaction with the laser light and contribute to the signal. Thus the high resolution is achieved with the maximum possible excitation efficiency. This is in contrast to the other Doppler-free techniques which are based on the selection of a certain velocity class of atoms.

So far one has dealt with singly-charged atomic ions. However, from the spectroscopy point of view it is often preferable to use neutral atoms, because their excited states are at lower energy and can be reached more easily with the radiation of narrow-band cw lasers². Neutral-atom beams are readily obtained by charge-exchange reactions. The primary ion beam (ions X^+) passes through a metal vapor cell containing neutral atoms (Y), and the reaction



occurs with a large cross-section of the order 10^{-14} cm² if the process is resonant in energy, meaning that the energy defect $|\Delta E|$ is below about 1 eV. In this situation the kinetic cross-section for momentum-changing collisions is negligible, and consequently the original beam quality and the kinetic energy spread are preserved [59, 60].

It is not only technically most convenient to use alkali metal vapors as a charge-exchange medium. Also the ionization energy of alkali atoms (4–5 eV) with the resonance condition $\Delta E \approx 0$ offers an interesting feature for laser spectroscopy: Atoms with inaccessible resonance lines from the ground state often have long-lived (metastable) states in this energy range, i.e. 4–5 eV below the ionization threshold. These are selectively populated and can be used for excitation with visible laser light [61]. The most remarkable examples for this kind of spectroscopy are the noble gases from neon to radon (see below).

Now all ingredients of collinear laser spectroscopy have been presented. At this point it is important also to consider the measuring procedure. Scanning the laser frequency over the hyperfine structure pattern of a particular isotope will be straightforward. However, the determination of an isotope shift is based on the measurement of a frequency difference between two isotopes of masses m_A and $m_{A'}$ which also involves a difference in the Doppler shifts. The Doppler-shifted frequencies for co-propagating (contra-propagating) beams

²The experiments require continuous wave (cw) laser radiation with a spectral linewidth and frequency stability of about 1 MHz. For the whole range of visible light (400–800 nm) this can be obtained from dye lasers. Ultraviolet wavelengths down to about 270 nm can be reached by frequency doubling. For the near infrared Ti:Sapphire lasers or inexpensive semiconductor diode lasers are most suitable.

are given by

$$\nu = \nu_L \frac{1 \mp \beta}{\sqrt{1 - \beta^2}}, \quad (35)$$

where the laser frequency is ν_L . Thus the calculation of a frequency difference between the resonances of two isotopes A and A' requires the knowledge of β_A and $\beta_{A'}$. With

$$\beta = \sqrt{1 - \frac{m^2 c^4}{(eU + mc^2)^2}} \approx \sqrt{\frac{2eU}{mc^2}}, \quad (36)$$

this means one has to accurately know the isotope masses m_A , $m_{A'}$ and the acceleration voltages U_A , $U_{A'}$.

As voltages or Doppler shifts can anyway not be eliminated from the evaluation of spectroscopic results, it is much more elegant to rely exclusively on voltage measurements and not to tune the laser frequency, but the Doppler-shifted frequency seen by the atoms [60]. This ‘‘Doppler-tuning’’ is easily done by post-accelerating or -decelerating the beam with the help of an electrical potential applied to the excitation region (in case of ions) or to the charge-exchange cell (in case of neutral atoms). At a fixed and stabilized laser frequency the spectra are taken as a function of a voltage in the few kV region which has to be controlled and scanned with a resolution better than 100 mV. The data evaluation on a frequency scale is based on (35) and (36), applied for different isotopes whose resonances appear for the same laser frequency at different values of β . One has to know the primary acceleration voltage (≈ 60 kV) and the post-acceleration voltage to a precision better than 10^{-4} , and the optical transition frequency [62] and the masses [63] of the isotopes of interest to about the tabulated accuracies.

For an illustration of the experimental setup we refer to the left-hand part of Fig. 7. The example of an experimental spectrum is shown in Fig. 8 for a measurement of dysprosium isotopes [64]. It includes the hyperfine structure of ^{151}Dy and the isotope shift plus Doppler shift between ^{151}Dy and the two stable reference isotopes ^{156}Dy and ^{158}Dy which exhibit no hyperfine structure.

Fluorescence Detection

The observation of fluorescence photons from the spontaneous decay of the excited state is a standard technique of detecting optical resonance excitation in laser spectroscopy experiments. For the weak beams of radioactive isotopes investigated by collinear laser spectroscopy one has to rely on single-photon counting using photomultipliers. As the laser frequency is very selective for a particular element, this ‘‘optical’’ detection is insensitive to contaminations of the beam by isobars. However, as the light-emitting section of the beam is pencil-shaped it is rather difficult to collect and detect the photons with

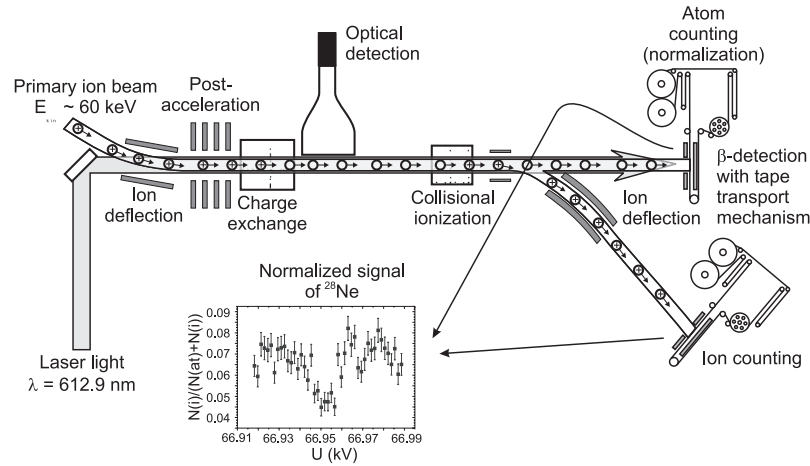


Plate 7. Experimental setup for collinear laser spectroscopy with collisional ionization and radioactivity detection

high efficiency. The main problem is the background of laser light passing through the interaction region and scattering into the direction of observation from diaphragms and other components of the apparatus, in particular the entrance and exit windows for the laser beam. A typical number for the efficiency of detecting an emitted fluorescence photon is about 10^{-3} with a background of some 1000 counts per second from scattered laser photons. Additional background can be produced by excitation in collisions of the 60 keV ions/atoms with the charge exchange medium or the rest gas. Finally, also the radioactivity collected in the apparatus may cause a considerable background level.

It is obvious that the sensitivity limit – i.e. the minimum beam intensity required for a measurement – depends on many parameters. As a rule of thumb, one can assume limits between 10^5 and 10^7 atoms/s, depending on the complexity of the atomic spectrum (see [64]). Only in the exceptional cases of so-called two-level systems, where the atoms can be excited many times without losses by a “branched decay” of the excited state, about 10^4 atoms/s have been sufficient.

Methods of background suppression, such as blocking of the laser light by filters and detecting on a different transition, have been favorably used for a few experiments. Their common disadvantage is a reduction not only of the background, but also of the signal. Only a very recent development promises to solve this problem more rigorously. The laser spectroscopy group at Jyväskylä has introduced a quadrupole cooler [65, 66, 67] to get at the same time a beam of low energy spread and a bunched time structure. Then the bunched observation is used to reject very efficiently the background from stray-light and radioactivity which is constant in time.

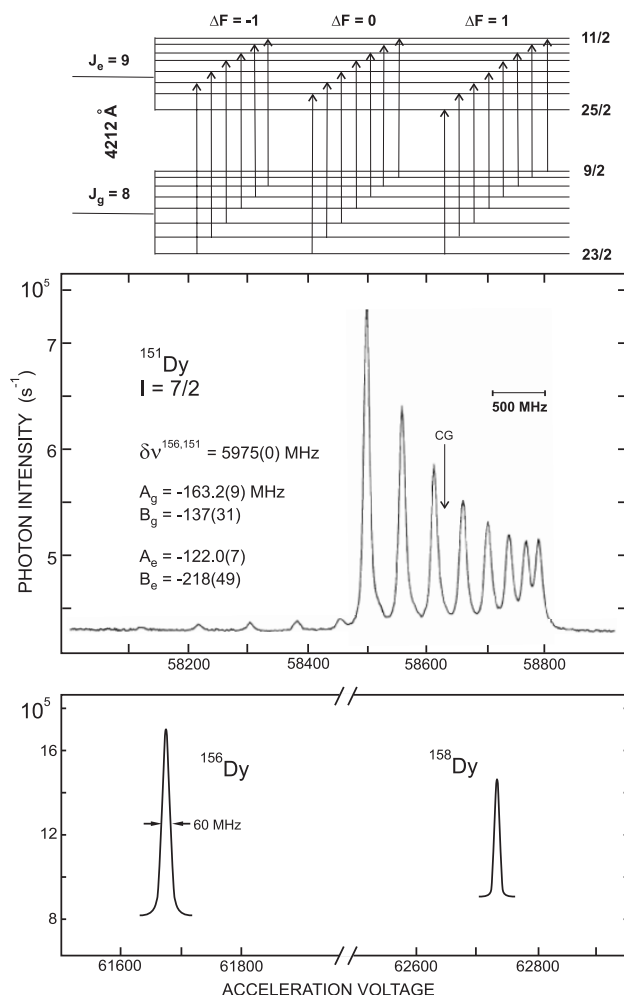


Plate 8. Spectra of ^{151}Dy , ^{156}Dy and ^{158}Dy obtained by collinear laser spectroscopy in the 421.2 nm line. The voltage scale translates into the frequency scale via (35) and (36)

Non-optical Detection and High Sensitivity

Alternatives to an improved sensitivity of optical detection have been sought in unconventional detection schemes. The idea is to introduce processes that can distinguish between atoms having or not having interacted with the laser light, and then to count either ions or atoms, or to detect their radioactive decay. The great advantage of the detection of massive particles is an efficiency close to 100% without the problem of background from the laser light. 60 keV ions or neutral atoms impinging on a surface can be counted by

secondary electron multiplication. For β -radioactivity one uses plastic scintillators with photomultipliers, and for α -particles silicon detectors which can even discriminate sharply between different decay energies.

The key to a non-optical detection scheme is efficient optical pumping, i.e. the transfer of population from the initial to a final atomic state via repeated excitation and decay. For the detection one needs a process discriminating between these states. Several such schemes were proposed and have been used very successfully in experiments. They are based on the state dependence of cross-sections of charge-changing collisions [68] or on the particular properties of the radioactive β -decay.³ The possible schemes depend very much on the properties of the atomic spectra. This is why they are rather specific for different classes of elements:

- The singly-charged ions of alkaline earth and related elements have $^2S_{1/2}$ ground states and metastable $^2D_{3/2,5/2}$ states below the excited $^2P_{1/2,3/2}$ states which are reached from the ground state. The transfer of population from the ground state to these metastable states can be detected by the use of a charge-exchange reaction that fulfills the resonance condition for the metastable state [69]. Then the neutralized fraction of the beam is detected as a function of the laser frequency. Experiments using this technique were performed mainly on Sr [70, 71] and Ca [72] isotopes.
- A somehow complementary process gives access to strongly bound neutral atoms, such as noble gases, for which the first excited states above the ground state can hardly be reached by cw lasers. The charge exchange according to (34) populates a metastable state about 4–5 eV below the ionization limit whose cross-section for being ionized in collisions is much larger than of the low-lying ground state [73]. Thus optical pumping from the metastable to the ground state gives rise to a decrease of the ionization rate when the (neutral) beam passes through the section of a thin gas atmosphere (“gas target”). This method has been applied very successfully in investigations of the unstable isotopes of all noble gas elements [74, 75, 76, 77]. This will be discussed below in more detail for the example of a recent experiment on short-lived Ne isotopes.
- The excitation by circularly polarized light transfers angular momentum from the light field to the atomic system. This can be exploited for polarizing atoms by optical pumping between the Zeeman components of a particular level. If a nuclear spin $I \neq 0$ is involved and coupled with J to a hyperfine structure level F , also the nuclear spin is polarized. Due to parity violation, the β -emission occurs preferentially along or against the spin direction. With the atoms implanted into a suitable host crystal, the β -decay asymmetry can serve as a detector for optical excitation. This

³The conceptually related method of resonance ionization spectroscopy (RIS), mostly used on thermal atomic samples or beams, will be treated separately in Sect. 4.2.

will be discussed in the last part of this section, while applications using the β -NMR technique will be presented in Sect. 5.3.

Ultra-sensitive Spectroscopy on Short-lived Noble Gas Isotopes

Collinear laser spectroscopy on the noble gas elements has been brought up to a sensitivity level at which nearly all known isotopes are within reach. The most recent example of an experiment on Ne isotopes may serve to explain the technique in more detail and to present some results.

The essential parts of the experimental setup are shown in Fig. 7 [9]. The 60 keV Ne⁺ beams of different isotopes from ISOLDE are neutralized by charge exchange in a Na-vapor cell, whereby the metastable $J = 2$ level of the configuration $2p^53s$ is preferentially populated. The atomic energy levels involved in this process and in the optical pumping cycle are displayed in Fig. 9. Laser light of the resonance wavelength 614.3 nm excites the metastable atoms to a $J = 2$ state of the configuration $2p^53p$. From there the decay to the $2p^6\ ^1S_0$ ground state occurs via the intermediate $2p^53s$ ($J = 1$) states. The detection makes use of the large cross-section for collisional ionization from the metastable level. The neutral beam passes through a thin differentially pumped Cl₂ gas target, with the resulting beams of singly-charged ions and neutral atoms being separated by an electrostatic deflector. Both these beams impinge on moveable tapes surrounded by scintillators which detect the β -decay. Optical resonance in a scan of the voltage applied to the

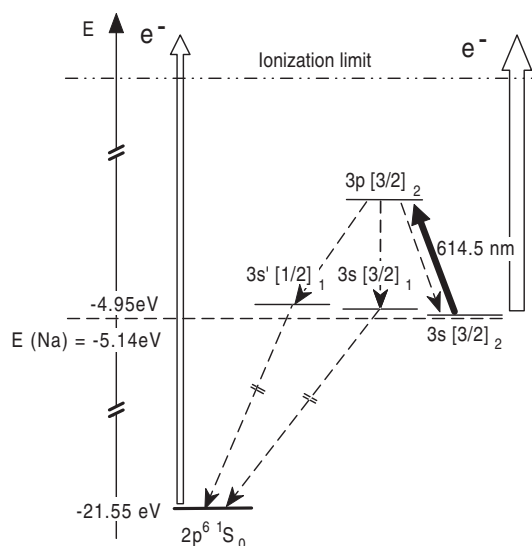


Plate 9. Atomic energy levels involved in the neutralization of a Ne^+ beam, optical excitation and decay to the ground state, and state-selective collisional ionization

charge-exchange cell reduces the ionization rate and gives rise to a drop of the ratio of both count rates.

For the weak beams of short-lived Ne isotopes the detection of radioactivity gives a very efficient rejection of the dominant background of stable-isobar beams. In this way a measurement of the resonance of the neutron-rich isotope ^{28}Ne was possible with the very low intensity of only 40 atoms/s.

In the general discussion of Sect. 4.1 it was pointed out that accurate voltage measurements are important for determining isotope shifts. This statement is too weak for isotopes of an element as light as neon because of the huge Doppler shifts. The accuracy of 1 MHz needed for the extraction of the small (~ 10 MHz) field shifts from the Doppler-shifted resonance positions requires a knowledge of the 60 keV beam energy to better than 1 eV. Such accurate voltage measurements are elaborate and not even sufficient, because for a plasma ion source the potential at which the ions are created is not well known. In the described experiment on Ne isotopes the required accuracy of the actual beam energy was obtained by an absolute Doppler shift measurement [78]. Two excitation lines from the metastable state have transition frequencies differing by about twice the Doppler shift for a beam energy of 60 keV. Both these transitions can be coincidentally excited by one laser beam which is retro-reflected at the end of the apparatus. If ν_1 is the frequency of the transition induced by the collinear and ν_2 the one induced by the anti-collinear beam, the condition for the laser frequency

$$\nu_L = \sqrt{\nu_1 \nu_2} \quad (37)$$

follows from (35). Given this condition, the beam energy becomes⁴

$$eU = mc^2 \frac{(\sqrt{\nu_1} - \sqrt{\nu_2})^2}{2\sqrt{\nu_1 \nu_2}}. \quad (38)$$

From the accurately known wave numbers the beam energy corresponding to a measured acceleration voltage can be calibrated with an accuracy of about 0.5 eV.

Some typical hyperfine structure spectra measured for three odd- A isotopes of neon are shown in Fig. 10. These data yield the nuclear spins, the magnetic moments and, for $I > 1/2$, also the quadrupole moments [79]. The experiment, including measurements of the isotope shifts, has been performed on all unstable Ne isotopes (except ^{27}Ne) from ^{17}Ne at the proton drip line to the neutron-rich ^{28}Ne . It addresses several interesting physics problems:

(i) ^{17}Ne , a nucleus at the proton drip line, has a weakly bound proton pair in the sd shell. Therefore, from theoretical considerations and from nuclear reaction cross-sections a proton halo has been postulated for this nucleus. The corresponding radial extension of the wave function of these “halo” protons

⁴For deducing this simple formula from (35) one has to use the relativistically correct version of (36).

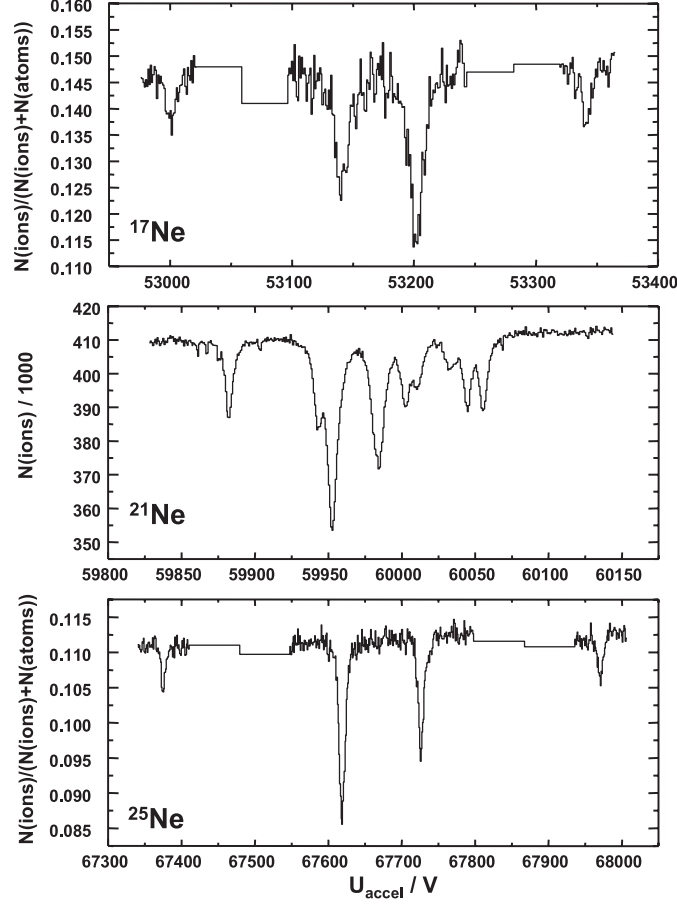


Plate 10. Hyperfine structure spectra of three odd- A Ne isotopes

would affect the radial nuclear charge distribution reflected in $\langle r^2 \rangle$ and can be probed very sensitively by the isotope shift.

(ii) Although the isotope shift is by far dominated by the mass shift, it is possible to extract field shifts by using muonic X-ray results on the three stable isotopes for calibration. Thus the dependence of the radii on the neutron number can be determined rather safely.

(iii) The measured moments [79] and radii give a valuable basis for comparisons with predictions from mean-field and shell-model calculations for the sd -shell nuclei and for ^{17}Ne where a $p_{1/2}$ neutron hole is coupled to sd -shell proton pair.

(iv) ^{17}Ne together with ^{17}N form one of the very few isospin $T = 3/2$ nuclear mirror pairs for which the magnetic moments are known [79]. Complementary information on a possible halo structure can thus be obtained

by investigating the mirror symmetry in the contributions of the *sd*-shell proton pair of ^{17}Ne and the corresponding neutron pair of ^{17}N to the respective magnetic moments.

Polarization by Optical Pumping

Polarization of a fast beam by optical pumping was introduced for the β -asymmetry detection of optical resonance in collinear laser spectroscopy [80]. However, it has turned out that most applications took advantage of the additional option to perform nuclear magnetic resonance spectroscopy with β -asymmetry detection (β -NMR) on a sample obtained by implantation of the polarized beam into a suitable crystal lattice. Whatever is the particular goal of such an experiment, it is important to achieve a high degree of nuclear polarization.

The most suitable transitions for polarization by optical pumping are found in alkali-like systems such as neutral alkali atoms or singly-charged alkaline earth ions. To be explicit, we take the example of Na to be optically pumped in the yellow D_1 or D_2 resonance line. The hyperfine structure for such a transition has been shown in Fig. 4. The ground state, $3s\ ^2S_{1/2}$, is split into two hyperfine structure components with $F = I + 1/2$ and $F = I - 1/2$. The excited states are $3p\ ^2P_{1/2}$ and $3p\ ^2P_{3/2}$, again with the corresponding hyperfine structure components. All hyperfine structure levels are $(2F + 1)$ -fold degenerate with respect to the M_F quantum number. If a weak magnetic field defines the quantization axis in the direction of the atomic and the laser beam, each absorption of a circularly polarized photon introduces one unit of angular momentum in the atomic system. This can be expressed by the selection rule

$$\Delta M_F = \pm 1 \quad \text{for } \sigma^\pm \text{ light,} \quad (39)$$

with σ^+ and σ^- being the conventional notations for the circular polarization of the light with respect to the direction of the magnetic field.

Figure 11 shows the optical pumping scheme for the example of a nuclear spin $I = 1$ which corresponds to the situation in ^{23}Na . Repeated absorption and spontaneous emission of photons results in an accumulation of the atoms in one of the extreme M_F states for which the total angular momentum $\mathbf{F} = \mathbf{J} + \mathbf{I}$, for an S state just composed of the electron spin and the nuclear spin, is polarized.

Now one has to consider the geometrical implications of a realistic experiment. Figure 12 displays the experimental setup which up to the optical excitation region is identical to the apparatus for collinear laser spectroscopy shown in Fig. 7. The optical pumping produces a longitudinal polarization of the beam. The implantation and detection occurs in a transversal magnetic field of about 0.5 Tesla which is used for the NMR experiments. This field is strong enough to decouple the electronic and nuclear spins of the alkali-like system. For a rotation of the spin from the longitudinal to the transversal

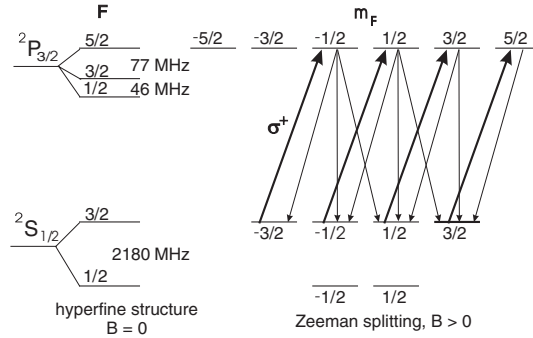


Plate 11. Optical pumping within the hyperfine structure Zeeman levels for polarization of the nuclear spin. The example shows the case of $I = 1$ for the case of ^{28}Na

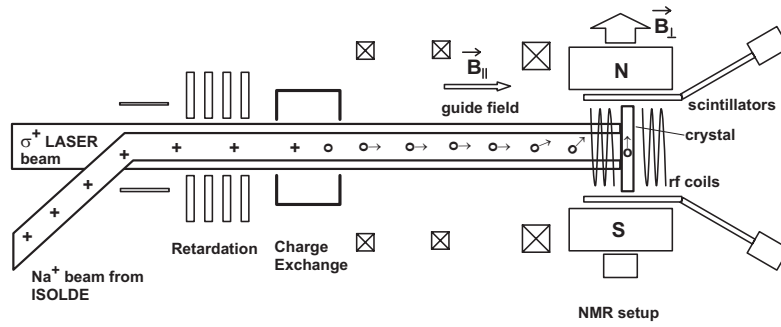


Plate 12. Experimental setup for in-beam optical polarization and β -NMR spectroscopy

direction, it is sufficient that the magnetic field changes slowly from longitudinal to transversal. The spins follow adiabatically, provided the rotation is much slower than the Larmor precession. This condition can be fulfilled only for the coupled system for which the g -factor g_F is determined by the electron. Therefore the rotation has to take place in an intermediate field region [81], before the electronic and nuclear spin are decoupled.

