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The Sun: New Challenges

Proceedings of Symposium 3 of JENAM 2011

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Springer

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Preface

This book contains the Proceedings of the Joint European and National Astronomical Meeting (JENAM-2011) held on July 4–8, 2011 at St. Petersburg.

The main topics discussed at the meeting were:

1. The unusual sunspot minimum—challenge to the solar dynamo theory
2. New observational window: terra-Hertz emission
3. Wavy corona
4. Space weather agents—initiation, propagation, and forecast.

Now and again, the Sun sets new problems before the astronomers. One of such problems is the abnormal behavior of solar activity during the past, 23rd cycle. Even now, it is not clear whether the anomalies have ceased with the beginning of the new cycle 24 or we are still facing a long period of low solar activity. The anomalies in question have manifested themselves in various parameters, such as the sunspots per se, the number and intensity of coronal mass ejections, extraordinary brightness distributions in the corona, solar wind parameters, and the persistent big low latitude coronal holes.

We discussed at the symposium the following problems:

- What are the characteristics of solar activity that display abnormal behavior? Is it possible that we are on the threshold of a strong decrease of solar activity? Were analogous episodes in the history of solar activity? What are the similar features and differences between the activity cycles in the Sun and stars?
- Are the present-day theories able to account for strong variations in the height of the cycles (up to an order of magnitude) on one and the same star? Is it possible to predict the heights and peculiarities of the cycles on the basis of the dynamo theory?
- Are there additional arguments for the influence of planets on solar activity?

Observations in the sub-THz range of large solar flares have revealed a mysterious spectral component increasing with frequency and hence distinct from the microwave component commonly accepted to be produced by gyrosynchrotron

(GS) emission from accelerated electrons. Evidently, having a distinct sub-THz component requires either a distinct emission mechanism (compared to the GS one) or different properties of electrons and location or both. It is interesting to discuss the complete list of possible emission mechanisms.

It is the magnetic field that determines the variations of the coronal brightness. However, the mechanism of the corona heating and, therefore, of the relationship between the corona brightness and magnetic field is unclear. This is, obviously, due to the fact that there are several heating mechanisms that play different roles in different areas (active regions, quiet Sun, coronal holes). So far, it is not clear whether the DC or AC mechanisms prevail in one or another object in the Sun. What is the role of different-scale magnetic fields in the heating of the solar corona and how does their relative contribution change with time? What kinds of manifold observational waves and oscillations are significant to understand the heating of upper solar atmosphere?

The progress in studying the key objects of the Space Weather problem—CMEs and high-speed solar wind streams will be, apparently, achieved owing to a wide use of stereoscopic data from the STEREO-A and spacecraft as well as the high-quality solar obtained in several channels corresponding to different plasma temperatures (SDO). The main topics to discuss were as follows:

- The nature of coronal mass ejections and their connection to various-scale fields
- The acceleration of the solar wind and connection between the solar wind and various features in the Sun
- 3D structure and physical parameters of CME sources, including the flares and filament eruptions
- The correlation between coronal hole characteristics and the parameters of the solar wind.

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The Unusual Sunspot Minimum: Challenge to the Solar Dynamo Theory

V.N. Obridko, Yu.A. Nagovitsyn, and Katya Georgieva

Abstract The last cycle 23 was low, long, complex, and very unusual. The “peculiarity” of the minimum was that the field was weak, but also that the morphology of the heliosphere was very complex. A large number of features of intermediate scale—neither global nor local—were observed. There are reasons to believe that the amplitude and the period of a cycle are determined by the large-scale meridional circulation which, in turn, may be modulated by planetary tidal forces. There are evidences that at present the deep meridional circulation is very slow, from which a low and late maximum of cycle 24 can be predicted. Calculations of the planetary tidal forces indicate that cycle 25 will be still lower, and therefore cycle 24 is the beginning of a secular solar activity minimum. Various prediction methods are summarized, all indicating that we are entering a period of moderately low activity, and the possibility of a Maunder-type minimum is very small. Arguments are also presented in favor of a near-surface dynamo.

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1 Introduction

We begin with the question why the paper is headed in this way, and in what properly is the challenge to the solar dynamo theory. The last minimum and the whole cycle 23 were quite strange. Does this strangeness really go beyond the limits of the formerly observed scatter in the activity cycles' characteristics? And can we at present even broadly identify the reasons for the significant variations in the solar cycles' characteristics? It should be noted that the main peculiarity of the last minimum was that the fields were weak and the morphology of the heliosphere was complex. A big number of features of intermediate scale—neither global nor local—were observed. Can't this be used to assess the amplitude of the next maximum? What are the recent models of solar dynamo, and can they provide a description of the characteristics of the 11-year cycle, in particular the observed variations in the sunspot magnetic field intensity, and what the level of activity in a given cycle depends on? Finally, what can we expect of cycle 24 which we have entered?

2 Some Characteristics of the 23/24 Solar Minimum

In Fig. 1 the monthly number of the spotless days in the last five solar cycles is shown based on data from the Mountain station of Pulkovo observatory [1]. It can be seen that cycle 23 which just ended is characterized by a big number of spotless days. However, during the last century, the number of spotless days was even bigger in 1913 minimum (Fig. 2).

The number of sunspots in 2008 was extremely low, but the heliospheric current sheet was not flat as it should be during the minimum when all other harmonics disappear except for the axial dipole. Figure 3 demonstrates the coronal structure during three consecutive high-latitude scans of the Ulysses spacecraft. It can be seen that the situation in 2008 minimum strikingly differs from the situation in 1996 minimum. The structure of the corona in 1996 is standard, with a pronounced streamer in the equatorial plane which is an evidence of a typical axial dipole structure. In 2008, on the other hand, numerous extra-equatorial streamers are observed which cannot be associated with the axial dipole.

Figure 4 demonstrates the cyclical variation of the effective multipolarity index introduced in [2]. This index is equal to 2 for a dipole and strongly increases in the cycle maximum.

In the 23/23 cycle minimum the situation is significantly different from 20/21 and 21/22 minima. While in the previous minima the index dropped to almost dipole values, the decreased which began during the declining phase of cycle 23 was later replaced by an increase to almost a maximum value. The increase of this index is an indication of an unusually large number of equatorial coronal holes. The

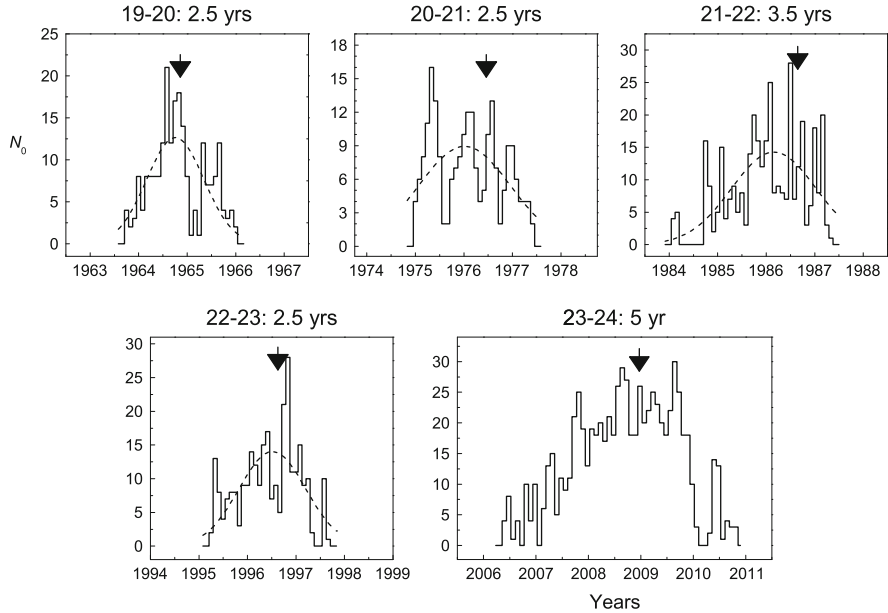


Fig. 1 Monthly number of the spotless days in the last five solar cycles is shown based on data from the Mountain station of Pulkovo observatory

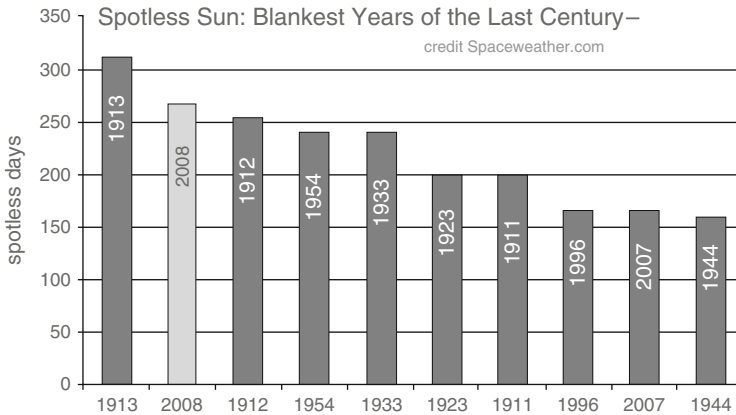


Fig. 2 Years with biggest numbers of spotless days in the last century. From [33]

calculations evidence that the global field even in 2009 was determined not only by the dipole, but it contained numerous small open magnetic field areas at all latitudes (Fig. 5). As a result, a large number of extra-equatorial solar wind sources were observed [3].

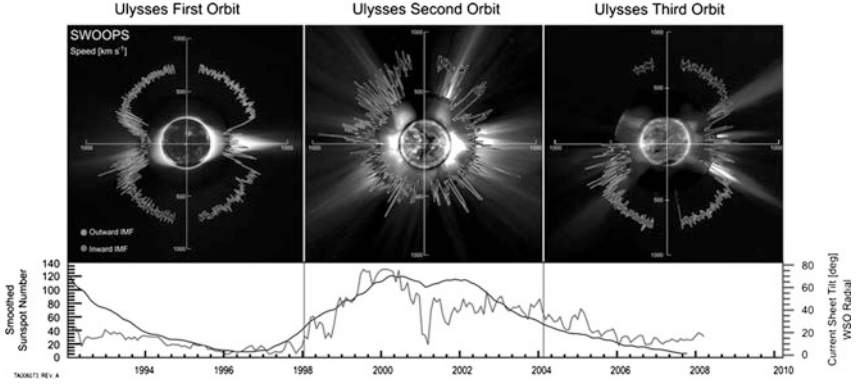


Fig. 3 *Upper three plots:* Polar plots of the solar wind speed for Ulysses' three polar orbits plotted over characteristic solar images for solar minimum for cycle 22 (8/17/96)—*left*, solar maximum for cycle 23 (12/07/00)—*middle*, and solar minimum for cycle 23 (03/28/06)—*right*; *Bottom plot:* Contemporaneous values for the smoothed sunspot number (*thick black curve*) and heliospheric current sheet tilt (*thin grey curve*), lined up to match the *upper three panels*. From [34]

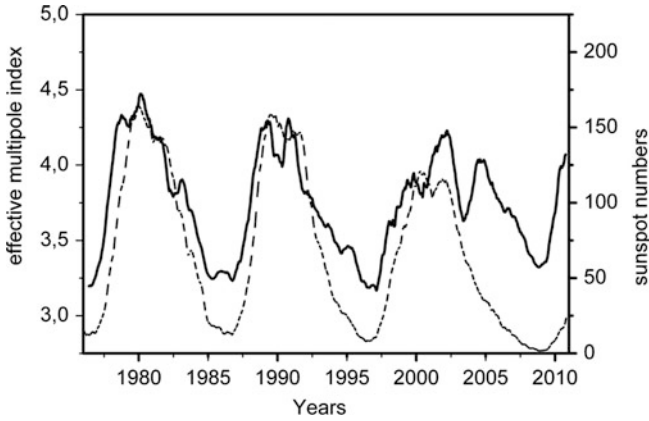


Fig. 4 Cyclic variations of the effective multipolarity index [2]

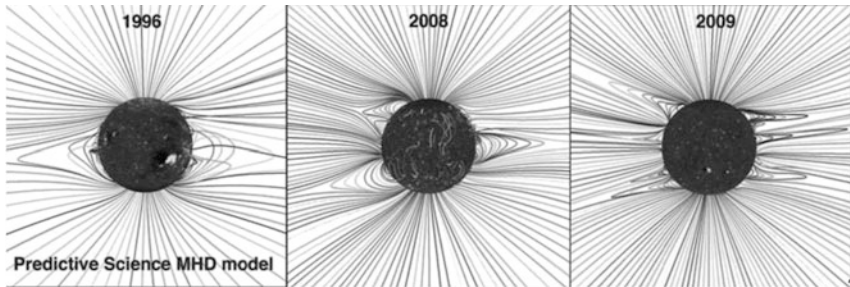


Fig. 5 The structure of the solar corona in 22/23 cycle minimum (*left*) and 23/24 cycle minimum (*middle and right*) as determined by the outward extension of the Sun's magnetic field. From [35]

3 Amplitude of the Solar Cycle and the Large-Scale Meridional Circulation

What determined such a low number of sunspots in the 23/24 cycle minimum? In general, what determines the amplitude of a cycle in both the global and local fields? There are reasons to believe that this is related to the variations in the large-scale meridional circulation.

It is known that the solar cycle consists of two processes. At the first stage the toroidal field is generated from the poloidal field (Ω -effect). In sunspot minimum the magnetic field has a quasi-dipolar (poloidal) structure. Due to the velocity gradient, the differential rotation at the base of the convection zone stretches the field lines of the poloidal field and deflects them in azimuthal direction, thus creating the toroidal component of the field. The magnetic buoyancy forces raise the toroidal field tubes to the surface, giving rise to standard bipolar solar groups. This process is well studied and raises no particular doubt. (But note by the way, that it is not clear whether it allows for local isolated flux tubes to be created with intensities of up to 3,000 G).

The reverse process of generation of the poloidal field from the toroidal field (α -effect) is much less clear. The most cited is the Babcock–Leighton mechanism [4,5]: due to the Coriolis force, the leading sunspots are closer to the equator than the trailing sunspots. In the end of the cycle when the sunspot pairs appear at very low heliolatitudes, the leading spots diffuse across the equator and their flux is canceled by the oppositely signed flux of the leading spots in the opposite hemisphere. The flux of the trailing sunspots and of the remaining sunspot pairs is carried to the poles. The excess trailing polarity flux cancels the poloidal field of the old cycle (with the polarity corresponding to the leading sunspots polarity), and accumulates to create the new poloidal field with a polarity opposite to the polarity of the previous cycle.

Though the dynamo α -effect branch is not clear in details, it is almost obvious that the meridional circulation itself is the key to the understanding of the cycle characteristics. The cycle's amplitude and period are determined by the speed of the meridional circulation [6–11], while the coefficient of the turbulent diffusivity determines the regime of operation of the dynamo [12–14].

4 Estimation of the Speed of the Meridional Circulation Based on Geophysical Data

We can suppose that the generation of the poloidal field is complete when the equatorial coronal holes merge with the polar coronal holes, forming giant coronal holes spreading from the pole to the equator and further in the opposite hemisphere.

An analogous area—a coronal hole of the opposite polarity—forms at antipodal latitudes and longitudes. The formation of these coronal holes corresponds to the geomagnetic activity maximum related to the outflow of high speed solar wind

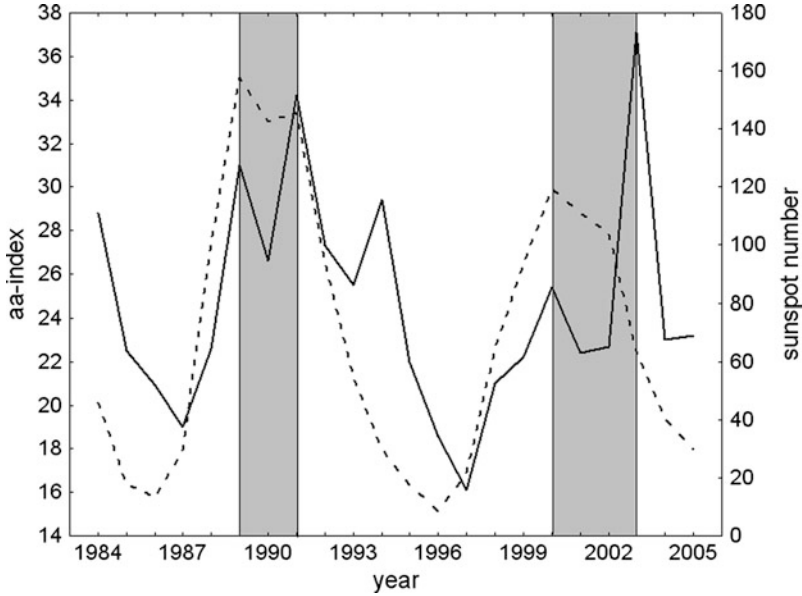


Fig. 6 Sunspot number (solid line) and aa-geomagnetic index (dashed line). The time between sunspot maximum and the following aa-maximum (grey shading) is used for estimation of the surface poleward meridional circulation, and the time between this aa-index maximum and the next sunspot maximum (white shading)—for estimation of the deep equatorward meridional circulation (see the text)

streams from such open magnetic field configurations. Therefore the time between the sunspot maximum and the geomagnetic activity maximum on the sunspot declining phase can be considered equal to the time it takes the flux to reach from sunspot latitudes to the poles (Fig. 6, grey shading). From it the speed of the surface poleward meridional circulation can be derived. It is found that the faster the poleward circulation, the lower the amplitude of the following sunspot maximum. The correlation coefficient is $r = -0.7$, $p = 0.03$ (Fig. 7; note the reversed scale). The negative correlation means that the advection time-scale is shorter than the diffusion time-scale and with a faster poleward flow there is less time for the leading polarity flux to diffuse across the equator and to cancel with the leading polarity flux in the opposite hemisphere, less uncanceled trailing-polarity flux reaches the pole to form the polar field of the next cycle. From the weaker polar field, a weaker toroidal field is generated in the base of the convection zone [6].

The speed of the reverse flow (the deep equatorward meridional circulation), can be estimated from the time between the geomagnetic activity maximum on the sunspot declining phase and the following sunspot maximum (Fig. 6, white shading). Depending on the diffusion coefficient, three regimes are possible: fully advection-dominated if the diffusion in the upper part of the solar convection zone is very low, intermediate with a higher diffusivity, and fully diffusion-dominated with

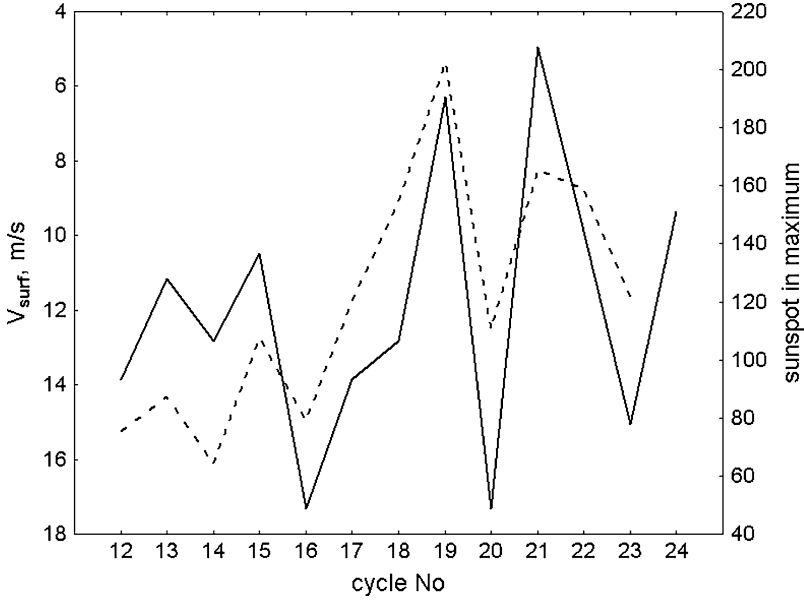


Fig. 7 Speed of the surface poleward circulation (*solid line*, note the reversed scale) and the amplitude of the following sunspot maximum (*dashed line*)

a very high diffusivity [13]. In the intermediate diffusivity regime which is probably observed in the Sun, a part of the flux diffuses to the base of the convection zone without reaching the pole, “shortcircuiting” the meridional circulation, another part completes the full cycle (Fig. 8). In this case the sunspot cycle is a superposition of the two surges of toroidal field, and a double peaked cycle maximum is observed. The correlation between the speed of the deep equatorward circulation and the amplitude of the following sunspot maximum indicates the regime of operation of the solar dynamo in the bottom part of the solar convection zone. The observed positive correlation (Fig. 9) is an evidence of a diffusion-dominated regime which means that diffusion is more important than advection near the tachocline, and with a faster flow, there is less time for diffusive decay of the flux during its equatorward transport along the tachocline. If advection were more important than diffusion, diffusive decay would be less efficient, and a faster flow would mean less time for toroidal field generation and therefore a lower sunspot maximum [12].

To summarize, there is a negative correlation between V_{surf} and the following polar field ($r = -0.8$), a negative correlation ($r = -0.75$, $p = 0.005$) between V_{surf} and the following V_{deep} , a positive correlation ($r = 0.81$, $p = 0.01$) between V_{deep} and the following maximum sunspot number. However, there is no correlation at all between the maximum sunspot number and the following V_{surf} .

Therefore, the sequence of relations is $V_{surf} \Rightarrow B_{pol}$, $V_{deep} \Rightarrow B_{tor}$. Here the chain breaks and it is not possible to forecast further.

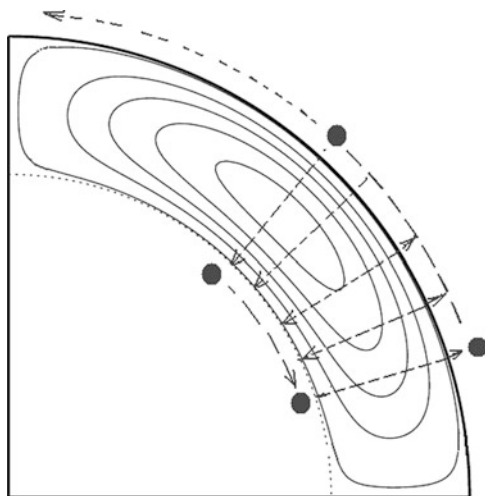


Fig. 8 Intermediate diffusivity regime in the upper part of the solar convection zone: a part of the flux diffuses to the base of the convection zone without reaching the pole, “shortcircuiting” the meridional circulation, another part completes the full cycle to the poles, down to the high-latitude tachocline, and equatorward to sunspot latitudes

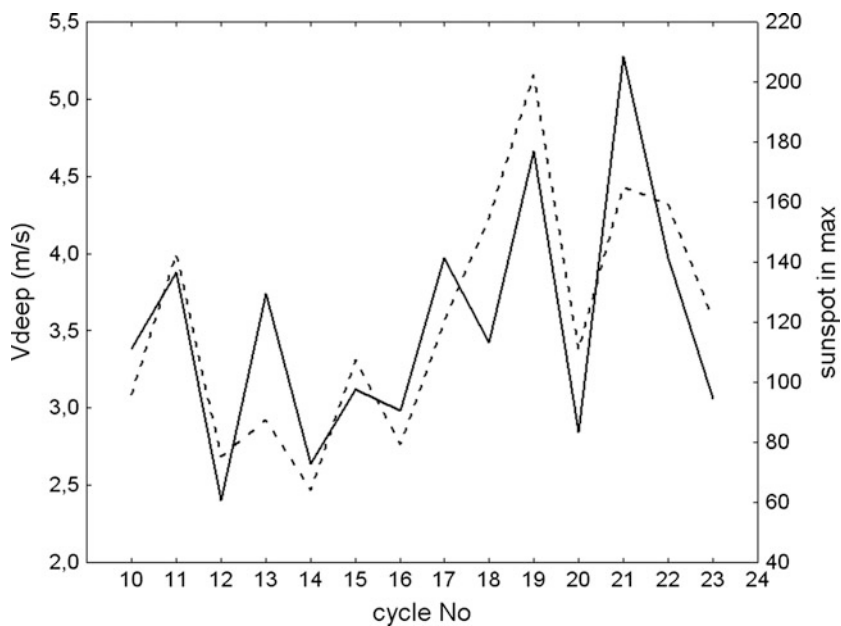


Fig. 9 Speed of the deep equatorward circulation (*solid line*) and the amplitude of the following sunspot maximum (*dashed line*)

5 Can Planetary Tidal Forces Modulate Solar Activity?

The dependence of the sunspot cycle amplitude on the speed of the meridional circulation makes it possible to understand the periodically debated but so far not physically supported correlation between the variations in solar activity and the orbital periods of the major planets. Of course, the planets themselves cannot cause solar activity, but it can be speculated that they can modulate solar activity by modulating the speed of the meridional circulation [15, 16].

The tidal forces are vertical and horizontal. The elevation caused by the vertical tidal force is very small (~ 1 mm) which is often cited as an argument against the reality of planetary influences on the solar activity. The horizontal tidal force causes acceleration in both zonal and meridional directions. The zonal acceleration can in principle change the rotation rate which is important for the generation of the toroidal magnetic field at the base of the solar convection zone ($0.7 R_S$). However, the background rotation rate is quite high, the linear rotation velocity is $\sim 2,000$ m/s; moreover, the tidal forces decrease with depth d as d^2 , while the density increases and at $0.7 R_S$ it is $\sim \text{gr/cm}^3$. Moreover, the sign of the tidal force depends on latitude so the latitudinal average is zero.

More significant can be the influence of the meridional tidal acceleration of the surface meridional circulation. The meridional tidal force is always directed toward the equator so its effect is to slow down the poleward meridional circulation, and a slower poleward meridional circulation results in a higher sunspot maximum. Let us estimate this effect. The tidal forces are biggest at the surface where the density is $\rho \sim 10^{-5} \text{ gr/cm}^3 = 10^{-2} \text{ kg/m}^3$. The speed of the surface poleward circulation is ~ 10 m/s. The acceleration is determined by $a = F/\rho$ where $F \sim 10^{-10} \text{ N/kg}$, therefore $a \sim 10^{-8} \text{ m/s}^2$. The characteristic time during which the tidal force acts upon the meridional circulation is the time for the transport of the flux from sunspot latitudes to the poles which is $\sim 2\text{--}3$ years, or $\sim 10^8$ s. This gives $dV_{\text{surf}} \sim \text{m/s}$, in agreement with the observed variations of the surface meridional circulation.

Figure 10 demonstrates the dependence of the amplitude of the sunspot maximum in consecutive solar cycles on the planetary tidal force modulating the poleward surface meridional circulation. This is a possible explanation of the mysterious correlations between the planets' orbital motions and the sunspot cycle amplitude.

An important question is whether the planetary modulation of solar activity can be used to predict the amplitudes of the future solar cycles. The planetary motions, respectively the planetary tidal forces, can be calculated with great accuracy for centuries ahead. To estimate the modulation of V_{surf} and therefore of the amplitude of the following sunspot maximum, it is necessary to calculate the tidal force from the major tidal-generating planets during the period in which the flux is being carried by the meridional poleward circulation from sunspot maximum latitudes to the poles—that is, in the interval between the previous sunspot maximum and the subsequent geomagnetic activity maximum on the sunspot declining phase. Knowing the dates of these two maxima in solar cycle 23, it can be predicted that

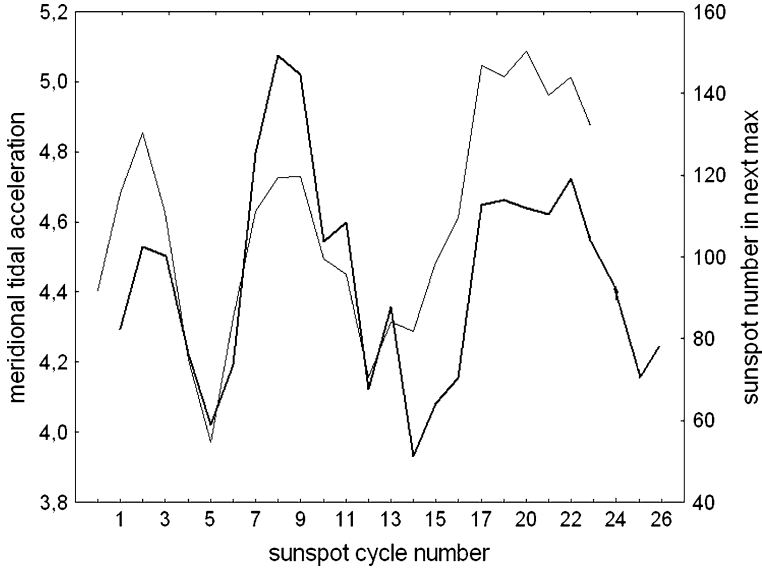


Fig. 10 Planetary tidal force between sunspot maximum and the following geomagnetic activity maximum modulating the poleward surface meridional circulation (*solid line*) and the amplitude of the following sunspot maximum in consecutive solar cycles

the maximum sunspot number in cycle 24 will be under 80 (Fig. 10). The times of the subsequent sunspot and geomagnetic activity maxima, and therefore of the sunspot cycles amplitudes, can be predicted with increasing uncertainty, but cycle 25 is expected to be even lower than cycle 24, and cycle 26 will probably be the beginning of a long-term increase in solar activity. Therefore, the beginning cycle 24 and the next two cycles, 25 and 26, are expected to constitute a secular minimum in solar activity.

6 Tachocline Dynamo as Compared to the Distributed/Near-Surface Dynamo

Most contemporary dynamo models assume that the toroidal magnetic field is generated at the base of the convection zone, at the so-called tachocline or just below it, and then emerges to the surface as sunspots [17–20]. The confidence in the preference of the deep-seated dynamo comes from the fact that this region is stable to allow for the accumulation of the magnetic flux, in spite of the magnetic buoyancy.

However, the observations of the rotation velocity of the emerging magnetic flux in different latitudinal zones seem to indicate a relatively small depth of the sunspots

[21], apparently rooted in the near-surface layer. This idea is also supported by the results of local helioseismology [22]. All active solar phenomena are the result of interaction of deep poloidal (axially symmetric quasidipolar) fields and non-axially symmetric (quadrupolar) fields. Therefore there are two regions of magnetic field generation: deep (tachocline) and near-surface layers [23].

A number of processes in which magnetic field generation is possible, occur in the near-surface layer, at the level of $0.995 R_S$ known as leptocline. An oscillatory regime of the seismic radius and a drastic change of the turbulent pressure are observed there, there are indications of the change of the radial gradient of the rotation velocity at 50° latitude, etc. Apparently that is the place where the sunspot magnetic field is rooted [24].

In 2005 Axel Brandenburg [25] formulated the arguments in favor and against the tachocline and the distributed near-surface dynamo. Among the disadvantages of the tachocline dynamo he pointed out that 1.3-year oscillations are observed in the tachocline, and there is no sign of the 11-year cycle; The tachocline dynamo cannot explain the generation of kG locally distributed magnetic fields. Among the advantages of the distributed near-surface dynamo, Brandenburg [25] mentioned the existence of topological pumping in the near-surface layer.

In a recent work [26] a model is presented with topological pumping in the near-surface layer. Such a dynamo model with a near-surface shear layer can satisfy all requirements for the generation of magnetic flux and in the same time is able to reproduce all known statistical features of the solar cycle, in general the relation between the cycle period and amplitude (Waldmeier rule) [26].

The downward turbulent pumping of the horizontal magnetic field (related to either the toroidal or the meridional component of the magnetic field) provides the best agreement of the characteristics of the theoretical dynamo with observations, increasing for the given turbulent diffusivity profile the magnetic cycle period. The model provides the asymmetry of the cycle growth and decay times (as well as the duration of the phases) of the toroidal magnetic field. The asymmetry is growing with growing gradient of the turbulent diffusion in the near-surface layer. From the calculations it also follows that in the beginning of the cycle the current helicity changes sign in the near-surface layers.

Therefore, the turbulent topological pumping is the mechanism transforming the weak diffusive field created by the deep tachocline dynamo into the strong toroidal field of the active regions.

7 What Can We Expect in the Coming One or Several Cycles?

To present time (winter 2011) there are a lot of predictions for the 24th solar cycle amplitude. Recent reviews of solar cycle forecasting methods and their results for cycle 24 are given in [27, 28]: see Table 1.

Table 1 A selection of forecasts for cycle 24 (from [28])

Category	Peak amplitude	Link	Reference (following [28])
Precursor methods			
Minimax	80 ± 25	Eq. (10)	Brown (1976); Braja et al. (2009)
Minimax3	69 ± 15	Eq. (11)	Cameron and Schssler (2007)
Polar field	75 ± 8	Sect. 2.2	Svalgaard et al. (2005)
Polar field	80 ± 30	Sect. 2.2	Schatten (2005)
Geomagnetic (Feynman)	150	Sect. 2.3	Hathaway and Wilson (2006)
Geomagnetic (Ohl)	93 ± 20	Sect. 2.3	Bhatt et al. (2009)
Geomagnetic (Ohl)	101 ± 5	Sect. 2.3	Ahluwalia and Ygbuhay (2009)
Geomagnetic (interpl.)	97 ± 25	Sect. 2.3	Wang and Sheeley Jr (2009)
Field reversal	94 ± 14	Eq. (12)	Tlatov (2009)
Extrapolation methods			
Linear regression	90 ± 27	Sect. 3.1	Braja et al. (2009)
Linear regression	110 ± 10	Sect. 3.1	Hiremath (2008)
Spectral (MEM)	90 ± 11	Sect. 3.2	Kane (2007)
Spectral (SSA)	117	Sect. 3.2	Loskutov et al. (2001)
Spectral (SSA)	106	Sect. 3.2	Kuzanyan et al. (2008)
Attractor analysis	87	Sect. 3.3.1	Kilcik et al. (2009)
Attractor analysis	65 ± 16	Sect. 3.3.1	Aguirre et al. (2008)
Attractor analysis	145 ± 7	Sect. 3.3.1	Crosson and Binder (2009)
Neural network	145	Sect. 3.3.4	Maris and Oncica (2006)
Neural network	117.5 ± 8.5	Sect. 3.3.4	Uwamahoro et al. (2009)
Model based methods			
Explicit models	167 ± 12	Sect. 4.3	Dikpati and Gilman (2006)
Explicit models	80	Sect. 4.3	Choudhuri et al. (2007)
Explicit models	85	Sect. 4.3	Jiang et al. (2007)
Truncated models	80	Sect. 4.4	Kitiashvili and Kosovichev (2008)

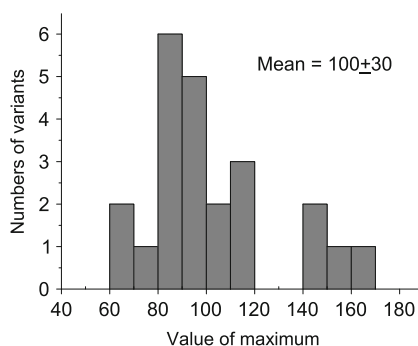
Fig. 11 Occurrence of different variants of forecast of amplitude of Solar cycle 24

Figure 11 illustrates the popularity of the different values of predicted amplitudes of cycle 24 according to Table 1. We see that the range of possible values of the maximum of the following cycle is very wide. Fortunately, the nearest 1–2 years must show the real level of activity of the cycle 24.

Table 2 Solar activity evolution in the next decades: scenarios (from [29])

Method	$W_{\max}(24)$	$W_{\max}(25)$	$W_{\max}(26)$	Nearest great minimum: year, type
1. Multiscale cloning	65 ± 20	80 ± 20	85 ± 20	$T \sim 2, 100$, Dalton
2. Multiscale autocorrelation	120 ± 25	75 ± 25	90 ± 25	Dalton
3. Statistics of states	100 ± 15	100 ± 15	–	Dalton
Mean	95 ± 10	90 ± 20	85 ± 20	Dalton

And what further? In [29] some forecasts and possible scenarios of solar activity evolution in the following nearest decades are presented—see Table 2. The main essences of these named methods are the following:

1. Using “the method of multi-scale cloning,” the time set of solar activity are split into wavelet components. Then for each of them various reasonable variants of future behavior are considered, constructed from their previous behavior (the “clones”). Carrying out the procedure for all frequency components and for all chosen ways of the cloning, thereafter we make the inverse wavelet-transform to obtain a forecast (scenario) of behavior of the function (solar activity variation) in the future.
2. Method of multi-scale autocorrelations: this method also implies the multi-scale approach and is based on precedents of behavior of a function in the past, but does it in the aspect of the nearest past from the point of view of the selected time scale.
3. In the method named “Statistics of states,” the solar activity evolution is presented in the form of a set of states (typical levels). Then the symbolical statistics is used and on the basis of the previous history of solar activity, the probabilities of its different subsequent states are calculated.

Following Table 2 one can conclude that after the high cycles of the twentieth century, the solar activity now turns to an average level (rather than to a low one, as some authors believe), and only by the latter half of twenty-first century we can wait for a great minimum of Dalton’s type. The probability of a Maunder’s type minimum is minimal.

An “alarming” fact has been found by Penn and Livingston [30]: a gradual decrease in the sunspot umbral magnetic field strength and the corresponding decrease in sunspot brightness. If this trend continues further, virtually no sunspots as dark features will be observed in cycle 25 (!). On the other hand, Pevtsov et al. [31] claim that the trend observed in [30], is actually a part of 11-year cyclical variations, and that in the nearest future the average sunspot magnetic fields will begin growing.

Another question is the date of the following sunspot maximum. It is known that at sunspot maximum the magnetic fields at the poles reverse. At this time the current sheets pass over the solar poles. In the previous three cycles there was no more than a year from the moment the current sheet passed the 70° latitude to the sunspot

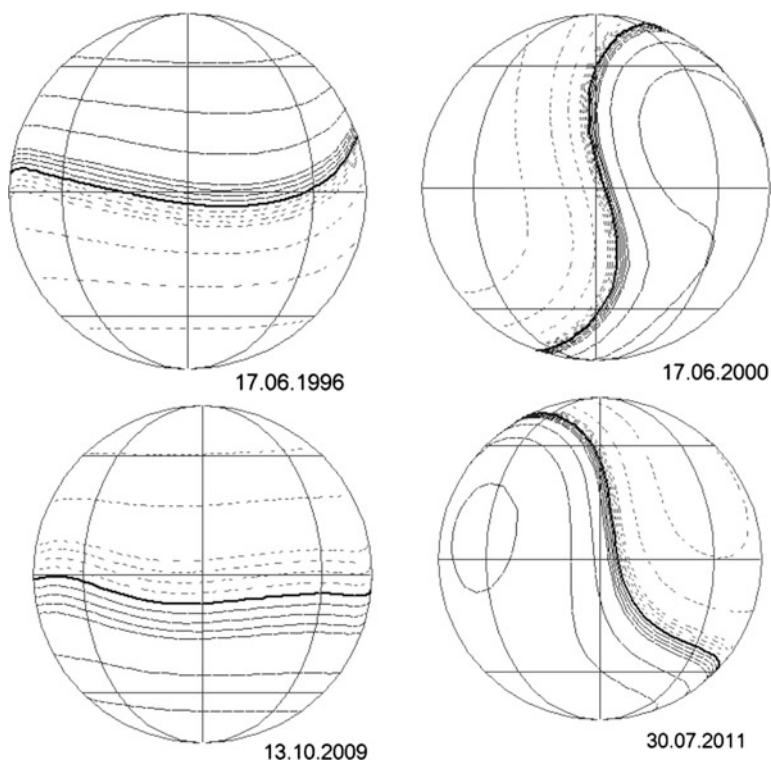


Fig. 12 Magnetic field in the 22/23 and 23/24 cycles minima (*left panels*, view from the Earth) and in cycle 23 maximum (*upper right panel*) and in 2011 (*bottom right panel*). Calculated on the base of the Stanford data

maximum. Recent observations indicate that already in July 2011 the inclination of the current sheet was 70° . This may mean that we are already on the eve of the maximum.

Figure 12 presents the solar magnetic field in the minima (*left panels*, view from the Earth) and at the North pole in the cycle 23 maximum (*upper right panel*) and at present (*lower right panel*). The situation in cycle 23 maximum and now is similar. The axial dipole is close to reversal, and the equatorial dipole is growing very slowly (<http://wso.stanford.edu/gifs/DipallR.gif>).

On the other hand, we should remind that the duration of the ascending branch of the cycle, as well as its amplitude, are determined by the speed of the deep circulation. There are indications that at present this circulation is very slow, which allows us to forecast a late occurrence and a low amplitude of the next sunspot maximum.

Finally, a very interesting consideration was presented by Altrrock [32]. Figure 13 shows coronal activity parameters, smoothed over seven solar rotations. In all

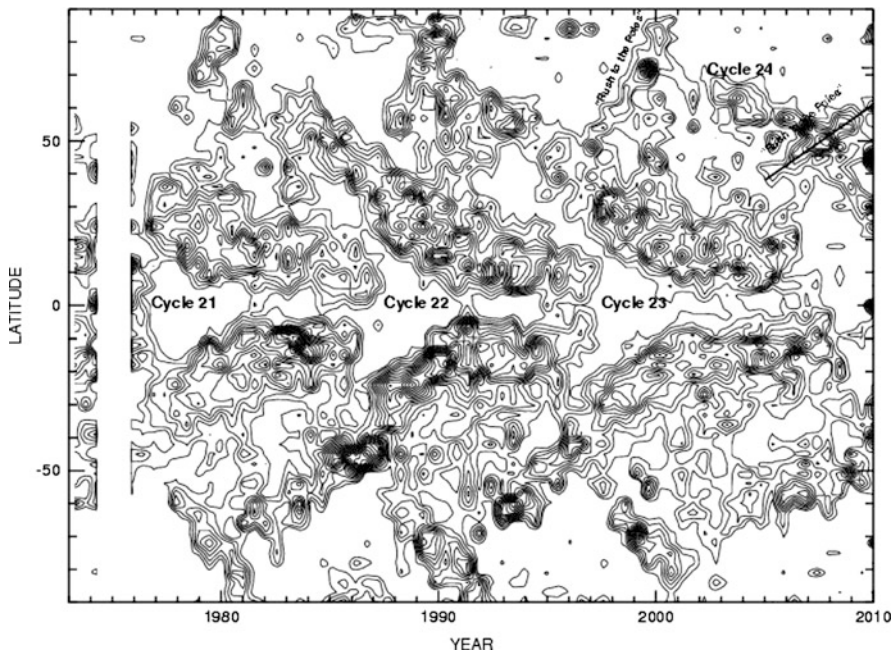


Fig. 13 Annual northern plus southern hemisphere averages of the number of Fe XIV intensity maxima from 1973 through 2009. The “Rush to the Poles” around 2000 is indicated, as well as the extended solar cycle 24, beginning in approximately 1999. From [33]

previous cycles shortly before the maximum a “rush to the poles” began. But at present there is no sign of this rush which could mean that cycle 25 will be also very low.

8 A General Conclusion

The unusual characteristics of cycle 23 proved very useful for studying the solar cyclicity. The unusual features of cycle 23 logically turned into unusual cycle 24. The peculiar characteristics of these two cycles turned out to be a convenient touchstone for understanding the nature and properties of solar cyclicity, the physical processes of its occurrence, the basis of its forecast.

Some “mathematical” scenarios testify that 20–30 years we shall be in conditions of average state of activity, but a diversity of various evidences indicate that in the near future we are entering a grand minimum similar in its characteristics to the Dalton minimum.

From the point of view of the dynamo theory, additional arguments appeared in favor of the distributed dynamo. The magnetic field generated in the tachocline is

diffusive, while the sunspot cycle itself with its specific properties (magnetic field intensity in the sunspots, cycle amplitude and duration, asymmetry of the ascending and descending branches, Waldmeier effect) is generated with the participation of processes in the near surface layer. Such a scheme seems to be the most acceptable one.

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The Evolution of Cyclic Activity of the Sun in the Context of Physical Processes on Late-Type Stars

Maria M. Katsova

Abstract Features of the solar cycle in the context of stellar activity are investigated. We discovered reliably differential rotation in chromospheres of some stars and presented the first stellar butterfly diagrams. These stars possess less regular variability and do not demonstrate excellent cycles. This is the first evidence for differences of the solar activity from processes on stars with Excellent cycles. We compare indices of the chromospheric activity of the Sun with that for above 1,300 northern and southern stars whose activity revealed during planet search programs. We argue the matter pro and con for two possible ways of an evolution of activity from a contraction phase to 10 Gyrs. When a young star brakes down, the chromospheric and the coronal activity weaken synchronously. The solar-like activity of the most main sequence F and early G stars does evolve by this path. The activity of the later stars from G5 to K7 after a definite level evolves by another way: the chromospheric activity diminishes up to the solar level, while coronae stay stronger than the solar one. Two possible paths of the evolution of activity are associated with the different depth of the convective zone of these stars. Physically this means that the relative input of small- and large-scale of magnetic fields differs for F–G and K stars.

1 Introduction

Changes of characteristics of activity in timethe evolution of activitybegan to be under investigation during the last 50 years. The HK Project, which was intended to reveal cyclic activity of stars in the solar neighborhood, became a basis for studies

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of the long-term variability of chromospheric activity. X-ray observations of late-type stars also made a major contribution to this field. Already the first results of studies of stellar activity demonstrated that the activity level is related to the axial-rotation rate. This means that angular momentum loss represents a basic factor determining the evolution of the activity. Recently, the rotational periods of stars in star-forming regions in the stage of gravitational contraction were determined directly from observations of the rotational modulation of their optical continuum emission. It was found that the rotational periods of young stars with masses from 0.8 to $1.2M_{\odot}$ vary from 7 days to about 1 day as the age varies from 1 to 70 Myr [1, Fig. 7]. The subsequent braking of the rotation occurs over significantly larger time scales of billions of years. The rotational period of a star in the deceleration stage is proportional to square root of the age as it follows from Skumanich's law [2], derived from chromospheric observations. This dependence was recently investigated in more detail in the projects "The Sun-in-Time" and "Living with a Red Dwarf" [3, 4]. An analysis of these results for G, K and M dwarf stars confirms Skumanich's law, without the need for any serious changes to it. This became the basis for the development of the method of gyrochronology that follows from combination of two relationships—"activity–rotation" and "rotation–age". This proved to be quite fruitful, since it enables to estimate of the age of the star with a given spectral type (for dwarfs of a given mass) using one parameter only, the level of its chromospheric activity [5]. However, the final correlation for "activity–age" leads to isochrone, which reflects only formally dependence of the level of the chromospheric activity on the spectral class of a star. In many respects this relation is determined by linear relations between indices of the coronal $\log L_X/L_{bol}$ and chromospheric $\log R'_{HK}$ activity. New observational data allow us to construct similar diagram for "chromosphere–corona" indices that can clear up physical aspects of the evolution of solar-like activity.

Current exoplanet-search programs provide as a by-product data on the rotation and activity levels of late-type stars. Indices of chromospheric activity S analogous to the corresponding index in the HK Project have been derived for several thousand F–M stars. The HK Project started in 1966 by O. Wilson at the Mount Wilson observatory was continued afterwards at several observatories, and new S values were reduced to a unified system. This index S is the ratio of the fluxes at the centers of the H and K lines of Ca II at 3,933.66 and 3,968.47 Å to the fluxes at the nearby continuum levels at 4,001–3,901 Å. The line width is 1 Å, and the width of the band at the continuum level is 20 Å. The S index takes into account the contributions of both the chromospheric and photospheric radiations. Later, to exclude the effect of the photospheric radiation, the quantity R'_{HK} —the flux of the stellar radiation in both Ca emission lines normalized to the bolometric luminosity—was introduced in place of S [6] (see [7] for details of the reduction from S to R'_{HK}). This transformation is well calibrated only for color indices $0.44 < B-V < 0.9$, whereas it becomes more uncertain for redder stars.

Thus, the chromospheric activity index R'_{HK} can be estimated only for F, G and K stars, and cannot be derived for M stars due to their weak blue (UV) continuum.

There are now more than 1,300 stars with reliable determined R'_{HK} indices. This dataset can be used to study the place occupied by solar activity among processes occurring on other late-type stars at a qualitatively new level, and to trace the evolution of solar-like activity from ages of 500 Myr to 10 Gyr. Considering these data together with current soft X-ray observations, we will analyze the relationship between the activity levels in the chromosphere and corona. In conclusion, we will discuss the consequences of our analysis of this larger amount of material for understanding of solar-like activity on late-type stars of various ages.

2 Chromospheric Activity of Late-Type Stars

Let us consider data on chromospheric activity obtained during exoplanet-search programs. The “California and Carnegie Planet Search Program” at the Keck Observatory has carried out a survey of the Northern Sky over 6 years. The HIRES echelle spectrometer mounted on the 10-m Keck Telescope operated at 3,850–6,200 Å with the high resolution of 67,000. The high signal-to-noise ratio of these data enabled the reliable detection of the H and K line fluxes not only for F, G and K stars, but also for M dwarfs [7]. The “The Magellan Planet Search Program” targeted at southern stars was started with the 6.5-m telescope of Las Campanas Observatory in Chili in the Autumn of 2002. These observations have been done using the MIKE echelle spectrograph yielding spectra at 3,900–6,200 Å with a resolution of 70,000 in the blue and 50,000 in the red. The first results of monitoring of the chromospheric radiation of several hundred late-type stars were recently published [8]. The final observational data for both surveys were reduced to a unified system corresponding to that adopted for the Mt. Wilson HK Project.

We combined both the northern and southern observations, and chose stars with trustworthy R'_{HK} values. All M stars were eliminated automatically by this approach, as well as duplicate targets. The final list of 1,334 stars contains F, G and K dwarfs with color indices $B - V \leq 0.9$.

Figure 1 presents the indices of chromospheric activity R'_{HK} for the entire set of observational data. Figure 1a shows that the vast majority of stars are G stars, with R'_{HK} indices between -4.9 and -5.1 . This is also illustrated by the histogram in Fig. 1b. A number of stars are characterized by higher activity levels, compared to most of the objects. Their chromospheric activity is close to the level typical of younger Hyades stars with ages of about 600 Myr. The corresponding isochrone is drawn in Fig. 1a. Similar results were obtained for northern stars in [9] and for southern stars in [8]. Several authors have noted the bimodal character of the activity distribution that can be traced in Fig. 1b. The corresponding centers of the distributions are $\log R'_{HK} = -5.01$ and -4.53 , and the distributions have widths of 0.25 and 0.31, respectively.

Such a large observational dataset enables study of the place occupied by the Sun among other late-type stars. For definiteness, we took the chromospheric index

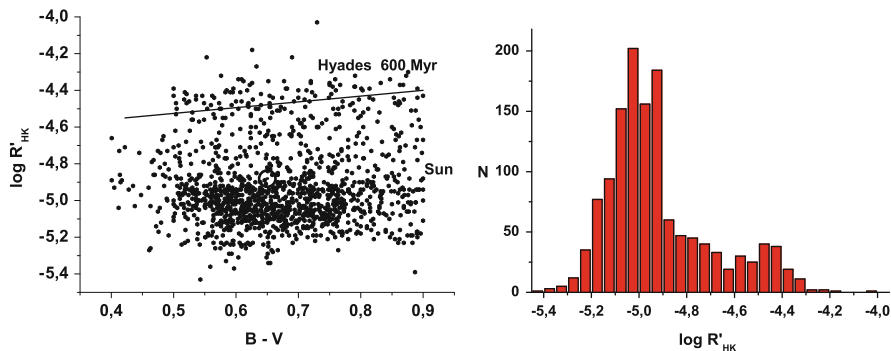


Fig. 1 Chromospheric activity of northern and southern stars recorded during exoplanet-search programs: (a)—chromospheric-activity index versus color index. The Sun is marked as a *circle* with $B - V = 0.65$. The isochrone is drawn for the age of the Hyades, 600 Myr; (b)—histogram of the number of stars with specified chromospheric-activity indices. The step in $\log R'_{HK}$ is 0.05

$\log R'_{HK}$ for the Sun to be -4.88 , which corresponds to epochs of relatively high activity (Wolf numbers of about of 80). The chromospheric activity of the Sun is clearly higher than for the overwhelming majority of stars in the solar neighborhood. Of course, differences in this parameter of a few hundredths are small, but can be significant when studying the formation of cyclic activity.

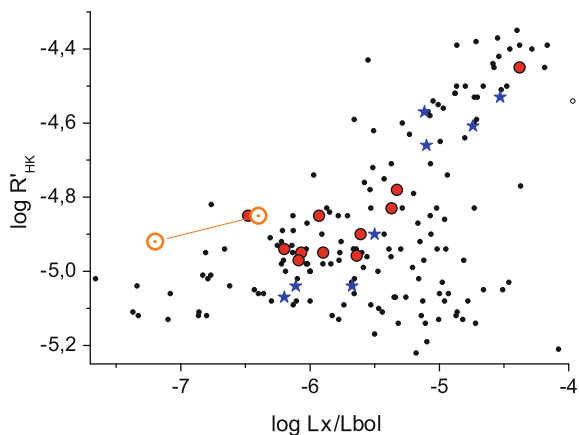
Note that our analysis reveals the existence of a certain number of F, and G stars whose activity is much weaker than the activity of the Sun, and even of most stars. This follows from both Fig. 1a, b.

3 Comparison of Chromospheric and Coronal Activity of Late-Type Stars

The appearance of a wider array of observational data for active late-type stars enables us to return to a comparison of the activity developing in different layers of the outer atmosphere. Observations in the soft X-ray can be used to compare data on the activity of late-type stars in the chromosphere and corona. We adopted ROSAT measurements for most of the stars (see, for instance, [10, 11], and used XMM-Newton data for several dozen stars [12, 13]. We adopted the X-ray-to-bolometric luminosity ratio, L_X/L_{bol} , as an index of coronal activity.

We obtained the coronal activity index $\log L_X/L_{bol}$ by reviewing the available soft X-ray data: XMM Newton observations, ROSAT data, and a few measurements of X-ray luminosities reported in [14]. Although the X-ray fluxes measured for the same star by two spacecraft were usually very close, in some cases, the difference exceeds an order of magnitude. This is primarily associated with variability of sources observed at different times. These few cases were analyzed in more detail,

Fig. 2 The diagram of the chromospheric and the coronal activity for the late-type stars. The stars of the basic data set are marked as *dots*. Accordingly to the type of a cycle, the stars of group “Excellent” are marked as *circles*, the stars of “Good” group are indicated as *asterisks*; the Sun at the maximum and at the minimum is denoted its own sign, connected by the *direct line*



and evidence was found for the possible development of flares on the star, or the presence of long-term changes in other energy ranges. As a result, we obtained a dataset for 172 stars (including the Sun) with certain indices of chromospheric and coronal activity (Table in [15]). Figure 2 shows a comparison of the chromospheric and coronal luminosities of the selected stars. Stars with Excellent and Good cycles detected in the HK Project are denoted by various symbols, as in our previous papers [16, 17]. Here, we also show the values for the Sun at its minimum and maximum activity. Adding new objects, most of which have lower levels of chromospheric activity, has significantly changed the general form of the dependence of $\log R'_{HK}$ on $\log(L_X/L_{bol})$. Of course, this reflects the basic fact that the activity level of the outer atmosphere is determined by the rate of axial rotation of the star. This is manifest most strongly in the existence of the previously discovered linear relation between the activity indices. However, in addition, some F and G stars with low chromospheric-activity levels (discussed in connection with Fig. 1) are located lower than the stars with cycles along the horizontal axis. These stars possess quite powerful coronas; i.e. their L_X/L_{bol} ratios are one to two orders of magnitude greater than that of the Sun at its maximum activity. Note that stars with detected exoplanets are characterized by fairly powerful coronas and very different levels of chromospheric activity. In [5], the linear relation for the entire interval of activity indices was given by

$$\log R'_{HK} = -4.54 + 0.289[\log(L_X/L_{bol}) + 4.92] \quad (1)$$

This reflects the behavior of points only in the upper branch of Fig. 2, where chromospheric activity increases with rise of the coronal radiation. The addition of active stars detected during exoplanet search programs increased the number of objects with chromospheric activity indices below $\log R'_{HK} = -5.0$. In accordance with the above formula, it was expected that the new points would fill the lower left corner of the diagram. However, they are located almost uniformly along the

horizontal axis at roughly similar low chromospheric activity levels. In other words, we have discovered a group of stars with chromospheric activities lower than the minimum level for the Sun, whose coronal radiation spans a broad range. Among them are stars with weak chromospheres and strong coronae. The appearance of stars with slowly varying chromospheric activity in the region where the X-ray luminosity changes substantially suggests the presence of a second branch on this diagram. The difference between these branches is such that the upper branch has a comparatively greater number of F stars and the lower branch a greater relative number of K stars. Figure 2 shows the presence of a number of objects with fairly powerful coronae (with about $\log(L_X/L_{bol}) = -5$) situated between these branches.

4 Features of the Cyclic Activity of the Sun

To understand the physical processes leading to the diversity of solar-like activity phenomena in late-type stars of different ages, it is helpful to consider some results of solar researches. We will briefly discuss the following issues: (1) the general character of the solar cycle and the relationship of phenomena to magnetic fields on various scales; (2) characteristics of the differential rotation; (3) the simultaneous manifestation of cycles of different durations.