The goal of polarizing the beam is to detect optical resonance by the asymmetry in the angular distribution of emitted β -decay electrons or positrons. For this purpose the beam is implanted into a suitable crystal surrounded by scintillation detectors. The polarization has to be preserved during the lifetime of the nuclei under consideration, which for typical spin relaxation times means that the method is suitable for isotopes with half-lives shorter than a few seconds.

The application of this polarization technique for β -NMR spectroscopy will be discussed in Sect. 5.3. This requires further consideration of the crystal properties. Implantation depths of low-mass 60 keV atoms are about 100 nm.

For clean crystals this is sufficiently deep to reach good lattice sites of the bulk material, which is of particular importance for studies of the quadrupole interaction.

4.2 Resonance Ionization Spectroscopy and Laser Ion Source

As was shown above, collinear laser spectroscopy is a rather flexible and widely applicable tool for experiments on unstable isotopes. Due to its high resolution this method is capable of dealing with the small hyperfine structure and isotope shift effects of lighter atomic systems. The sensitivity, however, depends very much on the particular properties of the atomic spectrum. This sensitivity problem is related to the fact that photon detection is inefficient and the particle detection schemes discussed so far are restricted to special groups of elements. A general method allowing efficient ion detection with very good background suppression is the stepwise excitation of neutral atoms to the ionization continuum. For collinear-beam spectroscopy, this was demonstrated in an experiment on neutron-deficient Yb isotopes [82], using narrow-band pulsed lasers with a high repetition rate. The power of pulsed lasers is needed to ionize efficiently, but for a continuous atom beam this also involves appreciable duty cycle losses.

On the other hand, it is not always necessary to achieve sub-Doppler resolution. Hyperfine structures and isotope shifts of heavy atoms are usually large enough to be accessible to Doppler-limited methods. Here one can choose a different high-sensitivity approach, based on the laser ionization of a thermal ensemble of atoms. This concept of resonance ionization spectroscopy (RIS) [83] is in many respects complementary to the methods described so far. On-line experiments on radioactive isotopes of a number of heavier elements were performed using techniques of sample collection and rapid atomic re-evaporation [84, 85]. RIS uses strong pulsed lasers for the stepwise excitation of an appreciable fraction of atoms to the ionization continuum. Moreover, ion counting, often combined with mass separation offers an efficient and background-free detection of the signals. When RIS is performed on thermal ensembles of atoms, the resolution is limited by Doppler broadening (see (29)). This does not represent a problem for the efficiency of the ionization process, because the spectral width of pulsed lasers is typically also in the GHz range. The modest resolution in measurements of large isotope shifts and hyperfine structures can even be an advantage, because it avoids scanning the laser frequency in small steps over a large ranges.

There is again an elegant way to combine the spectroscopic method with the concept of on-line isotope separation. The RIS principle forms the basis of the laser ion source [52, 57] which has become increasingly important for the production of clean beams of many elements. Using this ion source as a spectroscopic tool avoids the inevitable losses introduced by additional steps of sample preparation. This, together with the efficiency and the extreme selectivity of the ionization process yields the very high sensitivity which is

achieved when the laser ion source of an on-line isotope separator is directly used for spectroscopy [86, 87]. The detection consists simply in a measurement of the ionization rate as a function of the laser frequency. It is only necessary to control and accurately measure the laser frequencies.

Depending on the ionization energy and the available intermediate levels for the stepwise excitation process, two or three laser beams are used to ionize the atoms. In most practical cases, the first step is used for spectroscopy, yielding the desired information about the nuclear properties. This is demonstrated in Fig. 13 for the nuclear ground state and isomeric state of ^{185}Pb . The higher excited states can usually be chosen to be less sensitive to isotopic effects. This means that only the laser for the first step has to be tuned over a scanning range covering the hyperfine structures and isotope shifts.

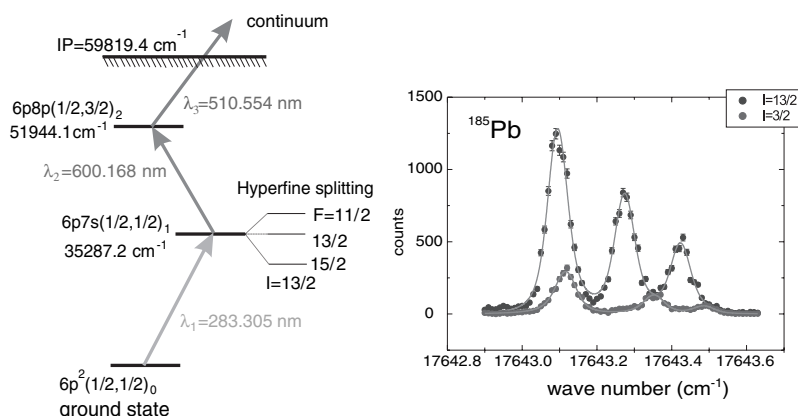


Plate 13. Three-step laser ionization scheme for Pb isotopes and hyperfine structure of two isomeric states in ^{185}Pb induced by scanning the laser frequency in the first step [88]; cf. Plate 14 in the Colour Supplement

Saturation of the optical transitions is required for maximum ionization efficiency. This is mainly a problem for the ionization step, because the photo-ionization cross-sections are small, typically 10^{-18} cm^2 . In favorable cases one can reach auto-ionizing states decaying to an ion and a free electron. These resonances in the continuum correspond to a two-electron excitation. Without such resonances one can rely on high laser power, e.g., use the strong pump laser beam to reach the continuum.

The laser ion source concept needs lasers with a high repetition rate. One has to make sure that any atom diffusing out of the target through the outlet tube serving as an ion source, is illuminated by at least one of the laser pulses. At ISOLDE, good results were obtained with dye lasers pumped by copper vapor lasers at a repetition rate of 10 kHz. More user-friendly alternatives

can be found in Ti:Sapphire lasers pumped by the frequency-doubled output of a Nd:YAG laser.

5 Methods Based on the Interaction of Nuclei with External Fields

Experimental techniques based on measuring the angular distribution of the radioactive decay are often more sensitive than the methods discussed so far, and in some cases also allow more precise measurements of the nuclear g -factor and quadrupole moment. This angular distribution is influenced by the interaction of the nuclear moments with externally applied magnetic fields and/or electric field gradients after implantation into a crystal (see Sect. 3.2). The radioactive decay intensity is measured as a function of time or as a function of an external variable, e.g., a static magnetic field or the frequency of an applied radio-frequency magnetic field. The former are called “time differential” measurements and the latter “time integrated” measurements.

5.1 Angular Distribution of the Radiation from Oriented Nuclei

The direction in which the β - or γ -decay of a radioactive nuclear state occurs, is determined by the direction of its spin. The angular distribution of radiation emitted by an ensemble of spin oriented radioactive nuclei with lifetime τ , in a direction (θ, ϕ) with respect to a chosen axis system (the LAB system), is given by [89]

$$W(\theta, \phi, t) = e^{-t/\tau} \sum_{k,n} \sqrt{\frac{4\pi}{2k+1}} A_k(\gamma, \beta, \dots) U_k B_k^n(I, \omega_L, \omega_Q, t) Y_k^n(\theta, \phi). \quad (40)$$

In this expression, A_k are the radiation parameters describing the type of radiation and its properties (β -decay asymmetry parameter, γ -decay of multipolarity M1, E2, ...), and Y_k are the spherical harmonics: they depend on the position of the detector (θ, ϕ) with respect to the LAB system. The orientation tensor B_k describes the spin orientation of the ensemble with respect to the LAB system, and its time-dependent change due to the interaction of the nucleus with the surrounding fields. The orientation of a nuclear spin ensemble is described easiest by the density tensor ρ_k^n , which is related to the orientation tensor $B_k^n = \sqrt{2k+1} \rho_k^{n*}$ and to the nuclear density matrix. It is this parameter that needs to be calculated to describe the influence of a particular interaction on the spin orientation and thus on the decay intensity. It is convenient to write the orientation tensor as a function of the initial orientation tensor $B_k^n(I, t=0)$ (the orientation before the interaction with the perturbing fields is applied) and a perturbation tensor $G_{kk'}^{nn'}(\omega_L, \omega_Q, t)$:

$$B_k^n(I, \omega_L, \omega_Q, t) = G_{kk'}^{nn'}(\omega_L, \omega_Q, t) B_k^n(I, t=0). \quad (41)$$

In order to observe an anisotropic radiation pattern, the initial spin orientation needs to be anisotropic, meaning that a spin-oriented ensemble of short-lived nuclei is required. Several experimental methods have been developed to produce such spin-oriented radioactive beams. One of these techniques has been described in Sect. 4.1: optical pumping on an ISOL beam by collinear polarized laser light. Another method is to use the spin orientation induced by the nuclear reaction that produces the exotic nuclei of interest (e.g., in fusion evaporation, projectile fragmentation, Coulomb excitation or transfer reactions). In this case, care needs to be taken that the spin orientation is maintained during the isotope selection process. More examples of methods that are used to produce exotic spin-oriented nuclei and the different formalisms which are used to describe the nuclear spin orientation can be found in [11] and references therein. In particular, we will not discuss in this lecture the low temperature nuclear orientation (LTNO) method [90]. In this case, spin orientation is not produced in the radioactive beam, but only after implantation. The ferromagnetic host crystal is cooled to very low temperature (mK) and the nuclei are submitted to a very strong magnetic field of typically more than 10 Tesla, such that the Boltzmann distribution causes a spin orientation.

5.2 Time-Differential Perturbed Angular Distribution (TDPAD) of γ -Decaying Isomers

Spin-oriented isomeric states implanted into a suitable host will exhibit a non-isotropic angular distribution pattern, provided the isomeric ensemble orientation is maintained during its lifetime. If an electric field gradient is present at the implantation site of the nucleus, the nuclear quadrupole interaction will reduce the spin orientation and thus the measured anisotropy. Also spin-relaxation effects (e.g., Korringa relaxation in metals [91] or quadrupole relaxation in transition metals [92]) can reduce or even fully cancel the spin orientation of the ensemble. If the implantation host is placed into a strong static magnetic field (order of 0.1–1 Tesla), the anisotropy is maintained. If the field is applied parallel to the symmetry axis of the spin orientation, the reaction-induced spin orientation can be measured.

Magnetic Interaction TDPAD

If a static magnetic field is placed perpendicular to the axial symmetry axis of the spin orientation, the Larmor precession of the isomeric spins in the applied field can be observed as a function of time [93], provided that the precession period is of the same order as the isomeric lifetime (or shorter). This method is called time-differential perturbed angular distribution (TDPAD).

A formal description of the perturbed angular distribution function can be derived starting from (40) and (41). Because γ -decay is not violating parity,

only even radiation parameters are non-zero. Furthermore, it is often assumed that the ($k \geq 4$) terms are negligibly small, reducing the angular distribution to

$$W(\theta, \phi, t) = e^{-t/\tau} \left[1 + \sqrt{\frac{4\pi}{5}} A_k U_k \sum_n B_2^n(I, \omega_L, t) Y_2^n(\theta, \phi) \right]. \quad (42)$$

The perturbation factor describing the Larmor precession of the nuclear spins around a static field, is given by $G_{22}^{nn}(t) = e^{-in\omega_L t}$. The time-dependent perturbed spin orientation, described in a reference frame with the z -axis along the magnetic field and the x -axis along the beam direction is then given by

$$B_2^0(t) = -\sqrt{\frac{1}{2}} B_2^0(t=0), \quad (43a)$$

$$B_2^2(t) = \sqrt{\frac{3}{8}} e^{-i2\omega_L t} e^{-i2\gamma} B_2^0(t=0), \quad (43b)$$

with $B_2^0(t=0)$ being the alignment with respect to the orientation symmetry axis at the time of implantation. This axis is along the beam direction for isomeric states investigated in-beam, while has an angle γ in the xy -plane with respect to the beam direction for isomers after in-flight mass separation (see e.g., [41] and [94]).

If the detectors are placed in a plane perpendicular to the magnetic field direction ($\theta = 90^\circ$) and at nearly 90° with respect to each other ($\phi_1 \approx \phi_2 + 90^\circ$), the $R(t)$ function in which the Larmor precession is reflected, is given by

$$R(t) = \frac{W(\phi_1, t) - W(\phi_2, t)}{W(\phi_1, t) + W(\phi_2, t)} = \frac{3A_2 B_2^0(t=0)}{4 + A_2 B_2^0(t=0)} \sin(\phi_1 + \phi_2 - 2\omega_L t - 2\gamma). \quad (44)$$

The TDPAD method has – since the seventies – been the most important method to measure g -factors of isomeric states produced and spin-aligned by fusion-evaporation and transfer reactions, with lifetimes spanning the range of 10 ns to about 100 μ s. Recently the method has also been applied to investigate the g -factor of neutron-rich isomeric states, which can not be produced by the former production methods. In this case the projectile fragmentation reaction was used in combination with a high-resolution doubly-achromatic spectrometer to produce and select a rather pure beam of spin-aligned neutron-rich isomers [94, 95]. Specific aspects of moment measurements on isomeric fragment beams will be discussed further below.

Quadrupole Interaction TDPAD

The TDPAD method can also be used to measure the quadrupole moments of these isomeric states, by implantation into a single crystal or a polycrystalline material with a non-cubic lattice structure providing a static electric

field gradient (no magnetic field is applied in this case). In a quadrupole interaction measurement, $2I$ 'quadrupole frequencies' occur as multiples of the basic coupling constant $\omega_0 = 3\pi\nu_Q/I(2I - 1)$ for half-integer spins and $\omega_0 = 3\pi\nu_Q/2I(2I - 1)$ for integer spins. They superimpose on each other in the $R(t)$ function, which makes it difficult to use this method for measuring quadrupole moments of high-spin isomeric states with lifetimes longer than 10 μ s. For in-beam experiments, the quadrupole $R(t)$ function is given by

$$R(t) = \frac{3A_2B_2^0(t=0)}{4 + A_2B_2^0(t=0)} \sum_n s_{2n} \cos(n\omega_0 t), \quad (45)$$

and the s_{2n} coefficients can be found in [96].

Examples of TDPAD Experiments

The in-beam TDPAD method continues to be a powerful tool to study g -factors and spectroscopic quadrupole moments of isomeric states. An example is the investigation of the rich nuclear structure in the neutron-deficient Pb isotopes [97, 98, 99]. In these isotopes with a magic proton number $Z = 82$, the spherical ground state is found to coexist with prolate and oblate deformed structures at very low excitation energies [100]. These deformed states are interpreted as arising from particle-hole excitations of protons across the $Z = 82$ shell gap into the $\pi h_{9/2}$ and $\pi i_{13/2}$ orbits. As a function of deformation, these orbits are "intruding" into the lower-shell orbits, giving rise to particle-hole excited states that occur at low energy. Therefore, these states are called "intruder states". The deformation-driving interaction is the proton-neutron interaction, which is enhanced when more neutron holes become available in going from the doubly magic ^{208}Pb ($N = 126$) down to neutron mid-shell ($N = 106$). Therefore, the energy of the intruder 2p-2h states is decreasing towards mid-shell, until in ^{186}Pb the lowest three states of the nucleus are found to be a spherical ground state, an oblate deformed first excited state and a prolate deformed second excited state [100]. Some of these intruder states are high-spin isomeric states, e.g., the $\pi(s_{1/2}^{-2}h_{9/2}i_{13/2})_{11-}$ isomers observed between ^{198}Pb and ^{188}Pb . In these nuclei, one has also observed a new type of nuclear rotation, called "magnetic rotation", because the properties of the observed rotational bands can be explained by the rotation of a magnetic dipole that forms an angle with respect to the symmetry axis of the deformed nucleus [101]. The band head of these rotational bands is based on a perpendicular coupling of the intruder 2p-2h proton to a high- j neutron-hole configuration, as proven experimentally by a g -factor measurement [97] (Fig. 14, left). This measurement was performed for the only "isomeric" band head in the region having a half-life of 9.4 ns, the $29/2^-$ isomer in ^{193}Pb . The quadrupole moment of the band head can reveal information about its deformation. A direct measurement for this very short-lived isomer has been performed recently (Fig. 14, right), revealing indeed a large quadrupole moment [98], similar to that of the intruder proton configurations [99].

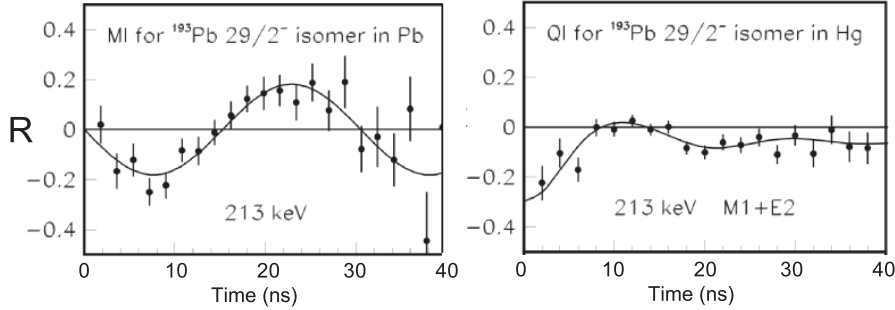


Plate 14. TDPAD spectra for the γ -decay of the $I^\pi = 29/2^-$, $t_{1/2} = 9$ ns isomeric rotational bandhead in ^{193}Pb , implanted respectively in a lead foil to measure its magnetic interaction (MI) and in cooled polycrystalline mercury to measure its quadrupole interaction (QI). Data are from Balabanski et al. [98]

In-beam Versus In-flight Experiments

In-beam fusion-evaporation reaction products are highly aligned and ideally suited for nuclear moment studies. However, this production method has some disadvantages for the study of moments of nuclei far from stability.

The first problem with in-beam experiments on nuclei far from stability is the low signal to background ratio in the photo-peak of interest. An example for techniques to improve this peak to background signal is the use of recoil separation of the isomers produced via a fusion-evaporation reaction in inverse kinematics. In this case it is important to be able to select the reaction products in charge states with a noble-gas configuration, in order to maintain the ensemble orientation during the in-flight separation process [102].

A second problem is that neutron-rich isomers are not produced in such reactions. However, the discovery of isomers [103] and of spin alignment [104] in the fully stripped radioactive beams from in-flight projectile fragmentation facilities, has very recently opened a new field of nuclear moments research, not accessible before. A first proof of principle, applying the TDPAD method to an isomeric projectile fragmentation beam, has been described in [105] and first experiments on exotic nuclei are described in [94, 95]. Figure 15(a) shows an example of some $R(t)$ curves obtained in the study of g -factors of $I^\pi = 9/2^+$ isomers in neutron-rich isotopes of nickel and iron. The isomers, with lifetimes of 13.3 μs and 250 ns, respectively, have been produced in a projectile fragmentation reaction at the LISE high-resolution in-flight separator at GANIL. The configuration of these isomers is suggested to be dominated by the $\nu 1g_{9/2}$ intruder orbit in these nuclei which have less than 40 neutrons. The g -factor measurement for the isomer in ^{61}Fe confirms this assignment, based on comparison with earlier measured g -factors for similar isomers closer to stability. For the isomer in ^{67}Ni , the experimental result deviates from the expected trend, as can be seen from Fig. 15(b). Further investigation on the

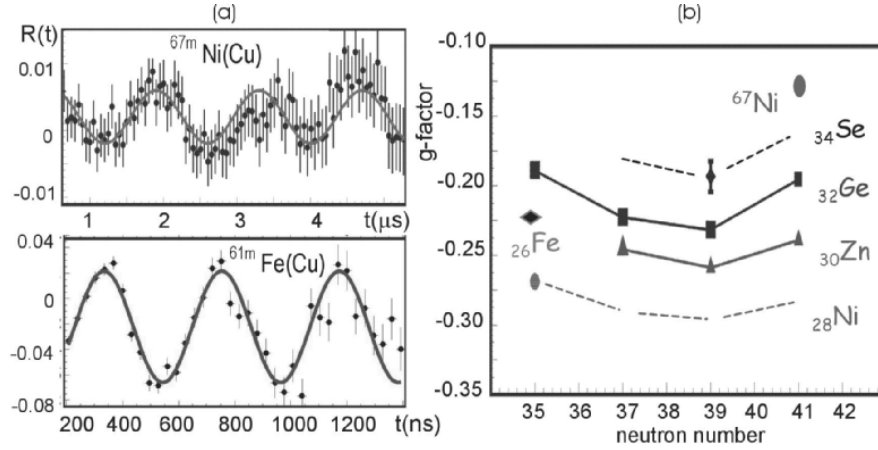


Plate 15. (a) TDPAD spectra for isomeric $I^\pi = 9/2^+$ states in the neutron-rich isotopes ^{61}Fe and ^{67}Ni , produced and spin-aligned via a projectile fragmentation reaction. (b) Experimental g -factors for $I^\pi = 9/2^+$ states in nuclei around $Z = 28$. The *dotted lines* are drawn to show the expected trend, based on comparison with the Zn and Ge trend lines

structure of this isomeric state near the suggested “doubly-magic” or “super-fluid” ^{68}Ni [106, 107, 108] is clearly needed.

5.3 Beta-Ray Detected Nuclear Magnetic Resonance (β -NMR)

Time-differential measurements as those discussed in the previous section are only suited for short-lived nuclear states, mainly because of relaxation effects causing a dephasing of the Larmor precession frequencies with time (typically in less than $100\ \mu\text{s}$) [89]. To measure nuclear moments of longer-lived isomeric states and also for ground states, a time-integrated measurement is required. Time integration of (44) and (45), taking into account the nuclear decay time, will lead to a constant anisotropy. Therefore, a time-integrated measurement of the angular distribution of this system will not allow one to deduce information on the nuclear moments. Hence a second interaction, which breaks the axial symmetry of the Hamiltonian, needs to be added to the system. For example, combining a quadrupole and a dipole interaction with their symmetry axes non-collinear (as described in Sect. 3.2), will give rise to resonant changes in the angular distribution at the magnetic field values where the nuclear hyperfine levels are mixing. The method based on this principle – the level mixing resonance (LMR) method – will be discussed in the next section.

Another way to introduce a symmetry breaking in the system, is by adding a radio-frequency (rf) magnetic field perpendicular to the static magnetic field (and to the spin-orientation axis). The similarity of this time-dependent way

of perturbing the system to the static case of non-collinear interactions, is discussed in [43]. If the nuclei are implanted into a crystal with a cubic lattice symmetry or with a noncubic crystal structure inducing an electric field gradient, respectively, one can deduce the nuclear g -factor or the quadrupole moment from the resonances induced by the applied rf field between the nuclear hyperfine levels.