The variations in the monthly average Wolf numbers for all 23 cycles show epochs of very low activity, an extreme example of which was the Maunder minimum, and of high activity, as in the late 1950s. It has been suggested to call such epochs Grand Minima and Grand Maxima [18]. In spite of the very different lengths of series of observations, some stars with Excellent cycles show more regular cyclic variations of their chromospheric emission.

Activity on the Sun is associated first of all with development of regions of local magnetic fields. But because the Sun is close to us, we are able to study in details some phenomena on the surface which do not effect on the brightness and other characteristics of the Sun as a star. Influence of the large-scale magnetic fields manifests itself in formation of coronal holes as well as in a tendency of appearance and existence of active regions near to definite longitudes at some epochs of the cycle. Besides, the evolution of the large-scale magnetic field determines the drift of active regions from high latitudes to the equator during the cycle (the butterfly diagram). Some non-stationary processes related to the large-scale magnetic field act as triggers for events occurring near spots, which are indicators of strong local fields. This has recently been observed during multiwavelength observations of solar flares.

We share the view that, although activity is associated with local magnetic fields, activity on the Sun is regulated primarily by large-scale magnetic fields.

When we began our studies of the differential rotation of HK Project stars, we expected the differential rotation of stars with Excellent cycles to be similar to that of the Sun. We selected 20 stars, including both such stars and others with less

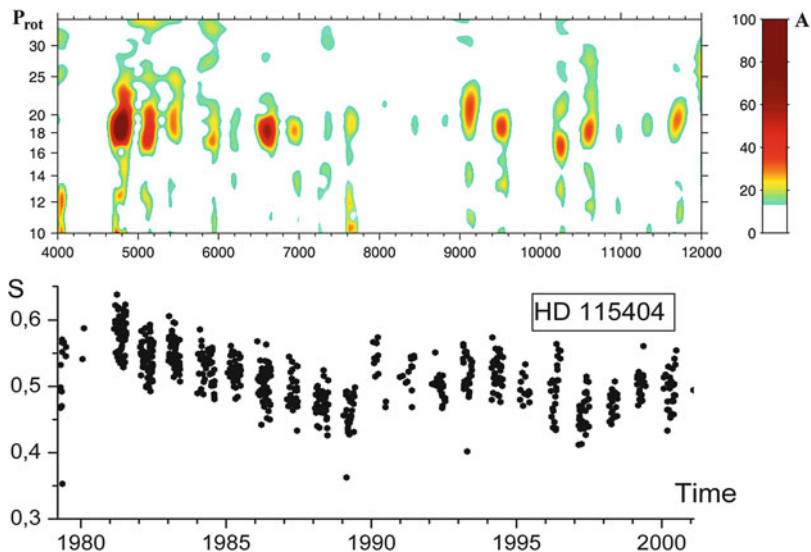


Fig. 3 (a) The “period-time” diagram for the star HD 115404. The scale of amplitudes of the wavelet transformation A on the right side is normalized to 100; the rotational period is given in days; (b) the temporal behaviour of the S index of the chromospheric activity. From [19]

regular activity (with Good cycles and others). Our wavelet analysis of the long-term variations of the S index of chromospheric activity for all of these stars gave a surprising result: spin-down of the rotation of large-scale inhomogeneities at epochs of high activity, repeating near the maximum of each cycle [19]. This is illustrated by a result for the star HD 115404 with the Good-pronounced 12-year cycle (Fig. 3). The mean period of rotation of surface inhomogeneities of this star of 18.5 days increase in 3–4 days in the epoch of the first observable maximum of the cycle near 1981 and this effect repeats 12 years later at the next maximum. The degree of the differential rotation is estimated as 0.14 that is close to the corresponding solar value. For another star with the Good cycle HD 149661 this value is twice smaller. Under certain assumptions it becomes possible to construct stellar butterfly diagrams.

Further we compare these results for stars with available data concerning solar differential rotation (see e.g., the review in [20] obtained with various techniques of measurements of differential rotation ranging from Doppler shift, Doppler feature tracking, magnetic feature tracking, and p-mode splitting. We confine ourselves to data of the brightness of the solar corona in the Fe XIV 5,303 Å coronal green line. These database collected by J. Sykora (Slovakia) for several solar cycles from 1939 to 2001 presents the daily measurements on the eastern (E) and western (W) limbs carried out for every 5° of latitude and recalculated to the central meridian.

A detailed study of the coronal rotation over the several last cycles is presented in [21] (see also references therein), which considers data on the brightness of the

Fe XIV 5,303 Å coronal green line. The rotation at various solar latitudes was determined from data averaged over six Carrington rotations. Among the main conclusions of [21] is that differential rotation occurs during the entire growth phase of the solar activity, while the rotation is close to rigid-body in the middle of the decay phase.

We proposed new representation for this database for calculation the total emission of the solar corona for each day of observations. We introduced as a new index of the solar activity GLSun (The Green-Line Sun) that is proportional to the total emission of the green corona behind the limb and on the visible disk. This index is purely observational and is free of the model-dependent limitations imposed on other indices of coronal activity. The GLSun index describes well both the cyclic activity and the rotational modulation of the brightness of the corona of the Sun as a star. The GLSun series was subject to a wavelet analysis similar to that applied to long-term variations in the chromospheric emission of late active stars. The brightness irregularities in the solar corona rotate more slowly during epochs of high activity than their average rotational speed over the entire observational interval [22].

Our research shows that solar activity differs from the activity of stars with well-pronounced Excellent cycles, primarily K stars. This could be related to either age differences or differences in the role of the large-scale magnetic fields. Returning to Fig. 2, we note that half the stars with Good cycles are adjacent to stars with well-defined cycles and the Sun, while others are characterized by higher chromospheric and coronal activities.

Solar investigations allow us to study in details the role of magnetic fields of different scales in formation of activity. So, observations of the large-scale solar magnetic field (synoptic maps) and measurements of the magnetic field of the Sun as a star (the total magnetic field) are used to determine the dipole magnetic moment and direction of the dipole field for three successive solar cycles [23]. Both the magnetic moment and its vertical and horizontal components vary regularly during the cycle, but never disappear completely. A wavelet analysis of the total magnetic field shows that the amplitude of the 27-day variations of this field is very closely related to the magnetic moment of the horizontal dipole.

The interval of slower rotation of the irregularities is close to the epoch when the Sun's field represents a horizontal magnetic dipole in each activity cycle, but is somewhat longer than the duration of the polarity reversal in both hemispheres [22]. The difference between the periods for the slower and mean rotation exceeds 3 days, as is typical for some stars with higher but more irregular activity than the Sun.

The largest-scale magnetic fields also affect the shape and development of coronal streamers. In the course of the cycle, the neutral line of the longitudinal magnetic field locates near the equator at the minimum and moves away up to 30–40° during other epochs. So, one can distinguish hills of an unipolar field or giant cells of diameters of 0.3–0.5 solar radius. Typically, active regions locate presumably close to the neutral line and the chromosphere above it is the brighter than averaged over the disk. Many coronal loops occur also above the neutral line.

In a certain sense we can speak about effect of magnetic fields of intermediate scales on activity.

As for the solar cycle, changes with a period of 10–11 years are expressed most clearly. However, there are some indications of other changes with shorter and longer periods. Apart from the 22-year magnetic cycle, we can note changes in various indices with a quasi-biennial period, as well as the secular 80–100-year cycles. The simultaneous existence of cycles of different durations on the Sun and some active late-type stars in the Good cycle group suggests that the solar cycle has not entered a strict asymptotic regime (see, for instance, [17] and references therein).

Thus, solar activity is due to a complex interaction of phenomena associated with the evolution of magnetic fields of different scales. Of course, this has only begun to be studied for stars. It may be that it is possible to observe effects associated with large-scale and fields on the Sun because the depth of the convective zone is small, and processes in the tachocline (the transition from the radiative to the convective zone) determine the appearance of the surface magnetic field of large scales. This field governs the development of solar activity. There is evidence that local fields can arise and be amplified directly beneath the photosphere. Thus, there are apparently two levels on the Sun where the dynamo process is realized. This new concept of a two-level dynamo may be useful for our understanding of solar-type activity on other stars.

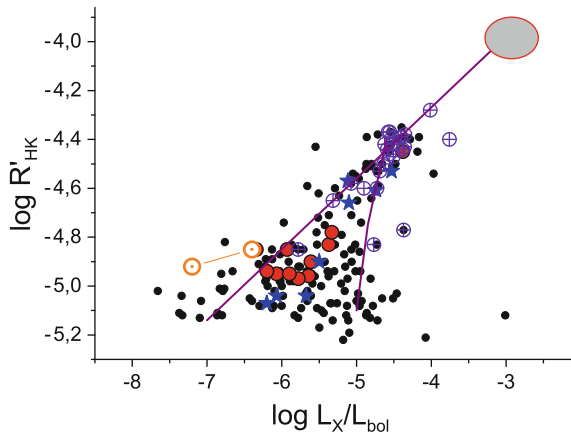
5 Possible Evolutionary Paths for Solar Activity

Thus far, we have considered stars with activity levels comparable to the solar level. The activity of young stars is substantially higher, and reaches even a saturation level for some stars. In soft X-rays this saturation corresponds to values of $\log(L_X / L_{bol})$ close to -3 . There is also the flux saturation in other energy ranges (e.g., the chromospheric emission). It is interesting to consider how the activity of young, low-mass stars evolves after the deceleration of their rotation (after ages of 70–100 Myr).

Let us consider the diagram in Fig. 2, adding the region of stars with saturated chromospheric and coronal activity, shown by the oval in Fig. 4. The linear relationship according to the formula given in Sect. 4 is also shown. This can be regarded as one evolutionary path for the activity as the rotation slows, that is characterized by simultaneous decrease in the chromospheric and coronal emission.

Further we add a set of late-type stars with registered Li I 6707 Å line to above considered stars of solar vicinity, including the HK Project stars, as well as to objects, whose activity had been studied in the course of the exoplanet search programs [7]. Based on observations carried out on the Haute Provence observatory [24] we chose G and K stars with measured lithium abundances, for which the index of the chromospheric activity has been determined earlier.

Fig. 4 Schematic illustrating two paths for the development of solar-like activity. Notation is the same as in Fig. 2. The investigated stars with measured lithium abundance and with direct measurements of $\log R'_{HK}$ indices are marked as *crosses* inside the *circles*. The area where the chromospheric and coronal fluxes are saturated is shown by the *grey oval*. The *straight line* corresponds to the expression (1)



In Fig. 4 it can be seen that the stars with heavy lithium are characterized by high activity level. Those newly added stars fill up the area, corresponding to the transition from the stars with the solar-like activity to the objects with saturated activity levels. The stars with saturation of the activity level were studied in detail by [25].

In Fig. 4 several stars with detected Li also deviate downwards the linear dependence. Such a trend can mean that there is another way for the evolution of solar-type activity beside a main path which is a basement of one-parameter gyrochronology (the straight line accordingly to the expression (1)). Namely, starting from a definite activity level of many late G and K stars, the chromospheres become weaken while coronae stay still powerful. Both paths shown conditionally in Fig. 4 can be considered as envelopes for all possible ways for the evolution of solar-like activity depending on masses and individual properties of stars. Note that stars with the Excellent cycles also deviate from the main linear relationship.

What is a cause of differences in activity of both groups of stars? The first group of stars adjacent to the straight line in Fig. 4 is characterized by a significant role of large-scale magnetic fields in addition to local magnetic fields in the formation of activity. Evidence for this includes various observational effects, such as the existence of active longitudes, analogs to the Maunder butterfly diagram, and the appearance a short cycle along with the main one. Solar-like activity appears in stars starting from spectral type F5, when the convective zone has a particular small depth of about 0.05 stellar radii. For stars of spectral type G4, the depth the convective zone begins to exceed 0.35 stellar radii. For stars with this range of parameters, the dynamo process probably develops in both the tachocline and sub-photospheric layers.

It is often supposed that large-scale magnetic fields occur near the lower bottom of the convective zone, whereas local fields are formed at relatively shallow depths beneath the photosphere. When the convective zone becomes sufficiently thick

(over 0.35 stellar radii), processes close to the bottom of this zone cease to affect the formation of solar-type activity at the surface.

For later-type stars, up to fully convective red dwarfs, the evolution of their activity is determined by processes occurring directly beneath the photosphere, where the dynamo mechanism generates local magnetic fields. For stars with spectral types later than G7 (color indices $B - V > 0.67$) and the thicker convective zone, effects of magnetic fields generated near its bottom are inconsiderable.

6 Discussion

The most extensive information on cyclic activity available for the closest to us star—the Sun. Our investigation of active late-type stars allows us to separate a new group of stars on the “chromosphere–corona” diagram that is an evidence for existence of two possible paths of the evolution of solar-like activity. How does this help to understand features of the cyclic activity of the Sun and to imagine what was activity of the young Sun and what happens with the Sun in the future?

The appearance of spots, plagues etc. on the solar surface is associated with local magnetic fields, but activity of the Sun as a whole govern by the large-scale magnetic field. This is inherent in the main group of stars where large-scale magnetic fields are observed directly, active longitudes exist, and activity expands to all atmospheric heights. The activity level of these stars is higher but is less regular than that of the Sun. Of course, the Sun is quite old star; its chromospheric activity is higher than that of stars of the similar age, while the solar corona is substantially weaker.

The conception about possible paths for the evolution of solar-type activity fits into those ideas that can explain features of effect of large-scale magnetic fields on activity and its relation with the depth of the convective zone. The large-scale magnetic field generated near the lower boundary of the convective zone governs activity both on the Sun and the most of stars of the main group. If these ideas for the Sun are formulated during last years thanks to success of the local helioseismology, so for stars we point out the first reasons for application of two-level dynamo.

What conclusions can we do from the proposed suggestion for the explanation of the physics of the cyclic activity of the Sun on various stages of its evolution? The basic features of activity of the Sun and the main group of stars, whose chromospheric and coronal radiation change simultaneously at large time scales, are similar in many respects. This relates to observable signs of large-scale magnetic fields, i.e. active longitudes, some instability of the cycle like different amplitudes and simultaneous presence of several periods of long-term variability. Despite the old age, the Sun does not yet reach that level of chromospheric activity that is characteristic for the most of late-type stars (Fig. 1). This can be associated with peculiar properties of the angular momentum loss after the age of 1–2 Gyr (in Fig. 4 this relates to the region of separation on two branches). A substantial argument for affiliation of the Sun to this group is general properties of their differential rotation [19, 22].

Fundamental difference between the Sun and the stars, whose activity evolves by another path, is distinction in filling factors, i.e. relative areas covered by active regions which are observed the best in the soft X-rays as well as in the area of spots. If the soft X-ray emission of the Sun out of flares variations changes at different time scales, so the later the star is then the corresponding variations are weaker. This reflects growth of number of weak non-stationary events. In Fig. 4 it is seen that coronae of stars located between two lines on the diagram, in particular, of stars with the Excellent cycles are significantly more powerful than the solar corona.

The regular cycle is not typical for the stars of the main group including the Sun, but this is characteristic for old low-rotating K stars. This can mean that access on the asymptotic regime of the dynamo with regular cycles occurs easier in the case when activity is governed by local magnetic fields.

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Long-Term Variations of the Solar Supergranulation Size According to the Observations in CaIIK Line

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Abstract This work contains analysis of distinctive size of chromospheric cells Kodaikanal (1907–1999) and Medon (1983–2010). At first the contrast of chromospheric grid was enlarged on image, by means of subtraction of the gradient from the solar disc intensity. This analysis was performed with the help of balanced wavelet transformation. It was discovered that distinctive size of chromospheric cells is close to 36 Mm but it has variations, in the phase of solar activity $\sim 1, 2$ Mm maximal size of the cell can be seen as a rule, in $\sim 1, 5$ year after maximum of the solar activity. There is a positive correlation ($R = 0, 83$) between the size of the chromospheric cells in maximum and the amplitude of the following activity cycle. Thus, the size of the supergranulation is connected with the solar activity and is ahead of it $\sim 8, 8$ of a year.

1 Introduction

Convective cell of different scale are registered on the surface of the sun. They correspond to granulation, meso-granulation, supergranulation and to gigantic cell. It is assumed that the first three cell types appear as a result of atomic hydrogen and helium ionization, and these cell have scale, comparable with the depths where ionization processes take place. There are descending substance streams with the speed of 1–2 km/s on boundaries of the supergranulas, these streams are concentrated in separate points, especially in places where several supergranulas are joined. It is considered that the on boundaries of supergranulas basic part of photospheric magnetic flux is concentrated [1].

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Granulation is clearly seen in the white light, has sizes about 1 Mm in diameter and average continuity for about an hour. Supergranulations, as a rule, have about 30–40 Mm in diameter and exist ~ 1 –2 days. Supergranulation was discovered by [2] in speed fluctuation. Doppler's visualization method for supergranulas was elaborated by [3, 17] and even nowadays remains a reliable means of supergranulas' observation. The presence of other cell is also assumed—they are meso-granulas (~ 7 Mm in diameter) and gigantic convective cell (~ 100 Mm in diameter) [4, 5], but there were no convective proofs after. The authors of the work [6] come to the conclusion, that meso-granulas are absent as a separate distinctive scale of the solar convection. It's quite possible that meso-granulas exist as a combination of the smaller supergranulas and bigger ones, unlike different components in convection itself [7].

Using the method of auto-correlation of dopplerogram [3] the authors got the values of 32 Mm size. Approximate size of the supergranulas cell was received during the analysis of filtergrams in line CaIIK in the work [8]. Using SOHO/MDI data of year 1996 it was discovered [6] that maximal distribution of the length of supergranulas was 36 Mm. In the work [9, 16] maximal distribution of the scale of supergranulas was also found—36 Mm in 2007.

Using SOHO/MDI dopplerograms in minimal 23/24 activity cycles [7] we discovered that in minimum of the 23rd cycle the size of supergranulas was $35.9 \pm 0,3$ Mm, and in the minimum of the 24th cycle— $35.0 \pm 0,3$ Mm. The changes of order 0,5 Mm in minimums of 23/24 cycles were also found in the work [10]. In [11], using two sets of data received in different levels of magnet activity, found decreasing of the size of supergranulation during the period of high activity. Analogous conclusion was made in the work of [12]. In [13] revealed lowering of the distinct cell size in the process of increasing field intensity in supergranulas, but pointed out that larger segments of supergranulas were linked with strong field grids on their boundaries. Therefore, it appears that negative or positive correlations can be received depending on the level of magnetic activity, it is defined in relation to interconnected fields or grids. In [13] also reported about absence of large supergranulation segments for supergranulas with intensive interconnected magnetic field, and indicated that interconnected magnetic fields have dynamic influence concerning supergranulas.

2 Method of Analysis

To analyze the supergranulas size we carried out preliminary processing of images, which included several stages. On the first stage the level of the quiet sun was defined, as a function of the distance from the disc's centre to position angle $I_Q = I_Q(r, \alpha)$ [14]. Then we corrected the darkening to the edge. Then the flocculi boundaries were specified [14]. For bright flocculi with their space more than $S > 1,000 \mu\text{hm}$, we replaced intensity value of every pixel with the values of background I_Q . Thus, we had the solar disc where the darkening to the edge was removed and flocculi spots were excluded.

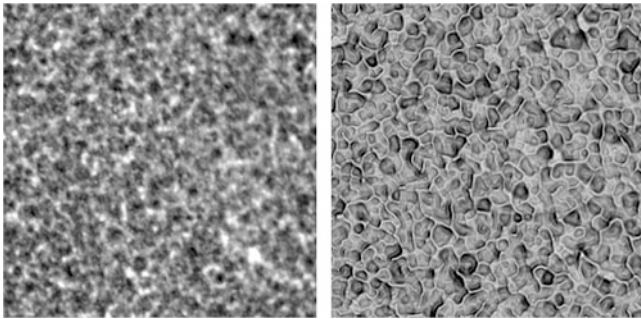


Fig. 1 The example of contrasting for singling out supergranulation cell for central area of observation 20.02.1997 according to Meudon

The second stage included contrasting of chromospheric grid, applying elements of the algorithm Canny for searching the boundaries [15]. The first step, as well as original method, contained the procedure of smoothing with the Gaussian filter 5×5 pixels, then we calculated the gradient by Sobel method and in the end we subtracted the gradient from the intensity. This procedure took place twice. The example of the original image and the result after processing is reflected in Fig. 1.

Then we calculated distinctive size of the chromospheric grid. To do this, we scanned images at a pitch of $0,2^\circ$ along the central meridian within the range of $\pm 10^\circ$ latitudes. For analysis we used the method of weighted wavelet Z-transform (WWZ), specially developed for irregular selected data [18]. We experimented with different values of parameter c , which defines the compromise between the spatial resolution and resolution frequency and chose $c = 0.005$. Acceleration spectral density of each day was summarized for a definite period. The the maximal distribution of acceleration spectral density was calculated.

As a test we conducted the processing of SDO doplerograms since June 2010 up to March 2011. Figure 2 reflects relative spectral power for doplerograms and data of observations in line CaIIK Kodaikanal in 1907–1999 and Meudon Observatory in 1983–2010. Maximal distribution is close to $2,98^\circ$ or $36,2 \text{ Mm}$, which is close to the values of distinctive supergranulation size, discovered by several authors [6, 7, 9]. Apart from the maximum $\sim 3^\circ$ in line CaIIK the local maximum $\sim 1,18^\circ = 14 \text{ Mm}$ is observed, which is close to the doubled size of meso-granulas cell.

3 Results of the Analysis

For searching the supergranulas cell we used sections along the central meridian on daily solar images. For each section acceleration spectral density was found. Daily data were summarized in 4 months intervals, then the maximal distribution of acceleration spectral density for the given altitude and time was calculated. Then this value was averaged in necessary range of altitudes.

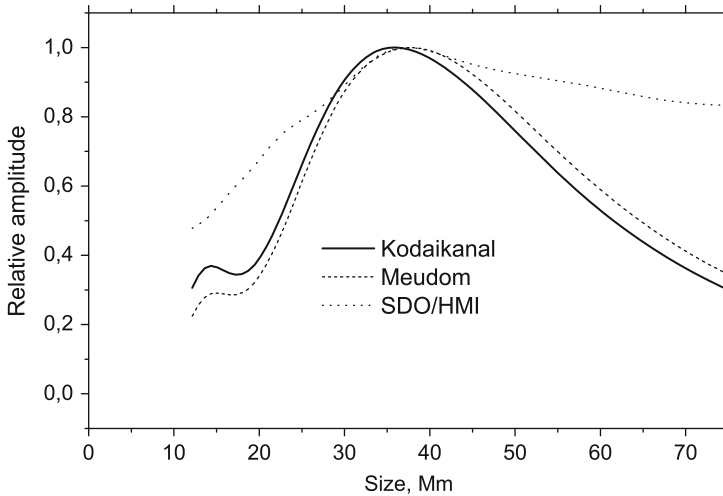


Fig. 2 Distribution of the spectrum power of the supergranulation size according to SDO doplerograms, and also observations in line CaIIK according to Meudon Kodaikanal. Maximal distribution is close to 36 m

Figure 2 demonstrates the values of maximal distributions for the data received at Kodaikanal Meudon. The average size of supergranulas cell, according to Kodaikanal in 1907–1999 was equal to $35,47(\pm 0,07)$ Mm, in accordance with Meudon in 1983–2010— $35,83(\pm 0,05)$ Mm. In spite of some differences, we can notice cyclical pattern of 11-year variations. Maximal supergranulation size was noticed in 1950–1953. Maximal size of the cell was attained during the phase of activity decay of the sun spots, in average, after ~ 1.8 year after minimal activity.

The size of supergranulas cell along the solar equator is a bit higher than along central meridian and amounts to $36,68(\pm 0,09)$ Mm, according to Kodaikanal.

4 Discussion

Cyclical motion of the supergranulation sizes can be received by methods of processing and the influence of floccules spots concerning small-scale harmonic components. But in this work we tried to minimize this influence by removing floccules and choosing wave-let transformation as a method of analysis. Also, there is a nonlinear connection between the solar activity and the size of the supergranulation cell. Figure 2 shows that during the maximal 19-th activity cycle local maximum of the supergranulation cell was relatively low, while larger size was observed during the 18-th activity cycle. There is a positive correlation $R = 0.83$ between supergranulation size and activity amplitude of the following solar spots cycle. Along with that supergranulation cycle passes ahead of the sun spots ~ 8.8 years (Figs. 3 and 4).

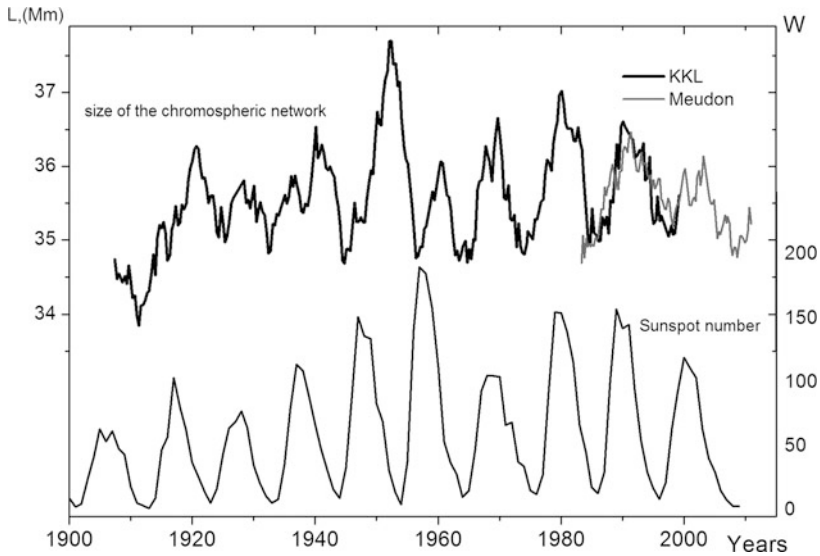


Fig. 3 The changes in the size chromospheric network smoothed with running window of 3 years. Presented data observatories Kodaikanal and Meudon. The Wolf sunspot number W is shown for comparison

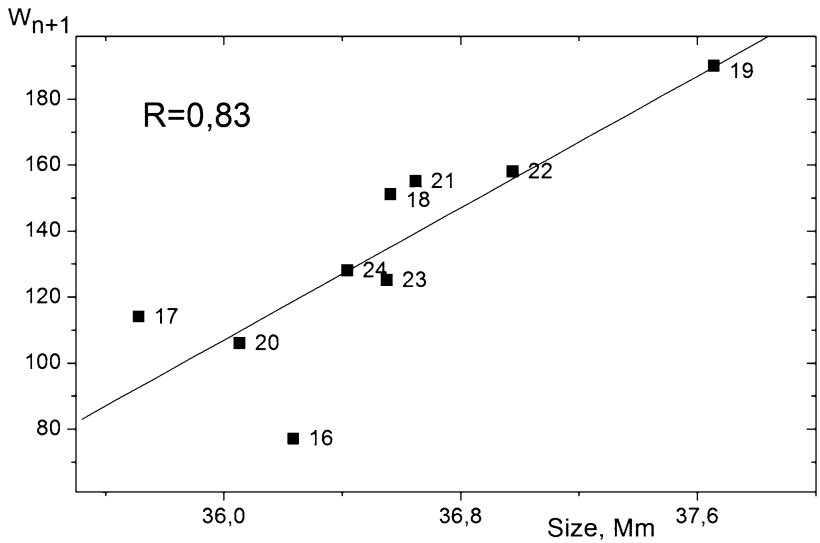


Fig. 4 Correlation between supergranulation size and activity amplitude of the following solar spots cycle according to Kodaikanal. Correlation coefficient is $R = 0.83$

Discovered correlation between the supergranulation size and activity amplitude of the following solar spots cycle can be linked with reconstruction of convection of the upper solar atmosphere, the consequences of which are present in the following activity cycle.

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On the Problem of Heat Transport in the Solar Atmosphere

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Abstract In the context of the problem of energy transport in the solar atmosphere, we present new results on physical properties of a transition region between the hot and cold plasma in quiet regions and in flares. In quiet regions, the transition region between the corona and chromosphere is shown to be a very thin layer, in which however the classical collisional approach is valid very well. A stability of the transition region is investigated. It is shown to be stable, moreover it is a stable consequence of the thermal instability in the condensation mode regime. We have developed also mathematical models for describing plasma heating in the corona and chromosphere by heat fluxes from a super-hot reconnecting current layer. It is shown that applicability conditions of classical heat conduction are not valid in the flare transition region and at the higher temperatures. Models, that account the effect of collisional relaxation, describe heat transport in flares better than Fourier's law. A possibility of comparing theoretical and observational data is discussed.

1 Introduction

In this paper we study the heat transport in quiet regions and solar flares. In quiet regions we investigate the chromosphere-corona transition region. We solve a stationary problem in which plasma heating by classical heat flux is balanced by radiative losses. In solar flares we consider plasma heated by a high-temperature turbulent-current reconnecting layer. In this case we solve a non-stationary problem in which internal thermal energy changes are balanced by heat conduction.

The aim of the work is determining mechanisms of heat transport in these two cases.

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2 Quiet Regions

In quiet regions we consider a magnetic tube which goes from the corona to the chromosphere. On the corona side the temperature is about 10^6 K. On the chromosphere side we assume that in the absence of heat flux from the corona to the chromosphere some other mechanism of chromosphere heating supports its time independent temperature.

We assume that in the described situation one can use the classical collisional approach. So we solve the balance equation where heat flux is due to Coulomb collisions between electrons:

$$\frac{d}{dx} \left(\kappa \frac{dT}{dx} \right) = L(T) n^2 - P_\infty. \quad (1)$$

Here κ is the coefficient of thermal conductivity. Neglecting the thermal conductivity of neutrals, we have the classical formula for the electron conductivity $\kappa = \kappa_0 T^{5/2}$ [1, 2]. The radiation power per unit volume of plasma $L(T) n^2 = P(T, n)$ where the function $L(T)$ is the radiative loss function, it describes the dependence of energy loss due to radiation on the temperature. $P_\infty = L(T_\infty) n^2$ is the power of the stationary heating of the chromosphere by an “external” source, partially, by the flux of waves from the convection zone. The dependence $L = L(T)$ is taken from the results of calculations [3] performed at the Potsdam Astrophysical Institute using the system of atomic data and CHIANTI programs (version 5.2). To define the dependence $n = n(T)$ without solving the full set of hydrodynamic equations, we consider two opposite limiting cases of pulse ($n = \text{const}$) and stationary ($p = \text{const}$) heating by the heat flow separately.

The calculated temperature distributions [4] along the depth $\xi = \int_0^x n(x) dx$ is shown in Fig. 1. As seen from Fig. 1, plasma is divided into high and low temperature parts, and this result does not depend on the heating regime. The stability of the obtained temperature distributions was checked in [5]. The temperature distributions were shown to be stable and, for the temperatures greater than 10^5 K, they were defined by the condensation mode of the thermal instability.

Our solutions allow us to compute the differential emission measure $DEM(T)$ of the plasma in the transition region. The theoretical dependence $DEM(T)$ was compared with the observations by SUMER/SOHO [6]. The result of this comparison is shown in Fig. 2. One can see that: (a) the observed points are located between two theoretical limits of fast and slow heating, (b) the observations are closer to the curve corresponding to the slow ($p = \text{const}$) heating.

So, the heating regime is rather close to the slow one in the transition region of the quiet sun. And what is the most important, the classical collisional approach is valid, because results obtained in this approach agree with observational data.

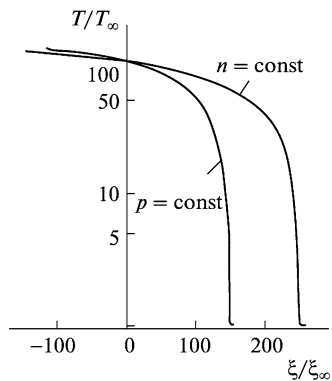


Fig. 1 The dependence $T = T(\xi)$ for the cases of fast ($n = \text{const}$) and slow ($p = \text{const}$) heating. The depth unit $\xi_{\infty} = 3.15 \times 10^{15} \text{ cm}^{-2}$

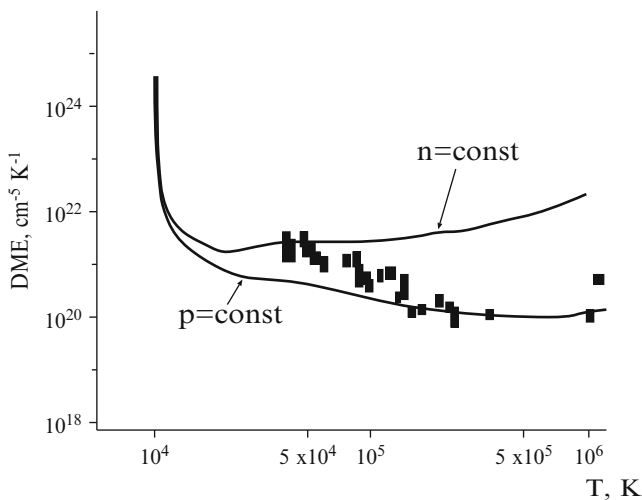


Fig. 2 The theoretical DEM(T) dependence for the cases of fast ($n = \text{const}$) and slow ($p = \text{const}$) heating. The points correspond to observations which were taken by the SUMER instrument on board of SOHO on 20 April 1997. From [6]

3 Solar Flares

In solar flares, plasma is heated by a high-temperature reconnecting current layer. Figure 3 shows magnetic field tubes that interact with the layer in the corona.

First, the tubes accompanied by coronal plasma move with relatively small velocities towards the layer, come into contact with it and penetrate into the layer. Then they reconnect in the center of the layer and move towards its edges, gradually accelerating to the high velocity. The plasma inside magnetic tubes heats

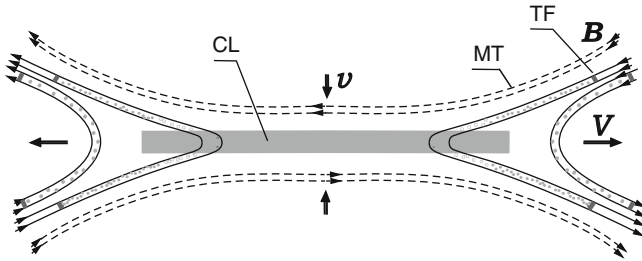


Fig. 3 Magnetic field tubes (MT) and a reconnecting current layer (CL). TF is a thermal front moving along the tube

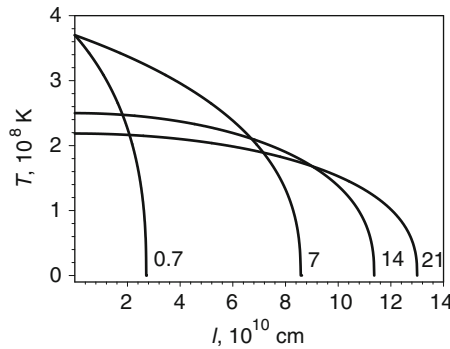


Fig. 4 Temperature distributions along the magnetic tube in the case of the classical heat conduction. The numbers near the curves denote the time (seconds) passed after the start of contact of the tube with the current layer

up to the temperature of the current layer. According to the self-consistent 2D model of the high-temperature turbulent current layer [7], its temperature is about 3×10^8 K. The characteristic time of tube contact with the layer is estimated about 7 s. After the tubes escape from the layer, they continue to travel in the corona. They are disconnected now from the heating source. The thermal energy redistribution continues in the form of a large-scale thermal wave passing along the magnetic tube.

We describe heat transfer in coronal plasma along the magnetic tube in terms of the simple heat conduction equation

$$\frac{\partial \varepsilon}{\partial t} = -\operatorname{div} \mathbf{F}. \quad (2)$$

Here ε is the internal energy of unit plasma volume, $\mathbf{F}$ is the vector of thermal energy flux density. The mathematical formulation of the problem is described in [8].

First, let us consider the classical heat conduction in the form of Fourier's law $\mathbf{F} = -\kappa_0 T^{5/2} \nabla T$. The problem has self-similar solutions [9]. We have applied this method to the conditions of solar flares in the vicinity of the current layer and have got the temperature distributions, presented in Fig. 4.

Note that the thermal wave front moves too fast. In the first second its mean speed is formally greater than the speed of light. Moreover the thermal wave propagates for a distance of 10^{11} cm in the first 20 s, which is comparable with diameter of the Sun. This considerably exceeds the size of solar active regions, $\sim 10^{10}$ cm. Note also that the classical heat flux in this thermal wave is greater than the maximum thermal flux at which the diffusion heat transfer is replaced by the convective transfer, i.e. the saturated flux [8].

These contradictions arise because the classical Fourier law is derived for the plasma that is very close to local thermodynamic equilibrium. This implies that the characteristic time scale for the process greatly exceeds the time scale on which electron collisions occur, and the characteristic length scale is far greater the thermal electron free path [10]. These conditions are not valid in the vicinity of the current layer. We deal with too fast process and with too steep wave front.

Another approach for describing the thermal energy transfer is based on 13 Grad's moment equations [11]. In contrast of the classical approach, this method does not imply that the distribution function of electrons is very close to the Maxwell one. The equation for heat flux is

$$\mathbf{F} = -\kappa \nabla T - \tau \frac{\partial \mathbf{F}}{\partial t}. \quad (3)$$

Here the first term describes the classical heat flux, and the second one is due to the relaxation. τ is certain characteristic time needed for maxwellization of distribution function. Its value is about 13 s (or even more) in solar flares [8]. It is comparable to the time of contact of the magnetic tube with the current layer (7 s). Therefore, taking this effect into account is absolutely necessary in solar flares.

Figure 5 presents the temperature distributions along a magnetic flux tube calculated with account for relaxation. It is seen that wave front becomes substantially steeper as compared to that in the classical case. Once the tube disconnected from the current layer, the wave shape significantly changed. The solution represents the wave with a steep forward front and gradual decay behind it. The speed of the front is almost time independent in the period under consideration and is about 2×10^9 cm/s. This is significantly less than the average speed in the classical case for the initial period. So, the collisional relaxation resolves the contradictions of the classical approach.

Our solution allows us to compute the differential emission measure $DEM(T)$ of the heated plasma surrounding the current layer. The method is similar to that described in [12]. Figure 6 presents $DEM(T)$ computed for different moments, and Fig. 7—integral emission measure $EM(t)$ as a function of time.

Using these results we can estimate the spectral line intensities of some high-charged ions such as Fe XXVI and Ni XXVII (Fig. 8). Their radiation power has a maximum at the temperatures $\gtrsim 10^8$ K. So this emission can be considered as an evidence of super-hot plasma in flares.

Fig. 5 Temperature distributions along the magnetic tube in the case of the relaxation heat conduction. The numbers near the curves denote the time (seconds) passed after the start of contact of the tube with the current layer. The scale is the same as in Fig. 4 for convenience

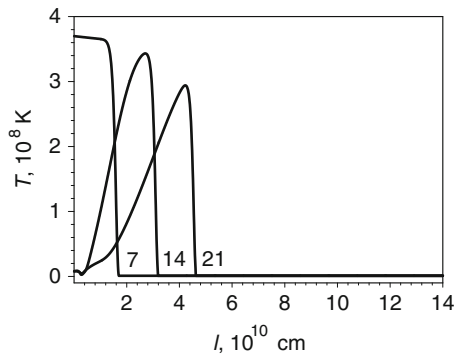


Fig. 6 Differential emission measure of the plasma heated by the current layer. Numbers near the curves denote time (seconds) passed after the start of the flare

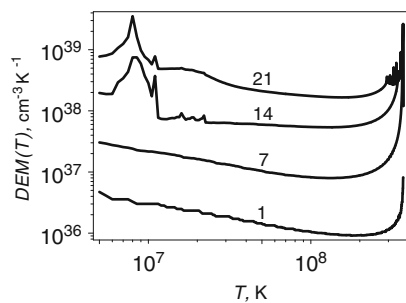


Fig. 7 Integral emission measure of the plasma as a function of time

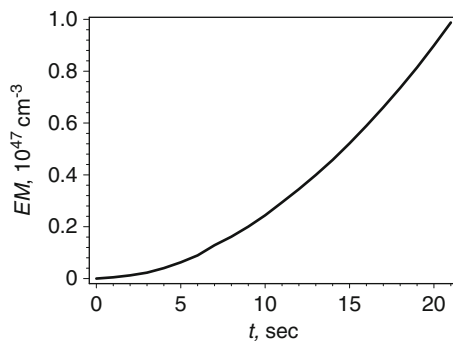
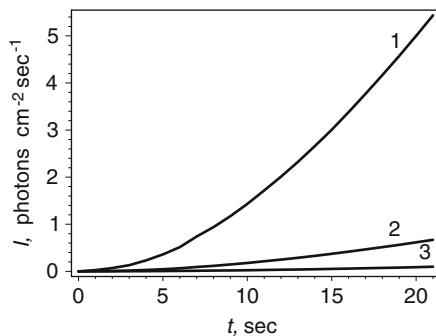


Fig. 8 Estimated intensity of spectral lines: 1, Fe XXVI ($\lambda = 1.78 \text{ \AA}$); 2, Fe XXVI ($\lambda = 1.51 \text{ \AA}$); 3, Ni XXVII ($\lambda = 1.59 \text{ \AA}$)



Our EM estimations are consistent by an order of magnitude with RHESSI observations in the beginning of a flare [13]. Intensity of spectral line Fe XXVI ($1.78 \text{ \AA}$) agrees with Yohkoh data [14].

4 Conclusion

We have demonstrated that the transition region between the corona and chromosphere in quiet regions is a very thin layer in which however the classical collisional approach is valid. The transition region is a stable consequence of the thermal instability in the condensation mode regime. The emission measure obtained from our temperature distribution shows a good agreement with observations.

In solar flares, the applicability conditions of the classical collisional approach are violated. Model, accounting the effect of collisional relaxation, describes heat transport in flares better than Fourier's law. Our estimations of emission measure and Fe XXVI ($1.78 \text{ \AA}$) spectral line intensity are in agreement with observations by RHESSI and Yohkoh. For more detailed comparison we need observations with high spectral resolution ($\sim 0.01 \text{ \AA}$ at $\lambda \sim 1 - 2 \text{ \AA}$) and simultaneously with high spatial one ($\lesssim 1''$). The detectors must be very sensitive because the emission measure of plasma in the vicinity of the current layer is small ($\lesssim 10^{47} \text{ cm}^{-3}$).

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Dynamics of the Electrical Currents in Coronal Magnetic Loops

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1 Introduction

We examined records of the radiation of three flare processes obtained in the Nobeyama observatory in 1992, 2001, and 2003. Observations with the Nobeyama radio heliograph at the frequencies 17 and 34 GHz were performed with a sufficient spatial resolution to determine the physical nature of the radiation source. In the first two cases, the radiation source appeared to be a coronal magnetic loop; in the third case, a group of sunspots. In the spectra of the first two events a characteristic modulation was detected, with the frequency that increased by a factor of 2–3 for a fairly short period immediately before the beginning of the flare and later decreased gradually to the initial value. In the spectrum of the event on June 17, 2003 (the radiation source were the sunspots), no modulation of such a type was seen. Here, we try to explain the origin of the modulation with the initial frequency of about $\nu \approx 0.005$ Hz at the pre-flare stage, and also to explain the reason for a strong increase in the modulation frequency directly before the flare.

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2 Observational Data

We studied low-frequency modulation of the microwave radiation of CML in the active regions AR9393 (March 30, 2001), NOAA7321 (November 02, 1992), and AR10386 (June 17, 2003). In all three cases, we analyzed the temporal profiles of microwave bursts recorded by Nobeyama radio polarimeters [1] at the frequencies 35, 17, 9.4, 3.75, 2, and 1 GHz. The spatial characteristics of the events were explored with the use of the Nobeyama radio heliograph at the frequencies 17 and 34 GHz [2]. In the first two cases, where the source was a coronal magnetic loop, the typical modulation of the radiation was detected. The algorithm that we used for the analysis of the low-frequency modulation is described in detail in [3].

In the first case (the active region AR9393), the observations covered the period 05:02–05:44 UT. The maximum intensity of the microwave radiation was observed at 05:14 UT. The source of the microwave radiation was a CML. The same active region was observed with the radio heliograph of the Siberian Solar Radio Telescope (SSRT) at frequencies near 5.6 GHz [4]. In the time interval 05:11–05:20 UT, a flare with the power M2.2 occurred in the mentioned region. The data of NoRH and SSRT radio heliographs indicate that the flare and the accompanying microwave radiation occurred in the coronal magnetic loop. Figure 1 presents an image of the loop in the emission measure of soft X-ray radiation, on which contours of the image of the microwave radiation source at 17 GHz were superimposed. It is seen that the intensity maxima of the microwave radiation coincide with the footpoints of the CML. The data on the emission measure in the soft X-ray radiation, as well as the SOHO/MDI data, make it possible to determine the parameters of the loop [4] for the event on March 30, 2001: the distance between the footpoints $d \approx 3.6 \times 10^9$ cm, height $h \approx 1.8 \times 10^9$ cm, thickness $w \approx 5 \times 10^8$ cm, length $l \approx 5.6 \times 10^9$ cm, plasma density $n \approx 1.4 \times 10^{11} \text{ cm}^{-3}$, temperature $T \approx 10^7$ K, and magnetic field at the loop footpoints $B_f \approx 400 - 600$ G.

The shape of the temporal profiles of the microwave radiation intensity at the frequencies 3.75, 9.4, and 35 GHz are consistent with the intensity profile at the frequency 17 GHz, which may indicate the general localization of radio emission sources at all these frequencies in the coronal magnetic loop (hereafter the curves of the Nobeyama spectrum polarimeter show the calibrated values in the solar flux units (s. f. u.)).

Figure 2a presents a part of the temporal intensity profile at the frequency 17 GHz, while Fig. 2b the dynamic spectrum of low-frequency modulation of the radiation obtained using the Wigner–Ville method. The Wigner–Ville transformation shows that before the flare energy was released, the loop radiation at the frequency 17 GHz had been modulated by a signal with the frequency $\nu \approx 0.005$ Hz (the period 200 s) and the relative bandwidth $\Delta\nu/\nu \approx 0.5$. About 50 s before the maximum of the microwave radiation was reached (which can be identified with the maximum of the flare energy release), the modulation frequency had sharply increased to $\nu \approx 0.035$ Hz (the period 28 s) and later went back to the initial value $\nu \approx 0.005$ Hz for a time of about 100 s, which coincides with the duration of the

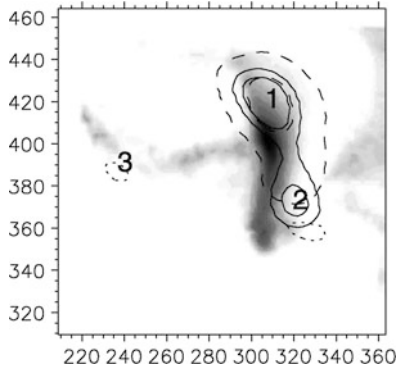


Fig. 1 Intensity of the soft X-ray radiation of the flare loop on March 30, 2001, based on the Yohkoh/SXT observations (gray background). The curves show the image obtained at 17 GHz using the Nobeyama radio heliograph. Solid curves correspond to the intensity levels of 0.25 and 0.5 of the maximum. Dotted-and-dashed curves indicate the polarization (at the levels -0.2 , -0.5 , 0.2 , and 0.5 of the maximum, respectively). The digits 1, 2, and 3 denote different radiation sources. Along the axes is the distance from the solar disk center in arcseconds

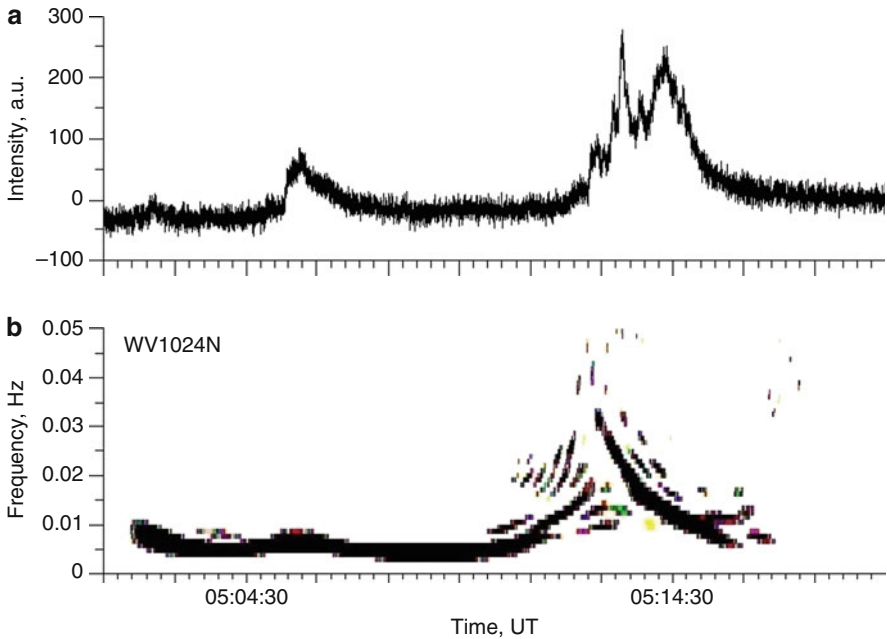
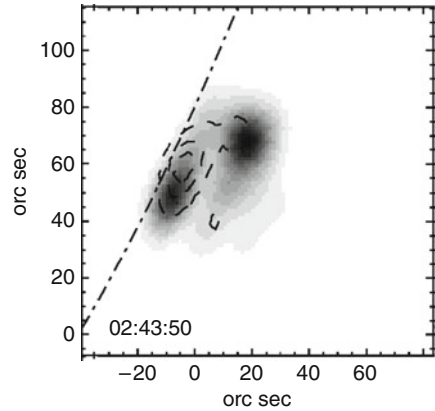


Fig. 2 The event on March 30, 2001: the microwave radiation intensity based on the Nobeyama spectrum polarimeter observations at the frequency 17 GHz (a); the spectrum of low-frequency modulation of the microwave radiation obtained using the Wigner–Ville transformation (b)

Fig. 3 The structure of the microwave source at 02:43:50 UT on November 2, 1992. Observations were made with the Nobeyama radio heliograph at 17 GHz. The intensity is shown by shades of gray, and the polarization by dotted lines. The distances on the axes are shown in arcseconds



explosive phase of the microwave burst at 17 GHz. After the end of the flare process, the modulation with the frequency $\nu \approx 0.005$ Hz almost disappeared at 17 GHz, but retained at the frequencies 2 and 3.75 GHz.

The second event that we analyzed is a limb solar flare observed on November 2, 1992 by the Nobeyama radio polarimeters at the same frequencies as those on March 30, 2001 in the active region NOAA7321. The total duration of the X9 class event was about 4 h (2:00–06:00 UT). We analyzed in detail two flare processes in the above interval (02:43–03:05 UT and 04:00–04:20 UT). The event on November 2, 1992 was also observed with the SSRT radio heliograph at the frequencies near 5.6 GHz [5]. According to the radio heliographic data, the radiation source was a CML. Using the available data, one can infer that the radiation at the time of the first flare (the radiation maximum at 02:49 UT) was produced by two coronal magnetic loops with roughly the same parameters [5]: the distance between the footpoints $d \approx 3.65 \times 10^9$ cm, height $h \approx 1.8 \times 10^9$ cm, thickness $w \approx 7.3 \times 10^8$ cm, length $l \approx 5.73 \times 10^9$ cm, plasma density $n \approx 2 \times 10^{10} \text{ cm}^{-3}$, temperature $T \approx 1.4 \times 10^7$ K, and magnetic field at the loop footpoints $B_f \approx 300$ G.

In the second flare with the radiation maximum at 04:09 UT the radiation source was a magnetic loop with the following parameters: the distance between the footpoints $d \approx 4 \times 10^9$ cm, height $h \approx 2 \times 10^9$ cm, thickness $w \approx 10^9$ cm, length $l \approx 6.3 \times 10^9$ cm, plasma density $n \approx 2 \times 10^{10} \text{ cm}^{-3}$, temperature $T \approx 1.4 \times 10^7$ K, and magnetic field at the loop footpoints $B_f \approx 300$ G.

The image of the radiating region corresponding to the instant 02:45:20 UT and obtained at the frequency 17 GHz is presented in Fig. 3 [5]. Figure 4 shows an image of the loop which was the radiation source during the flare at 04:09 UT [5].

Figure 5a shows the intensity of the microwave radiation at 17 GHz, Fig. 5b presents the corresponding spectrum of low-frequency modulation obtained using the Wigner–Ville transformation. A sharp increase in frequency ν from the initial value 0.01–0.025 Hz, i.e. by about a factor of 2.5, was observed directly before the

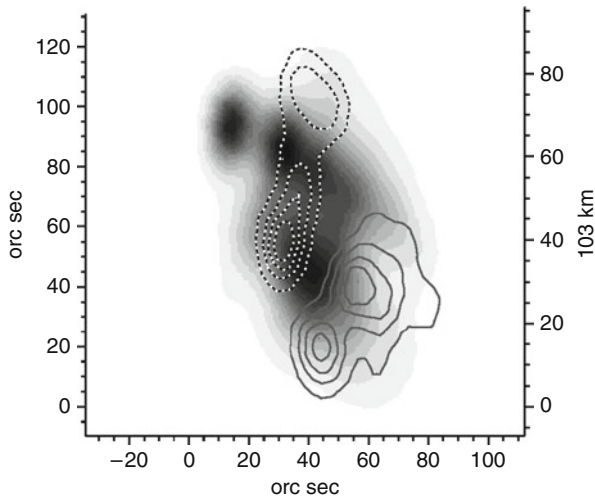


Fig. 4 The structure of the microwave source at 04:10 UT on November 2, 1992. The intensity is shown by shades of *gray*, and the polarization by *dotted curves*. The observations were made with the Nobeyama radio heliograph at 17 GHz. The distances on the axes are shown in arcseconds (the corresponding values in 1,000 km are shown on the right scale)

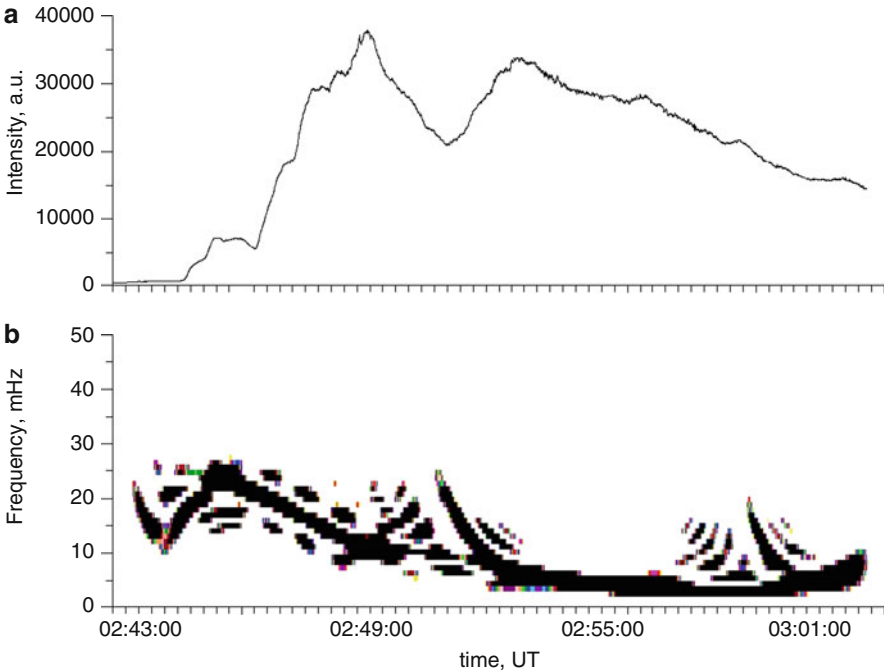
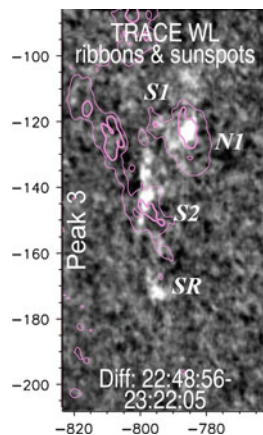


Fig. 5 The event on November 2, 1992: the microwave radiation intensity based on the Nobeyama spectrum polarimeter observations at the frequency 17 GHz (a); the corresponding spectrum obtained using the Wigner–Ville transformation (b)

Fig. 6 The flare region on June 17, 2003 at UV wavelengths of 160 nm (the difference of TRACE images [6]). The marks *S1*, *N1*, and *S2* correspond to the most significant sunspots, and *SR* to the southmost part of the region. The axes are graduated in arcseconds from the solar disk center



flare. Later, the frequency gradually decreased, while the radiation intensity began to rise. When the intensity reached its maximum the frequency was approximately equal to the initial value, $\nu \approx 0.01$ Hz. However, when the radiation intensity decayed after the first burst, the frequency ν began to increase again, reaching a value slightly less than 0.025 Hz at the maximum. Before the beginning of the second burst, the frequency decreased again and finally reached $\nu \approx 0.005$ Hz, as it was in the event on March 3, 2001. We should mention the following feature: in the first case, the frequency before the burst rose to a higher value than that in the second case, and one can see that a high maximum intensity was also reached in the first burst.

The spectra of the second event on November 2, 1992 show similar features.