Nuclear Magnetic Resonance: Features and Examples

Consider an ensemble of nuclei submitted to a static magnetic field B_0 and an rf magnetic field with frequency ν_{rf} and rf field strength B_1 . If the applied rf frequency matches the Larmor frequency of the nuclei:

$$\nu_{\text{rf}} = \frac{g_I B_0 \mu_N}{h}, \quad (46)$$

the orientation of an initially spin-oriented ensemble will be resonantly destroyed by the rf field [91]. For β -decaying nuclei that are initially polarized, this resonant destruction of the polarization can be measured via the change in the asymmetry of the β -decay, as demonstrated in Figs. 16(a) and 17(a). For an ensemble of nuclei with the polarization axis parallel to the static field direction, the angular distribution for allowed β -decay can be written as

$$W(\theta) = 1 + A_1 G_{11}^{10}(I, \omega_L, \omega_{\text{rf}}) B_1^0(t=0) \cos \theta, \quad (47)$$

with the NMR perturbation factor G_{11}^{10} describing the NMR as a function of the rf frequency or as a function of the static field strength (e.g., deduced in [43]). At resonance, the initial asymmetry is fully destroyed if sufficient rf power is applied [89], which corresponds to $G_{11}^{10} = 0$. Out of resonance we observe the full initial asymmetry and $G_{11}^{10} = 1$. Equation (47)

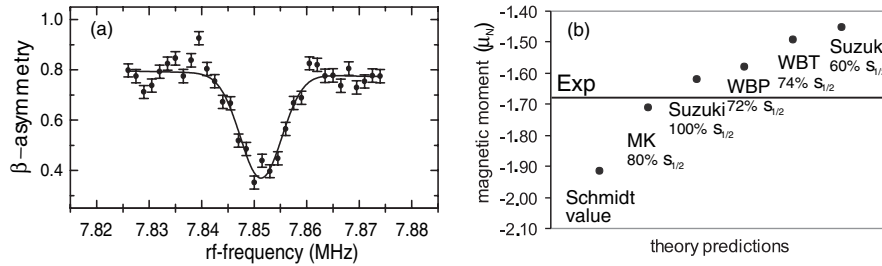


Plate 16. (a) NMR curve for ^{11}Be implanted in metallic Be at $T = 50$ K. At this temperature the spin-lattice relaxation time T_1 is of the order of the nuclear lifetime $\tau = 20$ s. (b) Experimental result compared to theoretical predictions obtained from different shell model approaches (see text for details and references)

can then be related to the well-known expression for the asymmetric β -decay, taking into account that $A_1 B_1 = AP$ and with the polarization $P = -\sqrt{((I+1)/3I)} B_1^0(t=0)$ [109]:

$$W(\theta, \beta^\pm) = 1 \mp AP \cos \theta. \quad (48)$$

The position of the NMR allows one to deduce the nuclear g -factor using (46). The amplitude of the NMR signal depends on:

- the β -decay asymmetry parameter A which can be estimated if the β -decay scheme is known [109]; if not, it is to be considered as a parameter, which can vary between -1 and $+1$ (and thus can even be close to zero, in which case NMR is very hard to detect);
- the amount of initial polarization P which is maintained after implantation in the crystal ($B_1^0(t=0)$); Crystals with a long spin-lattice relaxation time are therefore preferred;
- effects reducing the full breakdown of the asymmetry because not *all* the nuclei are at resonance with the applied rf field.

It is obvious that the statistical quality of the NMR signal increases linearly with the resonance amplitude and with $\sqrt{N}$, N being the number of detected β -particles. Thus it is important to maximize the NMR signal as much as possible, with a gain in the polarization having the largest effect. Full destruction of the nuclear polarization (and thus maximal effect in the measured NMR, taking into account a given A and P) is obtained provided that the following conditions are fulfilled:

- All nuclei in the crystal should interact with the same magnetic field B_0 . In some crystals the implanted nuclei are not all positioned in an unperturbed, substitutional lattice site. Such nuclei might interact with a slightly different magnetic field or their energy levels might be perturbed by a small quadrupole interaction. These nuclei will not be at resonance with the applied rf field and their spin orientation will thus not be affected.
- The rf interaction strength, $\omega_1 = g_I B_1 \mu_N / \hbar$, should be larger than the inverse nuclear lifetime (Heisenberg uncertainty principle). This puts a condition on the lower limit of lifetimes for which this NMR method can be applied, depending on the rf field strength B_1 , namely $\tau \gg \hbar / g_I B_1 \mu_N$. For a typical rf field of about 10 Gauss and a nucleus with g -factor $g_I = 1$, this lower limit is of the order of 100 μ s.
- To avoid relaxation-induced destruction of the spin orientation, the nuclear lifetime should be shorter than the spin-lattice relaxation time T_1 . This poses an upper limit on the lifetime of nuclei for which the NMR method can be applied. The spin-lattice relaxation time in isolators is typically of the order of several seconds or tens of seconds, while in metals it is of the order of a less than a second. In metals where the Korringa relaxation is the dominating relaxation process, the relaxation time T_1

can be increased by decreasing the lattice temperature T (according to $T_1 T = C_K$, the Korringa constant).

The NMR method can be applied to polarized beams of short-lived ($T_{1/2} < 10$ s) exotic nuclei, provided the spin relaxation time in a suitable host is longer than the nuclear lifetime. Polarization of the nuclear spins can be obtained e.g., by the optical pumping method, as described in Sect. 4.1, for an ISOL beam. A recent series of experiments has been performed at ISOLDE-CERN to investigate the nuclear moments of halo nuclei. The spin $I = 3/2$, the g -factor and the quadrupole moment of the two-neutron halo nucleus ^{11}Li have been measured some years ago [80, 110] and will now be improved on the basis of recent high-precision measurements of the moments of ^8Li and ^9Li [111].

A precision measurement of the g -factor of the one-neutron halo nucleus ^{11}Be [112] will be discussed here as an example (see Fig. 16(a)). A single halo neutron governs the main nuclear structure properties of this nucleus with 4 protons and 7 neutrons. Based on the positive parity assigned to its ground state, it was suggested that the unpaired neutron mainly occupies the intruder $\nu 2s_{1/2}$ orbit rather than the negative parity $\nu 1p_{1/2}$ orbit as would be expected for a nucleus with 7 neutrons. A measurement of the g -factor, $g_I = -3.3632(16)$, confirmed this $1/2^+$ assignment. The magnetic moment deduced for an $I = 1/2$ state is $\mu_I = g_I I \mu_N = -1.6816(8) \mu_N$, which is close to the Schmidt value of $-1.91 \mu_N$ for a $\nu 2s_{1/2}$ configuration (while for a $\nu 1p_{1/2}$ configuration this value is $+0.64 \mu_N$). The question is whether the halo neutron wave function is a pure s -wave (neutron in the $s_{1/2}$ orbital). To see how much admixture of a $d_{5/2}$ neutron coupled to the 2^+ state of ^{10}Be is present, we compare in Fig. 16(b) the experimental result to theoretical predictions, i.e. the Schmidt moment for a $\nu 2s_{1/2}$ configuration, an effective $\nu 2s_{1/2}$ moment, an empirical value assuming 40% admixture with a $\nu(1d_{5/2} \times 2^+)$ configuration (from [113], see also [114]) and two values calculated with the OXBASH shell model code using the WBT and the MK interactions for the $p - sd$ shell. The difference between these interactions is in the cross-shell matrix elements between the p and the sd shell. The calculations have been performed with free-nucleon g -factors. If an effective g_s value is used, the calculated magnetic moments deviate more from the experimental value. Both shell model calculations predict that about 80% of the neutron wave function is in the $\nu 2s_{1/2}$ configuration and the remaining part is in a $\nu(1d_{5/2} \times 2^+)$ state. A comparison of all these calculations with the experimental value shows that there is evidence for the halo neutron to occupy partly a d -wave state.

The high precision (5×10^{-4}) which has been obtained in the ^{11}Be g -factor experiment is due to the fact that the NMR method is a resonance technique. Of course such a high precision requires a very homogeneous magnetic field (homogeneity of 10^{-5} over the surface of the implanted beam spot) with a stable magnet power supply and a crystal of good quality (with very little

inhomogeneous line broadening). However, also the beam intensity plays an important role. At beam intensities of 10^5 – 10^6 ions/s implanted into the crystal (as was the case for ^{11}Be), several attempts for a resonance search can be made in a reasonable time. Experiments in different crystals can thus be performed to determine the one that gives the smallest resonance line width. Moreover, fine scans of the resonance are feasible, all within a few days of beam time.

The experimental situation becomes much less comfortable when the beam intensity is as low as 10^2 – 10^3 ions/s which is often the case if one studies nuclei far from stability. In experiments using projectile fragmentation reactions, the moments of exotic nuclei have been investigated with rates as low as a few 1000 ions/s. The advantage of using this reaction mechanism, apart from being faster than the ISOL production method, is that a spin-oriented ensemble is obtained from the nuclear reaction process itself. A spin-polarized beam can be selected for every fragment of interest⁵, whereas the optical pumping method is limited to a few elements with suitable atomic spectra. Furthermore, the projectile fragmentation reactions allow the production of a very broad range of isotopes of all elements, while the ISOL method is element dependent, although it often gives better yields for the isotopes that are accessible. The typical amount of polarization observed for projectile fragments is of the order of 5–15% [116], while the polarization produced with the optical pumping method varies typically between 20–60% [81].

Not only the low count rate, also the fact that in general very little is known on the ground state of exotic nuclei, makes g -factor measurements on such species very difficult. Sometimes not even the spin/parity is known, which implies that the g -factor range in which one needs to search for a resonance, is very broad. A way to scan a “broad” range of g -factors in a reasonable time is based on frequency modulation of the rf field. By modulating the rf frequency over a range $\Delta\nu$ around a central value ν_{rf} , one can scan a g -factor range according to

$$\Delta g_I = \frac{\Delta\nu h}{B_0 \mu_N}. \quad (49)$$

Using this measuring approach with a modulation range $\Delta\nu = \pm 10\%$ of the applied frequency, it is possible to scan a range of g -factors between $g_I = 0.6$

⁵The spin orientation is mainly related to the transfer of orbital momentum from the projectile to the fragment, assuming a peripheral collision [104]. If the fragment spin is dominated by spin momentum of a single nucleon, then it might happen that the spin orientation is very low. For example, no spin polarization was observed for the halo nucleus ^{11}Be , which has a spin $I = 1/2$ ground state dominated by a single particle $s_{1/2}$ wave function [112], in the fragmentation of a ^{18}O beam [115]. Generally, spin orientation has been observed for most fragment beams, albeit the interplay between nuclear structure and reaction mechanism aspects in the production of spin orientation has not been investigated systematically.

and $g_I = 1.9$, by varying the applied magnetic field over just a few values (as shown in Fig. 17(a1)). This reduces drastically the measuring time needed for locating the resonance. The uncertainty of such a g -factor measurement is given by the applied modulation width. In a subsequent scan, using a smaller frequency modulation, the range in which to search for the resonance can be reduced as well as the uncertainty of the deduced g -factor, as demonstrated in Fig. 17(a2).

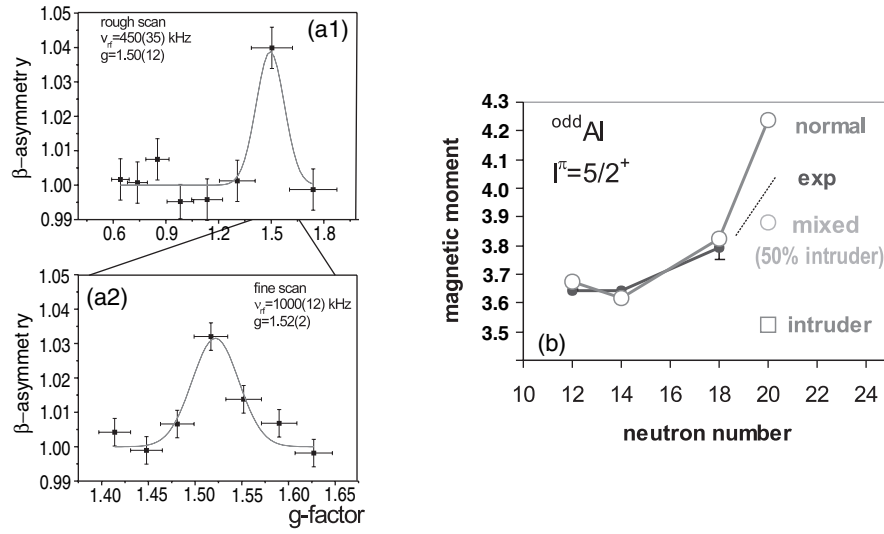


Plate 17. (a) NMR curves for ^{31}Al implanted in MgO . The scan performed as a function of the static field B_0 covers a broad range of g -factors, the modulation width of the rf frequency amounts to 8% (a1) and 1.2% (a2), respectively. (b) Comparison of experimental magnetic moments to different predictions for ^{33}Al , assuming a “normal” ground state, an “intruder” ground state or a “mixed” ground state [117]; cf. Plate 15 in the Colour Supplement

Such a frequency-modulated scan over a broad g -factor range needs an rf field strength B_1 that can only be achieved by integrating the rf coil in an RLC resonance circuit. This field strength varies with the applied frequency ν_{rf} and depends on the circuit properties. For the applied rf-generator voltage $V_{\text{gen}} \cos(\omega_{rf}t)$ the rf current depends on the impedance $Z = \sqrt{R^2 + (L\omega_{rf} - \frac{1}{C\omega_{rf}})^2}$ of the circuit:

$$I_{rf}(t) = \frac{V_{\text{gen}}}{Z} \cos(\omega_{rf}t - \Phi). \quad (50)$$

This impedance changes with the applied frequency. To avoid large variations of the rf field strength B_1 over a wide frequency scanning range, it is preferable

to perform the scan as a function of the static field B_0 , using a fixed rf frequency for which the RLC circuit is optimized. This procedure was used in the measurements shown in Fig. 17.

Frequency modulation has another important advantage which is related to the amplitude of the resonance. As discussed above, the condition for a full resonant destruction of the β -decay asymmetry is that *all* nuclei implanted in the crystal are at resonance with the rf field at exactly the same frequency. This requires that the Zeeman splitting of the nuclear levels is the same for all nuclei in the sample, which is not always the case. For example, a small EFG due to imperfections of the crystal, may induce small shifts of the Zeeman levels, or a magnetic field which is not perfectly homogeneous over the implanted beam spot may result in a different Zeeman splitting for different nuclei. With a modulated rf frequency, *all* the implanted nuclei will be simultaneously at resonance and contribute to the NMR signal. In order to reach saturation of the resonance, the rf field strength B_1 needs to be high enough, such that the homogeneous line broadening of ($\Gamma_{hom} \sim g_I B_1$) is larger than the inhomogeneous broadening. As the sensitivity of the method increases with the amplitude of the resonance squared, it is clear that a maximum signal is desired for experiments with low beam intensities.

At GANIL the g -factors and quadrupole moments of neutron-rich nuclei in the region of ^{32}Mg are presently under investigation. In this region of the nuclear chart it has been observed that some ground-state properties do not behave as expected for nuclei with a (near) magic neutron number $N \approx 20$: they have a deformed ground state [118]. This has been explained as due to excitations of neutrons from the sd shell into the pf shell. Such intruder particle-hole excited states become ground state due to the interplay between an enhanced proton-neutron interaction and a reduction of the $N = 20$ shell gap [119]. The region of nuclei that have ground states dominated by an intruder configuration, has been called the “island of inversion”. The goal of several experimental programmes is now to determine the borders of this island of inversion. At GANIL, spin-polarized and spin-aligned beams are obtained by selecting the secondary beam with a well-defined longitudinal and transverse momentum [42, 120]. Typical beam intensities of these nuclei vary from 5×10^4 down to 5×10^2 ions/s for the most exotic cases which are at the limit of still being accessible to such studies. The measurement of the ^{31}Al ($Z = 13$) ground-state g -factor has been reported recently [117], using the frequency modulation technique described above (see Fig. 17(a)). In a normal sd -shell model picture, the ground state of this odd- Z nucleus is expected to be dominated by the $\pi d_{5/2}^5$ configuration, with a ground state spin/parity $I^\pi = 5/2^+$. From the measured g -factor, which agrees very well with the calculated value, a firm spin assignment could be made. It was found that the ^{31}Al ground-state properties are well described in the sd -shell model with the USD interaction (see Fig. 17(b)). For the more neutron-rich Al isotopes, in particular ^{33}Al , more advanced interactions allowing excitation of

neutrons into the pf shell [121, 122], predict that these nuclei have some intruder components in their ground-state wave function. This affects the predicted g -factor, as shown in Fig. 17(b), and a measurement of the g -factor will thus provide information on the amount of intruder admixture.

Also at ISOLDE this region of nuclei has been investigated. Using optically polarized Na beams, the g -factors and quadrupole moments of neutron-rich Na isotopes have been measured [81, 123]. These results revealed that ^{29}Na , ^{30}Na and ^{31}Na , with the neutron numbers $N = 18, 19$ and 20 , have intruder components in their ground state wave functions [124].

In the chain of the Mg isotopes, which occur between the normal Al and the intruder dominated Na isotopes, the situation has become clear only very recently. ^{31}Mg , with $N = 19$, is expected to have a ground state spin/parity of $3/2^+$, due to the hole in the $\nu d_{3/2}$ orbital. This was also the tentative assignment based on earlier β -decay experiments [125].

The decisive experiment has been performed at ISOLDE using an optically polarized Mg^+ beam, for which the atomic hyperfine structure and the nuclear g -factor were measured independently. The hyperfine structure of the transition $3s\ ^2S_{1/2} \rightarrow 3p\ ^2P_{3/2}$ (D_2 line), observed in the β -decay asymmetry as a function of the Doppler shift (optical pumping section at a variable electrical potential ΔU), is shown in Fig. 18. The dominating ground-state ($J = 1/2$) hyperfine splitting is proportional to the product $g_I(I + 1/2)$. For each assumed spin one can thus determine to some accuracy the corresponding g -factor. With this knowledge, a β -NMR measurement does not only give the g -factor directly (inset of Fig. 18), but also decides firmly on the spin value. This combination of techniques has yielded the unexpected spin $I = 1/2$ and a precise value of the magnetic moment for the nuclear ground state of ^{31}Mg . Details including the theoretical interpretation are explained in [126]. The spin $I = 1/2$ can be understood in the Nilsson model if the nucleus is strongly prolate deformed. In the newly developed shell models, it is interpreted as an intruder-dominated state. A particle-hole excited state becomes the ground state due to the reduced $N = 20$ shell gap in this region of neutron-rich isotopes. The result will play a key role for the understanding of changes in the nuclear shell structure in the neutron-rich isotopes around the proton number $Z = 12$ and the disappearing magic neutron number $N = 20$.

NMR with Quadrupole Interaction: Features and Examples

The NMR method applied to nuclei implanted in crystals that have a non-cubic lattice structure involves the quadrupole interaction with an electric field gradient (EFG) at the lattice site of the nuclei. This will be explained by using the examples of ^8Li and ^9Li . Earlier experiments on these isotopes and on ^{11}Li at ISOLDE had provided first information on the spin and magnetic moment [80] as well as the quadrupole moment [110] for the prototype two-neutron halo nucleus ^{11}Li . Improved investigations of nuclear resonances in different cubic and non-cubic crystals have been performed recently [111],

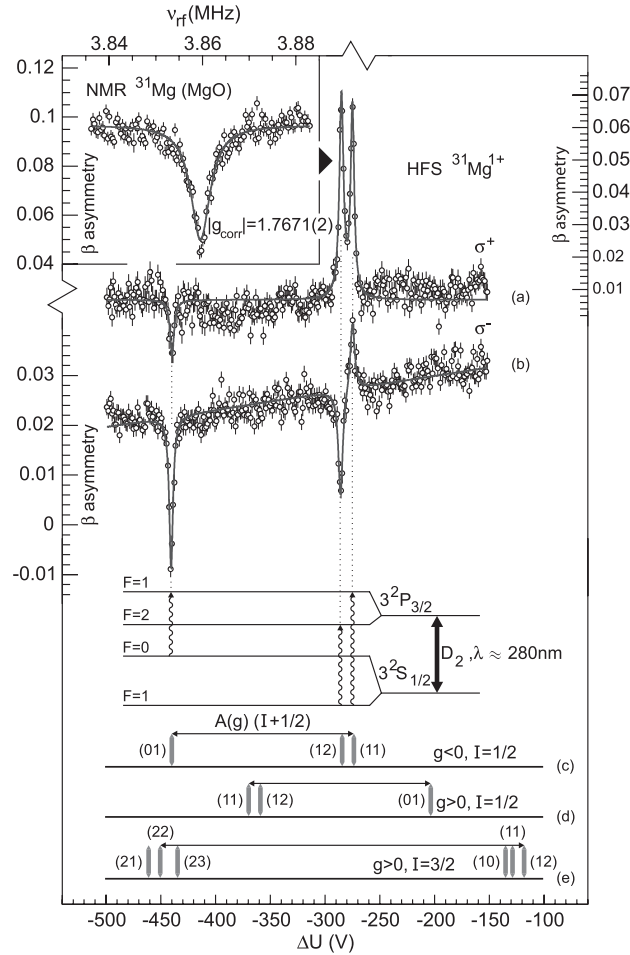


Plate 18. Atomic hyperfine structure in the D_2 line of $^{31}\text{Mg}^+$, measured for an optically polarized ion beam which is implanted into a MgO crystal for β -asymmetry detection. This, combined with a separate β -NMR measurement of the nuclear g -factor (see inset), determines unambiguously the nuclear spin and gives a very precise value for the nuclear magnetic moment. The lower part shows the hyperfine structure expected for different spins and positive/negative sign of the g -factor; cf. Plate 16 in the Colour Supplement

as a preparation for a high-precision measurement of the ^{11}Li quadrupole moment.

Beams of ^8Li and ^9Li atoms, with intensities of 10^7 and 2×10^5 atoms/s, respectively, were polarized via the optical pumping method and implanted into different crystals oriented with their c -axis along the static magnetic field direction. An rf field was applied perpendicular to the static magnetic field. The rf frequency was scanned around the Larmor frequency, giving rise to equidistant resonances in the asymmetry of the β -decay, as illustrated in Fig. 19 for ^8Li and in Fig. 20 for ^9Li . The distance Δ between the resonances depends on the quadrupole frequency ν_Q , and this dependence changes with the spin, as illustrated in Figs. 19 and 20 for nuclei with spin $I = 2$ and $I = 3/2$. By fitting the spectrum with equidistant resonances one can deduce from the resonance distance the quadrupole interaction frequency, $\nu_Q = 4\Delta$ for ^8Li ($I = 2$) and $\nu_Q = 2\Delta$ for ^9Li ($I = 3/2$). The parameter Δ contains the product of the nuclear quadrupole moment and the EFG at the site of the nucleus, $Q_s V_{zz}$.

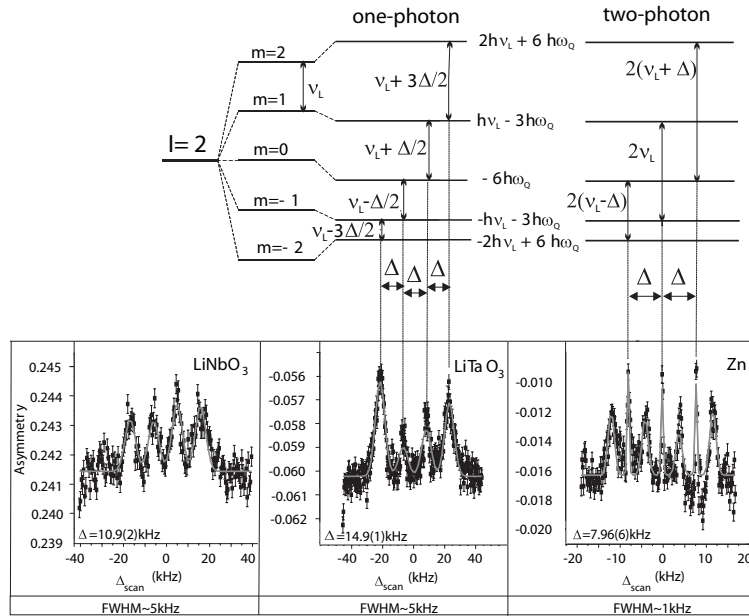


Plate 19. Nuclear magnetic resonances for ^8Li ($I = 2$) implanted into different non-cubic crystals. This illustrates the influence of the implantation host on the quadrupole frequency as well as on the resonance line widths. The nuclear level splitting for a nucleus with spin $I = 2$, submitted to a magnetic field and an EFG, and the corresponding transition frequencies are shown for one- and two-photon transitions. The five levels are non-equidistant, resulting in four equidistant one-photon resonances in the NMR spectrum; cf. Plate 17 in the Colour Supplement

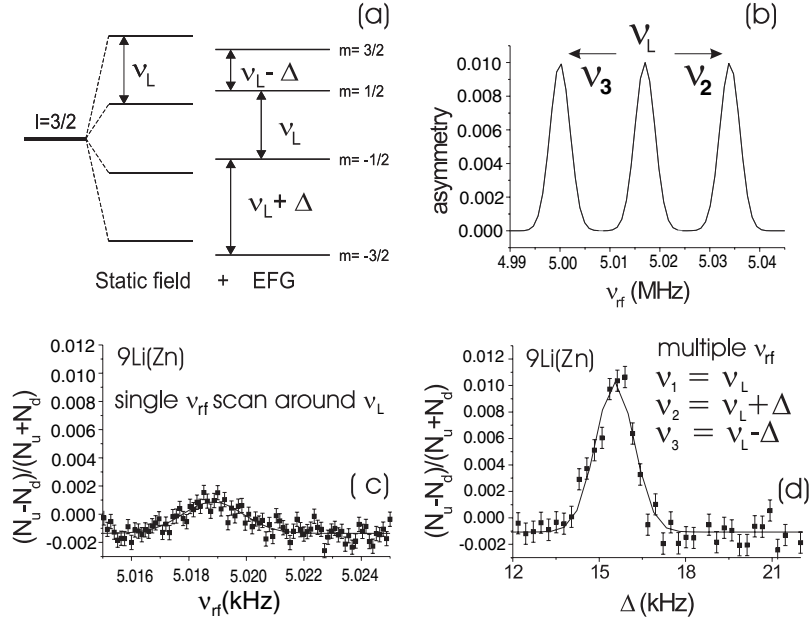


Plate 20. (a) Hyperfine levels of a nucleus with $I = 3/2$ submitted to a static magnetic field and an EFG. (b) Simulation of the three NMR peaks appearing in the β -decay asymmetry due to a resonant breakdown of the ensemble polarization. (c) Experimental result of a single-frequency scan around the Larmor frequency for $^9\text{Li}(\text{Zn})$. (d) Multiple-rf scan for $^9\text{Li}(\text{Zn})$.