The third event which we would like to present as an example is a class M6.8 flare observed at the Nobeyama observatory on June 17, 2003 in the active region AR10386. A radio image of the radiation source at 17 GHz is given in Fig. 6 [6]. In this case, the main sources of the microwave radiation were not the loops, but a set of sunspots of different polarity (*N1*, *S1*, and *S2* in the figure). The maximum magnetic field amounted to +3,080, -2,120, and -1,750 G for the sunspots *N1*, *S2*, and *S1*, respectively.

Figure 7 shows the initial realization of the event at 17 GHz and the corresponding Wigner–Ville low-frequency modulation spectrum. In this case a strong modulation line with a frequency of about $\nu \approx 0.003$ Hz was observed in the spectrum before the main burst, but a characteristic increase followed by a decrease in frequency, which was detected in the spectra of the events on March 30, 2002 and November 2, 1992, is not seen here. Possibly the presence of such a modulation in the low-frequency spectrum can be related to global oscillations of the solar photosphere, whose frequency is close to the observed value (the corresponding period is close to 5 min).

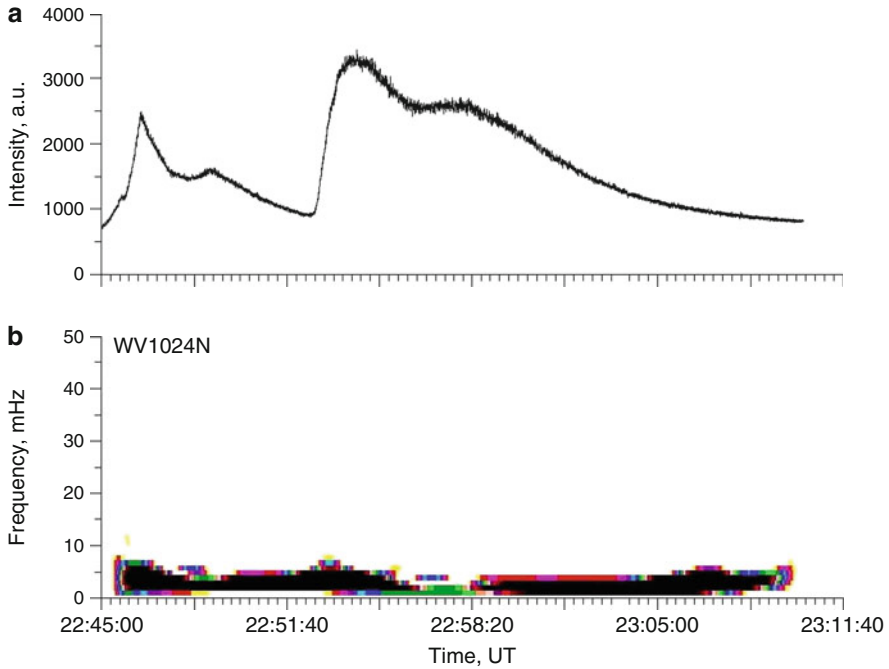


Fig. 7 The event on June 17, 1993 (the second flare). A part of the realization based on Nobeyama observations at the frequency 17 GHz (a), and the corresponding spectrum obtained using the Wigner–Ville transformation (b)

3 Analysis of the Origin of the Low-Frequency Modulation

Microwave radiation of coronal magnetic loops is usually interpreted as gyrosynchrotron radiation of some population of fast electrons at the gyrofrequency harmonics in the magnetic field of the loop. In the case of a power-law energy distribution of electrons $f(E) \propto E^{-\delta}$, the intensity of gyrosynchrotron radiation from an optically thin source is determined by the relationship [7]

$$I_\nu \propto B^{-0.22+0.9\delta} (\sin \theta)^{0.43+0.65\delta} \quad (1)$$

For typical values of the index of the energy spectrum of electrons $2 < \delta < 7$, (1) yields a relatively strong dependence of the intensity on the magnetic field and significant angular anisotropy of the radiation: $I_\nu \propto B^{1.58...6.08} (\sin \theta)^{0.87...4.12}$. Thus, oscillations of the magnetic field in a coronal magnetic loop or the loop oscillations leading to variations in the angle θ between the magnetic field and the line of sight may result in the modulation of the intensity of the received radio emission.

Oscillations of coronal magnetic loops can be studied in terms of a uniform plasma cylinder of the radius $r = w/2$ and length l , whose ends are loaded by

a dense chromosphere and can therefore be assumed frozen-in. The plasma inside the cylinder has a density ρ_i , a temperature T_i , and a magnetic field B_i . Outside the cylinder, the corresponding parameters are ρ_e , T_e , B_e . The magnetic field variations can be caused by several types of MHD eigenoscillations of the plasma cylinder.

3.1 Fast Magneto-Sonic (FMS) Oscillations

In the case of a thin ($r/l \ll 1$) and dense ($\rho_e/\rho_i \ll 1$) cylinder, the oscillations display the frequency [8]

$$\omega_+ = (k_{\parallel}^2 + k_{\perp}^2)^{1/2} (C_{si}^2 + C_{Ai}^2)^{1/2}, \quad (2)$$

where $k_{\parallel} = \pi s/l$, $s = 1, 2, 3 \dots$; $k_{\perp} = \lambda_i/r$, (λ_i are the zeros of the Bessel function $I_0(\lambda_i) = 0$); $C_{si} = (k_B T_i/m_i)^{1/2}$ is the sound velocity (k_B is the Boltzmann constant), and $C_{Ai} = B_i/(4\pi\rho_i)^{1/2}$ is the Alfvén velocity inside the cylinder. If the arch is sufficiently “thick”, namely, $(l/r) < 1.3(\rho_i/\rho_e)^{1/2}$, the global fast magneto-sonic mode has the frequency [9]

$$\omega_{+G} \approx k_{\parallel} C_{Ae} \approx \frac{\pi}{l} C_{Ae}, \quad (3)$$

where C_{Ae} is the Alfvén velocity outside the cylinder.

We assume that the magnetic field in a coronal magnetic loop is not potential and decreases with the altitude slower than the potential field. This is indicated, in particular, by a weak variation in the thickness of the coronal magnetic loop with the altitude, as is seen in the soft X-ray radiation. Assuming that the average magnetic field in the flare loop is $B_i \approx 300$ G and the plasma density is $n_i \approx 1.4 \times 10^{11} \text{ cm}^{-3}$ inside and $n_e \approx 2 \times 10^{10} \text{ cm}^{-3}$ outside the tube, we obtain an estimate for the FMS oscillation periods for the event on March 3, 2001: $P_+ = 2\pi/\omega_+ \approx 4$ s, $P_{+G} = 2\pi/\omega_{+G} \approx 50$ s.

Assuming the plasma density $n_e \approx 2.1 \times 10^{10} \text{ cm}^{-3}$, $n_i \approx 2 \times 10^{11} \text{ cm}^{-3}$ and the magnetic field $B_i \approx 300$ G for the event on November 2, 1992, we obtain the following estimates: $P_{+G} = 2\pi/\omega_{+G} \approx 25$ s for the first flare with the maximum at 02:49 UT and $P_{+G} = 2\pi/\omega_{+G} \approx 27$ s for the second flare with the maximum at 04:09 UT.

Actually, each of these values can be somewhat smaller since we take the maximum value of the plasma density observed in this event at the top of the loop [4, 5].

3.2 Kink Oscillations of a Coronal Magnetic Loop

These oscillations [10] can lead to variations in the angle θ between the magnetic field direction and the line of sight. The frequency of kink oscillations is expressed by the formula

$$\omega_{kink} = k_{||} \left(\frac{\rho_i C_{Ai}^2 + \rho_e C_{Ae}^2}{\rho_i + \rho_e} \right)^{1/2}. \quad (4)$$

If the loop is sufficiently dense i.e. $\rho_i \gg \rho_e$, then for $B_i \approx B_e$ the frequency of kink oscillations is by the factor of $\sqrt{2}$ greater than that of the global FMS mode, i.e. $P_{kink} \approx 44$ s in the event on March 3, 2001 and $P_{kink} \approx 14$ s, $P_{kink} \approx 15$ s for the first and second flares on November 2, 1992, respectively.

3.3 Alfvén Oscillations

The frequency of these oscillations is

$$\omega_A = k_{||} C_{Ai} \quad (5)$$

The corresponding periods for the event on March 30, 2001 and the flares on November 2, 1992 at 02:49 UT and 04:09 UT are equal to $P_A \approx 62$ s, $P_A \approx 20$ s, and $P_A \approx 22$ s, respectively

3.4 Slow Magneto-Sonic (SMS) Oscillations

According to [11], the period of these oscillations can be estimated as

$$P_{GSM}(s) \approx 13 \times l[Mm] / \sqrt{T[MK]}, \quad (6)$$

where l is the length of the loop and T is the temperature.

Slow magneto-sonic oscillations can modulate the microwave radiation due to the Razin effect. In this case, an increase or a decrease in the plasma density due to the wave propagation in the loop leads to the corresponding increase or decrease in the radiation intensity. For the event on March 30, 2001, (6) yields $P_{GSM} \approx 230$ s and for two flares on November 2, 1992 at 02:49 and 04:09 UT, the values $P_{GSM} \approx 199$ s and $P_{GSM} \approx 218$ s, respectively, which are close to the observed period (200 s). However, when such a mechanism is used for the events on March 30, 2001 and November 2, 1992, it seems difficult to explain the observed linear frequency modulation, a severalfold increase in frequency directly before the flare and the subsequent decay during the event.

3.5 LRC-Oscillations of the Coronal Magnetic Loop

These are eigenoscillations of the loop as an equivalent electrical circuit [12]. Convective flows of photospheric plasma interacting with the magnetic field at the loop footpoints generate an electrical current flowing from one footpoint to the other through the coronal part and closing up in the subphotospheric layers where the conductivity becomes isotropic. Thus, the loop with a subphotospheric current channel is similar to a coil with electrical current, for which the equation of an equivalent electrical circuit for small deviations of the current I from the stationary value I_0 can be written [12]:

$$\frac{1}{c^2}L \frac{d^2 I}{dt^2} + \left[R(I_0) - \frac{|V_r l_1|}{c^2 r_1} \right] \frac{dI}{dt} + \frac{1}{C(I_0)} I = 0 \quad (7)$$

where L is the inductance of the loop and r_1 is the loop radius in the region of the footpoints. If the loop is approximated by a coil of a length l and a small radius r , so that $r \ll l$, then the inductance can be described by the well-known expression [13]

$$L \approx 2l \left(\ln \frac{4l_2}{\pi r_2} - \frac{7}{4} \right), \quad (8)$$

where $C(I_0)$ is the effective capacitance dependent on the current flowing in the circuit,

$$\frac{1}{C(I_0)} \approx \frac{I_0^2 l_2}{c^4 \rho_2 \pi r_2^2} (1 + b^{-2}). \quad (9)$$

Parameter b is determined by the values of the magnetic-field components and the pressure on the tube axis and outside the tube:

$$b = \frac{B_{\phi 0}(r_2)}{B_{z0}(r_2) - B_{z0}(0)} \approx 6 \frac{B_{\phi 0}(r_2) B_{z0}(0)}{8\pi [p(\infty) - p(0)]}. \quad (10)$$

Here, r_2 is the tube radius in the coronal part of the loop, $B_{\phi 0}$ and B_{z0} are the azimuthal and axial components of the magnetic field of the coronal magnetic loop. The effective resistance of the loop is determined by the formula

$$R(I_0) \approx \frac{F_1^2 I_0^2 l_1}{(2 - F_1) c^4 n_1 m_i v_{ia}^{\prime} \pi r_1^4} (1 + b^{-2}), \quad (11)$$

where l_1 , r_1 , n_1 , F_1 are the altitude length of the photospheric electromotive force area, the tube radius, the electron number density, and the relative mass of neutrals in this area, respectively, and $v_{ia}^{\prime}$ is the effective frequency of electron–atom collisions. In (7), V_r is the radial component of the velocity of the convergent convective flow of photospheric plasma at the footpoints of the CML. The resistance of the circuit is mainly contributed by the footpoints of the coronal magnetic loop with a

relatively low conductivity stipulated by ion–atom collisions (the so-called Cowling conductivity). It can be assumed that $b^2 \gg 1$ in the coronal part of the loop; hence, for the frequency of LRC oscillations of the loop we obtain

$$\nu_{LRC} = \frac{c}{2\pi \sqrt{LC(I_0)}} \approx \frac{1}{(2\pi)^{3/2} \sqrt{\Lambda}} \frac{I_0}{cr_2^2 \sqrt{n_2 m_i}}. \quad (12)$$

Here, $\Lambda = \ln \frac{4l_2}{\pi r_2} - \frac{7}{4}$, n_2 is the electron number density in the coronal part of the loop.

Equation (7) assumes that the electrical-current oscillations are in-phase at all points of the loop as an equivalent electrical circuit. On the other hand, the current variations propagate along the loop with the Alfvén velocity. Hence, the in-phase condition requires that the Alfvén time $\tau_A = l_2/C_{Ai}$ be substantially smaller than the oscillation period $T_{LRC} = 1/\nu_{LRC}$. Since $I_0 \approx cr_2 B_{\phi 0}(r_2)/2$, the in-phase condition takes the form

$$\frac{B_{\phi 0}(r_2)}{B_{z0}(0)} \ll \pi \sqrt{2\Lambda} \frac{r_2}{l_2}. \quad (13)$$

It should be emphasized that the magnetic loops observed in the solar corona have, as a rule, a very slight twisting of the magnetic field; hence, (12) is valid for at least not too long magnetic loops. For the considered loop in the event on March 30, 2001, the in-phase condition is fulfilled for the current $I_0 \leq 9 \times 10^{10}$ A. For the event on November 2, 1992, we obtain the estimate $I_0 \leq 1.9 \times 10^{11}$ A for the flare loops at 02:49 UT and $I_0 \leq 3 \times 10^{11}$ A for the flare loop at 04:09 UT. The average magnetic field in the coronal part of the loop was assumed to be $B_{z0} \approx 300$ G in all these cases.

Assuming that for the event on March 30, 2001 the modulation frequency of the microwave radiation at the pre-flare stage was $\nu \approx 0.005$ Hz, from (12) we obtain the current $I_0 \approx 10^{10}$ A for the observed parameters of the loop. Similar estimates for the flares on November 2, 1992 at 02:49 and 04:09 UT yield $I_0 \approx 10^{10}$ A and $I_0 \approx 5 \times 10^{10}$ A, respectively.

4 Discussion

The analysis of the low-frequency modulation by the Wigner–Ville method shows that quasistationary oscillations with the frequency $\nu \approx 0.005$ Hz (the corresponding period $T \approx 200$ s) exist at the pre-flare stage in the CML. Such oscillations cannot be explained by the eigenoscillations of the loop as an MHD resonator since fast magneto-sonic, Alfvén, and kink oscillations have periods substantially shorter than the observed modulation period. Acoustic oscillations with the period of about 400 s cannot be responsible for the observed modulation, either. We assume that in this case the modulation source is the oscillations of the CML as an equivalent electrical circuit. The oscillation frequency depends on the electrical

current flowing along the magnetic loop and generated as a result of interaction between the convective flows of atmospheric plasma and the magnetic field at the loop footpoints. Our estimates for those different flows show that the modulation of the microwave radiation with the frequency $\nu \approx 0.005$ Hz in the pre-flare phase arises if the current in the coronal magnetic loops has the values $I_0 \approx 10^{10}$ A, $I_0 \approx 10^{10}$ A, and $I_0 \approx 5 \times 10^{10}$ A (for the flare on March 30, 2001 and two flares on November 2, 1992, respectively).

It follows from the dynamic spectra presented in Fig. 2 that about 50 s before the flare the modulation frequency increased relatively rapidly from $\nu \approx 0.005$ Hz to $\nu \approx 0.035$ Hz, i.e., by the factor of seven. Hence, according to our interpretation, the current flowing in the coronal magnetic loop increased by the same factor. After the maximum was reached, the modulation frequency and, therefore, the electrical current decreased to the pre-flare values for a time of the order of the flare duration ($\tau f \approx 200$ s). The pattern is also similar for two flares of the event dated November 2, 1992.

Thus, our data show that the flare process is preceded by a sharp rise in the electrical current in the coronal magnetic loop. It follows from (11) that the Joule heating power is proportional to the fourth power of the current:

$$\frac{dW}{dt} = R(I_0)I_0^2 = \frac{F_1^2 l_1 I_0^4}{(2 - F_1)c^4 n_1 m_i v_{ia}' r_1^4} \text{ erg s}^{-1}. \quad (14)$$

Here, $v_{ia}' = 2, 25 \times 10^{-11} F_1(n_1 + n_a)\sqrt{T}$, $n_1 = (1 - F)(n_1 + n_a)$. The subscript 1 relates to the energy release region. The rate of the energy release depends on the tube radius at the footpoints of the magnetic loop. For a stationary tube with the magnetic field $B_z(0) = 500$ G on the axis and with the radial component of the photospheric convection velocity $|V_r| = 10^4$ cm·s⁻¹ at an altitude of 500 km above the photosphere, the standard photospheric model [14] leads to the estimate of the magnetic tube radius at the footpoints $r_1 \approx 2.5 \times 10^6$ cm.

For the event on March 30, 2001, the current increases to $I_0 \approx 7 \times 10^{10}$ A, for the flare at 02:49 UT on November 2, 1992, to $I_0 \approx 5 \times 10^{10}$ A, and for the second flare at 04:09 UT on the same day, to $I_0 \approx 1.5 \times 10^{11}$ A.

Optical flares usually occur for the following values of the density and temperature [15, 16]: $3 \times 10^{12} < n_1 + n_a < 5 \times 10^{13}$ cm⁻³ and $T \approx 10^4$ K; in this case, $F_1 \approx 0, 44$. The Joule heating power amounts to $(dW/dt) \approx 6 \times 10^{27}$ erg s⁻¹, and the total energy input into the plasma heating reaches about 6×10^{29} erg for the flare process duration. During the energy release, the electrical current dissipates, its value decreases, and the LRC modulation frequency returns to its pre-flare value.

The electrical current increase before the flare can be caused by the development of the balloon mode of the flute instability at the footpoints of the coronal magnetic loop which play the role of a flare trigger [17]. At the chromosphere base, the loop radius increases with the altitude because of the decrease in the external gas pressure; hence, a curvature of the magnetic field appears, so that the outer plasma is affected by a centrifugal force directed into the tube creating conditions for the

development of the flute instability. The instability criterion in this case has the form [17]

$$2 \frac{n}{n + n_a} - \cos \theta > 0, \quad (15)$$

where n and n_a are the number densities of electrons and neutrals, respectively, at the loop footpoints, θ is the angle between the direction of the curvature radius and the vertical line. The instability develops if the degree of polarization is not too small, which requires a certain preheating of the footpoints of the magnetic tube. The characteristic rise time of the balloon mode of the flute instability is given by

$$\tau_B \approx 2 \times 10^{-3} \sqrt{H(cm)} s, \quad (16)$$

where $H = \kappa_B T / \mu m_i g$ is the reduced altitude of the non-uniform atmosphere and μ is the average atomic mass. If the temperature of the outer plasma surrounding the loop footpoints is $6 \times 10^3 \text{ K}$, $\tau_B \approx 7 \text{ s}$. The development of the flute instability leads to the invasion of the additional plasma surrounding the loop footpoints into the interior of the magnetic tube and, as it follows from the induction equation, to the corresponding amplification of the electrical current. If, for example, the azimuthal component of the magnetic field before the development of the flute instability depended on the radius, i.e. $B_{\varphi 0}(r) = B_{\varphi 0} r / r_1$, then after the invasion of a plasma “tongue” with the velocity $V_r(r, t) = -V_r(t) r / r_1$ the electrical current in the tube would increase as $I(t) = I_0 \exp[2 \cdot \int_0^t V_r(t') / r_1] dt'$. A sevenfold increase in the current in the event on March 30, 2001 corresponds to the exponent index 195.

As the plasma is heated inside the tube, the “tongue” is stopped by an increase in the internal pressure, and the flute instability disappears.

Thus, the performed analysis leads to the following conclusions:

1. In the cases in which the source of microwave radiation was a coronal magnetic loop, we detected a characteristic modulation of the radiation, with the frequency increasing severalfold for a few tens of seconds before the flare and then decreasing to the initial value as the flare proceeds. It was noted that the intensity profile often reproduces, to a certain extent, the profile of the modulation line.
2. In the analysis of a flare, whose source was the active region with several sunspots, only a modulation line with a constant frequency corresponding to 5-min photospheric oscillations was observed. The different character of the spectra is the evidence for different mechanisms of the flare development in the cases in which the source is a coronal magnetic loop or sunspots.
3. The frequency closest to the observed displays a type of oscillations of the coronal magnetic loop as an equivalent electrical circuit. In this case, the oscillation frequency is proportional to the electrical current in the loop, so that a sharp increase in the frequency before the flare is possibly due to a dramatic increase in the current strength.
4. A possible reason for the increased current is the flute instability that develops in the loop.
5. The main results of the work were published in the paper [18].

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Observations of Solar Flares from GHz to THz Frequencies

Pierre Kaufmann

Abstract The discovery of a new solar burst spectral component with sub-THz fluxes increasing with frequency, simultaneous but separated from the well known microwave component, brings serious constraints for interpretation. Suggested explanations are briefly reviewed. They are inconclusive indicating that further progresses on the understanding of nature of the emission mechanisms involved require the knowledge of GHz to THz continuum burst spectral shapes. New 45 and 90 GHz high sensitivity solar polarimeters are being installed at El Leoncito high altitude observatory, where sub-THz (0.2 and 0.4 THz) solar flare flux data are being obtained regularly since several years. Solar flare THz photometry in the continuum should be carried in space or at few selected frequency windows at exceptional ground-based sites. A dual photometer system, operating at 3 and 7 THz, is being constructed to be flown in a long duration stratospheric balloon flight in Antarctica (summer 2013–2014) in cooperation with University of California, Berkeley, together with GRIPS experiment. One test flight is planned for the fall 2012 in USA. Another long duration balloon flight over Russia is considered (2015–2016), in a cooperation with Moscow Lebedev Physics Institute.

1 Introduction

The radio emission of solar bursts was well studied in the past century covering three frequency decades, from tens of MHz up to tens of GHz. Models of interpretation were believed to be solidly founded, despite of the complete ignorance of flare emissions on the wide frequency gap between tens of GHz and the visible. Radio

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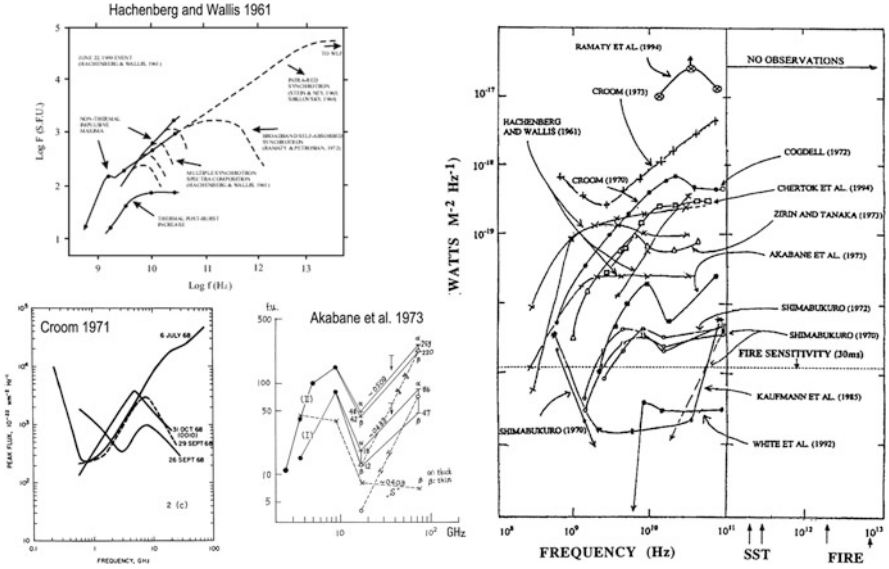


Fig. 1 Various examples of solar burst spectra with fluxes increasing with frequency, up to 100 GHz (after [25]). The *inner panels* show examples of time-variable spectra for single bursts [1, 9, 19]

emissions in the MHz to few GHz range are produced in the solar corona, being usually attributed to plasma waves excited by flare generated shock waves and/or beams of fast electrons accelerated [2, 4, 61]. Emissions in the few GHz to tens of GHz frequency range originate from a distinct mechanism, usually attributed to gyrosynchrotron radiation of mildly relativistic electrons (up to 1 MeV) in magnetic field at solar active centers [2, 11].

Radio burst emission temporal and spectral features at cm to mm-waves (GHz to tens of GHz) have been reviewed [25]. It has been shown that although the majority microwave bursts exhibit maximum emission up to about 10 GHz [2, 39], there were various observations of burst spectra with fluxes increasing with frequencies up to 100 GHz [1, 8, 9, 19, 23, 47, 52, 60, 63]. Several examples of solar burst spectra with fluxes increasing in the 10–100 GHz range are illustrated in Fig. 1. One solar hard X-ray impulsive solar burst occurred on May 21, 1984, produced emission in the 30–90 GHz range only, with rapid time structures [23], shown in Fig. 2. Higher radio frequency emission spectral components might be, in fact, quite common in the impulsive phase of solar flares. They cannot be ignored, adding new parameters for interpretation [3, 24, 35, 43, 46, 48, 57, 59]. There were earlier suggestions associating the high frequency spectral trend to synchrotron radiation due to extremely high electrons energies, peaking somewhere in the infrared-visible range [53, 56]. These results indicate that observations only in microwave spectral range are not sufficient to describe the emission processes in solar bursts. Observations at frequencies larger than 100 GHz, with enough sensitivity and time resolution, become essential to understand the physical processes at the origin of flares.

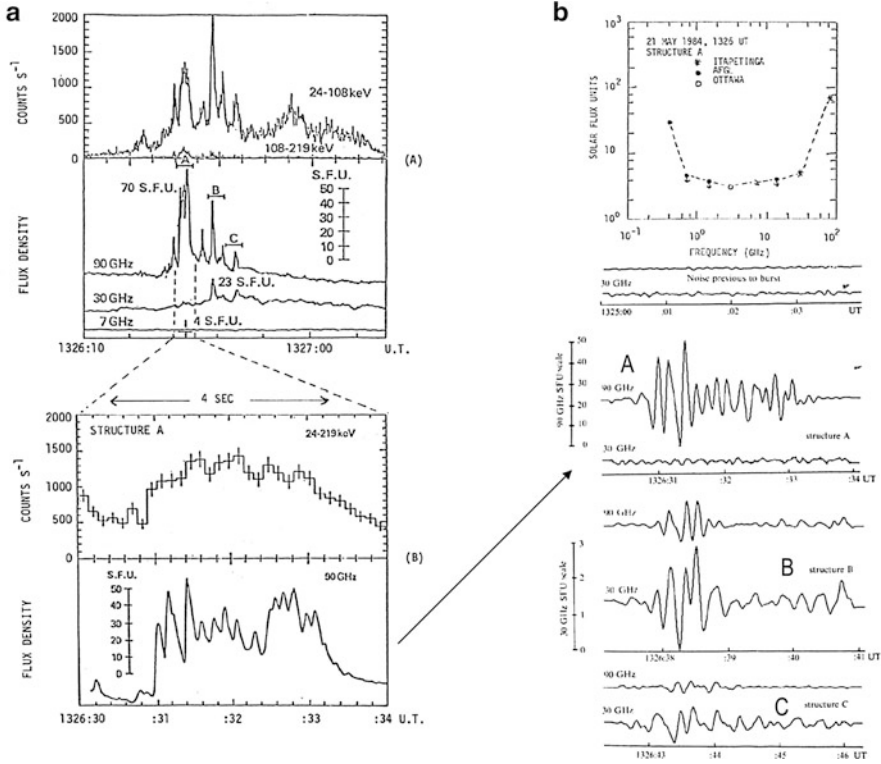


Fig. 2 The solar burst of May 21, 1984, exhibited emission observed at 90 GHz only, in (a), well correlated in time with SMM hard X-rays [23]. Inner panels at left show in (b) the spectrum for the main time structure, in (c) the rapid superimposed time structures observed at 90 and 30 GHz

2 Solar Activity at Frequencies Larger Than 100 GHz

The first sub-THz solar observations were carried out in 1968 [7], with a 250 GHz bolometer on the Queen Mary College 1.5 m optical telescope, UK. They found unexpected 100 K brightenings on active regions, on time scale limited to 1 min, producing an important impact in interpretation [3]. Similar observations in 1975 produced inconclusive results [21]. Nearly three decades elapsed before the Brazilian Solar Submillimeter Telescope (SST) was installed, optimized, and operated in the Argentina Andes within cooperation with Complejo Astronomico El Leoncito, San Juan, Argentina, at a 2,550 m altitude site with more than 300 days good for sub-THz observations [29]. It consists in a 1.5 m Cassegrain antenna on an alt-azimuth positioner, enclosed in a 3 m GoreTex radome (Fig. 3). Independent radiometers and feed horns are placed in the focal plane, four at 0.2 THz and two at 0.4 THz, producing three 0.2 THz partially overlapping beams, each one with 4 arc min beam widths (2, 3 and 4 in Fig. 1) with one 0.4 THz, 2 arc min beam in the

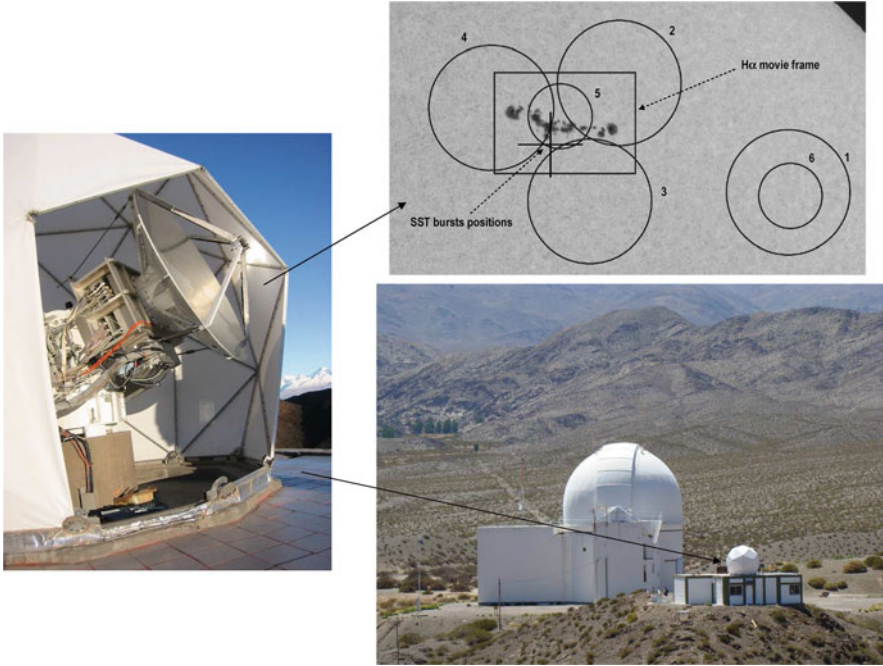


Fig. 3 The Solar Submillimeter Telescope (SST) 1.5 m Cassegrain antenna at El Leoncito Astronomical Complex, San Juan, Argentina Andes. The *top right panel* shows the six antenna beams pointing to a solar active region (on February 8, 2010) [34]

center (beam 5), and two beams displaced by 8 arc min, at 0.2 and 0.4 THz (beams 1 and 6 respectively). The comparison of antenna temperatures at beams 2, 3 and 4 permits the determination of the burst position in space and time needed for flux calculations [17, 18].

The first SST observations revealed the presence of numerous sub-second time structures exhibiting occurrence rates correlated with the intensity of X-rays burst emissions [25]. A relationship between the onset of sub-THz pulsations and the acceleration phase of CMEs has been suggested [26]. Quiescent active center observations suggested the presence of two thermal sources, one hot at microwaves and another cold at sub-THz [55].

The most important result obtained by SST was the discovery of a new spectral component with increasing fluxes at sub-THz frequencies that appears along with, but separated from, the well-known emission component seen at microwave frequencies [27, 30, 31, 55] (see examples in Figs. 4 and 5 [30, 31]).

A number of emission processes were suggested to explain the sub-THz spectral component. They include emission by free-free collisions of thermal electrons, synchrotron produced by high energy electrons [27, 55], synchrotron from relativistic positrons [55, 58]; emission by Langmuir waves excited by beams of electrons and protons at denser regions of the solar active centers [49, 50] by inverse-Compton

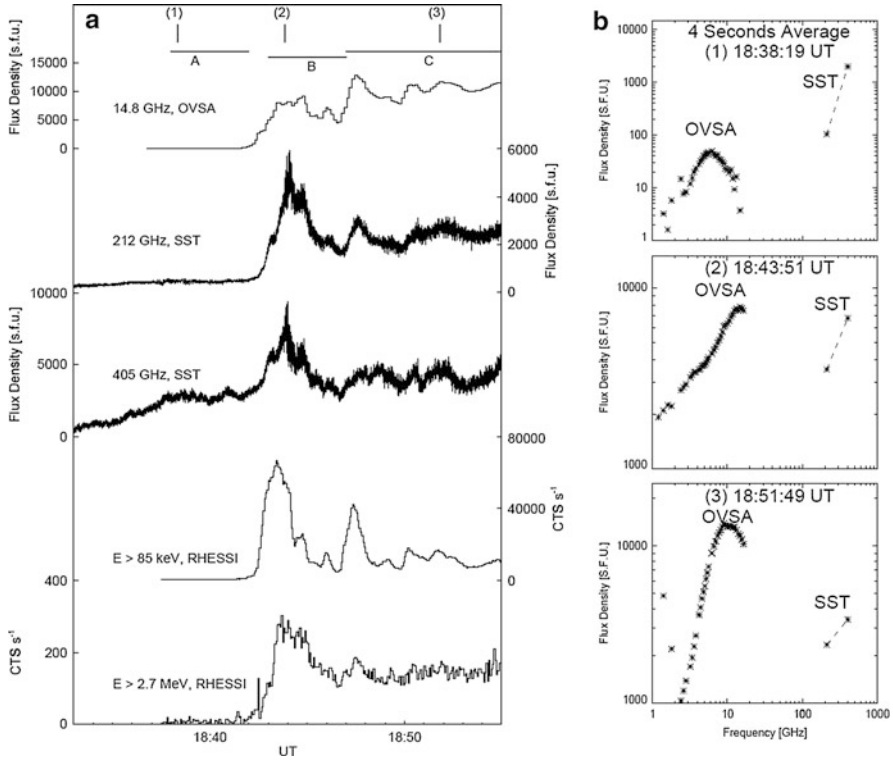


Fig. 4 Example of a solar burst occurring on December 6, 2006 showing in (a) the time profiles, from top: microwaves, 0.2 and 0.4 THz, two energies hard-X-rays; in (b) the new spectral component increasing with frequency in the sub-THz range, together with the well known microwave spectral component at different times, labeled (1), (2) and (3) in (a) [30]

effect on the field of synchrotron electrons [24], on the field of turbulent Langmuir waves [13] and the Vavilov–Cherenkov emission by high energy electrons on an assumed partially ionized chromosphere gas [13].

The double spectral feature (i.e., one component at microwaves, and another at THz) poses serious difficulties for the suggested models. Another possibility recently suggested assumes that a mechanism observed in laboratory accelerators might become important also in solar flares [28, 62]. The high frequency emissions are attributed to incoherent synchrotron radiation (ISR) produced accelerated beams of high energy electrons with intensity peaking at THz frequencies while certain wave-particle instabilities may set in the electron beam, giving rise to bunching of the electrons which radiate powerful broadband coherent synchrotron radiation (CSR) in the microwave spectrum peaking at wavelengths comparable to the size of the bunching [62]. The acceleration of a distinct ultra-relativistic electron beams in solar flares is suggested by the sudden reduction of photon energy spectral index at hard X- and gamma-rays, observed in certain events [40]. Simulations of the

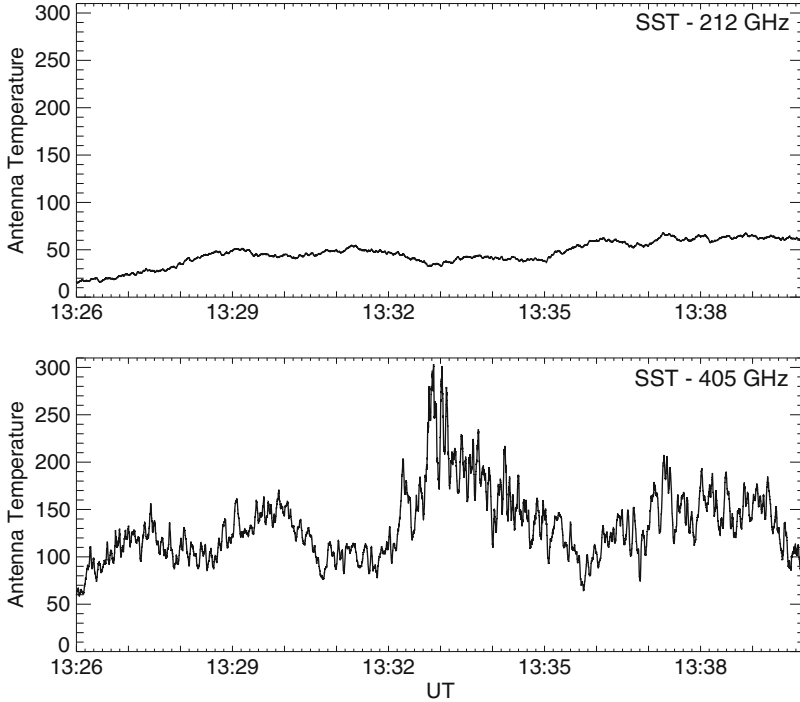


Fig. 5 A small complex impulsive burst produced emission at 0.4 THz only, on February 8, 2010 [34]

ISR and CSR radiation spectral components, using typical active region physical conditions [36, 37] have shown that the mechanism may be extremely efficient in solar flares (Fig. 6).

It is likely that more than one of the above mentioned mechanisms might be acting at the same time, at distinct proportion for different bursts, with the free-free contribution always being present to a certain level. To fully understand the nature of the high frequency emission in flares it became clear enough that it is necessary to measure the complete continuum spectra at higher THz frequencies.

3 Solar Burst Observations at THz Frequencies

Solar activity observations in the THz range of frequencies requires the use of detecting systems placed outside the terrestrial atmosphere, such as at high altitude aircraft [12], stratospheric balloon [10], satellites [45], or through few atmospheric THz transmission “windows” at exceptionally good high altitude ground based locations [6, 15, 41].

Detection of excess solar continuum radiation in the THz range poses several technological challenges [5, 20, 32, 54]. They include the need of efficient

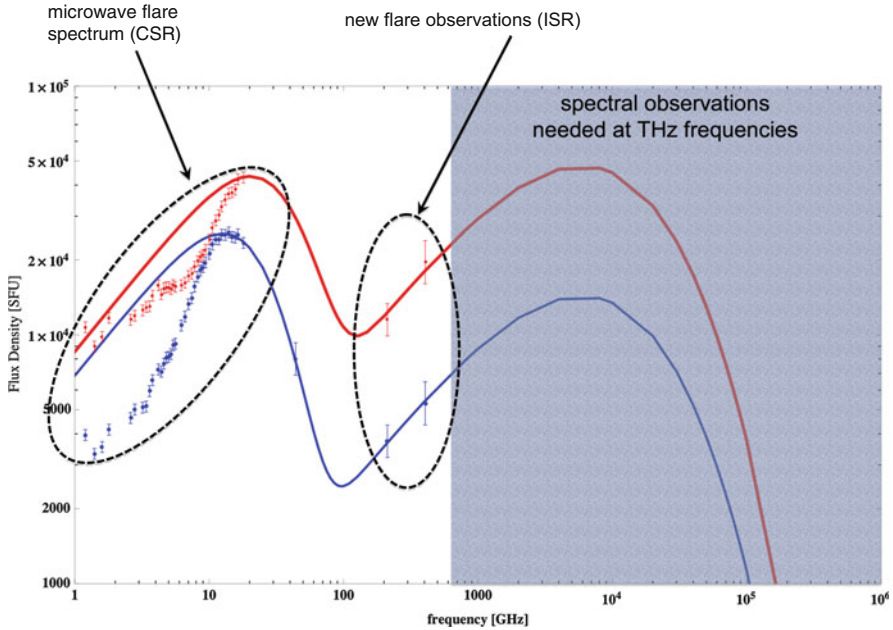


Fig. 6 The simulation of microbunch instability in a solar laboratory accelerator electron beam [36,37], adopting typical solar flare parameters, reproduces well the solar burst of November 4, 2003, observed at 0.2 and 0.4 THz by SST and at microwaves by OVSA. The THz component is due to incoherent synchrotron radiation. The microwaves are the result of microbunching instability which generates large amount of coherent synchrotron radiation with only a very small fraction of accelerated electrons. The steeper observed microwave spectral index for frequencies smaller than 10 GHz is due to self absorption and Razin suppression of the medium where the flare occurred, while the laboratory radiation is produced in vacuum

suppression of the powerful visible and near infra-red background component of emission, the band-pass filtering and the appropriate selection of uncooled detector system and optical setup design to obtain enough sensitivity and time resolution [14, 31–33]. A system has been designed utilizing a new concept that is capable to observe the whole Sun with sufficient sensitivity to detect flares subtended by much smaller angular sizes [32, 33, 42]. It operates at central frequencies of 3 and 7 THz (see Fig. 6). The detection systems will be sensitive to fluxes larger than 60 SFU (one solar flux unit = $10^{-22} \text{ W m}^{-2} \text{ Hz}^{-1}$), time resolution of 100 ms, at any location in the solar disk.

The 3–7 THz photometers system utilizes two identical 76 mm Cassegrain telescopes producing a solar disk image on non-imaging photon concentrator placed in front of the Golay cell sensitive surface (Fig. 7). The Cassegrain primary surface has been roughened to diffuse a significant fraction of radiation with $\lambda < 30 \mu$ [38]. The Golay cell sensor is preceded by a 20 Hz resonant fork chopper, 3 and 7 THz metal mesh band pass filters and TydexBlack low-pass membranes, one for each telescope.

The SOLAR-T experiment is planned to fly on board of a long-duration stratospheric balloon flight, coupled to the GRIPS (Gamma Ray Imaging Polarimeter

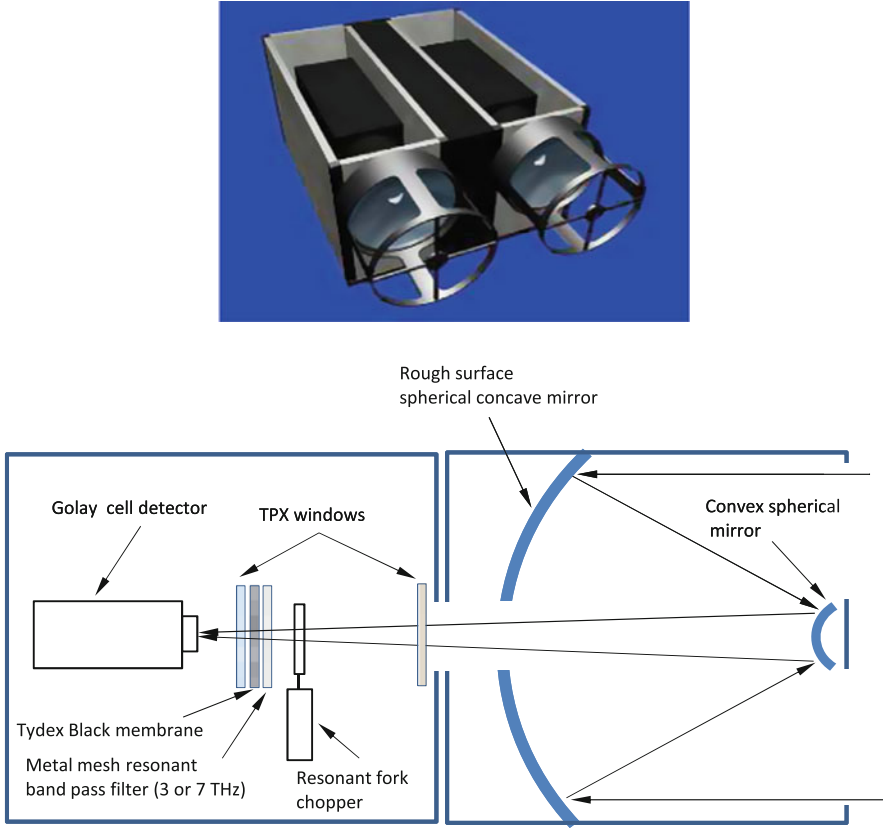


Fig. 7 Schematic diagram showing the principal components of the THz photometer, using a Cassegrain optical configuration. The *upper right panel* shows the dual 3 and 7 THz photometers conceptual configuration

for Solar flares) gamma-ray experiment [51] in cooperation with University of California, Berkeley, USA. One engineering flight is scheduled for fall 2012 in USA, and a 2 weeks flight over Antarctica in 2013–2014. Another long duration stratospheric balloon flight over Russia (1 week) is planned in cooperation with Lebedev Physics Institute, Moscow (2015–2016).

4 Completing Flare Spectra at Frequencies Lower Than 100 GHz

Solar burst observations in the frequency gap between 15–18 GHz (the upper limit of OVSA [22] and RSTN observatories [16]) and the lower SST limit (200 GHz) is only partially completed by Nobeyama 35 GHz polarimeter and 80 GHz



Fig. 8 The solar patrol polarimeters completed for installation at El Leoncito Observatory. The 90 GHz 22 cm aperture at the *left side* and the 45 GHz 44 cm at the *right side* produce a 1° beam for full sun regular observations

radiometer [44], which cannot observe simultaneously with SST. Observations in this frequency gap became very important after the discovery of the new independent THz spectral component. The knowledge of burst emission spectral features in the range 18–200 GHz is essential to identify the transition of microwaves to sub-THz components and how they evolve with time during the burst.

To complete the 18–200 GHz spectral gap two solar polarimeter patrol telescopes have been built, to operate at 45 and 90 GHz at El Leoncito Astronomical Complex together with SST. The polarimeters have full Sun half power beams, sensitivity of ten solar flux units, with time resolution of 10 ms (Fig. 8). They are scheduled to begin regular solar observations in December 2011.

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On the Interaction of Solar Rotational Discontinuities with a Contact Discontinuity Inside the Solar Transition Region as a Source of Plasma Heating in the Solar Corona

S.A. Grib and E.A. Pushkar

Abstract We consider the self-similar MHD problem of the oblique interference of a solar rotational (Alfven) discontinuity A and a stationary contact discontinuity C . The interaction between A and C is studied for typical conditions in the solar corona. Since solar Alfven waves observed in the solar plasma are numerous, prerequisites exist for the formation of a solar rotational discontinuity that propagates from the chromosphere through the transition region to the corona. Dissipative slow MHD shock waves with insignificant variation of the magnetic field also appear due to the refraction of the solar non-dissipative rotational discontinuities against a contact discontinuity inside the transition region. It is supposed that a real source of plasma heating may exist in the high solar corona due to the well-known mechanism of Landau damping of the dissipative slow MHD shock waves. Frequently observed explosive events may also be triggered in the solar chromospheric plasma. Thereby, we suggest a new model of the coronal plasma heating.

1 Introduction

Dynamics of the upper chromosphere, transition region, and lower solar corona can play an important role in the heating of the coronal plasma. The parameters of the heat flux in the solar corona are still unknown [5]. Thus, at present the main goal of solar physics is to estimate the heating processes in the solar atmosphere.

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In numerous papers it was supposed that the coronal heating must be connected with kinetic Alfvén waves or with the so-called magnetic reconnection and turbulence [11]. However, finding the location for the kinetic Alfvén wave resonance constitutes a complex problem. Currently, no exact theory is available to describe such processes.

The existence of non-dissipative MHD Alfvén waves coming from the photosphere is confirmed by numerous data. Due to the velocity shear instability and the nonlinear evolution, it is precisely these waves that may generate rotational discontinuities [13], frequently observed in the solar wind.

It was suggested [7] that solar rotational discontinuities may create dissipative coronal shock waves, which heat the coronal plasma in low-beta regions [15]. If the magnetic pressure is not very high (close to gas-kinetic pressure), slow MHD shock waves may attenuate due to the Landau damping [2].

Some studies [5] noted the importance of the interface region between the photosphere and corona for the energy transfer.

In the solar atmosphere, some plasma boundaries can be described as solar magnetohydrodynamic (MHD) discontinuities of different types: stationary contact and tangential, nonstationary fast S , slow shock waves S , rotational (Alfvén) discontinuities [4].

On a contact MHD discontinuity C the normal component of the plasma velocity is equal to zero, and the density across the contact discontinuity displays a jump: $v_n = 0$, $\{v\} \neq 0$, $\{\rho\} \neq 0$, $M_A = 0$, where the braces $\{\dots\}$ denote a jump in a parameter across the discontinuity.

MHD fast and slow shock waves S can be described by the relations with the effective Mach number greater than unity:

$$v_n \neq 0, \{\rho\} \neq 0, \mathbf{B}_1, \mathbf{B}_2, \mathbf{n} \in (S), \{\mathbf{B}_t\} \neq 0, M_{eff} > 1$$

Rotational discontinuities A have a normal component of the plasma velocity equal to the Alfvén velocity and the Alfvén–Mach number equal to unity:

$$v_n = \frac{\pm B_n}{\sqrt{4\pi\rho}} \neq 0, \quad M_A = 1$$

The plasma density and the magnetic field increase in the fast shock waves, while the magnetic field strength decreases in the slow shock waves:

$$S_+ : \rho \uparrow, |\mathbf{B}| \uparrow; \quad S_- : \rho \uparrow, |\mathbf{B}| \downarrow$$

Numerous data indicate the existence of magnetohydrodynamic (MHD) directional (tangential and rotational) discontinuities both in the solar corona and in the solar wind plasma [4]. Rotational (or Alfvén) and tangential discontinuities belong to directional discontinuities, where the first type is the surface of a non-stationary discontinuity that propagates through an unperturbed flow, and the second type is

at rest relative to the unperturbed region. For both types of MHD discontinuity the integral relations expressing the conservation laws must be satisfied [12].

It was shown by Grib et al. [8] that MHD tangential discontinuities appearing on the boundaries of coronal holes and inside the streamers may help to slow shock dissipative waves generated as a result of refraction of the solar fast shock waves.

2 Analysis

Let us consider regular interaction between a solar rotational (Alfven) discontinuity A and a contact discontinuity C which simulates a strong perturbation of the medium density. The interaction between A and C is studied in the case of a strong magnetic field $\mathbf{B}$, when the magnetic pressure is higher than the gas-kinetic by a factor of ~ 8 (plasma beta-parameter $\beta \sim 0.13$) and the density drops on C by the factor of ten. These conditions are typical for the solar corona. We will vary the angle ψ of inclination of $\mathbf{B}$ to the front of C and study the flow developed assuming that the interaction between A and C is regular and represents a fan of MHD shock and rarefaction waves and rotational discontinuities divergent from the line of intersection of A and C . The rotational discontinuity A is assumed to be plane-polarized, so that the tangential component of the magnetic field changes its sign and maintains its absolute value. In such a rotational discontinuity the magnetic field rotates around the normal $\mathbf{n}_A$ by the angle 180° and remains in the plane of the flow (the developed flow is plane-polarized). Compared with the non-plane-polarized rotational discontinuities, the plane-polarized rotational discontinuity has the maximum intensity $|A|$ since the magnetic field in it rotates by the maximum possible angle.

The flow developed as a result of the interaction will be sought as the exact solution of the problem of breakdown of the discontinuity between two states: the state downstream of the impinging rotational discontinuity and the state on the other side of the contact discontinuity. We will evaluate the possible role of the solar rotational discontinuity in the perturbation of plasma inside the transition region from the chromosphere to the solar corona.

Suppose that an Alfven wave is propagating on the photospheric side along a magnetic flux tube. Its amplitude will increase on the way to the corona and we will obtain [18]

$$\frac{v_\perp}{v_A} \sim \frac{B_\perp}{B} \sim \rho^{\frac{1}{4}} B^{-1} \cos^{-1} \theta, \quad (1)$$

where v_A is the Alfven velocity, B is the magnetic field strength, ρ is the density, and $v_\perp$ is the amplitude.

For unperturbed conditions, $\rho_{tr}/\rho_{ph} \approx 10^{-5}$ and $B_{tr}/B_{ph} \approx 10^{-2}$, we find that $v_{\perp tr}/\rho_{\perp ph} \approx 20$ and $(v_\perp/v_A)_{tr}/(v_\perp/v_A)_{ph} \approx 5$.

Thus, a linear MHD Alfven wave may become a nonlinear wave or a rotational discontinuity when a magnetic flux tube broadens significantly.

Therefore, it is reasonable to consider an interaction of a rotational discontinuity A travelling from the chromosphere downside with the transition region described in a frame of a contact discontinuity C : $\underline{A} C$; across C , the plasma density varies dramatically and the gas-kinetic pressure is constant. Therefore, we obtain the relations:

$$\{nkT\} = 0, \{n\} \neq 0, v_n \neq 0, B_n \neq 0 \quad (2)$$

where the braces $\{\dots\}$ denote sharp variation across the surface of the discontinuity, v_n is the normal component of the plasma velocity, B_n the normal component of the magnetic field intensity, and n the concentration of plasma.

For the rotational discontinuity, we have the well-known conditions

$$v_n = \frac{\pm B_n}{\sqrt{4\pi\rho}} \neq 0, \quad \{\rho\} \neq 0, \quad B_n \neq 0, \quad v_t = \frac{B_t}{\sqrt{4\pi\rho}} \neq 0 \quad (3)$$

where v_t and B_t are tangential components of the plasma velocity and magnetic field strength, respectively.

The main effect of the rotational discontinuity is the rotation of the tangential component of the magnetic field on the front of a strong discontinuity without variation in its value and the self-consistent variation in the tangential velocity induced by the electrical current flowing inside the rotational discontinuity.

Without considering in detail the generation of a rotational discontinuity due to nonlinear processes in a magnetic field tube, let us analyze the interaction of the discontinuity with the transition region approximated by a contact discontinuity.

It is known that a sharp jump-like increase in the electron temperature and a sharp decrease in the electron concentration exist inside a region with the width of 500 km [10]. Using the data obtained by Gabriel [6], we suppose that the plasma density and temperature follow the relation $n = \text{const}/T$ for the sharp temperature variation from 5×10^4 to 5×10^5 K.

The application of the model of the contact discontinuity to the transition region was reported only by Shibata [18].

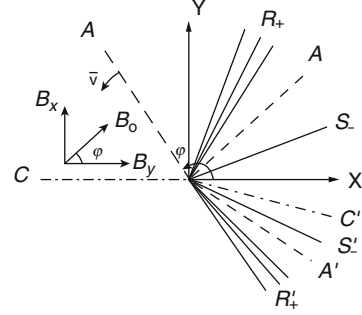
It should be noted that the rotational discontinuity can be adequately characterized by the angle φ at which it impinges on C and by the parameters of the undisturbed region ahead of the discontinuity [1] (Fig. 1). Also, for the rotational discontinuity, the relation between the Mach number and the plasma parameter β is:

$$M_*^2 = \frac{2}{\gamma\beta} \sin(\varphi - \psi) \quad (4)$$

where φ is the angle of incidence of A , ψ is the angle of inclination of the magnetic field to the front of C , β is the ratio of the gas-kinetic and magnetic pressures, and γ is the polytropic index.

The numerical solution of the problem of the interaction of a solar rotational discontinuity A with a contact discontinuity C in the transition region may be found on the basis of the conservation laws (the conditions of dynamic correspondence) similar to considerations carried out by Barmin and Pushkar [1] and Grib and

Fig. 1 A schematic diagram of the peculiar interaction of a solar rotational discontinuity with a contact discontinuity in the transition region



Pushkar [9] when analyzing oblique interactions of MHD strong discontinuities. In this case the oblique interaction between a plane-polarized rotational discontinuity and a contact discontinuity is reduced to a steady-state self-similar MHD problem. The generalized polars are used in accordance with [1], so that all the flow parameters downstream of new MHD shock and/or self-similar rarefaction waves can be found for the given initial conditions.

Thus, for the plasma parameter $\beta < 1$, for the angle $\psi = \mathbf{B}_0 C = 30^\circ$ and for the angle 175° between C and the X axis, we obtain $\overrightarrow{AC} \rightarrow \overleftarrow{R_+ AS_- C'} \overrightarrow{S_- AR_+}$, where R_+ is the fast rarefaction wave, S_- is the slow shock wave, and C' is the modified contact discontinuity.

Therefore, a slow dissipative shock wave directed to the upper solar corona may appear as a result of the interaction of a rotational discontinuity with the transition region. In the case of $\psi = 15^\circ$, we have a refracted fast shock wave.

3 Conclusion

Let us present the results of our consideration for various values of ψ .