Figure 19 shows the NMR scans for ^8Li implanted in LiNbO_3 , in LiTaO_3 and in metallic zinc. The crystal plays an important role in the quality of the resonances: not only the EFG is different, but also the resonance amplitude as well as the resonance widths can differ significantly from crystal to crystal. This has implications on the accuracy with which the quadrupole frequency, and consequently the quadrupole moment, can be measured. Notice that in the case of zinc, in which the line width is smallest and the applied rf field strength highest, one observes not only the $2I$ one-photon resonances $\Delta m_I = \pm 1$ with frequencies separated by Δ , but also the $2I-1$ two-photon resonances inducing transitions with $\Delta m_I = \pm 2$. These latter appear between the one-photon resonances at the same frequency distances Δ , all being symmetric with respect to the Larmor frequency. This is illustrated in the top part of Fig. 19.

Due to the non-equidistant level splitting, only two m_I quantum states can be at resonance with the applied rf frequency. The resonance amplitudes therefore reflect the destruction of asymmetry due to equalizing two level populations only, and thus are much smaller than in the purely magnetic case where all level populations are equalized simultaneously at the Larmor

frequency ν_L . The small resonance amplitudes often make it difficult to measure a full NMR spectrum as it was shown for the case of ^8Li . Figure 20(a) and (b) illustrates the levels and a simulated spectrum of three resonances for the $I = 3/2$ case of ^9Li . A scan of one of these resonances is shown in Fig. 20(c). For determining the quadrupole moment, one needs at least the distance of two such resonances, which for statistics reasons may be impossible for very exotic nuclei.

A solution to this problem was found already in the earlier measurements on ^{11}Li [110] (see also [127]), namely the simultaneous application of several correlated rf frequencies. In the spin $3/2$ case, all four m_I levels will be at resonance with the rf field for the correct value of Δ , when three frequencies ν_L , $\nu_L + \Delta$ and $\nu_L - \Delta$ are applied simultaneously. This is explained in Fig. 20(a). A search for a resonance can be performed by scanning the 3 rf frequencies simultaneously: one is fixed to the Larmor frequency, the other two are scanned symmetrically with respect to ν_L , as demonstrated in Fig. 20(b). The result of such a scan is a single resonance in which all levels are contributing to the NMR effect, as shown in Fig. 20(d). In this case the resonance amplitude is a factor of 6 higher than for the single-rf scan. From the resonance position the quadrupole interaction frequency can be obtained very accurately, and with a known EFG the nuclear quadrupole moment can be deduced. A requirement for the efficient application of this multiple-rf method is that the Larmor frequency ν_L is known very accurately (typically to a precision better than 10^{-3}) from the measurement in a cubic crystal.

5.4 Beta-Ray Detected Level Mixing Resonance (β -LMR)

The resonances observed in a LMR experiment are not induced by the interaction with a rf field, but by misaligning the magnetic dipole and electric quadrupole interactions [41]. This experimental technique does not need an additional rf field to induce changes of the spin orientation. The change of the spin orientation is induced by the quantum mechanical “anti-crossing” or mixing of levels, which occurs in quantum ensembles where the axial symmetry is broken. From the similarity between a dynamic and a static symmetry breaking it can be shown [43] that the observed resonances have similar features. There are however a few important differences. The β -LMR method can be applied to purely spin-aligned ensembles, while the β -NMR method requires a spin-polarized ensemble [128]. In combination with projectile fragments, this difference is important: a spin-aligned fragment beam is obtained by selecting the fragments in the forward direction (i.e. with the highest fragment yield), while the secondary beam needs to be selected asymmetrically with respect to the primary beam to obtain a spin-polarized ensemble. In general, this leads to a reduction of a factor of 5 in the intensity of the selected fragment beam. For exotic nuclei, this can make the difference between a feasible and a non-feasible experiment.

From a LMR curve, one can deduce the ratio of the nuclear moments, μ_I/Q_s . In order to deduce from this result the quadrupole moment, one needs to measure the nuclear g -factor with another method, e.g., the NMR method on a polarized fragment beam implanted in a crystal with cubic lattice symmetry. Another option is to induce the spin polarization by applying a resonant rf field to a spin-aligned ensemble implanted in a crystal with non-cubic symmetry [43]. In Fig. 21(a) the nuclear hyperfine levels are plotted for ^{12}B ($I = 1$) implanted in a Mg single crystal with hcp lattice structure. At the level anti-crossing field, a β -LMR is observed in the β -decay asymmetry as a function of the magnetic field strength (see middle resonance in Fig. 21(b) and (c)). The static fields for which the applied rf field is at resonance with the ensemble is shown in Fig. 21(a), once for a lower and once for a higher rf frequency applied. At these fields, β -NMR resonances occur, positioned symmetrically around the β -LMR. These data have been obtained by using a polarized ^{12}B beam produced via a $^{11}\text{B}(\text{d},\text{p})$ reaction and selected at about 40° with respect to the deuteron beam. Using this method, it has been possible to measure the ground-state moments of spin-aligned ^{18}N fragments, produced in the fragmentation of a ^{22}Ne beam at the LISE fragment separator at GANIL [42, 129].

Another feature of the β -LMR method, is that it allows to measure directly the nuclear spin. In Sect. 3.2 and in Fig. 6 it is demonstrated that the position of the resonances, as well as the distance between two resonances, is directly related to the nuclear spin. Fitting the experimental resonances with a model that has the nuclear spin, the ratio of the nuclear moments and the initial spin orientation as variables, one can directly deduce the nuclear spin. This has, however, not been applied till now.

6 Conclusions

In order to measure the static moments of exotic nuclei, whether in their ground state or in one of their excited isomeric states, it is necessary to apply complementary experimental techniques that cover a wide range of nuclear lifetimes and spins. In this lecture it has been shown that β -NMR methods allow the determination of high-precision values for the g -factor and quadrupole moment of short-lived nuclear ground states. Examples of such studies on exotic nuclei in neutron-rich isotopes of Be, Na, Mg and Al have been discussed. The Be, Na and Mg isotopes have been produced at ISOLDE-CERN, where they are spin-polarized using the resonant interaction with a collinear laser beam. The Al isotopes are more efficiently produced and spin-polarized using a projectile-fragmentation reaction followed by in-flight separation, as it was done at GANIL. Nuclear moments from atomic hyperfine structure and isotopic changes of the mean square nuclear charge radii can be measured by high-resolution optical spectroscopy, using a variety of techniques based on the resonant interaction of atoms or ions with a laser

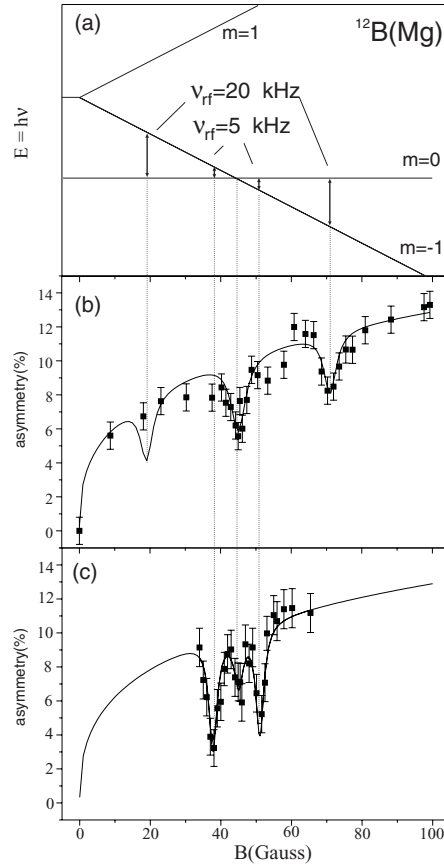


Plate 21. (a) Hyperfine levels for ^{12}B implanted in a Mg single crystal. (b) and (c) Resonances observed in the β -decay asymmetry due to nuclear level mixing and due to interaction with a resonant rf field

beam. These techniques are limited to ISOL-type production schemes. As an example, recent collinear laser spectroscopy experiments on short-lived neutron-deficient and neutron-rich Ne isotopes have been discussed. For the study of excited isomeric states, the spin orientation of isomeric projectile fragment beams will open a new field of research, e.g., the investigation of the moments of neutron-rich isomeric states.

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Spallation Reactions in Applied and Fundamental Research

J. Benlliure

Universidade de Santiago de Compostela, 15782 Santiago de Compostela, Spain

Abstract. Spallation reactions have recently gained new interest not only due to their application as neutron or radioactive nuclear beam sources but also for their implications in understanding cosmic ray abundances or investigating the dynamics of nuclear matter. The purpose of this lecture is to discuss the role of these reactions in different areas of interest, the modern experimental techniques currently being used for their investigation and finally some fundamental underlying physics.

1 Introduction

Spallation reactions are defined as interactions between relativistic light projectiles, mostly hadrons, and heavy target nuclei which are smashed into many fragments. In the relativistic energy domain, the wave length associated with the incoming projectile is such that the interaction can be described as a sequence of nucleon-nucleon collisions referred to as intra-nuclear cascade. The simplest illustration of this reaction is presented in Fig. 1. This process leads to the fast emission of some nucleons, identified as pre-equilibrium emission. The properties of these nucleons are basically determined by the kinematics of the nucleon-nucleon collisions, and the mass and charge distribution of nucleons inside the target nucleus.

As a consequence of this intra-nuclear cascade, the residue of the target nucleus gains thermal excitation energy and angular momentum. In a second step of the reaction, this target nucleus remnant or pre-fragment is expected to equilibrate all its excited degrees of freedom according to the compound-nucleus hypothesis. A subsequent statistical de-excitation process then leads to a final target residue in its ground state. This de-excitation mechanism has been explained as γ -ray, nucleon or cluster evaporation in competition with fission or even multi-fragmentation.

Nuclear reactions induced by light-relativistic projectiles were investigated for the first time more than 70 years ago using cosmic rays [1]. The observed high multiplicities of cascade particles had already been attributed to the spallation mechanism. Later on, these reactions led to the discovery of the pion [2]. High-energy particle accelerators came into operation at this time, and experimental programs to investigate these reactions were first undertaken using cyclotrons such as the one at Berkeley [3], which



Plate 1. Illustrative representation of a spallation reaction; cf. Plate 18 in the Colour Supplement

covered the energy range up to a few hundred MeV, and later on using proton synchrotrons like the Brookhaven Cosmotron [4] which reached the GeV energy range. At that time, particle detection techniques were based on radiochemical or photographic methods used to characterise the residual nuclei produced in these reactions. The main results of these programs were the observation of high-energy protons emitted in forward direction, understood as pre-equilibrium emission during the intra-nuclear cascade [5], and the identification of two de-excitation mechanisms leading to final residue formation, nucleon evaporation and fission [6].

With the advent of more advanced detection techniques, light particle emission in spallation reactions could be investigated more accurately during the late '60 and the '70s using magnetic and time-of-flight measurements with plastic scintillators [7] or ΔE -E telescopes based on semiconductor detectors [8]. These measurements permitted an accurate determination of the angular and energy distribution of light nuclear species (p, d, t, ^3He and ^4He) together with π and K mesons produced in these reactions. Two components were observed in the measured spectra: mesons and fast-forward emitted light-charged particles, attributed to the pre-equilibrium mechanism, and, the isotropic emission of light-charged particles, identified as the evaporation process from the thermalised target pre-fragment. At the same time, on-line mass separators were also introduced to investigate the nature of target residues produced in these reactions [9, 10].

These experimental investigations validated the already mentioned two-stage model proposed by Serber to describe spallation reactions [11]. According to this picture, these reactions are understood as an intra-nuclear cascade followed by the statistical de-excitation [12] of the resulting pre-fragments. The first attempts to compute the intra-nuclear cascade using the Monte Carlo technique were made by Goldberger [13], Bernardini et al. [14],

Morrison et al. [15], McManus et al. [16] and Meadows [17]. In addition, Meadows [17] and Jackson [18] used these calculations as a starting point for evaporation estimates to explain the yields of spallation residues observed in radiochemical studies. At that time, Metropolis and collaborators [19] and Bertini [20] provided the most advanced calculations describing both, particle emission and residual nuclei production in spallation reactions.

In the '80s, spallation reactions were identified as an optimum reaction mechanism for investigating the dynamics of hot nuclei, particularly the so-called multi-fragmentation process, which was thought to indicate the phase transition between nuclear and hadronic matter. Proton induced reactions were believed to heat the target nucleus without introducing a large amount of angular momentum or density variation. Consequently, this was a clean reaction mechanism for investigating thermal instabilities leading to multi-fragmentation. Such investigations required exclusive measurements of all reaction products including intermediate-mass fragments [21, 22]. Experiments of this kind were made possible in part thanks to the new BEVALAC heavy-ion synchrotron at Berkeley, which allowed the investigation of spallation reactions in inverse kinematics [23].

The parallel progress in computer hardware and numerical methods helped to develop new model calculations for spallation reactions. Some examples of this are intra-nuclear cascade models [24, 25], dynamical models based on kinetic equations including a collision term, like the Boltzmann-Uhling-Uhlenbeck (BUU) approach [26], or models such as Quantum Molecular Dynamics [27] which provided a full description of many-body nuclear dynamics.

In the last decade spallation reactions have gained new interest in several fundamental and applied research fields. They are considered as optimum neutron sources for solid-state physics or material-science investigations [28, 29, 30, 31], or for energy production and nuclear-waste transmutation in accelerator-driven systems [32, 33, 34]. These reactions have also been used during the last 35 years to produce and investigate nuclei far from stability and are currently being proposed as a production mechanism for next-generation exotic-beam facilities [35]. In addition, cosmic rays are known to undergo spallation reactions with interstellar medium during their propagation which affect the abundances observed on earth. In fact, due to the highly accurate cosmic-ray data provided by new observation techniques, most of the remaining uncertainties in the interpretation of these data are due to an inadequate description of the spallation process [36].

In this lecture the main emphasis will be put on discussing the present interest on spallation reactions, the current techniques used for their experimental investigation and the underlying physics. Moreover, some examples of basic research on the structure and the dynamics of the atomic nucleus using spallation reactions will be presented.

2 Fields of Application

2.1 Astrophysics: Cosmic Rays

Cosmic rays constitute one of the most important sources of information in Astrophysics. They consist of massive particles (protons and nuclei), photons and neutrinos. Cosmic rays can be of either galactic or extragalactic origin and it is well established that they undergo reactions (mainly spallation) with interstellar medium or with our atmosphere, which lead to secondary cosmic rays.

By studying the composition and energy distribution of cosmic rays one hopes to gain information about nucleosynthesis processes, galactic mass density or their lifetime/flight path. Since the motion of the particles involved (protons, nuclei and leptons) is roughly randomised by the galactic magnetic field, they provide very little information about the direction of the source. Their energy distribution peaks in the range of 100–1000 MeV, where the intensity of cosmic rays with an energy of 1 GeV or greater is about $1 \text{ cm}^{-2}\text{s}^{-1}$. The corresponding energy density is about 1 eV/cm^3 , which can be compared to the energy density of stellar light of 0.3 eV/cm^3 .

The chemical composition of cosmic rays is shown in Fig. 2. This distribution is almost independent of energy, at least in the dominant energy range between 10 MeV and several GeV. This composition has been measured by instruments mounted on balloons, satellites and spacecrafts. The figure also shows the chemical distribution of these elements in our solar system (dashed line), which in some cases is markedly different from that of the cosmic rays. The observed differences between cosmic-ray and solar system abundances are explained as being due to the reactions of energetic nuclei with the interstellar medium during transport through the galaxy from their source to the observation point. Since the interstellar medium is largely composed of hydrogen and helium, most of these reactions are spallation reactions. Consequently, the observed overabundance of elements in the region of lithium-beryllium-boron and scandium-manganese are due to the spallation of carbon, oxygen and iron, respectively. The enhanced production of odd-Z nuclei in the cosmic ray abundances is also thought to be caused by spallation reactions.

Measuring the abundances of cosmic rays can provide valuable information about their typical path length and propagation time. Assuming that the spallation process is sufficiently well-known, the path length of cosmic rays can be determined from the observed differences in the abundances in the regions of lithium-beryllium-boron and scandium-manganese. According to those calculations the typical path length of cosmic rays is of the order of 5 g/cm^2 . Using this value for the path length, the mass density within intergalactic space (1 proton/cm^3) and a propagation velocity close to the velocity of light one can crudely estimate that the cosmic-ray lifetime is around $3 \cdot 10^6 \text{ y}$.

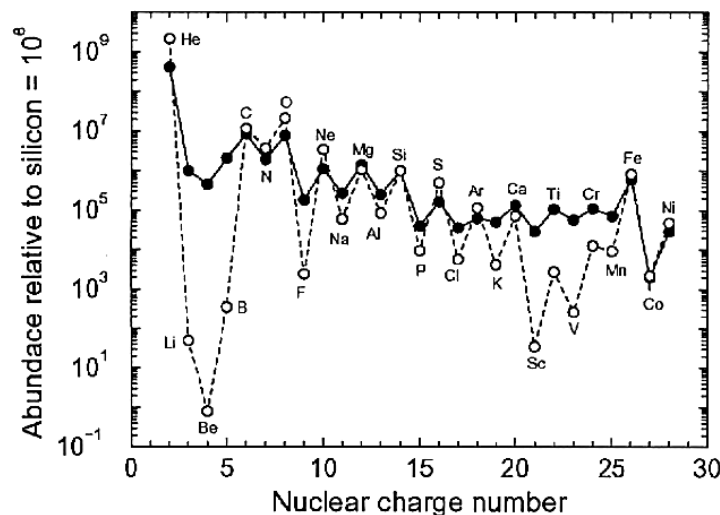


Plate 2. Chemical composition of galactic cosmic rays (*closed circles connected by a solid line*) compared to the abundances of these elements in our solar system (*open circles connected by a dashed line*) normalised to Si = 10^6

A more accurate determination of cosmic-ray lifetime requires information about the isotopic distribution of cosmic-ray abundances. Since 1997, the Cosmic Ray Isotope Spectrometer (CRIS), launched aboard the Advanced Composition Explorer (ACE), has provided this information. Abundances of long-lived radioactive secondary cosmic rays such as ^{10}Be , ^{26}Al , ^{36}Cl or ^{54}Mn can be used to derive the confinement time of cosmic rays in the galaxy. These radioactive “clock” abundances reflect the balance between production by spallation in the interstellar medium and decay or escape during the propagation time. ^{10}Be is an optimum radioactive “clock” since its lifetime ($1.5 \cdot 10^6$ y) is of the same order as the estimated cosmic-ray lifetime ($3 \cdot 10^6$ y). The observed abundance of ^{10}Be is only 20%–30% of what was expected according to cosmic-ray propagation calculations relative to other lithium-beryllium-boron isotopes. From this, the cosmic-ray lifetime is estimated to be of the order of $2\text{--}3 \cdot 10^7$ y.

Other radioactive isotopes such as ^{49}Vd and ^{51}Cr produced during cosmic-ray transport through the galaxy show evidence of decay due to electron capture. This decay happens only at energies low enough for the stripped nuclei to catch an electron from the interstellar medium. If cosmic rays were accelerated incrementally by multiple shocks over a period of time, the effects of decay during the lower energy stages would be observed in the parent/daughter ratios at higher energies.

All these techniques used to describe cosmic rays rely on the description of spallation reactions in cosmic-ray propagation models. Today, measurements of cosmic-ray isotopic abundances are so accurate that the main source of

uncertainty in the cosmic-ray propagation models comes from the description of spallation reactions. Though qualitatively well-known, the present understanding of spallation reactions is insufficient for the degree of accuracy now required in cosmic-ray interpretation.

2.2 Production of Radioactive Nuclear Beams

Nowadays, one of the major challenges of Nuclear Physics is to extend the present limits of the nuclides chart towards the driplines. The production of nuclides far from stability opens the possibility of investigating the isospin degree-of-freedom over a wide range. During the last decades, experimental programs devoted to investigating the structure of nuclei far from stability have made an enormous progress [37]. Some outstanding discoveries are the modification of shells with neutron excess [38], the discovery of halo or skin matter distributions in loosely bound nuclei close to the driplines [39], the observation of new resonant modes in neutron-rich nuclei [40] or the identification of new radioactive decays such as proton and two-proton radioactivity [41]. Nuclei far from stability also have strong implications in Nuclear Astrophysics [42], investigations of fundamental interactions [43], solid-state physics [44] and nuclear medicine [45].

However, the production of nuclides approaching the driplines still represents a technical challenge. Seventy years ago, I. Curie and J.F. Joliot artificially produced the first radioactive nuclear species by bombarding boron and aluminium foils with α particles emitted by radioactive sources. Later on, neutron-induced fission and reactions induced by energetic particles produced in particle accelerators made it possible a rapid increase in the number non-stable nuclear species being investigated. Nowadays about 3600 exotic nuclides are known. With these nuclides one has been able to explore the proton dripline up to the region of lead. The neutron dripline, however, has only been explored up to neon. Since mass models predict that around 6000 different nuclides could be bound by nuclear force, most of the 2400 nuclides not yet observed are assumed to be on the neutron-rich side of the β -stability valley.

Along with some technical issues related to high-current particle accelerators, ion sources or identification and separation methods, the choice of the reaction mechanism plays a major role in optimising the intensity of exotic beams. While different reaction mechanisms can be used to produce these nuclei, heavy-ion collisions at low energies such as fusion, deep inelastic or multi-nucleon transfer can only be applied with thin targets, thus limiting the intensity of exotic beams. Fragmentation or spallation reactions at high energies seem better suited for this task. These two reaction mechanisms also cover large ranges in excitation energy and N-over-Z ratio, leading to the production of a large variety of final nuclides, particularly neutron-rich ones [46].