1. $\psi = 30^\circ$. A regular solution exists only for $\varphi_A = 5^\circ$, where $\varphi_A = \pi - \varphi$ is the angle between the fronts of C and A . Then the angle θ between the normal $\mathbf{n}_A$ and $\mathbf{B}$ is equal to 55° and the intensity $|A|$ is large. Since $\theta > 0$, the vector of the magnetic field strength rotates counterclockwise in A and the magnetic field lines behind A form an obtuse angle with the X axis directed along the initial front of C (Fig.1). Since in the gas below C the direction of the magnetic field strength does not vary the difference between these directions leads to generation of strong rotational discontinuities in the media both above and below C . In these rotational discontinuities the magnetic field must rotate, so that the conditions of continuity of the magnetic field could be satisfied on the resultant C' .

The flow developed as a result of the $A \rightarrow C$ interaction displays the wave pattern $(R_+ AS_-) \leftarrow C' \rightarrow (S'_- A' R'_+)$, where the symbol $\leftarrow$ denotes the group of waves reflected from C (above the transformed C' in Fig.1) and the symbol

$\rightarrow$ denotes the group of waves refracted through C (below the transformed C in Fig.1). The refracted discontinuities and rarefaction waves have the specific hatching. All the waves are of the large intensity greater than $|A|$. The flow is entrained by A , but the velocity of the resultant stream after the interaction between A and C is lower than that downstream of the initial A . In the fast rarefaction waves R_+ and R'_+ the density and the magnetic field strength $|\mathbf{B}|$ decrease by $\sim 20\%$ and $\sim 10\%$, respectively, while in S_- and S'_- the density increases by the factor of 2.3 and $|H|$ varies only slightly. On the resulting contact discontinuity C' the density and the pressure of the medium increase by the factors of 2 and 4.6, respectively, while the magnetic field strength $|B|$ decreases by $\sim 20\%$. Thus, as a result of the interaction between A and C , the medium is strongly compressed and intense slow MHD shock waves are generated.

2. $\psi = 45^\circ$. A regular solution exists only for $\varphi_A = 5^\circ$. Then the angle θ between the normal $\mathbf{n}_A$ and $\mathbf{B}$ is equal to 40° ; the intensity $|A|$ is large and approximately equal to that for $\psi = 30^\circ$.

The flow developed as a result of the interaction $A \rightarrow C$ has the same wave pattern $(R_+AS_-) \leftarrow C' \rightarrow (S'_-A'R'_+)$, where C' is a modified contact discontinuity similar to that in the case $\psi = 30^\circ$. However, the intensities of almost all the waves are significantly smaller than $|A|$, and only $|R_+| \approx |A|$. The flow is entrained by A and the velocity of the resultant stream after the interaction between A and C is lower than that downstream of the initial A , but the deceleration is weaker, compared with the case $\psi = 30^\circ$. In the fast rarefaction waves R_+ and R'_+ the density decreases by $\sim 10\%$, while in S_- and S'_- it increases by the factor of 1.7. On the resulting contact discontinuity C' the density and the pressure of the medium increase by the factors of 1.7 and 2.55, respectively, while the magnetic field strength $|B|$ decreases by $\sim 9\%$. Thus, as a result of the interaction between A and C , the medium is compressed and fairly intense slow MHD shock waves are generated.

3. $\psi = 60^\circ$. Regular solutions exist for $\varphi_A = 5^\circ$ and 10° , when the angles θ between the normal $\mathbf{n}_A$ and $\mathbf{B}$ are equal to 25° and 20° , respectively. The intensity $|A|$ is significantly smaller than that for $\psi = 30^\circ$.

For $\varphi_A = 5^\circ$, the flow developed as a result of the interaction $A \rightarrow C$ displays the same wave pattern $(R_+AS_-) \leftarrow C' \rightarrow (S'_-A'R'_+)$ as those for $\psi = 30^\circ$ and 45° , while for $\varphi_A = 10^\circ$ the wave flow pattern is radically different: $(S_+S_-) \leftarrow C' \rightarrow (S'_-S'_+)$, and there are no rotational discontinuities. In the latter case, the intensities of S_+ and S'_+ are rather large due to an increase in the magnetic pressure $B^2/8\pi$. The flow is entrained by A and the velocity of the resultant stream after the interaction between A and C is lower than that downstream of the initial A , but the deceleration is weaker compared with the previous values of ψ .

For $\varphi_A = 5^\circ$, the fast rarefaction waves R_+ and R'_+ are weak, whereas for $\varphi_A = 10^\circ$ the fast shock waves S_+ and S'_+ are of small intensity; the density and the pressure increase by the factors of 1.14 and 1.3, respectively. When $\varphi_A = 5^\circ$, the density and the pressure in S_- and S'_- increase by the factors of 1.3 and 1.55, respectively; for $\varphi_A = 10^\circ$, this growth is significantly smaller, by the factors of

1.12 and 1.21. Thus, as a result of the interaction between A and C , the medium is compressed only slightly and slow MHD shock waves of the low intensity are generated.

4. In the case $\psi = 75^\circ$, regular solutions exist for $\varphi_A = 5^\circ$ and 10° , when the angles θ are equal to 10° and 5° , respectively. The intensity $|A|$ is fairly small. For both values of φ_A the wave flow pattern consists of only shock waves: $(S_+ S_-) \leftarrow C' \rightarrow (S'_+ S'_+)$, and contains no rotational discontinuities. It is of interest that the intensities of all the shocks developed in the interaction are smaller than those of A , $|S_-|$, and $|S'_+|$, being substantially smaller than $|S_+|$ and $|S'_+|$.

For $\theta = 10^\circ$, the compression of the medium on C' is very weak and the density and pressure variations do not exceed 6 and 10%, respectively. For $\theta = 5^\circ$, the variations are smaller approximately by the factor of two. Thus, a weak rotational discontinuity interacting with C , in which the density drops by the order of magnitude, generates slow MHD shock waves of a very low intensity and virtually disappears, initiating fast shock waves.

5. When $\psi = 90^\circ$, regular solutions exist for $\varphi_A = 15^\circ$, 10° , and 5° , so that the angles θ are equal to -15° , -10° , and -5° , respectively. Thus, compared with the above cases of the smaller ψ , the angle of inclination of the magnetic field strength to $\mathbf{n}_A$ changes its sign, while the magnetic field strength $\mathbf{B}$ rotates in the opposite direction (clockwise). This leads to restructuring the wave flow pattern. For all φ_A the wave flow patterns are of the same form: $(R_+ S_-) \leftarrow C' \rightarrow (S'_- R'_+)$. In these flows the rarefaction occurring in R_+ is compensated by the compression in S_- . For example, when $\varphi_A = 15^\circ$ ($\theta = -15^\circ$), the density decreases by $\sim 20\%$ in R_+ and then increases by $\sim 20\%$ in S_- . The magnetic field strength varies only slightly. Similarly to the case $\psi = 75^\circ$, $|S_-|$ and $|S'_-|$ are substantially smaller than $|R_+|$ and $|R'_+|$ (fast waves are now rarefaction waves). In the cases of $\varphi_A = 10^\circ$ and 5° ($\theta = -10^\circ$ and -5°) the slow shock waves display a very low intensity. Thus, a weak rotational discontinuity, in which $\mathbf{B}$ rotates clockwise, interacting with C generates slow MHD shock waves of very low intensity and disappears initiating fast rarefaction waves.
6. When $\psi = 105^\circ$, regular solutions exist for $\varphi_A = 15^\circ$, 10° , and 5° , so that the angles θ are equal to -30° , -25° , and -20° , respectively. For all φ_A the wave flow patterns have the same form $(R_+ S_-) \leftarrow C' \rightarrow (S'_- R'_+)$. In these flows the compression in S_- is stronger than the rarefaction occurring in R_+ , so that the density on the resultant C' is higher than that in the initial state. For example, when $\varphi_A = 15^\circ$ ($\theta = -30^\circ$) the density decreases by the factor of 1.59 in R_+ and then increases by the factor of 1.93 in S_- ; therefore, the density on C' increases by the factor of 1.22. In R_+ the magnetic field strength decreases by the factor of 0.65.

The slow shock waves are of hydrodynamic nature and $|\mathbf{B}|$ varies only slightly in them. The gas is accelerated in the initial A and thereafter in R_+ , while in S_- the gas velocity does not vary. Thus, as a result of the interaction between A and C , the medium is compressed and fairly intense slow MHD shock waves are generated.

7. When $\psi = 120^\circ$, regular solutions exist for $\varphi_A = 10^\circ$ and 5° ; the angles θ are equal to -40° and -35° , respectively. The wave flow patterns have the same form as that for $\psi = 90^\circ$ and 105° : $(R_+ S_-) \leftarrow C' \rightarrow (S'_- R'_+)$. The rarefaction occurring in R_+ is weaker than the compression in S_- and on the resultant C' the density is higher than that in the initial state. For example, when $\varphi_A = 10^\circ$ ($\theta = -40^\circ$) the density decreases by the factor of 1.74 in R_+ and then increases by the factor of 2.45 in S_- , so that the density on C' increases by the factor of 1.41. In R_+ the magnetic field strength decreases by the factor of 0.6. The slow shock waves are of hydrodynamic nature and $|\mathbf{B}|$ varies only slightly in them. The gas is accelerated in the initial A and thereafter in R_+ , while in S_- the gas velocity does not vary. Thus, as a result of the interaction between A and C , the medium is compressed, the magnetic field decreases significantly, and very intense slow MHD shock waves are generated.
8. Finally, when $\psi = 135^\circ$, a regular solution exists for $\varphi_A = 5^\circ$ ($\theta = -50^\circ$). The wave flow pattern is $(R_+ S_-) \leftarrow C' \rightarrow (S'_- R'_+)$. The density decreases by the factor of 1.84 in R_+ and then increases by the factor of 2.73 in S_- , so that the density on C' increases by the factor of 1.48.

In R_+ , the magnetic field strength decreases by the factor of 0.57. The slow shock waves are of hydrodynamic nature and $|\mathbf{B}|$ varies by 3–7% in them. The gas is accelerated in the initial A and thereafter in R_+ . Thus, as a result of the interaction between A and C , the medium is compressed, the magnetic field decreases significantly, and very intense slow MHD shock waves are generated.

When impinging on a density jump simulated by a contact discontinuity, the Alfvén discontinuity is transformed into a fan of MHD shock or rarefaction waves and rotational discontinuities divergent from the line of intersection of A and C . It has been found that a regular solution exists only for almost a plane impingement of A on C , when the angle between A and C does not exceed 10 – 15° . When the magnetic field is inclined to C at an acute angle, the magnetic field rotates counterclockwise and the gas entrained by the impinging A begins to move in the direction of the propagation of the latter. When the magnetic field and C form an obtuse angle, the magnetic field rotates clockwise and the gas is accelerated in the direction opposite to the motion of A . When the rotational discontinuity is fairly weak and the angle between its normal and the magnetic field does not exceed 20° , the flow consists of fast shock or rarefaction waves and slow MHD shock waves for the intensities of the slow shocks commensurable with those of the impinging A . At certain angles between the normal to A and the magnetic field there are no rotational discontinuities in the induced flow.

When the density on C drops significantly, as in the case considered, the MHD slow waves generated in the interaction are always shock waves. Despite the fact that slow shock waves are inherent only in MHD flows and their origin is associated with the MHD model of infinitely conducting gas with a frozen-in magnetic field, the slow shocks generated in the considered interaction, as a rule, display gas-dynamic variations of the parameters of the gas (the density and the gas-kinetic pressure increase by the factors of 2.6 and 10, respectively), while the

magnetic field varies only slightly in them. Thus, the MHD slow shock waves are rather strong.

Variations in the magnetic field can be significant only for strong impinging A . At any angles of inclination of the magnetic field to the initial contact discontinuity, this angle varies so that it becomes close to the right angle on the resulting contact discontinuity in the flow developed in the interaction.

It is shown that solar shock waves heating the coronal plasma may appear as a result of the collision of a solar rotational discontinuity with a contact discontinuity in the transition region, providing the source of dissipation of the magnetic field energy and causing spicules and explosive events, such as those described by [14].

In the picture of the upper corona with the flow of the solar wind ejected from it, we may indicate the effect of collisionless wave dissipation through the Landau damping and cyclotron damping associated with wave-particle interactions. In particular, Landau damping is an important mechanism, through which MHD slow shock waves heat and accelerate coronal and space plasma. The solar slow shock wave was observed by Helios-1 at 0.31 a.u. [16].

These main points are consistent with the results obtained by simulations [3] for the conversion of transverse to longitudinal modes of MHD waves and with the observations obtained with SOHO and TRACE spacecrafts and discussed in connection with shock formation and microflares [17].

Note that the solar activity may be indicated by rotational discontinuities appearing also as secondary discontinuities in the case of the oblique interaction between strong MHD discontinuities in the solar plasma.

4 Concluding Remark

Solar slow shock waves are always generated if there is a 10 times plasma density drop in the solar transition region and are sufficiently strong if the falling rotational discontinuity A is not too weak and the angle between the magnetic field and the normal to A exceeds 20° . These waves have the dissipative nature and may damp inside the upper corona due to Landau damping, thereby partially heating the coronal plasma.

5 Summary

1. A solar MHD rotational discontinuity creates a MHD slow shock wave as a result of its interaction with a contact discontinuity inside the transition region.
2. Dissipative in their nature, the slow shock waves heat the coronal plasma and are damped in the upper corona and solar wind due to the Landau damping.
3. Nonreversible slow MHD shock waves refracted to the solar corona can partially heat the corona.

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Complex Magnetic Evolution and Magnetic Helicity in the Solar Atmosphere

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Abstract Solar atmosphere is a single system unified by the presence of large-scale magnetic fields. Topological changes in magnetic fields that occur in one place may have consequences for coronal heating and eruptions for other, even remote locations. Coronal magnetic fields also play role in transport of magnetic helicity from Sun's subphotosphere/upper convection zone to the interplanetary space. We discuss observational evidence pertinent to some aspects of the solar corona being a global interconnected system, i.e., large-scale coronal heating due to new flux emergence, eruption of chromospheric filament resulting from changes in magnetic topology triggered by new flux emergence, sunspots rotation as manifestation of transport of helicity through the photosphere, and potential consequences of redistribution of energy from solar luminosity to the dynamo for solar cycle variations of solar irradiance.

1 Introduction

Solar atmosphere is not simply a collection of individual features. It is a single system unified by the presence of large-scale magnetic fields. As magnetic field emerges through the photosphere into the corona, it expands significantly forming a canopy of relatively strong magnetic fields overlying field-free or weaker field areas. X-ray and EUV images show a “network” of loops interconnecting neighboring and distant active regions even across the solar equator (e.g., [21, 30]). In some respect, at any given moment the solar corona is completely filled by the magnetic fields at different scales and field strengths. Because of $\nabla \times \mathbf{B} = \mathbf{0}$ condition, there are no “free” magnetic polarities: every magnetic “pole” is connected to

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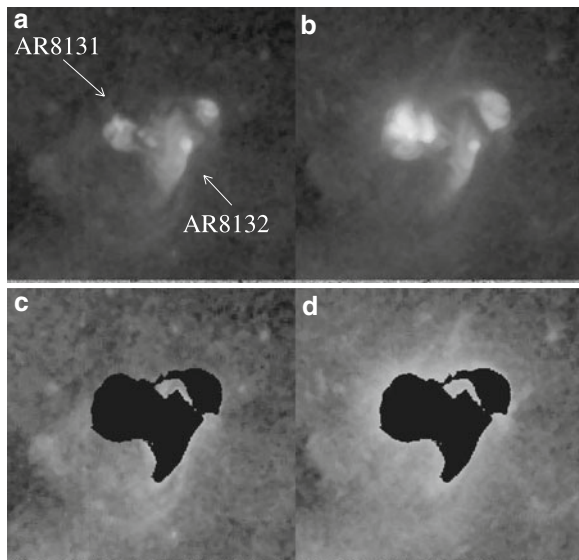
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somewhere else. Still, observations show that shortly after its emergence, new magnetic flux establishes new connections with its neighbours, which implies that other previously existed connections would inevitably change. Thus, a seemingly localized flux emergence may lead to readjusting magnetic topology over much larger area, potentially causing additional heating and/or destabilizing distant coronal flux systems. Due to page limitations, this article is restricted to a discussion of effects of (localized) change in magnetic topology on coronal heating and remote triggering of eruptions. In addition, we also discuss the nature of sunspot rotation as possible indication of transport of helicity from below the photosphere and present consideration of the role of energy diverted to operate the solar dynamo in the total solar irradiance variations.

2 Enhanced Coronal Heating in Response to a Remote Emergence of a New Flux

A simple look at solar images taken in EUV or X-ray wavelength bands leaves no doubt that magnetic fields are present almost everywhere in the corona. By its nature, the coronal fields maintain a constant dynamic equilibrium: changes in magnetic connectivity in one part of the corona, may lead to changes in other parts. When a new magnetic flux emerges through the photosphere, it does not emerge in magnetically empty corona; its magnetic field will interact with pre-existing large-scale field. On smaller spatial scales, such interaction may transfer flux between closed and open fields, leading to formation of coronal jets [17] or coronal bright points [15]. If a new active region develops underneath a large scale magnetic field, the coronal flux system will re-adjust to accommodate the new flux. This re-adjustment may include development of new connections (for example, see Fig. 6 in [14] showing a new loop developing between emerging active region AR9574 and existing active region AR9570, and Fig. 4 in [21] showing development of transequatorial loops between emerging and pre-existing active regions). Moore et al. [18] observed episodic increase in brightness of coronal loops in the vicinity of a new flux emergence site, and have contributed these variations to the reconnection events associated with interaction between the emerging and existing magnetic flux systems. Pevtsov and Acton [22] reported increase in brightness of solar corona over a large fraction of solar disk, associated with the emergence of a single active region. Shibata et al. [28] have suggested that the reconnection between emerging and pre-existing magnetic systems may result in heating of large-scale corona above the emerging flux. More realistic reconnection in 3D geometry also shows formation of area of enhanced heating above the emerging flux [10]. Pevtsov and Kazachenko [23] studied the emerging active region AR8131 and its interaction with the existing region AR8132. They found a significant increase in brightness of solar corona in areas adjacent to the emerging flux even though there was no corresponding change in magnetic flux in the same area in the photosphere. Thus,

Fig. 1 X-ray images of ARs 8131 and 8132 taken by SXT on Yohkoh with Al.1 filter on 11 January 1998, 16:08:01 UT (panels **a** and **c**), and 12 January 1998, 00:12:55 UT (panels **b** and **d**). In *lower panels*, the area of brightest corona is masked to demonstrate the increase of brightness in extended area surrounding emerging active region. Adopted from [23]



for example, Yohkoh soft X-ray images showed about 485% increase in X-ray intensity over the area encompassing two active regions, but excluding the emerging flux region itself (Fig. 1). The change occurred over the 8 h time interval. Change in the photospheric flux over the same area and same time interval was about 8%. Pevtsov and Kazachenko [23] have estimated the total amount of thermal energy deposited in the corona as the result of interaction between the emerging and existing flux systems (but excluding the coronal loops directly associated with the emerging flux) as $2.4\text{--}3.3 \times 10^{30}$ erg. The rate of total thermal energy was found to be nearly constant during the early stages of emergence of the active region, which suggests a continuous heating.

3 Changes in Large-Scale Magnetic Connectivity and Eruption of a Filament from a Distant Location

New flux emergence is often considered as a potential trigger for coronal mass ejections (CMEs), filament eruptions, and flares (e.g., [5]). Past statistical studies found that as much as two-thirds [4] to three-fourths [9] of quiescent filaments were de-stabilized by the birth of a nearby active region. Wang and Sheeley [32] used potential field solar surface (PFSS) model to demonstrate that a newly-emerged magnetic flux may reconnect with the magnetic fields of arcades overlying chromospheric filament. Weakening the arcade may destabilize the filament and lead to its eruption.

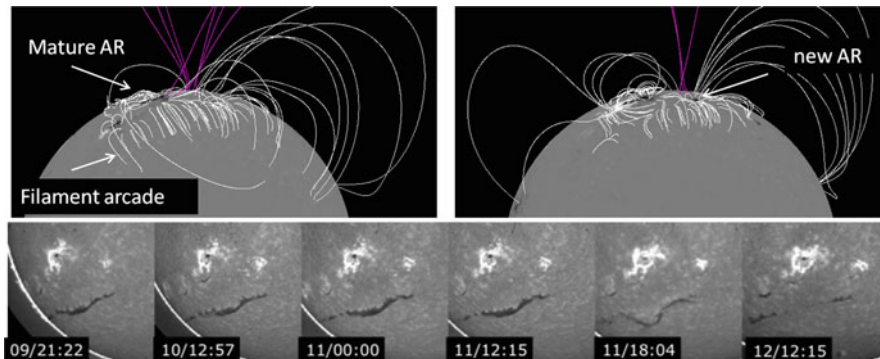


Fig. 2 Magnetic field lines of filament arcade, mature active region (AR10830), and emerging region AR10831 from PFSS model (*upper panel*). *Lower panel* shows quiescent filament on different stages of its evolution on June 9–11, 2003. Rapid rise of central part of filament prior to its eruption can be seen on a panel corresponding to June 11 at 18:04 UT. Magnetic field of emerging AR reconnects with mature AR, which in its turn, “steals” field lines from magnetic arcade above the filament

However, the connectivity change may be indirect and more complex. Balasubramaniam et al. [2] have presented case when a filament eruption was the result of multi-step reconnection. First, a newly-formed active region 10381 had developed new magnetic connectivity with existing region 10380. This new connectivity disrupted the previously existing connections between two fluxes of opposite polarity comprising AR10380, which, in its turn, led to establishing new magnetic connections between AR10830 and neighbouring magnetic fluxes. The latter re-configuration weakened the magnetic arcade above the filament channel next to AR10830 and resulted in a filament eruption. Figure 2 shows overall magnetic topology (based on PFSS model extrapolated from SOHO/MDI magnetograms) and evolution of filament as observed by the ISOON H α telescope [1].

4 Sunspot Rotational Motions as Indication of Helicity Transport

Sunspot rotational motions, when a sunspot exhibits a clockwise/counter-clockwise (CW/CCW) rotation relative to its geometric center, have been first reported more than a century ago [11]. Later studies by several researchers (e.g., [7, 8, 13, 16, 19, 25, 29]) established typical properties of sunspot rotation including their average angular rotation rate ($17 \pm 15^\circ \text{ day}^{-1}$, e.g., [25]). In bipolar active regions, sunspots of leading and following polarity were observed to rotate in phase either in the same or opposite direction (see Fig. 1 in [25]). Some sunspots exhibited change in the direction of their rotation, which was doubted “torsional oscillations of sunspots”

(e.g., [8, 25, 29]). The periods of the torsional oscillations were found to be of order of a few days, although much shorter periods (of a few hours) had also been reported (e.g., [6]). Some sunspots exhibited torsional oscillations with decreasing or increasing amplitude (e.g., [13, 25]). Amplitude of sunspot torsional oscillations was found to show a solar cycle dependency [12]. Torsional oscillations of sunspots were used to estimate the depth to which sunspot rotational motions penetrate below the photosphere (10,000 km— [29]; and 7,500 km— [25], accordingly). Renewed interest in sunspot rotational motions came with high-cadence data from TRACE (e.g., [3]).

It has been suggested that sunspot rotational motions may be an indication of helicity (twist) transport across the photosphere. According to one scenario, sunspot rotation may “pump” helicity to the corona leading to flares and CMEs. Numerical estimates indicate that the amount of helicity transported by a typical rotating sunspot is in agreement with the amount of helicity ejected by CME (e.g., [31]). Kinetic energy of sunspot rotation is about 10^{31} erg [25], which is comparable to energy of a typical flare.

Alternatively, one can hypothesize that sunspot rotation is a response to a removal of magnetic twist (helicity) from the corona by flare or CME. In this latter scenario, subphotospheric portion of magnetic flux tube serves as a reservoir of helicity for the coronal portion. Prior to eruption, both parts are in equilibrium, but removing helicity from the corona disturbs the equilibrium and causes helicity to be transported from below the photosphere until a new equilibrium is established. Such evolution of twist (helicity) is observed in emerging active regions (e.g., [24]). Active regions with strong kinetic helicity below the surface are found to be more flare productive [26].

These two scenarios can be distinguished by the timing of sunspot rotation and flare/CME eruption. If the rotating sunspot twists the coronal magnetic field, the flares should occur at/near the maximum of twist (i.e., when the sunspot rotation is strong). If the rotation starts after a flare/CME eruption, this might indicate that it is a response to helicity removal from the corona (our second scenario). Figure 3 shows that a period with several large flares in NOAA AR9236 is followed by increase in amplitude of sunspot rotation in this active region. This seems to be in agreement with our hypothesis that sunspot rotation is a response of magnetic field on helicity removal from the corona. However, to verify the commonality of such scenario requires study of additional cases of sunspot rotation. It is worth noticing that the direction of sunspot rotation maybe hemisphere dependent. In a recent study, R. Nightingale (private communication) had found that about 70% of rotating sunspots show counter-clockwise rotation in Northern hemisphere. For the Southern hemisphere the asymmetry is weaker, with about 56% of sunspots rotating in clockwise direction. About 15% of sunspots in both hemispheres had shown change in the direction of rotation (earlier referred to as torsional oscillations of sunspots). The hemispheric preference in rotation of sunspots is in agreement with well-known hemispheric helicity rule [20], which provides an indirect support for our second scenario (sunspot rotation as transport of helicity).

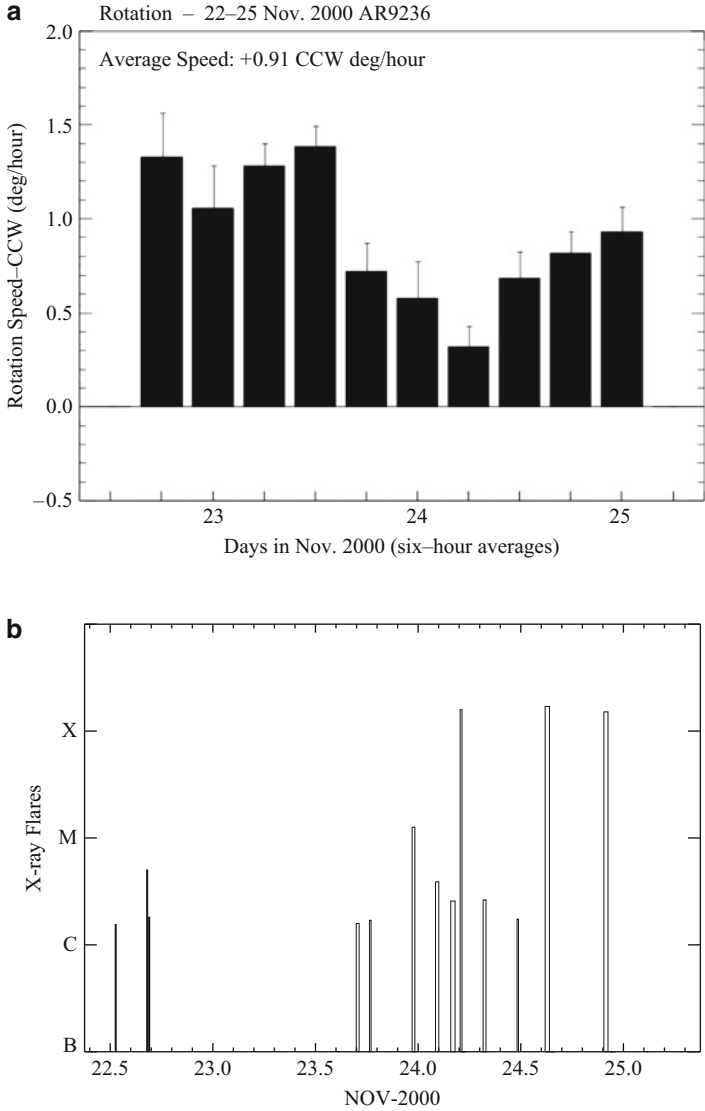


Fig. 3 Rate of sunspot rotation (**a**) and flare activity (**b**) of active region NOAA AR9236. Maximum X-ray flare flux and total duration of flares are shown on panel (**b**). Panel (**a**) is courtesy of R. Nightingale

5 Solar Dynamo and Luminosity

Magnetic field on the Sun is generated by the processes collectively called solar dynamo. In a nutshell, motions of highly conductive plasma in presence of seed magnetic field creates electromotive effect that further amplifies magnetic field.

Energy that drives these flows comes from nuclear reaction in the core of the Sun—same source that powers total solar luminosity. Thus, this energy spent on generation of magnetic field is taken out of energy going to luminosity. If there is no phase-shift between the production of the magnetic field in the convection zone and its emergence through the photosphere, bulk of magnetic field should be generated at/near solar maximum. Therefore, one can expect a dip in solar luminosity when the dynamo operation is at its maximum because more energy is diverted to the dynamo action. How significant is the effect? Rempel [27] has estimated that the total energy of magnetic fields (E_m) stored at the base of the convection zone over 10 year solar cycle is about $E_m \approx 10^{38}$ – 10^{39} erg. In comparison, total thermal energy emitted by Sun over same period is $3.9 \times 10^{33} \text{ erg} \cdot \text{s}^{-1} \times 10 \text{ years} \approx 10^{42}$ erg or $E_L \approx 10^{41}$ erg per year.