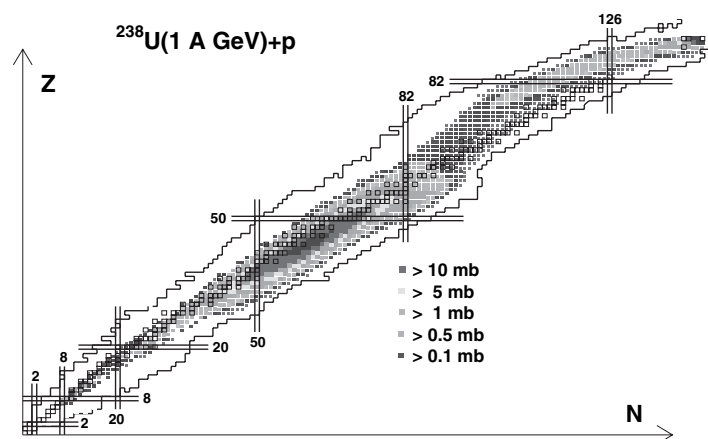


Plate 3. Residual nuclei produced in the interaction of 1 GeV protons with ^{238}U with a cross section larger than $10\ \mu\text{b}$. Measured cross sections are represented in form of chart of the nuclides and are coded by different grey hatching according to five lower limits between 0.1 and 10 mb. Data taken from [47]; cf. Plate 19 in the Colour Supplement

Spallation reactions induced by relativistic protons on various targets have been used for the last 35 years to produce nuclei far from stability at Isolde/CERN. Figure 3 represents the distribution of production cross sections that have been measured for reactions induced by 1 GeV protons on ^{238}U with a cross section larger than $10\ \mu\text{b}$ [47]. Thanks to the recent investigations on spallation reactions, one now has a good understanding on the reaction mechanisms leading to the production of a given nucleus and can thus optimise the reaction parameters in order to increase the intensity of exotic beams of interest.

In Fig. 3 three groups of nuclei can be distinguished. The high- Z group corresponds to residues produced in a nucleon or cluster evaporation process. These residues are mostly neutron-deficient covering a large range in atomic number from the target nucleus down to a charge of 65. On average, for a given element, the isotopic distribution covers around 20 different nuclides with a cross section larger than $10\ \mu\text{b}$ and the production cross section decreases with increasing distance from the β -stability line. The final isotopic residue distribution is basically determined by the neutron to proton evaporation ratio and the amount of excitation energy induced in the first stage of the collision. Since proton and neutron evaporation is governed not only by the respective binding energies but also, in the case of protons, by the Coulomb barrier, the equilibrium between these two processes is reached on the neutron-deficient side of the β -stability line. This is why most of the evaporation residues lie on this side of the valley of stability, defining the so-called “evaporation corridor”. The large fluctuations in excitation energy induced

by the collision also facilitate the production of extremely neutron-rich and neutron-deficient residual nuclei [46]. The initial excitation energy induced in the collision defines the length of the evaporation chain and consequently the mass loss of the final residues with respect to the initial target nucleus.

In the intermediate-mass region in Fig. 3, one can identify a second group of residues with atomic numbers between 23 and 65, covering a large range in neutron excess. In this case, the isotopic distributions are broader, populating on average some 25 isotopes of a given element. The distribution is centred to the right of the stability line, and the production cross-sections are much larger than those observed for most of the evaporation residues. These nuclei are interpreted as being produced by fission of the initial pre-fragments. The neutron excess of the fissile target nucleus is preserved in the fission process. This neutron excess can be partially lost when fission takes place at high excitation energies, namely by pre- and post-fission neutron evaporation which broadens the final isotopic distribution of fission residues [48]. The observed mass distribution of fission residues is mostly symmetric, in contrast to the well-known asymmetric distribution of residues produced in the low-energy fission of ^{238}U , indicating that fission occurs at high excitation energy. Nevertheless, very peripheral collisions induce low-energy fission, populating the most neutron-rich residues [49].

Finally, the third group of nuclides that can be distinguished in Fig. 3 is located below $Z = 23$. Here the isotopic chains are narrower due to the fact that the driplines are closer to the stability line. On average, the isotopic distributions are centred on the neutron-rich side of the stability line and the production cross sections increase with decreasing Z . Recent investigations have shown that in the case of reactions induced by 1 GeV protons on heavy target nuclei like ^{238}U , these residues are produced in a binary evaporation process. In fact, a detailed analysis of the kinematic properties of these nuclides has revealed that the most neutron-rich residues are produced in very asymmetric fission processes [50]. For more violent spallation reactions this binary de-excitation mechanism has been found to be replaced by a multi-fragmentation decay, leading to the production of light residues [51] (see Sect. 5.4 for a detailed discussion).

Conclusions drawn so far on spallation reactions are also valid for projectile fragmentation. The differences in the production rates from these two reaction mechanisms are mostly of technical nature. A detailed discussion of the Isol and in-flight techniques can be found in references [52] and [53], respectively.

2.3 Spallation Neutron Sources

Neutron beams are a fundamental tool in solid-state and material physics investigations, but also in applied research and industrial applications. In addition, some years ago it was proposed that an external neutron source could be used to feed sub-critical reactors, in order to burn long-lived nuclear

waste [32, 33] or to produce energy [34]. These ideas led to the concept of the Accelerator-Driven System (ADS) which is presently being considered worldwide as a possible solution for the nuclear material management problem, including proliferation and as an alternative to geological nuclear waste disposal.

A common issue in all these applications is the neutron source. Although nuclear reactors have been used to produce huge neutron fluxes, they cover a limited energy range and are not suitable for many of the present applications, in particular for ADS. Other possible neutron sources are listed in Table 1. According to this comparison, relativistic ion-ion collisions have a rather poor neutron economy. Muon-catalysed fusion, bremsstrahlung induced fission or deuterium on beryllium reactions produced similar neutron multiplicities per energy unit. However, the best neutron economy is obtained in spallation reactions induced by relativistic protons on a lead target which produce an average of 30 neutrons per incident proton at 1 GeV. This is why several spallation sources are currently under construction or study in the USA (SNS [28]), Europe (SINQ [29], ESS [30]) and Japan (NSP [31]).

Table 1. Neutron economy in different neutron sources

Reaction	Neutrons Per		Energy	Neutrons Per GeV
	Projectile	Beam Intensity		
Neutron-induced fission	2.5	$10^{14} \text{ n/cm}^2/\text{s}$		
$\text{d}+\text{C} \rightarrow \mu^- + \text{d} + \text{t}$	24	10^{16} s^{-1}	1.5 GeV	15
$\text{U}+\text{U}$	1000	10^{12} s^{-1}	238 GeV	4.2
$\text{e}^- + \text{W} \rightarrow \gamma + \text{U} \rightarrow \text{fission}$	0.1	$3 \cdot 10^{14} \text{ s}^{-1}$	0.1 GeV	15
$\text{d}+\text{Be}$	10	10^{16} s^{-1}	1 GeV	10
$\text{p}+\text{Pb}$	30	10^{16} s^{-1}	1 GeV	30

In the case of ADS, the required proton beam intensity of 10 mA for the spallation neutron source is so high that the spallation target assembly design becomes a real challenge. Among other things, activation, long-term radio-toxicity or corrosion problems in the spallation target due to the production of spallation residues should be considered. Structural materials, in particular the window between accelerator and target, will also suffer due to high intensities and long irradiation periods.

Indeed, two main factors that influence the design and construction of such a target are neutron yields and the nature and kinematic properties of residual nuclei produced in the reaction. Neutron production has to be described in terms of neutron multiplicity and its spatial and energy distribution. Neutron multiplicity determines both, the current and the beam energy of the proton-driver accelerator, while their energy and spatial distri-

bution will shape the geometry of the spallation target and the shielding to high-energy neutrons. As already mentioned, the production and kinematic properties of residual nuclides are of interest with respect to activation and radiation damage.

For these reasons, a detailed design of these technological applications requires powerful and reliable computational tools that can accurately predict particle and nuclide production in spallation reactions. Figure 4 gives as an example the partial and total simulated activity induced in a cylindrical lead target by a 1 mA, 1 GeV energy proton beam during one year of irradiation and 10^6 years of cooling time. It is evident that the cooling times and the total activities induced in the target are neither negligible nor easy to handle.

However, the reliability and predictive power of the present simulation codes describing spallation reactions are still limited, as can be seen from Fig. 5. Here, we illustrate the activities ratio calculated by using two different intra-nuclear cascade simulation codes, Bertini [20] and INCL4 [25] for the same target and proton beam as in Fig. 4. The results of these calculations

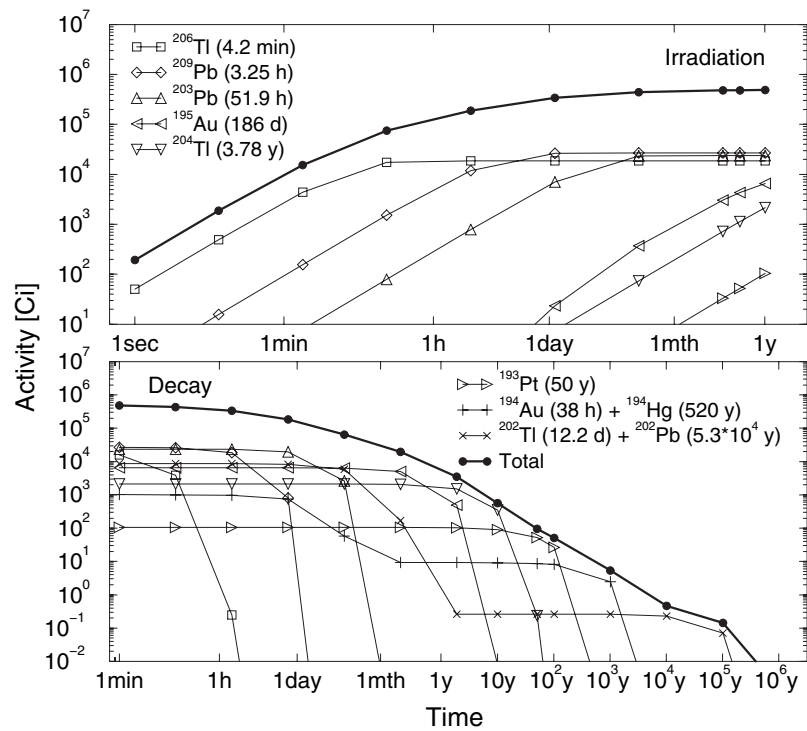


Plate 4. Total and partial activities induced in a ^{nat}Pb target ($L = 100$ cm, $R = 10$ cm) by a 1 GeV, 1 mA intensity proton beam as a function of the irradiation (*upper panel*) and cooling time (*lower panel*). The simulation has been performed with the LAHET transport code [54] and taken from [55]

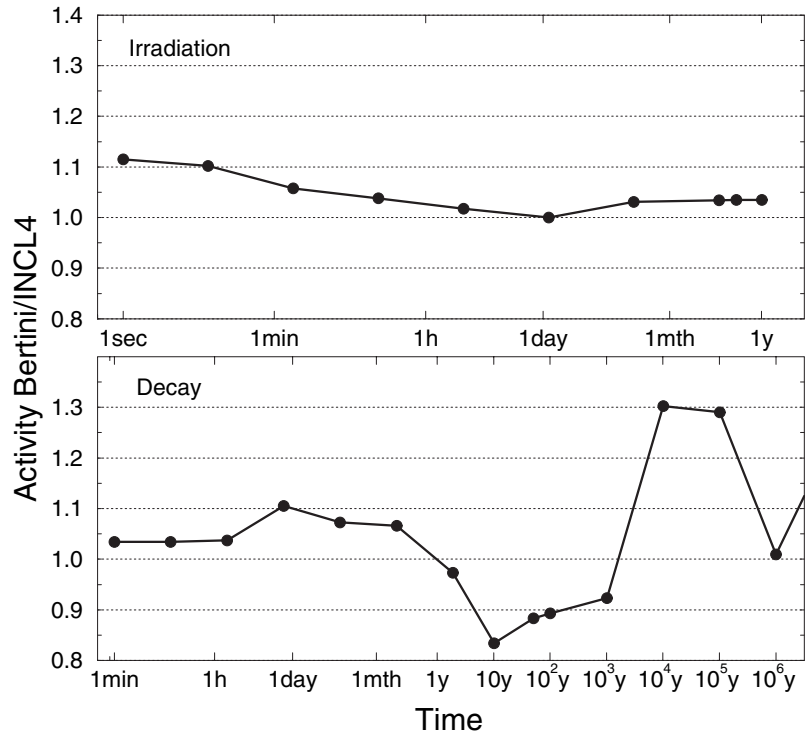


Plate 5. Ratio of the total activities induced in a ^{nat}Pb target under the same irradiation conditions as in Fig. 4, given by two different intra-nuclear cascade codes Bertini [20] and INCL [25], during the irradiation (*top*) and decay (*bottom*) phases. Data taken from [55]

show that the discrepancies between different codes can reach 30%, which are unacceptable for technological applications such as the ones we are discussing here.

From these results one can appreciate the good qualitative understanding of spallation reactions that one has gained from more than 50 years of research. However, this knowledge is clearly insufficient for some new technological applications or even to fully interpret the new accurate data obtained on cosmic rays. This insufficiency has catalysed experimental programs around the world, which currently investigate spallation reactions in an effort to develop codes that are able to reliably predict particle and nuclide production in these reactions.

3 Experimental Techniques

3.1 Neutron Production

The neutron flux produced in spallation reactions depends strongly on the projectile-target combination. As a general rule, the heavier the target nucleus the larger the neutron excess and the neutron yield. The gain factor between heavy and light targets is approximately a factor of five. However, the radio-toxicity induced in the spallation target could be drastically reduced by using lighter targets [56]. In addition to the neutron yields, reliable information on the energy and spatial distribution of the neutrons is required. Different experimental devices are needed to describe the neutron production in spallation reactions.

Measurement of Neutron Yields

Neutron multiplicities can be investigated using liquid scintillator-based detectors with a large angular acceptance. The BNB (Berlin Neutron Ball) [57] and ORION [58] detectors used by the NESSI collaboration (Berlin-Ganil-Jülich) may serve as examples. This collaboration has performed a large experimental program to determine the neutron yields produced in thin and thick targets for a large range of primary projectiles and energies. Experiments were done at GANIL (France) [58], COSY (Germany) [57] and CERN [59].

Figure 6 shows representative results obtained by this collaboration at Jülich with the BNB detector. For different target materials, neutron multiplicity saturates at a given target thickness which increases with the proton energy. Considering that the mean free path of relativistic protons on lead is of the order of 15 cm, the observed saturation at about 30 cm indicates that these relativistic protons produce on average two nuclear collisions in lead.

Energy and Spatial Distributions

Specific experimental setups are needed to measure the spatial and energy distribution of the neutrons produced in spallation reactions, the experiments performed by the Transmutation collaboration at Saturne (France) exemplify this. Such data were obtained by using two different experimental techniques to cover the full energy range of the neutrons produced in the reaction. The detection of neutrons with energies below 400 MeV was based on a measurement of their time of flight, i.e. the time difference between the incident proton, tagged by a plastic scintillator, and a signal provided by a neutron-sensitive liquid scintillator [60]. Neutrons with higher energies were measured using (n,p) scattering on a liquid hydrogen converter and reconstruction of the proton trajectory in a magnetic spectrometer [61]. An additional collimation system was used to determine the angular distribution of the neutrons.

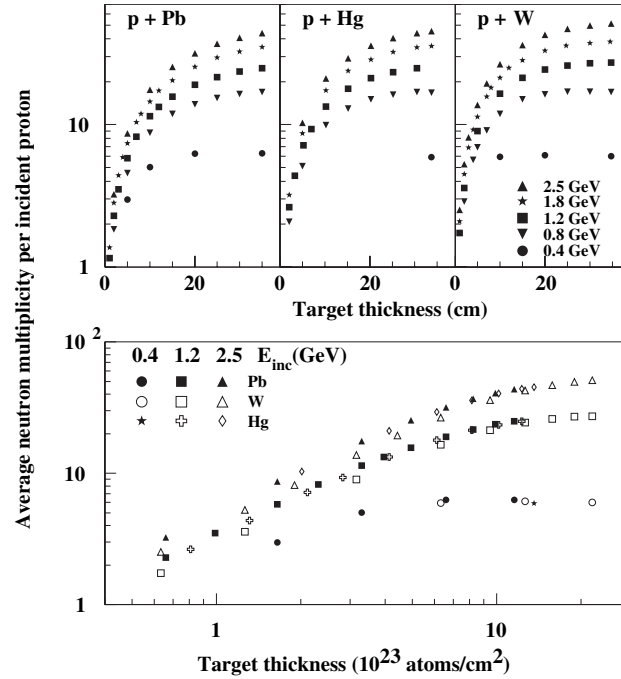


Plate 6. Average neutron multiplicity per incident proton as a function of target thickness and beam energy for lead, mercury and tungsten targets obtained by the NESSI collaboration [57]

Neutron production was studied in reactions induced by protons and deuterons with energies between 0.8 and 1.6 GeV on thin and thick aluminium, iron and lead targets [62]. Figure 7 shows the results obtained on a 2 cm thick lead target with a 1.2 GeV proton beam, simulating a spallation neutron source. High-energy neutrons emitted at small angles are produced during the first stage of the collision while low-energy neutrons emitted isotropically correspond to the evaporation phase. Measurements done with thicker targets provided information about the inter-nuclear cascade.

3.2 Production of Light Charged Particles

Neutron production in spallation reactions is always accompanied by the emission of light charged particles. Their nature and kinematic properties provide information on the dynamics of the hot nuclei, from which they are emitted, but also play a major role causing damages in target and structural materials used in spallation targets, the production of hydrogen and helium isotopes being particularly relevant.

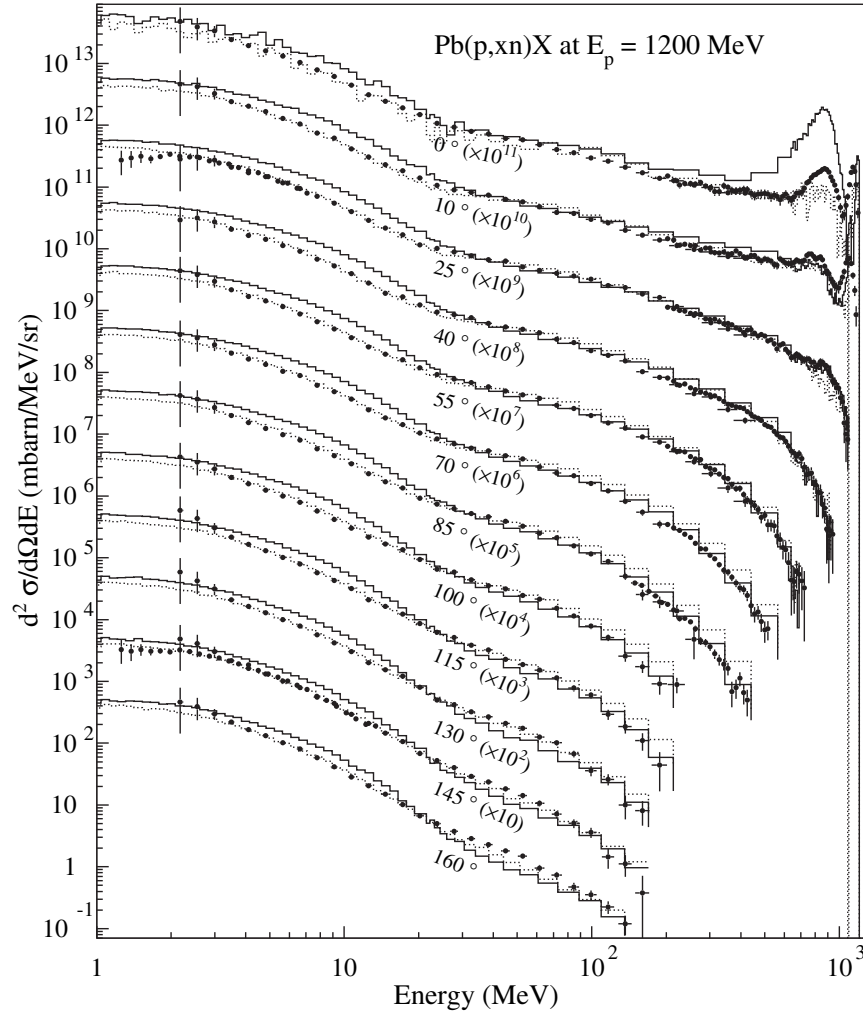


Plate 7. Double-differential cross sections for neutron production in reactions induced by a 1.2 GeV proton beam on a 2 cm thick lead target [62]. The histograms show calculations using the Bertini INC code [20] while the dotted lines correspond to calculation done with the Cugnon INCL code [25]

Although the emission of light charged particles in spallation reactions has been investigated since the sixties [7, 8], new experimental programs have recently been initiated. They aim at obtaining more accurate data, in particular concerning correlations between charged particles and neutrons. Some of the most outstanding experiments are those performed by the NESSI [63] and PISA [64] collaborations at the proton accelerator facility COSY in Jülich, Germany.

3.3 Production of Residual Nuclides

Residue production in spallation reactions can be investigated using two different experimental approaches. In the standard approach, the reaction is induced by light energetic projectiles on a heavy target. As the recoil velocity of the residues produced in this reaction is not sufficient to make them leave the target, the spectroscopy of γ -rays or masses is used to identify them. The main limitation of this technique is that for most of the residues the measurement is done after β -decay and consequently only isobaric identification is possible.

Investigation of the reaction using inverse kinematics is better suited to unambiguously identify spallation residues. In this case a heavy nucleus is accelerated to relativistic energies and impinges on a light target. Due to the kinematic conditions, the reaction residues leave the target easily and can be identified in flight.

Inverse Kinematics

Outstanding experiments were performed by a German-French-Spanish collaboration at GSI (Germany). The technique used in these experiments made use of inverse kinematics and provided full identification of the reaction residues with respect to their mass number (A), atomic number (Z) and kinematic properties.

These experiments were performed at the SIS synchrotron at GSI. Primary beams of ^{56}Fe [51], ^{136}Xe [65], ^{197}Au [68, 69], ^{208}Pb [66, 67] and ^{238}U [70, 71] were accelerated up to an energy of 1 A GeV and impinged on a liquid hydrogen or deuteron target. At this energy all residues of the reaction are predominantly fully stripped, i.e. bare ions. The achromatic high-resolution magnetic spectrometer FRS [72] equipped with an energy degrader, two position sensitive scintillators and a multi-sampling ionisation chamber allowed one to identify all reaction residues with half lives longer than 200 ns according to their atomic and mass number. Resolving powers of $A/\Delta A \approx 400$ and $Z/\Delta Z \approx 200$ were achieved with this technique, and production cross sections were obtained with nearly 10% accuracy. In addition, the high resolving power of the magnetic spectrometer makes it possible to determine the recoil velocity of the reaction residues. This information is relevant for the characterization of the damages induced by the radiation in the accelerator window or the structural materials of an ADS. More details about these experiments can be found in references [66, 67, 68, 69].

Figures 3 and 8 compile the data obtained at GSI for the reactions $^{238}\text{U}(1 \text{ A GeV})+\text{p,d}$, $^{208}\text{Pb}(1 \text{ A GeV})+\text{p,d}$ and $^{56}\text{Fe}(1 \text{ A GeV})+\text{p}$ are presented at the top of different charts of the nuclides. About 4000 nuclei were identified in these reactions. As can be seen from this figure, the spallation residues produced with heavy projectiles populate two different regions of the chart of the nuclides. The high- Z region corresponds to spallation-evaporation

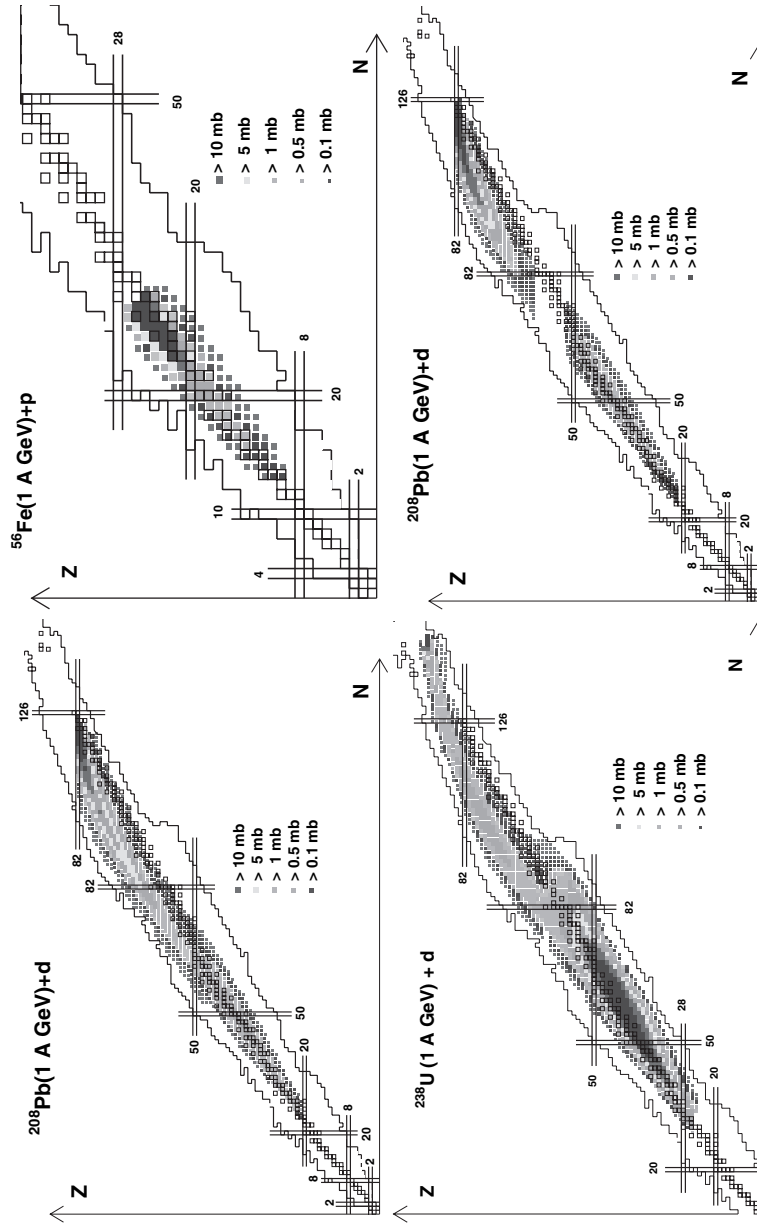


Plate 8. Cross sections for the production of spallation residues, measured at GSI for the reactions $^{238}\text{U}(1 \text{ A GeV})+d$, $^{208}\text{Pb}(1 \text{ A GeV})+p$, and $^{56}\text{Fe}(1 \text{ A GeV})+p$. The experimental data are presented in form of a chart of nuclides and are coded by different grey hatching according to five lower limits between 0.1 and 10 mb. Data taken from [47, 51, 67]; cf. Plate 20 in the Colour Supplement

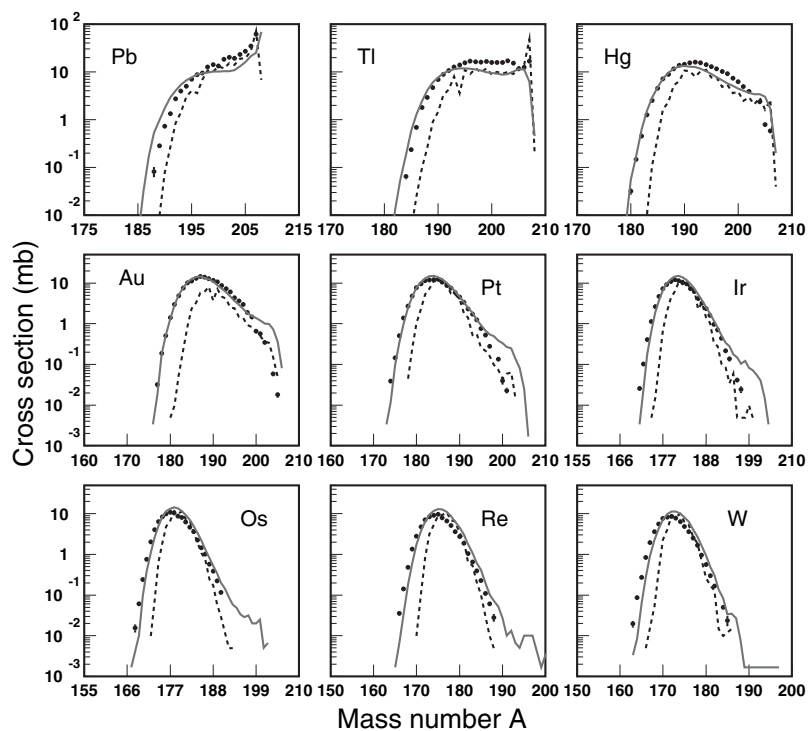


Plate 9. Isotopic distributions of the cross sections for the production of isotopes of some heavy elements, measured at GSI for the reaction $^{208}\text{Pb}+p$ at 1 A GeV [67]. The data are compared with two model calculations. The *dashed lines* correspond to results obtained with the Lahet code [54] while the *solid line* curves represent results obtained with the intra-nuclear cascade model of Cugnon [86] coupled to the evaporation-fission code ABLA [73, 74]; cf. Plate 21 in the Colour Supplement

residues which populate the so-called evaporation-residue corridor. The second, medium-mass region corresponds to residues produced in spallation-fission reactions. Though inherently different, both fission and evaporation reactions contribute to the production of residues.

The measured isotopic production cross sections for some selected elements are presented in Fig. 9, which clearly shows the quality of the measured data that can be used to benchmark any model calculation.

Figure 10 shows the average kinetic energy in the centre-of-mass frame of fragmentation and fission residues produced in the reaction $^{208}\text{Pb}+p$ at 1 A GeV as a function of their atomic number [67]. These results show an increase of the recoil velocity of the fragmentation residues for the most violent collisions, leading to the production of lighter residues. The large kinetic energies of the fission residues are a key parameter used to evaluate the heat load of the spallation target.

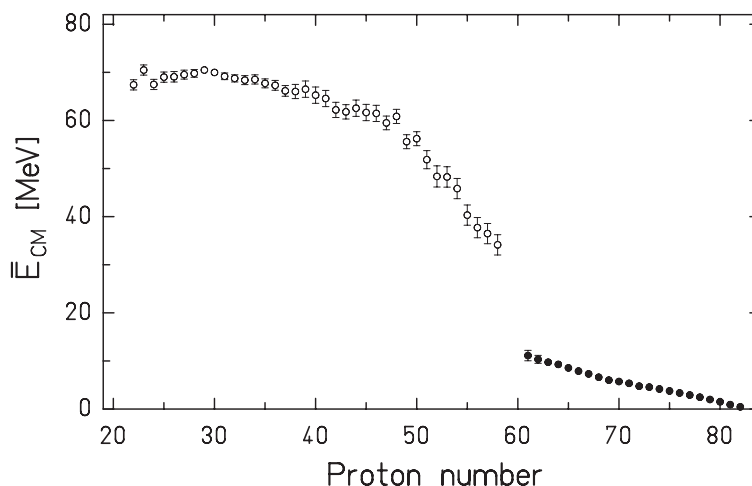


Plate 10. Average kinetic energy in the centre-of-mass frame of fragmentation (*closed symbols*) and fission (*open symbols*) residues produced in the reaction $^{208}\text{Pb}+\text{p}$ at 1 A GeV as a function of their atomic number [67]

Direct Kinematics

Gamma-ray spectroscopy makes it possible to investigate the production of spallation residues in direct-kinematic reactions. Although this method is restricted to isobaric identification after β -decay, it is possible to determine the primary production cross-sections for some shielded isotopes. This technique consumes less beam time than the method based on inverse kinematics, and production cross-sections can be obtained over broad energy ranges for selected isotopes, as shown in Fig. 11. The method can also be applied to thin and thick targets at low and high energies and complements the inverse-kinematic technique in this aspect.

Presently in Europe there are two main experimental programs using this technique. Michel and collaborators at the University of Hannover (Germany) perform experiments mainly at Saturne (France), TSL (Sweden) and PSI (Switzerland) [75]. In some of these experiments different target material were irradiated with protons in the energy range 20–2600 MeV. Figure 11 shows excitation functions for the production of some nuclei in collisions induced by protons on lead [76]. These results clearly indicate that low-energy reactions mainly produce residues close to the target nucleus, while most of the reaction residues further away from the target are produced by energetic particles. Similar experiments are also performed at the ITEP in Moscow (Russia) [77].

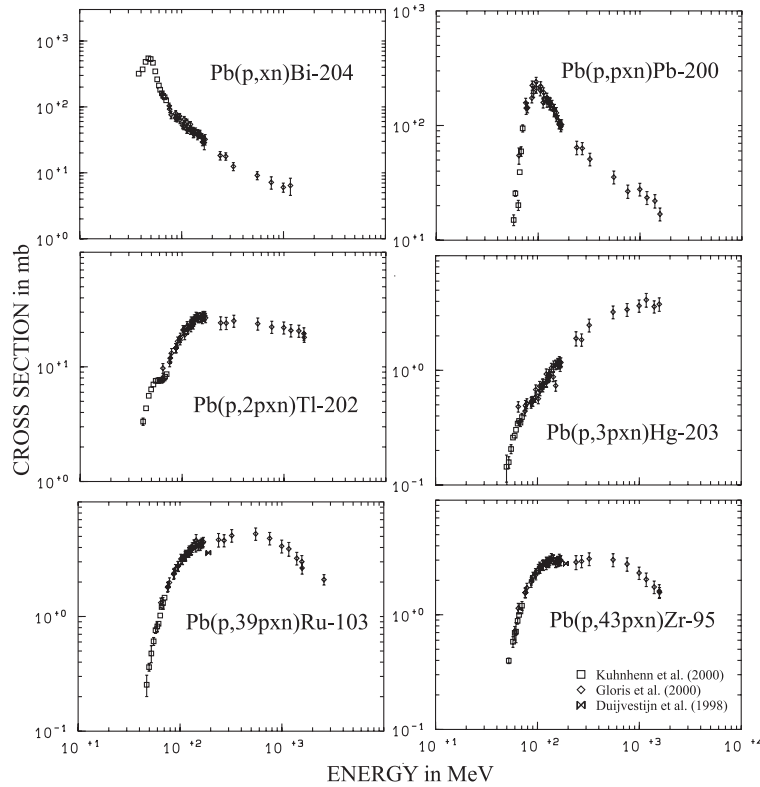


Plate 11. Excitation functions of the production cross sections for some selected nuclides produced in the interaction of protons with lead, measured using γ -spectroscopy techniques [76]

3.4 Reactions Induced by Projectiles of 20–200 MeV Energy

Reactions induced by neutrons and light charged particles in the energy range between 20 and 200 MeV simulate the inter-nuclear cascade in the spallation target. These reactions play a major role in the multiplication and moderation of the neutrons. Only a few nucleons are emitted due to the small amount of energy dissipated in these reactions. Consequently, the resulting mass and atomic number of the residual nuclei is close to that of the target nucleus. The experiments performed in this energy range are intended to measure the double-differential cross sections for production of neutrons and light charged particles in reactions induced by protons and neutrons with energies in the range 20–200 MeV. Fission is also investigated in such reactions. Moreover, measurements of the total fission cross sections and mass distributions of fission residues for these projectile energies are in progress. These measurements will make it possible to extend the Evaluated Nuclear Data File libraries (ENDF) up to 200 MeV.

A significant number of European laboratories contribute to this experimental program, taking advantage of a large network of European facilities that deliver protons and neutrons in this energy range: KVI (Netherlands), Louvain la Neuve (Belgium) and Uppsala (Sweden). Most of them contributed to the HINDAS project of the Fifth Framework Program of the European Commission [78]. This European project also included the measurement of residue production in inverse kinematics done at GSI and light charged-particle production experiments at COSY, both in the 200–2000 MeV region.

Another important program supported by the European Commission is the n_TOF project. This project is based on a time-of-flight neutron facility recently set up at CERN. After moderation, the neutrons produced in spallation reactions induced by 20 GeV protons, from the proton synchrotron (PS), on a lead target cover an energy range between 1 eV and 200 MeV [79]. The experimental program developed at this facility includes a large number of experiments related to neutron capture and neutron-induced fission reactions. The neutron flux produced in this facility will make possible the use of radioactive targets for some of the measurements. Results from these experiments are expected in the next few years.

4 Reaction Physics and Model Description

Reactions induced by relativistic light projectiles on heavy target nuclei are generally interpreted on the basis of the two-step mechanism proposed by Serber [11]. According to this model, in the collision between projectile and target, a remnant of the heavy target nucleus gains excitation energy and angular momentum while it loses some nucleons through pre-equilibrium emission. This fast stage of the reaction is followed by a slower stage where the target remnant or pre-fragment equilibrates all its excited degrees of freedom and de-excites emitting γ -rays, nucleons, clusters or even fissioning. The following sections will deal with the main physics concepts involved in these two processes as well as the corresponding model description.

4.1 Nucleon-Nucleus Interaction

Full microscopic calculations of the interaction between heavy ions exceed by far the present computational possibilities due to the many constituents and degrees of freedom involved in these collisions. Therefore different approaches are used to simplify the description of heavy ion collisions according to the energy regime. At low energies, the reactions are governed by the nuclear potential and quantum mean-field calculations, such as the time-dependent Hartree-Fock approach (TDHF), are used [80, 81]. At relativistic energies, the wave-length of the incoming particle and subsequent collision products is of the order or smaller than the average inter-nucleon distance within the

nucleus ($\approx 10^{-13}$ cm). Consequently, classical models based on a sequence of particle-particle collisions inside the nucleus known as intra-nuclear cascade can be used.

More advanced models are based on both, a description of the nuclear mean field and the intra-nuclear collisions. These dynamical models are based on kinetic equations, for example the Vlasov equation, which describe the motion of averaged classical quantities, such as one-body distribution function, in the phase space $f(r, p, t)$. Additionally, a collision term including the Pauli blocking is accounted for in the so-called Boltzmann-Uhling-Uhlenbeck (BUU) equation [26].

Although the BUU approach has been widely used to describe the first stage of heavy-ion collisions, it does not account for many-body correlations. These correlations are fundamental to the explanation of certain processes that require symmetry breaking as in the case of multi-fragmentation. Two methods are being used to overcome this limitation. The first, known as the Boltzmann-Langevin approach, seeks to simulate the correlations using a stochastic term in the BUU equation [82, 83]. The second one is based on the solution of many-body nuclear dynamics using either a classical approach where all nucleons are localised [84, 85] or a more sophisticated one that explicitly accounts for the fermionic character of the nucleons through the antisymmetrization of the corresponding wave functions (Quantum Molecular Dynamics) [27].

The dynamical evolution of the mean field plays a minor role in spallation reactions, due to the fact that only a few nucleons are affected by the collision and consequently the nuclear density as well as the mean field are only slightly and temporarily modified. These considerations lead one to conclude that the classical approach based on the intra-nuclear cascade model is well-suited for describing this kind of collisions. Its simplicity and the correspondingly short computer time needed for its applications have made it very popular indeed.

Intra-Nuclear Cascade Models

Intra-nuclear cascade models describe the interaction of high-energy (>100 MeV) nucleons with complex nuclei as a sequence of two-body interactions between the incident particles and the individual nucleons in the target nucleus. In a similar way, the collisions of the struck nucleons with the remaining target nucleons are also considered as a sequence of two-body interactions.

These models are based on the principle that incident nucleons follow straight trajectories until they collide with a target nucleon under a given criteria determined by the nucleon-nucleon interaction cross section. Relativistic kinematics is used to describe the trajectories, and refraction and reflexion effects due to the nuclear potential are accounted for. The elastic or inelastic character of the collision is determined by in-medium nucleon-nucleon cross sections. The upper energy limit for the validity of these models is a few GeV,

and in the case of inelastic collisions, pion production and absorption are supposed to occur from the following reactions:

$$NN \rightleftharpoons N\Delta, \quad \Delta \rightleftharpoons \pi N \quad (1)$$

The momenta of the reaction products are obtained from the angular dependence of the cross section of the corresponding reaction channel and from energy and momentum conservation. Moreover, Pauli blocking is implemented using different approximations according to the model. The main idea behind these assumptions is that collision products can not populate an area of the phase space that is already occupied. Hadrons reaching the nuclear surface with an energy larger than their binding energy will escape through pre-equilibrium emission. The propagation of this cascade is followed until the total energy (kinetic plus potential plus mass) of the cascade particles drops below a certain cutoff energy or the time of the cascade propagation reaches a limiting value after which thermalisation is assumed. The final excitation energy and angular momentum induced by the cascade is obtained from particle-hole excitations due to cascade particles bound by the nuclear potential and holes produced in the initial Fermi distribution of target nucleons.

The properties of such calculated cascades were first investigated by Goldberger [13] using the Monte Carlo technique and a two-dimensional model of the nucleus. In these calculations it was assumed that the characteristics of the nucleon-nucleon collisions within the nuclear matter are the same as those in free space, except for the Pauli exclusion principle. Several similar investigations of this problem followed, the most detailed being those of Metropolis et al. [19] and Bertini [20]. The most recently developed intra-nuclear cascade codes are ISABEL by Yariv and Fraenkel [24, 88] and the Liège code (INCL) by Cugnon [25, 86, 87].

ISABEL is a generalisation of the VEGAS model [89] and can describe hadron-nucleus and nucleus-nucleus collisions up to energies of 1 A GeV. INCL describes hadron-nucleus and nucleus-nucleus collisions up to energies of 2.5 A GeV and for nuclei with mass number of 4 or less. The major differences between these two codes are related to the treatment of the nuclear medium and the criterion used to stop the cascade.

In the case of ISABEL, the nucleus is considered as a Fermi sea of nucleons. The interaction of the excited particles with the nucleons in the Fermi sea can produce new cascade particles. After each interaction the nuclear density is readjusted and the trajectory of the new excited particle is computed. The cascade stops when the energy of every cascade particle falls below a pre-defined cut-off energy given by the Coulomb barrier plus two times the binding energy. After this point is reached, the energy of all the excited particles that remain bound in the nucleus is assumed to be distributed among all the nucleons of the nucleus, leading to a thermalised pre-fragment.

By contrast, in the INCL code, the nucleus is considered to be a set of compact balls (nucleons) whose trajectories can be followed at any moment

during the intra-nuclear cascade. At each time step, two nucleons can collide if the distance between their associated trajectories becomes smaller than a pre-defined minimum. The calculation continues until the computed time reaches a given value τ_{eq} which determines the end of the pre-equilibrium stage. The latter parameter is determined from the analysis of the time evolution of the excitation energy of the system. During the thermalisation process, the kinetic energies of the emitted pre-equilibrium particles cause an abrupt reduction of the excitation energy of the nucleus. As thermal equilibrium is approached the excitation-energy decrease becomes smoother. The parameter τ_{eq} can thus be defined as the transition between these two regimes. Finally, the internal energy of the nucleus is obtained as the difference between the sum of the kinetic energies of all the excited nucleons and the energy of the ground state in a Fermi gas. Both quantities refer to the bottom of the potential well describing the nuclear interaction.

Illustrative Results

Figure 12 presents INCL predictions of excitation-energy and mass distributions of pre-fragments produced in the first stage of the reaction $p(1 \text{ GeV}) + {}^{208}\text{Pb}$. According to this calculation, the average number of pre-equilibrium nucleons emitted during the first stage of the collision is five, with the corresponding number of protons and neutrons being scaled to the neutron to proton ratio of the target nucleus. This process results in an average excitation energy of 200 MeV. This means that the average excitation energy induced by the abrasion of one nucleon is 40 MeV. The same calculation gives an average induced angular momentum of $10 \hbar$ for the pre-fragments.

From this calculation one can conclude that the number of abraded nucleons represents only a small fraction, and consequently spallation reactions

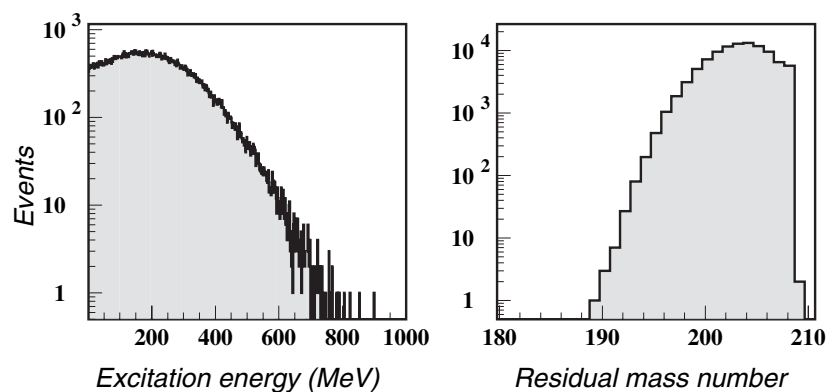


Plate 12. Distributions of excitation energy (*left panel*) and mass (*right panel*) of the pre-fragments produced in the first stage of the reaction $p(1 \text{ GeV}) + {}^{208}\text{Pb}$ calculated by using the code INCL

are not expected to cause large modifications in the target nucleus in terms of deformation. Though the average angular momentum induced is quite small, the final pre-fragment gains a large amount of excitation energy. Consequently, the degree of freedom more affected by the collision is the internal energy of the nucleus. This is why spallation reactions are thought to be an optimum tool for investigating the dynamics of nuclear matter under extreme conditions of excitation energy or temperature.

4.2 Pre-fragment De-excitation: Statistical Model

The statistical model describes the de-excitation of the pre-fragment produced in the first stage of the collision assuming the compound-nucleus hypothesis of Bohr [90]. According to this hypothesis, at the end of the pre-equilibrium phase the pre-fragment reaches statistical equilibrium defined by a given value of the excitation energy and the total angular momentum. This assumption of equilibrium implies that all possible decay channels open for the equilibrated system have, on the average, the same probability of being populated.

In the case of spallation reactions, the excitation energy induced during the first stage of the collision is so high that many possible decay channels are available. In particular, emission of nucleons or light clusters, fission, emission of intermediate-mass fragments and γ radiation, are possible decay channels. Since the amount of energy release in any of these emissions is limited, the statistical model assumes that the initial pre-fragment will undergo a sequential decay process, or evaporation chain, until the final residual nucleus reaches its ground-state configuration. The model assumes that statistical equilibrium is reached after every evaporation step.

Using this hypothesis and the principle of detailed balance, Weisskopf and Ewing [12, 91] introduced the first formulation of the decay rate of the compound nucleus. The principle of detailed balance for two systems a and b in statistical equilibrium and with level densities ρ_a and ρ_b states that the depletion of the states of system a by transitions to b equals their increase by the time-reversed process $b \rightarrow a$. Thus if $R_{ab} = \Gamma_{ab}/\hbar$ is the decay rate (probability per unit time) for transitions from a to b and $R_{ba} = \Gamma_{ba}/\hbar$ is the decay rate for the inverse process, one has:

$$\rho_a \Gamma_{ab} = \rho_b \Gamma_{ba} \quad (2)$$

Weisskopf's statistical model was improved by Hauser and Feshbach [92] who introduced a proper quantum-mechanical treatment of the angular momentum. Another formulation of the statistical model is based on the transition state method that Bohr and Wheeler proposed to describe the fission decay channel [90]. According to this formalism, the decay rate of a fissioning nucleus is not determined by the statistical configuration of the final state but rather by the configuration of states above the potential barrier that governs this process. The following sections will present the formalism used to

describe the decay rates of the different de-excitation channels a hot nucleus can follow.

Particle and γ Emission

If we consider an excited nucleus in thermal equilibrium which decays into a final nucleus emitting a nucleon or light nucleus, the final state of the system is characterised by the fraction of the phase space occupied by the emitted particle and the number of states of the daughter nucleus that can be populated in this transition. The number of final states can be statistically calculated as a level density. In this particular case, the decay width of the time-reversed process can be obtained from the cross section of the capture process (fusion) and the velocity of the emitted particle [93, 94].

Taking these considerations into account and assuming an initial nucleus characterised by an excitation energy U_i and total angular momentum J_i , which decays into a final nucleus (U_f, J_f) emitting a nucleon or a light nucleus ν with a given kinetic energy ϵ_ν , spin s_ν and orbital angular momentum l , the decay width for this process can be calculated according to:

$$\Gamma(U_i, J_i; U_f, J_f, s_\nu) = \frac{(2s_\nu + 1)k_f^2}{2\pi^2} \sigma_{fi}(U_i, J_i) \frac{(2J_f + 1)\rho(U_f, J_f)}{(2J_i + 1)\rho(U_i, J_i)} \quad (3)$$

The energies U_i and U_f are related by $U_i = U_f + \epsilon_\nu + S_\nu + B_\nu$, where S_ν represents the separation energy and B_ν the Coulomb barrier for the emitted particle. The cross section for the time-reversed process, fusion, can be calculated in term of the transmission coefficients as:

$$\sigma_{fi}(U_i, J_i) = \frac{\pi}{k_f^2} \frac{2J_i + 1}{(2s_\nu + 1)(2J_f + 1)} \sum_{S=|J_f-s_\nu|}^{J_f+s_\nu} \sum_{l=|J_i-S|}^{J_i+S} T_\nu^l(\epsilon_\nu) \quad (4)$$

where $S = J_f + s_\nu$ is the channel spin. Spin-orbit coupling is neglected in this expression. The quantity T_ν^l incorporates the effects of the Coulomb and the centrifugal barriers together with the nuclear potential. By combining the previous equations one obtains the average decay rate R_ν of a nucleus emitting a particle ν :

$$R_\nu(U_i, J_i; U_f, J_f) = \frac{1}{h} \frac{\rho(U_f, J_f)}{\rho(U_i, J_i)} \sum_{S=|J_f-s_\nu|}^{J_f+s_\nu} \sum_{l=|J_i-S|}^{J_i+S} T_\nu^l(\epsilon_\nu) \quad (5)$$

Finally, the decay width of a system characterised by an excitation energy U_i , an angular momentum J_i and the nature of the emitted particle ν , results from summing over all possible final values of the angular momentum J_f and integrating over all possible kinetic energies ϵ_ν of the emitted particle:

$$\begin{aligned}
\Gamma_\nu(U_i, J_i) &= \hbar \int_0^{U_i - B_\nu} \sum_{J_f} R_\nu(U_i, J_i; U_f, J_f) d\epsilon_\nu \\
&= \frac{2m_\nu(2s_\nu + 1)}{2\pi^2 \hbar^2 \rho(U_i)} \int_0^{U_i - B_\nu} \epsilon_\nu \sigma_{fi} \rho(U_f) d\epsilon_\nu
\end{aligned} \tag{6}$$

where m_ν is the mass of the emitted particle. This formulation of the decay process of an excited nucleus in thermal equilibrium indicates that the whole process is governed by the transmission coefficient for the emitted particle and the level densities of the initial and final nuclei involved in the decay. An accurate determination of the decay widths requires a correct evaluation of these quantities. The transmission coefficients T_ν^l can be obtained solving the corresponding Schrödinger equation with an optical potential describing the particle–nucleus interaction as discussed in [42]. However, the high excitation energies involved in spallation reactions lead to an enormous number of possible decay channels and very long evaporation chains that make it impossible to perform calculations within a reasonable short time. For this reason, realistic approximations are frequently used to simplify the calculations. One of the most commonly used approximations for evaluating the inverse cross-section is to approach the nuclear potential by a black disk. In this case, the transmission coefficient is represented by a step function with a maximum value determined by the geometrical interaction cross-section. In addition, if the excitation energy of the decaying nucleus is sufficiently high, the upper integration limit in (6) can be extended to infinity. Correspondingly, the following equation can then be used to calculate [95] the decay width for a given particle ν emission as:

$$\Gamma_\nu(U_i) = \frac{1}{2\pi \rho(U_i)} \frac{4m_\nu R^2}{\hbar^2} T_f^2 \rho(U_i - S_\nu - B_\nu) \tag{7}$$

where R is the radius of the nucleus and T_f the temperature of the final residue. The differences between the exact calculation (6) and the approximated one (7) can be accounted for by using effective Coulomb barriers B_ν for the emitted particles [96].

In analogy to the particle decay, the average decay rate of an excited nucleus emitting γ rays can be written as [97]:

$$R_\nu(U_i, J_i; U_f, J_f) = C_\lambda(\epsilon_\gamma) \epsilon_\gamma^{2\lambda+1} \frac{\rho(U_f, J_f)}{\rho(U_i, J_i)} \tag{8}$$

where the two first terms represent the matrix element of the transition, $C_\lambda(\epsilon_\gamma)$ is the radiative strength function, and the last term corresponds to the phase-space ratio. As the statistical emission of γ radiation occurs predominantly through the giant dipole resonance, its decay width can be written as:

$$\Gamma_\gamma(U_i) = \sum_{J_f} \int_0^{U_i - \epsilon_\gamma} \epsilon_\gamma^3 C_\lambda(\epsilon_\gamma) \frac{\rho(U_f, J_f)}{\rho(U_i, J_i)} d\epsilon_\gamma \tag{9}$$

Comparing (6) and (9), one sees that the emission of particles is favoured over dipole γ rays by a factor of $\approx 10^{-7}$. Thus, γ -ray emission is important only in the later stages of the evaporation chain at energies around or below the particle separation energy. Taking $\epsilon_\gamma = S_\nu$ and using the power approximation for the radiative strength function [98] and the constant temperature model [99], (9) can be parameterised as [100]:

$$\Gamma_\gamma(\epsilon_\gamma = S_\nu) = 0.624 \cdot 10^{-9} A^{1.6} T^5 \text{ MeV} \quad (10)$$

Here, A is the mass of the mother nucleus and $T = 17.6/A^{0.699}$ MeV is the nuclear-temperature parameter of the constant-temperature model.

Fission

Fission is the consequence of a large-scale collective motion of nucleons inside the nucleus that splits the initial nucleus into two fragments. This process can occur spontaneously for certain nuclei, or can be induced in nuclear reactions where part of the internal or excitation energy gained by the fissioning nucleus is transformed into collective motion that provokes a significant deformation of this nucleus. As the nucleus deforms (at constant volume), the magnitude of the negative contribution of the surface energy to the binding energy increases and opposes deformation. At the same time, the increasing separation of the nuclear charges (the protons) reduces the Coulomb energy and favours further elongation. The competition between these two energy changes has the effect of creating an angular-momentum-dependent potential energy barrier in the deformation coordinate. The top of the barrier corresponds to the “no return” or saddle point, the height of which is known as the fission barrier B_f . Beyond the saddle point, the system descends to the point of scission where it separates into two fragments.

The fission barrier, which has a magnitude of only a few MeV, results from the partial cancellation of two quantities (the surface and Coulomb energies), which are both of the order of several hundred MeV. This indicates that fission is very sensitive to small energy variations such as those resulting from shell effects or dissipation. Correspondingly, fission is considered to represent a laboratory for investigating the structure and dynamics of the atomic nucleus. However, these issues also identify the major difficulty found in the model description of this process. In fact, 65 years after Bohr and Wheeler provided a qualitative understanding of this process based on the statistical liquid-drop model [101], a well defined and universally accepted model describing fission is not yet available.

The statistical model of fission is based on the transition-state method. According to Bohr and Wheeler and in contrast to the decay modes considered so far, the decay rate for fission does not depend on the level density of the residual nuclei, which are the fission fragments at infinite separation. It depends rather on the properties of the “transition states” of the compound

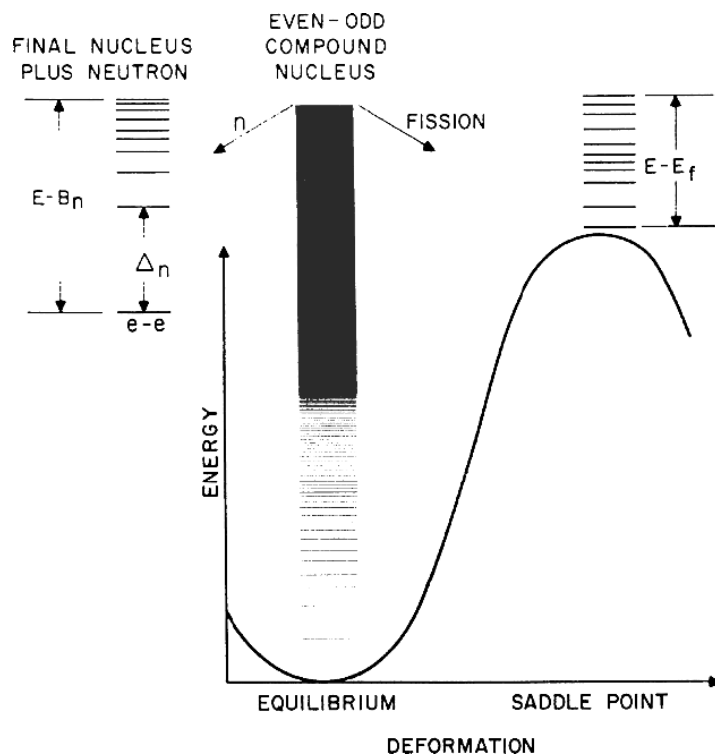


Plate 13. Illustrative representation of the intrinsic states of an even-odd nucleus at its equilibrium and saddle-point deformations (*right-hand side*) and of the residual nucleus following neutron emission (*left-hand side*). Figure taken from [102]

nucleus at the saddle point (see Fig. 13). These states are characterised by two factors. The first one is the density of states above the fission barrier due to intrinsic excitations that can be written as $\rho(U_i - B_f - \epsilon)$, where B_f is the fission barrier and ϵ is the kinetic energy associated with the motion of the fission fragments at the saddle point. The second factor is related to the fraction of the phase space occupied by the movement of the nascent fission fragments along the deformation axis. If we call R the deformation coordinate, the phase space defined by the fission movement will be $\Delta R dp/h$. Since one is only interested in those states around the saddle point (transition states), the fraction of the total phase space that corresponds to the transition states is $v\Delta t/\Delta R$, where v is the velocity of the fission motion and Δt the time the system stays around the saddle point. The fission decay rate R_f is then defined by the ratio of densities for transition and initial states, normalised to Δt :

$$R_f(U_i, J_i) = \frac{1}{h\rho(U_i, J_i)} \int_0^{U_i - B_f(J_i)} \rho(U_i - B_f(J_i) - \epsilon) d\epsilon \quad (11)$$

where the relation $vdp = d\epsilon$ has been used. Equation (11) does not account for the transmission through the barrier. Using the Hill-Wheeler parabolic barrier approximation the following generalisation can be deduced:

$$R_f(U_i, J_i) = \frac{1}{h\rho(U_i, J_i)} \left[\sum_i \frac{1}{1 + e^{2\pi(B_f - \epsilon_i)/\hbar\omega_f}} + \int_0^{U_i - B_f(J_i)} \rho(U_i - B_f(J_i) - \epsilon) d\epsilon \right] \quad (12)$$

ϵ_i being the kinetic energies of the discrete states with energies below the fission barrier, and ω_f the frequency of the oscillator describing the fission barrier around the saddle point. Extending the integration limit to infinity in (11) [95], the fission decay width Γ_f can be calculated according to the transition state model of Bohr and Wheeler:

$$\Gamma_f^{BW}(U_i, J_i) = \frac{1}{2\pi\rho(U_i, J_i)} T_{sadd} \rho(U_i - B_f) \quad (13)$$

Here T_{sadd} represents the nuclear temperature at saddle. Together with the description of the level densities, the main ingredient of this model is the fission barrier. Fission barriers can be calculated accurately using the finite-range liquid-drop model [103] taking into account microscopic contributions from ground-state shell structure [104].

The transition-state method can also be used to describe the emission of light particles by replacing the fission barrier by the separation and Coulomb energy that characterise this disintegration mode. It has actually been shown that both approaches of describing particle emission, the transition state method and the statistical evaporation model, correspond to the same unified description of the process [112].

The present formulation of the fission model does not account for the mass asymmetry degree of freedom. Consequently it yields fission probabilities but not the mass or charge distribution of the fission process. Several attempts have been made to introduce additional degrees of freedom into the transition state fission model. In the context of the adiabatic approximation, Moretto [105] proposed that a mass asymmetry degree of freedom q with conjugate momentum p would introduce an additional factor $dqdp/h$ in the description of transition states around the saddle point. The sum of the potential and kinetic energy ($\epsilon = p^2/M$, with M being the inertia parameter corresponding to this new degree of freedom) must then be subtracted from the total energy U_i in the calculation of the level density:

$$\frac{1}{h} \frac{d^3\Gamma_f}{d\epsilon_f dp dq} d\epsilon_f dp dq = \frac{\rho(U_i - B_f(q) - \epsilon_f - \epsilon)}{\rho(U_i)} \frac{d\epsilon_f}{h} \frac{dq dp}{h} \quad (14)$$

where $B_f(q)$ represents the potential energy at the saddle point as a function of the mass asymmetry and ϵ_f is the kinetic energy along the fission degree

of freedom (deformation). If one considers a parabolic potential $B(q) = B_o + \alpha q^2$, and expands $\rho(U_i - B_f(q) - \epsilon_q - \epsilon)$ one gets:

$$\frac{1}{h} \frac{d^3 \Gamma_f}{d\epsilon_f dp dq} d\epsilon_f dp dq = \frac{\rho(U_i - B_o)}{\rho(U_i)} e^{-\beta(\alpha q^2 + \epsilon_f + \epsilon)} \frac{d\epsilon_f}{h} \frac{dq dp}{h} \quad (15)$$

with β being the inverse of the nuclear temperature at the saddle point. Integrating over ϵ_f and p and extending the integration limits to infinity yields the decay width for a given mass asymmetry as:

$$\frac{1}{h} \frac{d^3 \Gamma_f}{dq} dq = \frac{(2\pi M)^{1/2}}{h^2 \beta^{3/2}} \frac{\rho(U_i - B_o)}{\rho(U_i)} e^{-\beta \alpha q^2} dq \quad (16)$$

Benlliure and collaborators improved this formalism by introducing a more realistic description of the mass-asymmetry dependence of the energy potential at saddle point, based on a semi-empirical description of shell effects [74]. Using experimental results and the model of Brosa and collaborators [106], the authors defined three main components of the mass-asymmetric potential as a function of the neutron number at the saddle point. The first one is the symmetric component (V_{mac}) described by a parabolic function as defined by the liquid-drop model. This parabola is assumed to be modulated by two neutron shells, located at mass asymmetries corresponding to neutron numbers $N = 82$ ($V_{sh,1}$) and around $N = 88$ ($V_{sh,2}$) in the nascent fragments. Shell effects are represented by Gaussians as a function of the mass asymmetry. The influence of shell effects in the light fission fragments and in proton number are neglected. The total potential energy is thus given by the sum of five contributions:

$$\begin{aligned} V(N) = & V_{mac}(N) \\ & + V_{sh,1}(N) + V_{sh,1}(N_{CN} - N) \\ & + V_{sh,2}(N) + V_{sh,2}(N_{CN} - N) \end{aligned} \quad (17)$$

where the potential energy is symmetric around $N_{CN}/2$, N_{CN} being the neutron number of the fissioning nucleus. The curvature of the parabolic potential describing the symmetric component as well as the parameters of the Gaussian functions corresponding to the shell effects are obtained from experimental results. The probability for a given splitting of the fissioning nucleus is obtained by the statistical weight of transition states above this mass-asymmetric potential. The neutron to proton ratio of the final fission fragments is given by the unchanged charge density assumption of the fissioning nucleus with a width calculated in a touching-sphere configuration for a symmetric split.

A different approach is taken in scission-point models such as those proposed by Fong [107] or Wilkins, Steinberg and Chasman [108]. These models use the properties of the scission-point configuration to describe the nature of fission fragments, based on the assumption of a total or partial thermal equilibrium of the degrees of freedom at the scission point. In this approach, the

coupling between collective and intrinsic degrees of freedom is strong enough to transform all collective energy from saddle to scission into intrinsic energy shared by both fission fragments. These models are markedly different from the Bohr and Wheeler model since they assume strong coupling rather than adiabaticity for the saddle to scission path.

In the simplest version of the Fong model, the scission configuration consists of two touching spheres with a given intrinsic energy, where the final density of states is determined by the convolution of the density of states of the two fission fragments according to:

$$\rho_{12}(U) = \int_0^U \rho_1(U_1) \rho_2(U - U_1) dU_1 \quad (18)$$

where U is the available excitation energy at scission that can be calculated as $U = U_i + Q_f - B_C$, Q_f being the mass energy balance of the fission process and B_C the Coulomb energy of both fission fragments. Given the Fermi-gas level density one can derive the level density at scission [94]:

$$\rho_{12}(U) = K \frac{(a_1 a_2)^{1/2}}{(a_1 + a_2)^{5/4}} U^{3/4} e^{2\sqrt{(a_1 + a_2)U}} \quad (19)$$

with a_i representing the level-density parameter. The main limitation of this model concerns the exact description of the scission configuration where shell effects, pairing and deformation strongly influence the result of the calculation.

In the model of Wilkins, Steinberg and Chasman, the level density at the scission point is obtained by convoluting the level density of both fission fragments, the latter being approximated by the maximum value of the integrand in (18) according to [94]:

$$\rho_{12}(U - V) = \rho(U) e^{-V/T_{coll}} \quad (20)$$

where T_{coll} is referred to as the “collective temperature” and V is taken as the difference between the ground state energy of the fissioning nucleus and the lowest energy of any deformed configuration at the scission point. This model also includes a detailed description of the scission potential energy where shell effects, pairing and deformation are accounted for.

As will be further discussed in Sect. 5.3, this statistical description of fission was questioned by Kramers [109] just a year after its publication. Kramers considered fission to be a dynamical process that cannot be described in using purely statistical terms.

Intermediate-mass Fragment Emission and Multi-fragmentation

At high excitation energies, such as those reached in spallation reactions, the statistical emission of intermediate-mass fragments (IMF) has to be considered as a competing decay channel. From a formal point of view, IMF

evaporation can be considered as a very asymmetric fission process. Thus, the transition state and the scission point models discussed in the previous section, can be used to describe this process. The only additional difficulty is the description of the potential energy for extremely asymmetric configurations at high excitation energies. Another problem could be finding a consistent treatment for the evaporation of light particles and IMFs. The extension of the Weisskopf or the Hauser–Feshbach formalisms to IMF emission would correspond to the scission point models of Fong or Wilkins, Steinberg and Chasman. However, other authors such as Moretto [105, 110] and Swiatecki [111] prefer to use the transition state model to describe IMF evaporation and extend it to the emission of light particles and clusters. In the latter approach, defining the saddle point for these very asymmetric configurations seems to be a difficult task. A detailed discussion of this topic can be found in [112].

Spallation reactions induced at energies above 1 GeV could lead to the formation of nuclear matter at temperatures above 4 MeV. Under these conditions, a new decay channel called multi-fragmentation is expected to occur [113]. This decay channel is characterised by the complete disintegration of the nucleus in several simultaneously emitted IMFs. The description of this decay channel is beyond the sequential evaporation approach described in previous sections. Although multi-fragmentation represents a small fraction of the cross section in spallation reactions, they offer optimum conditions for investigating this decay channel, as will be discussed in Sect. 5.

Level Density

Level density is one of the key parameters in the statistical description of the de-excitation process of a hot nucleus. This topic has been discussed extensively in a previous series of lectures notes by Langanke, Thielemann and Wiescher [42], or in review papers such as those of Ericson [99] or Huizenga and Moretto [114]. Here the reader is reminded of its basic formulation and particularities at high excitation energy. According to the adiabatic formalism [116], the level density of an excited nucleus can be written as:

$$\rho(U, J) = K_{coll}(U)\rho_{int}(U, J) \quad (21)$$

where the internal ($\rho_{int}(U, J)$) and collective (K_{coll}) degrees of freedom are completely decoupled. The most widely used model for calculating level density due to intrinsic excitations is the non-interacting Fermi–gas model proposed by Bethe [115], according to:

$$\rho_{int}(U, J) = \frac{1}{2}F(U, J)\rho(U) \quad (22)$$

with

$$F(U, J) = \frac{2J+1}{2\sigma^2} \exp\left(\frac{-J(J+1)}{2\sigma^2}\right)$$

$$\rho(U) = \frac{\pi}{12} \frac{e^S}{a^{1/4} U^{5/4}} \quad (23)$$

Here $\sigma^2 = \Theta_{rigid} \sqrt{U/a}/\hbar^2$ is the so called spin cut-off parameter, with $\Theta_{rigid} = 2m_u A R^2/5$ being the moment of inertia of a nucleus with radius R and mass $m_u A$ and a the asymptotic level-density parameter that includes corrections for the nuclear surface, as discussed in [117]. The entropy S can be calculated using the back-shifted Fermi-gas prescription [118] in order to account for shell effects δU and pairing corrections δP , according to:

$$S = 2\sqrt{a\bar{U}} \quad (24)$$

where $\bar{U}$ represents an effective excitation energy that includes shell and pairing corrections. In the particular case of highly excited nuclei, the energy dependence of the shell and pairing corrections plays an important role in the last steps of the evaporation chain when the nucleus has lost most of its initial excitation energy. A detailed description of these excitation-energy dependent corrections to the level density will be given in Sect. 5.1.

Finally, the enhancement of the intrinsic level density due to rotation and vibrations ($K_{coll}(U)$) has to be taken into account for the case of the de-excitation of nuclei with collective properties. This topic will be presented in Sect. 5.2.

Evaporation Codes and Results

Statistical de-excitation codes are based on a sequential evaporation mechanism which neglects the probability of simultaneous emission of two or more particles or γ rays. In most of the codes it is assumed that the emitted particles are in their ground states and that the initial nucleus undergoes an evaporation chain (multi-step process) until a final nucleus in its ground state is reached. Furthermore, it is assumed that at each evaporation step the initial compound nucleus and final daughter nucleus are in statistical equilibrium. The compound nucleus assumption is then applied at each step of the de-excitation chain. The mass and atomic numbers, excitation energy and angular momentum which define each evaporation step are obtained by conservation laws.

Multi-step evaporation codes follow two different methods. The first one is based on a mass and atomic number grid with a population distribution of each nucleus in a two-dimensional space of excitation energy and angular momentum. Given the initial distribution in excitation energy and angular momentum of the compound nucleus, the population of the various daughter nuclei is calculated. The disadvantage of these codes is that they generally do not calculate angular and energy distributions of the emitted particles and do

not preserve correlations between the emitted particles. The second method uses Monte Carlo techniques to follow the decay of individual compound nuclei until the residual nucleus can no longer decay. The advantage of this method is that it can predict energy spectra, angular distributions and multi-particle correlations.

The existing evaporation codes also differ in the possible decay channels or the definition of statistical model parameters such as transmission coefficients, level densities or binding energies. In Table 2 Stokstad's list [112] of the most commonly used evaporation codes is updated, including the decay channels taken into account by each code and the formalism used to calculate the decay rates.

Table 2. List of most used statistical evaporation codes. Together with the main references, the models used to evaluate the different de-excitation channels are characterised by the following abbreviations: W.E. (Weisskopf–Ewing), H.F. (Hauser–Feshbach), T.S. (transition state model), F.P. (Fokker–Planck, dynamical description of the fission flux), PAR (parameterisation), MC (Monte Carlo) and GR (grid)

Code	Type	γ	n,p,He	IMF	Fission	Reference
ABLA	MC	yes	W.E.	no	F.P.	A. Junghans et al. [73]
ALICE	GR	no	W.E.	no	T.S.	M. Blann et al. [119, 120]
CASCADE	GR	yes	H.F.	no	no	F. Pühlhofer [121]
DRESNER	MC	yes	W.E.	no	T.S.	I. Dostrovsky et al. [122]
GEM	MC	no	W.E.	W.E.	PAR	S. Furihata [123]
GEMINI	MC	no	H.F.	T.S.	T.S.	R.J. Charity et al. [124]
LANCELOT	MC	yes	H.F.	no	no	A.J. Cole [125]
LILITA	MC	no	H.F.	no	no	J. Cómez del Campo [126]
PACE	MC	yes	H.F.	no	T.S.	A. Gavron [127]

It is beyond the scope of this lecture to benchmark all these codes. However, it is clear that the reliability of the predictions of these codes has increased with the quality of the available data. An example is shown in Fig. 9 where accurate measurements of the isotopic distribution of residual nuclei produced in spallation reactions induced by ^{208}Pb on hydrogen are compared to model calculations. Recent codes such as the INCL [25] coupled to the ABLA evaporation code [73] (dark line) provide a much better description of the data than older codes such as the Bertini cascade [20] coupled to the Dresner evaporation code [122] (grey line). A detailed discussion on some of these evaporation codes can be found in Chaps. 4 and 5 of [94].

5 Recent Investigations of Structure and Dynamics of Atomic Nuclei by Using Spallation Reactions

Spallation reactions are an important tool for studying the dynamics of hot nuclei because they allow one to investigate, in principle, nuclear properties as a function of temperature. These reactions lead to relatively high excitation energies, while the induced angular momentum or deformation in the pre-fragment remains quite small. At high temperature other decay channels become accessible to the nucleus, such as the emission of intermediate-mass fragments, in addition, evidences for a hindrance of the fission width with respect to the value predicted by statistical models have been observed [128]. At temperatures above a critical value, the emission of intermediate-mass fragments evolves toward a simultaneous break-up of the nucleus into many fragments, referred to as “multi-fragmentation” [129]. The onset of multi-fragmentation and the dynamics of fission at high excitation energies are topics which have been intensely studied in recent years.

In addition, it has been shown that spallation reactions in inverse kinematics or projectile fragmentation reactions can be used to investigate different nuclear-structure features. It is well established that even-odd or shell closure effects disappear at high temperature and consequently, are not expected in high energy reactions. However, the accurate measurement of the production yields of residual nuclei in these reactions has revealed complex nuclear-structure phenomena [130]. These observations indicate that the final residual nuclei are produced in high-energy reactions through long evaporation chains that originate from very hot nuclei. The final isotopic composition of the residues is thus determined during the last steps of the evaporation chain where the nuclei are sufficiently cold and their decay widths are sensitive to structural effects. Consequently, the investigation of the production yields from highly excited nuclei could be a rich source of information on nuclear-structure phenomena in slightly excited nuclei found at the end of their evaporation process. Two manifestations of nuclei structure have already been observed in the production yields of spallation and fragmentation residues, namely those related to pairing and collective excitations, as will be discussed in the following two sections.

5.1 Pairing

Figure 14 shows the production cross sections of light projectile-like residues from the reaction $^{238}\text{U}+\text{Ti}$ at 1 A GeV [130]. The data are sampled according to the neutron excess $N - Z$ for even-mass and odd-mass nuclei. This representation reveals a complex structure. For even-mass nuclei a clear even-odd effect is present, being particularly strong for $N = Z$ nuclei. Odd-mass nuclei show a reversed even-odd effect with enhanced production of odd- Z nuclei. This enhancement is stronger for nuclei with larger values of $N - Z$. However,

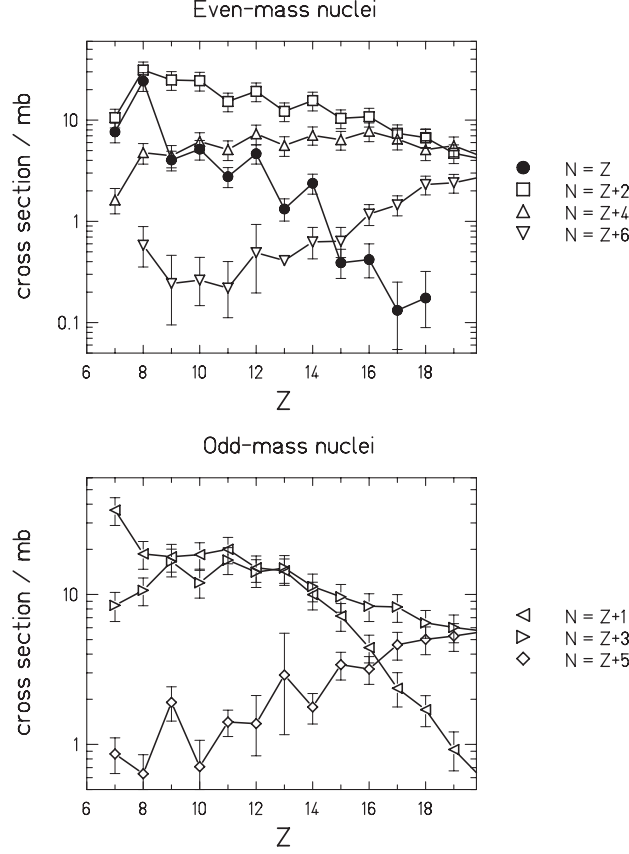


Plate 14. Formation cross sections of the projectile-like residues from the reaction $^{238}\text{U}+\text{Ti}$, 1 A GeV. The data are given for specific values of $N-Z$. The chain $N=Z$ shows the strongest even-odd effect, while the chain $N-Z=5$ shows the strongest reversed even-odd effect. Data taken from [130]

for nuclei with $N-Z=1$ the reversed even-odd effect vanishes at $Z=16$, and an enhanced production of even- Z nuclei can again be observed for $Z>16$. All observed structural effects seem to vanish as the mass of the fragment increases.

To interpret these effects, a consistent description of level densities and binding energies is required, including shell and pairing corrections according to the back-shifted Fermi-gas prescription. Following this model, the effective excitation energy $\overline{U}$ in (24) can be written as:

$$S = 2\sqrt{a\overline{U}} = 2\sqrt{a[U_{eff} + k(U_{eff})\delta S + h(U_{eff})\delta P]} \quad (25)$$

The asymptotic level-density parameter a can be calculated according to Ignatyuk [117] as:

$$a = 0.073A + 0.095B_s A^{2/3} \quad (26)$$

where A is the mass of the nucleus and B_s the ratio of the deformed nucleus surface to the spherical one. In (25), δS is the ground-state shell correction, calculated as the difference between the experimental ground-state mass and the corresponding value predicted by the liquid-drop model [103]. At the saddle point, shell correction can be assumed to be negligible. The function $k(U_{eff})$ describes the damping of the shell effect with the excitation energy, and can be calculated according to [131] as $k(U_{eff}) = 1 - \exp(-\gamma U_{eff})$, with the parameter γ being determined by $\gamma = a/(0.4A^{4/3})$ [132].

In the back-shifted Fermi-gas model, the effective pairing energy shift δP is calculated as:

$$\delta P = -\frac{1}{4}\Delta^2 g + 2\Delta \quad (27)$$

where the average pairing gap is $\Delta = 12/\sqrt{A}$, the single-particle level density at the Fermi energy is $g = 6a/\pi^2$ and $h(U_{eff})$ parameterises the super-fluid phase transition [133] at the critical energy $U_{crit} = 10$ MeV [134] according to:

$$h(U_{eff}) = \begin{cases} 1 - \left(1 - \frac{U_{eff}}{U_{crit}}\right), & U_{eff} \leq U_{crit} \\ 1, & U_{eff} > U_{crit} \end{cases} \quad (28)$$

The effective energy U_{eff} is shifted with respect to the excitation energy U to accommodate for the different energies of even-even, odd-mass and odd-odd nuclei:

$$\begin{aligned} U_{eff} &= U & \text{odd - odd} \\ U_{eff} &= U - \Delta & \text{odd - mass} \\ U_{eff} &= U - 2\Delta & \text{even - even} \end{aligned} \quad (29)$$

Using this model most of the observed even-odd effects shown in Fig. 14 can be understood when pairing and shell effects are consistently accounted for in both binding energies and level densities of the parent and daughter nuclei. In the case of odd-mass residues, the even-odd staggering in the production cross sections can be interpreted as being due to pairing correlations in the particle separation energies. The number of particle-bound states in the final residual nuclei follows the observed staggering in the production yields. However, in order to describe even-odd staggering in even-mass nuclei, the states available in both the daughter and mother nuclei must be taken into account in the evaporation chain. In both cases and for heavier nuclei, γ emission becomes a competitive decay channel in the final de-excitation steps, causing even-odd staggering to decrease as mass increases.

In spite of this success, the particularly strong staggering observed in the final yields on the $N = Z$ chain could not be reproduced. Although this question is not completely solved, some phenomena like Wigner energy, α clustering or neutron-proton correlations could explain this strong staggering. A detailed discussion of these results can be found in [130].

5.2 Collective Excitations

As is apparent from (23) and (24), shell effects modify binding energies and level densities. Analogous to pairing correlations, they should also manifest themselves during the last steps of the evaporation chain for those nuclei near a shell closure.

By using reactions induced by ^{238}U projectiles it has become possible to produce nuclei across the neutron shell $N = 126$. In this region of the chart of the nuclides, the production cross-sections of such nuclei are governed by the competition between neutron or proton evaporation and fission. In fact, the production cross sections of evaporation residues are a measure of the survival probability against fission. Near the shell closure $N = 126$, the higher fission barriers are expected to enhance the production cross-sections of such nuclei. However, such an effect has not been observed, as can be seen from Fig. 15.

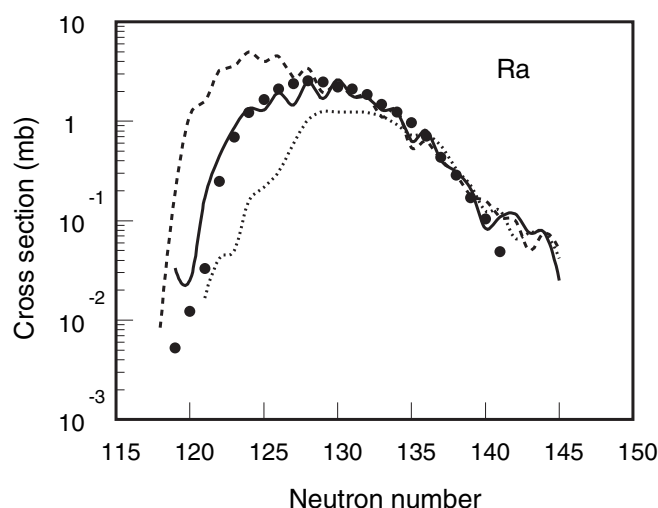


Plate 15. Production cross sections of radium isotopes measured for the reaction $^{238}\text{U}(1\text{ A GeV})+d$ (*full circles*). The experimental data are compared with model calculations using a Fermi-gas level density (*dotted line*), a level density including ground-state shell effects (*dashed line*) and one based on ground-state shell effects and collective excitations (*solid line*) [135]

The astonishing lack of stabilisation against fission for magic or near magic nuclei is interpreted as a signal of level-density enhancement due to the presence of rotational and vibrational collective excitations $K_{coll}(U)$ [73]. The large deformation of the mother nucleus at the saddle point favours the appearance of rotational bands above the fission barrier. These collective levels are added to the intrinsic levels of the nucleus, leading to an increase of the level density at saddle that favours the fission decay channel.

Junghans and collaborators [73] calculated the contribution of collective excitations to the level density $K_{coll}(U)$ according to:

$$K_{coll}(\bar{U}) = \begin{cases} (\sigma_{\perp}^2 - 1)f(\bar{U}), & \sigma_{\perp}^2 > 1 \\ 1, & \sigma_{\perp}^2 \leq 1 \end{cases} \quad (30)$$

where $\sigma_{\perp}$ is the spin-cutoff parameter and $f(\bar{U})$ describes the damping of the collective modes with increasing excitation energy. For nuclei with quadrupole deformation ($|\beta_2| > 0.15$), the collective enhancement corresponds to rotational excitations. This is described by a spin-cutoff parameter $\sigma_{\perp}^2(rot)$ defined as:

$$\sigma_{\perp}^2(rot) = \frac{\mathfrak{S}_{\perp} T}{\hbar^2}, \quad \mathfrak{S} = \frac{2}{5} m_o A R^2 \left(1 + \frac{\beta_2}{3} \right) \quad (31)$$

T being the nuclear temperature, $m_o A$ the nuclear mass and R the nuclear radius. At small deformations ($|\beta_2| \leq 0.15$), the collective enhancement is due to vibrations. In this case, the rigid-rotor moment of inertia is replaced by the irrotational-flow moment of inertia ($\mathfrak{S}_{irr} = \beta_{eff}^2 \mathfrak{S}_{\perp}$). The corresponding spin-cutoff parameter $\sigma_{\perp}^2(vib)$ is:

$$\sigma_{\perp}^2(vib) = S \beta_{eff}^2 \sigma_{\perp}^2(rot) \quad (32)$$

where S is the strength of the vibrational excitation and β_{eff} a dynamical deformation parameter that accounts for the variation in the energy of the vibrational bands according to the distance in neutron and atomic number $\Delta N(\Delta Z)$ to the nearest closed shell according to:

$$\beta_{eff} = 0.022 + 0.003\Delta N + 0.005\Delta Z \quad (33)$$

Following this formulation, the strength of the vibrational excitations (S) can be investigated along with the dumping of rotations and vibrations with temperature ($f(\bar{U})$). In [73, 135], different sets of data have been compared with these model calculations. As shown in Fig. 15, a consistent description of the data was obtained for a value of $S = 50$. Moreover, it was deduced that the damping of the collective enhancement can be represented by a Fermi function $1/(1 + \exp[-(\bar{U} - U_c)/d_c])$ with parameters $U_c = 40$ MeV and $d_c = 10$ MeV which, contrary to [136, 137] do not depend on deformation. These results are a good example of investigation of structural effects in the atomic nucleus that cannot be addressed by using spectroscopic techniques.

5.3 Fission Dynamics

According to Kramers [109], fission should be considered as a diffusion process above the nuclear potential in the deformation coordinate that cannot be properly described within the framework of Bohr and Wheeler's statistical model [101]. This dynamical picture of fission can be modelled by using a

Fokker–Planck or Langevin equation where the diffusion process is governed not only by the nuclear potential but also by a dissipation coefficient. This coefficient represents the coupling between the intrinsic and collective degrees of freedom populated in fission. In his pioneering work, Kramers showed that the stationary solution of this Fokker–Planck equation can be analytically described in terms of the statistical fission width Γ_f^{BW} proposed by Bohr and Wheeler:

$$\Gamma_f = K \cdot \Gamma_f^{BW} \quad (34)$$

Here K is the Kramers factor defined as:

$$K = \left\{ \left[1 + \left(\frac{\beta}{2\omega_o} \right)^2 \right]^{1/2} - \frac{\beta}{2\omega_o} \right\} \quad (35)$$

where β is the reduced dissipation coefficient, and ω_o represents the frequency of the harmonic oscillator describing the inverted potential at the fission barrier. Since K is smaller than one, the dynamical description of fission leads to a smaller fission width than the predicted by the statistical model.

The complete time-dependent solution of the Fokker–Planck equation was introduced in the eighties by Grangé and collaborators [138]. They showed that the fission-decay width across the barrier needs time to reach its asymptotic value defined by Kramer’s stationary solution of the Fokker–Planck equation. The main consequence of the work of Grangé et al. is that the hindrance of the fission width with respect to its statistical value is due to both stationary and transient effects in the fission flux across the barrier. During the transient time, i.e. before the stationary regime of the fission width is reached, other de-excitation channels are favoured. Consequently, the nuclear system cools down, reducing the fission probability even more with respect to other channels.

The work of Grangé and collaborators was triggered by the experimental observation of anomalously enhanced pre-scission neutron multiplicities in fission induced by heavy-ions [141]. The large pre-scission neutron multiplicities were interpreted as an indication of the delay of fission at high excitation energies. Meanwhile, other evidence for fission hindrance induced by transient and stationary dissipative effects was obtained from the analysis of γ -rays emitted during the de-excitation of the giant dipole resonance [142] or by directly measuring the fission time using crystal blocking techniques [143].

Recently, a novel technique based on a detailed investigation of fission products in projectile fragmentation reactions at relativistic energies [145] or spallation reactions [144] has been introduced. As already mentioned, the advantage of this reaction mechanism is that the excited fissioning nucleus is produced with well defined initial conditions that can be easily described using the Serber model [11]. In addition, these reactions lead to almost undeformed nuclei covering a large range of excitation energy. In these works,

the fission cross-sections, the charge distribution [145] or the isotopic distribution [144] of fission residues have been used as indications of fission dynamics.

An example of the results obtained in these investigations is shown in Fig. 16 which represents the isobaric distribution of residual nuclei produced in the reaction $^{197}\text{Au}(800 \text{ A MeV}) + \text{p}$ [144]. In this figure the hump centred around $\Delta A \approx 85$ corresponds to the fission residues. Comparison to model calculations demonstrates the sensitivity of the data to the different approaches describing fission. The predictions of the Bohr and Wheeler's statistical model clearly overestimate the fission cross section. The calculation based on the dynamical picture including stationary and transient effects provides a better description of the data, in particular if a value of $\beta = 2 \cdot 10^{21} \text{s}^{-1}$ is assumed for the reduced dissipation coefficient. In these calculations the time dependence of the fission width was described using a step function. As illustrated here, these new experimental data allow one not only to estimate the value of the reduced dissipation coefficient that describes the coupling between intrinsic and collective excitations in nuclear matter, but also to investigate the transient effects that manifests itself in the time dependence of the fission width [140]. However, very interesting questions concerning,

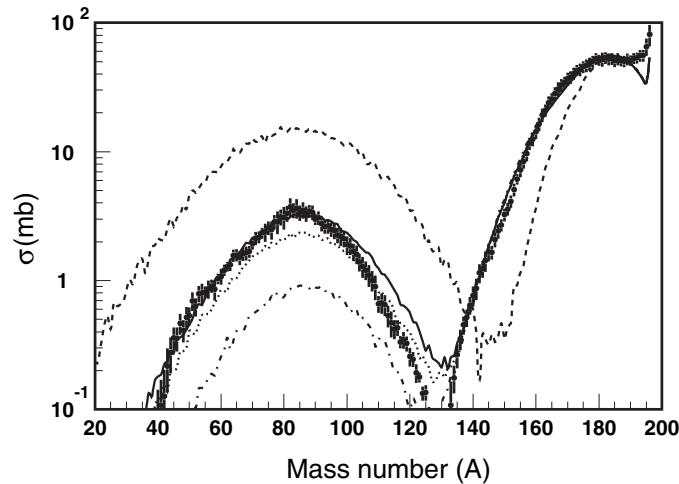


Plate 16. Measured isobaric distribution of nuclei produced in the reaction $^{197}\text{Au}(800 \text{ A MeV}) + \text{p}$ [144]. The experimental data (*black points*) are compared with the predictions obtained with different fission model calculations. The *dashed line* represents the results obtained with the statistical model of Bohr and Wheeler. The solid line corresponds to a dynamical calculation (Γ_f^K) with a reduced dissipation coefficient of $\beta = 2 \cdot 10^{21} \text{s}^{-1}$, while the dotted and *dash-dotted lines* are the results obtained with reduced dissipation coefficient of $\beta = 3 \cdot 10^{21} \text{s}^{-1}$ and $\beta = 5 \cdot 10^{21} \text{s}^{-1}$, respectively

e.g., the temperature or deformation dependence of the reduced dissipation coefficient are still open.

5.4 Multi-fragmentation

The investigation of the decay modes of extremely hot nuclei constitutes one of the approaches followed to characterise nuclear matter under extreme temperature conditions and thus, the equation of state of nuclear matter. The emission of intermediate-mass fragments competes with fission and light-particle evaporation during the de-excitation of highly excited nuclei. In principle, this process can be understood as an extension of neutron or light-charged particle evaporation to heavier clusters, for which, a sequential binary decay picture described by the statistical model can be used (see Sect. 4.2). At extremely high excitation energies, the nucleus becomes unstable and undergoes, instead of sequential binary decays, a simultaneous break-up process called multi-fragmentation [129, 146, 147]. Due to the similarities between the nuclear potential and the Lennard–Jones molecular potential, multi-fragmentation has been interpreted as a liquid-gas phase transition [113, 148, 149].

Experimentally, heavy-ion collisions in the Fermi–energy regime have been used to investigate this reaction mechanism. In this case, the result of the collision is a complex interplay between dynamic effects (compression, deformation and rotation degrees of freedom) and thermal excitation. On the contrary, peripheral heavy-ion collisions at relativistic energies allow for the production of projectile remnants heated mainly by thermal energy. In reactions induced by relativistic protons, the dynamical effects of the collisions have even less importance. For this reason spallation reactions are considered as an optimal choice to investigate “thermal” multi-fragmentation, allowing one to characterise intermediate-mass fragment emission and its relation to a phase transition [22, 150, 151].

The investigation of spallation reactions in inverse kinematics using a high resolution magnetic spectrometer has revealed new observables. In particular, it has been shown that the kinematic properties and the isotopic composition of light fragments produced in spallation reactions are sensitive to the decay mechanisms of the hot remnants produced in the reaction. Figure 17 illustrates how sensitively the kinematic properties of light residues produced in the reaction ^{56}Fe on proton at 1 A GeV depend on the decay mechanism [51]. The observed velocities are clearly compatible with a prompt break-up process. The isotopic composition of projectile remnants produced in peripheral collisions of relativistic heavy ions have also been used to determine the temperature of the nuclear system at break-up [152]. These investigations highlight high-resolution measurements as complementary source of information respect to exclusive measurements.

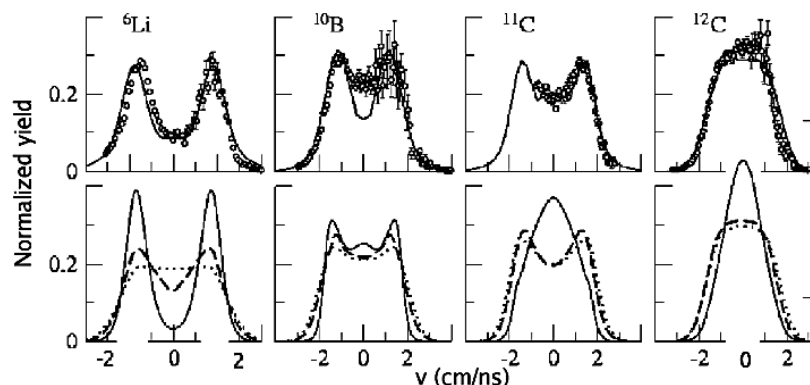


Plate 17. Upper row: measured (*circles*) and reconstructed (*lines*) velocity distributions of projectile remnants produced in the reaction ^{56}Fe on proton at 1 A GeV, represented in the frame of the average velocity of each fragment. Lower row: velocity distributions calculated with a sequential binary model (*solid line*) and prompt break-up model (*dashed and dotted lines*). For details see [51]

6 Summary and Outlook

This lecture addressed the increasing importance that spallation reactions have gained during the past few years. The implications of these reactions were discussed, concerning basic research such as the interpretation of the observed cosmic ray abundances, or applications such as neutron sources or the production of radioactive nuclear beams. Even though these reactions have been qualitatively understood for a long time, the present knowledge is not sufficient for most of these applications. The interpretation of accurate new data on cosmic ray abundances or the design of spallation targets for neutron or radioactive beam production requires a better understanding of these reactions that would enable one to significantly improve the predictive power of the present model calculations. This situation has inspired numerous efforts around the world, particularly in Europe, to obtain accurate data on these reactions using advanced experimental techniques. Some of these experimental programs were also briefly discussed in this lecture.

Moreover, this lecture reviewed the present understanding of this reaction mechanism using the two-step model. Intra-nuclear cascade models were presented as an optimal choice for describing the first stage of the reaction. The underlying physics of the subsequent statistical de-excitation of the remnant was then discussed in terms of light-charged particle evaporation, γ -ray emission, fission and multi-fragmentation.

The last section of this lecture highlighted some of the most interesting and recent investigations of the structure and dynamics of the atomic nucleus, obtained by using spallation or fragmentation reactions in inverse kinematics.

The accurate measurement of the isotopic composition and kinematic properties of residual nuclei have been proven to be new and powerful observables.

Acknowledgements

The contributions of all the colleagues who participated in the spallation experiments at GSI providing many of the results shown in this lecture are gratefully acknowledged. In particular I thank S. Leray and P. Napolitani for providing some figures and K.-H. Schmidt and M.V. Ricciardi for the careful reading of the manuscript and for many enjoyable and fruitful discussions.

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