Assuming that during solar maximum dynamo produces ten times more magnetic field as compared with solar minimum, one can arrive to two estimates of magnetic energy produced in solar minimum and maximum: $E_M(\text{minimum}) = 1.5 \times 10^{37}$ erg and $E_M(\text{maximum}) = 1.5 \times 10^{38}$ erg. By comparison with total radiative energy of the Sun, magnetic energy is only about 0.03% in solar minimum, and it reaches 0.15% in solar maximum. Although the magnetic energy makes such a small fraction of radiative energy, it is comparable in amplitude with cycle variation of total solar irradiance and may need to be taken into consideration.

Of course, this decrease in solar luminosity due to dynamo action is in anti-phase with cycle variation of TSI. However, if this expected decrease in luminosity is real, potentially it may offset even larger variations in TSI than have been observed.

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On Our Ability to Predict Major Solar Flares

Manolis K. Georgoulis

Abstract We discuss the outstanding problem of solar flare prediction and briefly overview the various methods that have been developed to address it. A class of these methods, relying on the fractal and multifractal nature of solar magnetic fields, are shown to be inadequate for flare prediction. More promise seems delivered by morphological methods applying mostly to the photospheric magnetic configuration of solar active regions but a definitive assessment of their veracity is subject to a number of caveats. Statistical and artificial-intelligence methods are also briefly discussed, together with their possible shortcomings. The central importance of proper validation procedures for any viable method is also highlighted, together with the need for future studies that will finally judge whether practically meaningful flare prediction will ever become possible, if only purely probabilistic.

1 Introduction

Recent major advances in heliophysics have made it clear that solar eruptive phenomena have multiple ramifications to the heliosphere, in general, and the geospace, in particular. Minutes after a solar flare occurs in the earthward solar hemisphere, hard X-ray (and γ -ray, in case of strong flares) photons shower Earth, while the first flare-accelerated energetic particles reach Earth within ~ 20 min, if suitable solar-terrestrial magnetic connections exist (see, e.g., [1] and references therein). Any flare detectors at 1 AU, such as the NOAA/GOES 1–8 Å X-ray solar flux monitors, observe these photons when they are already at Earth. In case a shock-fronted coronal mass ejection (CME) couples with the flare, the first shock-accelerated solar energetic particles (SEP) reach Earth within hours while the interplanetary CME

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(ICME) itself, including the remaining SEP population, encounters geospace within $\sim 1\text{--}4$ days after the occurrence of the eruption (e.g., [2, 3] and references therein).

Contrary to CMEs, therefore, in case of flares there is absolutely no early warning. Both high-energy photons and flare particulate, especially protons, can jeopardize sensitive space-borne equipment [4] and unprotected personnel in extra-vehicular activities [5]. For adequate protection from large solar flares, one ideally needs to know when a flare of a given magnitude will occur *before* it actually occurs.

The aim of this paper will be to briefly summarize the main research avenues to solar flare prediction and to assess which of these methods seem to hold better promise for the future, and why. A tentative classification of flare prediction methods is given in Sect. 2. An assessment of the predictive capability of multiscaling methods is included in Sect. 3, while Sect. 4 discusses morphological prediction methods and highlights one of them for clarity. Section 5 discusses caveats, shortcomings, and the importance of a proper validation for every flare prediction method, while Sect. 6 concludes the discussion, summarizing the main points and findings.

2 A Classification of Solar Flare Prediction Methods

It seems beyond doubt that solar flares of class C1.0 and higher ($1\text{--}8\text{ \AA}$ X-ray flux $\geq 10^{-6}\text{ W m}^{-2}$ at peak) stem from solar *active regions*. Recently, [6, 7] plotted the heliographic positions of $\sim 24,100$ microflares (the weakest of which were much smaller than C-class) observed by the RHESSI mission. These flare locations clearly outline the solar active-region belt. For major M- and X-class flares, therefore, one should use applicable flare prediction tools in active regions.

Although an exhaustive list of flare prediction methods has not yet been compiled, we attempt here a preliminary classification of techniques into four general categories:

1. *Monoscale* (fractal) and *multiscale* (multifractal) methods, applying mainly to solar photospheric magnetic field measurements [8–19].
2. *Morphological* methods, applying mostly to photospheric active-region (vector or not) magnetograms [20–33] but also to coronal emission, mainly X-rays [34, 35].
3. *Combinatorial*, *statistical*, *machine-learning*, and *logistic* methods, applying either to active-region magnetograms [27, 36–44] or to historical records of solar flares [45, 46].
4. *Helioseismic* methods, monitoring sub-surface kinetic helicity and applying exclusively to photospheric magnetograms processed by means of standard helioseismology techniques [47, 48].

In the first category above, the known fractal and multifractal nature of magnetic fields in active regions is utilized. Unfortunately, in the next section we show that multiscaling techniques *cannot* be used for flare prediction, most likely because fractality and multifractality are so widespread in the turbulent solar atmosphere,

quiet and active alone [49], that one cannot use them to actually distinguish between the two [14, 51]. The second category identifies efficient, physics-based parameters and metrics that exemplify flaring active regions. Subject to a number of caveats (Sect. 5), we explain why some of these metrics might indeed advance the flare-prediction cause. Among the morphological methods is the calculation of the *unsigned magnetic flux* $\Phi_{tot} = \int_{\mathcal{S}} |B_n| d\mathcal{S}$ in an active region, where B_n is the normal field component of an active region's magnetogram, integrated over the field of view $\mathcal{S}$ of the magnetogram. The unsigned magnetic flux is a traditional flare prediction method that has known disadvantages and weaknesses. As such, it is viewed as the first criterion of quality of a given flare prediction method—if the method does not score better than the unsigned flux it should not be considered promising. Many methods of the third category rely on combinations of readily calculated parameters exploring the idea that an enhancement of flare-predictive skills may be found into coupling the individual predictive capability of these parameters. In this category we have included statistical processing of historical flare records and artificial-intelligence, machine-learning techniques that aim to automate the flare-prediction process. Despite multiple and multifaceted efforts, current progress is not exactly what one might have wished. The last category includes a single method in which one assesses the sub-surface kinetic helicity by helioseismic vorticity proxies before active regions emerge in the solar atmosphere. In this way one obtains an advance (2–3 day) knowledge of active regions' eruptive potential. This is a very recent technique and is still under scrutiny, but initial validation efforts seem promising.

3 Assessing Multiscaling Flare Prediction Methods

A sizable body of literature discusses fractality and multifractality in the solar atmosphere, from photosphere to corona [50, 51, and references therein]. In brief, though, fractality and multifractality, including intermittency in the spatial distribution of solar magnetic fields and the dynamical (flaring) response of solar active regions (see Fig. 1 for an example) are a manifestation of the self-similarity acting either globally or within certain scale ranges in active regions. Self-similarity is most likely due to self-organization, leading to a hierarchical selection of a few *important* degrees of freedom out of a large array of possible such degrees in a nonlinear dynamical system far from equilibrium [52]. An important sub-class of self-organization is self-organized criticality (SOC) for which there is also a vast body of literature exploiting the concept for active regions and solar flares (see [53] for a detailed account). Quite possibly, self-organized or SOC behavior stem from the widespread turbulence in the solar atmosphere.

From a theoretical viewpoint, multiscaling behavior in a nonlinear dynamical system implies *spontaneity* in the system's dynamical response, namely lack of predictability on whether and when an instability of a given size will be triggered in the system. In this sense, one might expect that multiscaling methods cannot, and should

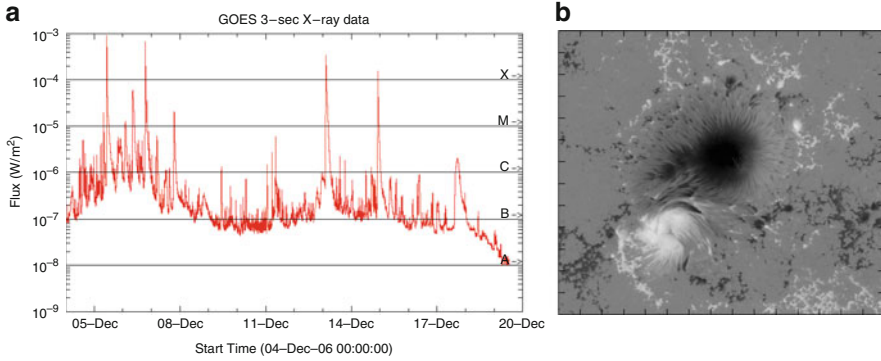


Fig. 1 Example of distinct intermittency in both space and time for solar active region NOAA AR 10930. **(a)** The 1-8 Å X-ray flux of the region between 4 and 20 December 2006. 4 X-class, 3 M-class, and dozens of C-class flares can be seen. This complicated dynamical response stems exclusively from NOAA AR 10930 as it was the only active region present in the earthward solar hemisphere. **(b)** A snapshot of the intermittent line-of-sight magnetic field component of the region recorded by the Solar Optical Telescope's (SOT) Spectropolarimeter (SP) onboard the Hinode spacecraft at 04:30 UT on 2006 December 13. Tic mark separation is 10''

not, be used for flare prediction. Nonetheless, multiple authors [11, 13, 15, 16] have expressed hopes that multiscaling methods indeed hold promise to flare prediction. This author [14] also concluded that for a *single* flare case a certain multifractal parameter exhibited a significant decrease just before the flare that persisted after the event. To further investigate this, but also in order to assess the predictive capability of multiscaling methods, [51] selected three of the reportedly most promising fractal and multifractal parameters and applied them to a comprehensive active-region magnetogram sample. These parameters were

- The *fractal dimension* D_0 , given by the scaling

$$N(\varepsilon) \propto \varepsilon^{-D_0}, \quad (1)$$

where $N(\varepsilon)$ is the number of boxes of variable linear size ε that include part of an active region's boundary.

- The multifractal *structure function spectrum*, $S_q(r)$, given by

$$S_q(r) = \langle |\Phi(\mathbf{x} + \mathbf{r}) - \Phi(\mathbf{x})|^q \rangle, \quad \text{with } S_q(r) \propto r^{\zeta(q)}, \quad (2)$$

where $\Phi(\mathbf{x})$ is the local magnetic flux at a given location $\mathbf{x}$, $\mathbf{r}$ is a given displacement, q is an integer called the *selector* and spatial averaging ($\langle \rangle$) corresponds to a circle of radius r around $\mathbf{x}$.

- The *turbulent power spectrum* $E(k)$, given by the scaling

$$E(k) \propto k^{-\alpha}, \quad (3)$$

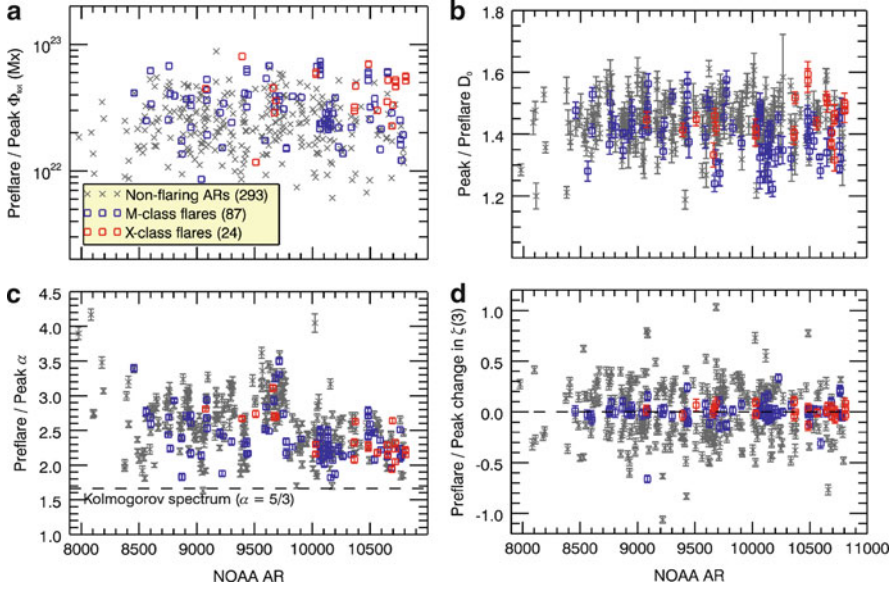


Fig. 2 Preflare vs. peak values of the unsigned magnetic flux (a), fractal dimension (b), turbulent inertial-range power-law index (c), and change of multifractal-spectrum scaling index (d) for 17,733 SoHO/MDI active-region magnetograms (from [51])

where k is the wavenumber and α is the turbulent *inertial-range* exponent. A value of $\alpha \simeq 5/3$ signals a Kolmogorov-type turbulence in the system.

Regarding the above parameters, [13] reported that on a comprehensive set of $\sim 10,000$ active-region magnetograms from the Michelson Doppler Imager (MDI) [54] onboard SoHO, a necessary but not sufficient condition for active regions with major flares was $D_0 \geq 1.20$. Moreover, active regions with larger D_0 were found statistically more flaring. Regarding $\zeta(q)$, [14] found a single case of distinct decrease before and after a flare for $q = 3$. Finally, [11] concluded that stronger departure from $\alpha = 5/3$ implies statistically stronger flares in the active region under study. Moreover, it was reported that α reflects future flare probability at the time the active-region magnetic configuration becomes fully evolved.

In [51] the above claims were tested on a set of 370 active-region magnetogram timeseries (293 non-flaring, 60 M-class flaring, and 17 X-class flaring) captured in $\sim 17,700$ SoHO/MDI magnetograms. The preflare values of the unsigned magnetic flux Φ_{tot} , fractal dimension D_0 , inertial-range exponent α , and the change of the multifractal scaling index $\zeta(3)$ are shown in Fig. 2. For non-flaring active regions we plotted the peak values of the above parameters over the recorded timeseries. At first glance, one notices a considerable mixing of preflaring and peak parameter values that implies an *inability* of these parameters to clearly distinguish flaring from non-flaring regions. Per previous claims, we find that (a) $D_0 \geq 1.2$ for *all* active regions, regardless of flaring, (b) $\alpha > 5/3$ for *all* regions, as well, and (c) $\zeta(3)$ *does*

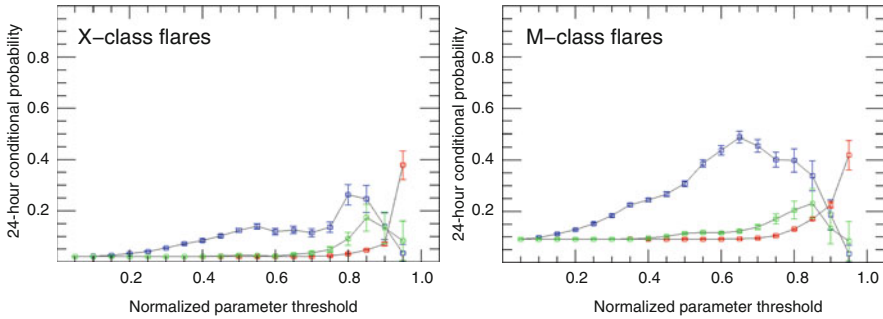


Fig. 3 Conditional flare probabilities over a 24-h window, inferred from the data shown in Fig. 4a–c for X-class (*left*) and M-class (*right*) flares, as a function of different (normalized against their maximum value) parameter thresholds. The probabilities obtained by the unsigned magnetic flux, the turbulent power-law index, and the fractal dimension are shown top-down respectively (from [51])

not show a clear preference for increase or decrease in the course of major flares. The respective predictive probabilities for a 24-h window for M- and X-class flares are shown in Fig. 3. There we notice immediately that the unsigned magnetic flux (blue curve) provides more significant flaring probabilities than the other fractal and multifractal parameters. Therefore, these parameters must be deemed insufficient for flare-prediction purposes.

4 Morphological Flare Prediction Methods and an Example

The driving motivation behind morphological methods of flare prediction is that flaring active regions tend to exhibit some tell-tale features. Most prominent of these is the formation of one or more distinct magnetic polarity inversion lines (PILs) in their photospheric configurations (i.e. in line with the “standard” CSHKP flare model) or of a conspicuous X-ray sigmoid in their corona [34]. Strong, flux-massive photospheric PILs do not simply decay without *at least* one eruption, i.e., major flare coupled by a CME. The only exception in active regions with major flares but without conspicuous PILs is regions with intense flux emergence coupled by strong photospheric flows (e.g., [55]). In these cases, flares (eruptions, in general) are interpreted by means of the interaction between newly emerged and pre-existing opposite-polarity magnetic flux, perhaps realized via one or more flux cancellation episodes [56]. As all active regions by definition show a phase of intense flux emergence, major flaring in some of them most likely relates to the sub-surface magneto-kinematic arrangement of magnetic flux tubes. In our view, flares in these cases are the hardest to predict, unless sub-surface techniques (e.g., [47, 48]) are successfully employed.

The majority of the reportedly promising morphological methods aim to quantify, in different ways, the strength of photospheric PILs (e.g., the WL_{SS} and WL_{SG} parameters [20–22, 25, 32], the R - [29], B_{eff} - [28], or $GWILL$ - [33] parameters. A detailed comparison between them exceeds the scope of this work. As an example, we discuss here the *effective connected magnetic field strength* (B_{eff}) parameter and report on preliminary results of its application to a comprehensive data set of active-region magnetograms. B_{eff} relies on a $N_+ \times N_-$ connectivity matrix Φ_{ij} that includes the magnetic fluxes committed to connections between a positive-polarity photospheric flux partition i ($i \equiv \{1, \dots, N_+\}$) and a negative-polarity partition j ($j \equiv \{1, \dots, N_-\}$). Flux partitioning has been accomplished as first described by [57]. If the respective matrix giving the distances between the flux-weighted centroid positions of each partition is L_{ij} , then B_{eff} is given by

$$B_{eff} = \sum_{i=1}^{N_+} \sum_{j=1}^{N_-} \frac{\Phi_{ij}}{L_{ij}^2}, \quad (4)$$

so it has magnetic-field units. Any connectivity matrix will give rise to an independent B_{eff} -value. In an effort to emphasize photospheric PILs, however, we infer Φ_{ij} by means of a simulated annealing technique [59] that provides a solution which absolutely minimizes a given functional. In this case the functional is

$$\mathcal{R} = \sum \left(\frac{|\mathbf{r}_l - \mathbf{r}_m|}{R_{max}} + \frac{|\Phi_l + \Phi_m|}{|\Phi_l| + |\Phi_m|} \right), \quad (5)$$

where Φ_l , Φ_m are the flux contents of photospheric partitions l and m with centroid positions $\mathbf{r}_l$ and $\mathbf{r}_m$, respectively, R_{max} is a *fixed* maximum length scale of the studied magnetogram (typically its diagonal length) and the sum refers to all possible connectivities. Evidently, Eq. (5) preferably connects opposite-polarity fluxes that are as close to each other as possible, thus minimizing L_{ij} and hence maximizing B_{eff} (Eq. (4)). A magnetogram enclosing an intense PIL will then yield a larger B_{eff} -value than that of a magnetogram lacking a strong PIL. Most flaring active regions, then, should statistically show larger B_{eff} -values than non-flaring ones.

The inference of B_{eff} was criticized by [31]. First, these authors argued that the connectivity matrix Φ_{ij} should depend on the choice of the coordinate system origin since the fixed length R_{max} now used in Eq. (5) was $|\mathbf{r}_l| + |\mathbf{r}_m|$ in the original work of [28]. Although tests have shown only a minimal change of B_{eff} for different origins of the cartesian coordinate system, this drawback is now completely alleviated by the use of R_{max} that is insensitive to the coordinate system. Second, they maintained that simulated annealing yields a connectivity matrix that reflects neither the potential connectivity, that has a clear physical meaning, nor the *true* coronal connectivity. However, the latter is obviously unknown, so one cannot comment on its similarity, or difference thereof, with any given connectivity. The potential connectivity, on the other hand, should not reflect the connectivity of a

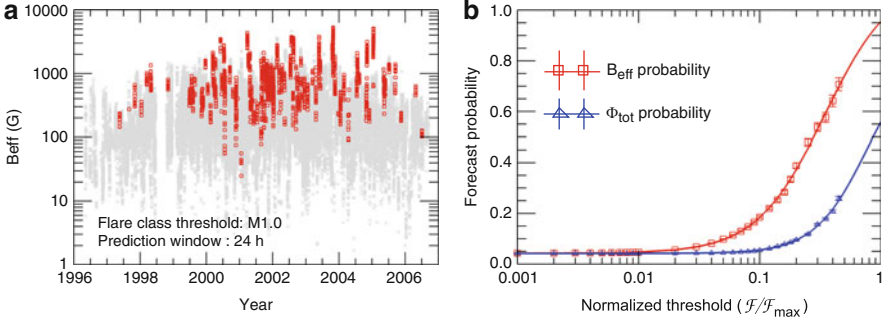


Fig. 4 (a) B_{eff} -values for 55,691 SoHO/MDI magnetograms of solar active regions. The values for active regions that flared with flare-class of at least M1.0 are shown with *squares*; the remaining B_{eff} -values are shown with gray crosses. (b) The respective 24-h conditional probability for B_{eff} (*squares*) and for the unsigned magnetic flux Φ_{tot} (*triangles*), plotted against the normalized B_{eff} - and Φ_{tot} -thresholds over their maximum values $\mathcal{F}_{max}$

strongly non-potential active region. Most importantly, the chosen construction of Φ_{ij} emphasizes PILs so the resulting B_{eff} -values are to be viewed in this respect. For these reasons we continue to use the simulated annealing method to calculate Φ_{ij} , and hence B_{eff} , and conclude that the values of B_{eff} calculated and tested by [31] are very different than the B_{eff} -values that should be considered instead.

In a work currently in preparation, B_{eff} -values have been calculated for a large dataset of 1,416 active-region magnetogram timeseries from SoHO/MDI observed in $\sim 55,700$ magnetograms that were recorded over the entire solar cycle 23. A subset of these regions gave rise to 66 X-, 623 M-, and 3,885 C-class flares. In Fig. 4a we show *preliminary* results of B_{eff} -values for all non-flaring regions plotted against the respective values of regions that gave flares of class at least M1.0 within the next 24 h. Both the flare-class threshold and the prediction window can be adjusted at will. The respective 24-h predictive probability compared to that of the unsigned magnetic flux is given in Fig. 4b. Both curves are fitted by a sigmoidal function of the form

$$P(\mathcal{F}) = A_2 + \frac{A_1 - A_2}{1 + \exp(\log(\mathcal{F}/\mathcal{F}_0)/W)}, \quad (6)$$

where A_1 , A_2 , $\mathcal{F}_0$, W are fitting constants and $\mathcal{F} \equiv \{\Phi_{tot}, B_{eff}\}$. Clearly, the B_{eff} -inferred probabilities are more significant than the Φ_{tot} -inferred ones. Notice that this test was not passed successfully by the multiscaling parameters discussed in Sect. 3. Therefore, the simulated-annealing B_{eff} -values seem to be viable parameters for flare prediction. Additional steps to take and challenges to overcome are briefed in the next section.

5 Challenges and Validation of Flare Prediction Methods

The MDI full-disk magnetograph onboard SoHO has provided us with an unprecedented, homogeneous dataset of line-of-sight active-region magnetograms over the entire solar cycle 23. This wealth of data allows us to build sufficient statistics to construct, evaluate, and validate flare-prediction methods. However, the MDI has now been succeeded by the Helioseismic and Magnetic Imager (HMI) [58] onboard the Solar Dynamics Observatory (SDO) mission. To be able to use MDI-built flare prediction statistics for SDO/HMI data, one has to fully understand how the values of a given parameter calculated on MDI data translate to the respective HMI parameter values. This may turn out to be a highly nontrivial task. First, inter-calibration issues between the two magnetographs need to be understood. Second, a chosen parameter must be sufficiently insensitive to the different spatial resolution of the observing instrument, or at least one should understand how a parameter value behaves for different spatial resolution. In [51] it was found that the multiscale parameters discussed in Sect. 3 show a variability with spatial resolution that cannot be modeled easily; only the fractal dimension D_0 showed a relative insensitivity but this parameter is not adequate for flare prediction (Figs. 2, 3). For B_{eff} and GWILL an effort is currently underway to understand the corrections necessary to translate MDI- to HMI-inferred values [60]. Preliminary results seem encouraging and will be reported in a future publication.

Besides susceptibility to varying spatial resolution, any photospheric flare-prediction parameter is subject to the relatively slow evolution of photospheric magnetic configurations that is due to the line-tied nature of the photospheric magnetic fields (e.g., [62]). Otherwise put, the values of a parameter just before and just after the flare will be fairly similar, thus returning similar flare probabilities. Repeatedly flaring regions may justify this but, clearly, this is not always the case. Combining morphological parameters with a proxy of the flaring history in an active region seems a promising tactic for this problem. This might be done by utilizing, say, mean flaring rates or waiting-time distributions of flares [63, 64].

An equally important task is to validate the results of a given prediction method. A variety of indices and skill scores may be used here: discriminant analysis was applied [31] to compare the prediction capabilities of the unsigned flux, a proxy of the magnetic free energy in active regions [27], the R -parameter [29], and the potential-connectivity B_{eff} —a similar analysis has been applied for the kinetic helicity of sub-photospheric active regions [48]. At least for B_{eff} the fact that potential connectivity was used by [31] dictates a redoing of the analysis with simulated-annealing B_{eff} . Adequate skill scores include the Heidke skill score (see [61] and references therein), the climatological skill score [65], a superposed epoch analysis similar to that of [47], the reliability diagram of, say, [66], or the typical 2×2 contingency table including correct positive prediction, correct negative prediction, missed positives, and false alarms, and the resulting indices. Not all validation methods are equally applicable to all datasets and care must be taken to avoid misleading conclusions. For B_{eff} -values relying on SoHO/MDI data

a complete validation effort is being examined and will be reported in a future publication [67].

6 Conclusion

This work briefly outlines the current state-of-the-art in solar flare prediction. Numerous attempts in this front have relied on the multiscaling nature of solar active-region magnetic fields and subsequent flaring activity. Although we have shown evidence that multiscaling techniques may not be our best asset for flare prediction, studies on the subject have exemplified fractality and multifractality in active regions. This multiscaling behavior, caused by self-organization and giving rise to self-similarity, may be viewed as evidence that solar flare prediction may remain inherently probabilistic.

Among the array of techniques used for flare prediction there are morphological methods that highlight and utilize tell-tale characteristics of flaring active regions. The vast majority of these methods apply to photospheric active-region magnetic fields. There may be promise in several of these methods, provided that (a) statistics gained in datasets of a given magnetograph can be properly translated to the data of another magnetograph, (b) the slow photospheric evolution that does not reflect the actual changes in the course of flares is accounted for, and (c) adequate validation procedures are established.

Combinatorial and machine-learning techniques may also advance the cause of flare prediction. Point taken, care should be exercised to (a) use terms that are independent of each other, rather than terms that can be inferred from each other, to maximize efficiency, and (b) train algorithms to make use of as much physics of the actual system as possible; flare prediction is ultimately a physics-based problem.

Only further studies with homogenous datasets and hands-on experience will judge whether flare prediction is a tenable problem to tackle. Research in this front is definitely worthy, nonetheless, given the profound societal benefits a viable solution will bring, and this explains the ever-increasing independent and synergistic efforts worldwide.

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Chromospheric Evaporation in Solar Flares

Zongjun Ning

Abstract Chromospheric evaporation implies the mass flow from chromosphere to corona along the loop legs in the solar flares. From observations, radio emissions show the high-frequency cutoff and with a drift toward the low frequency, and the coronal lines display a strong blueshift on the Doppler diagram, and hard X-ray emissions tend to rise up the double footpoint sources along the loop legs and finally merge together around the top at the same position as the loop top source. In this paper, we briefly review the radio, EUV and X-ray evidences of the chromospheric evaporation from the documents in recent years.

1 Introduction

Based on the standard flare model, magnetic reconnection is thought to be the primary energy release mechanism that heats the plasma and accelerates the bi-directional particles high in the corona. These particles, guided by magnetic field lines, not only travel to interplanetary space (as traced by radio type III bursts or solar energetic particle events), but also precipitate into the lower corona and upper chromosphere (as traced by reverse-slope type III radio bursts), where they lose energy through Coulomb collisions with the denser medium. This has become known as the “thick-target” model for the hard X-ray (HXR) emission. In the framework of this model, only a small fraction of the energy is lost by radiation. The bulk of the energy will heat the local chromospheric material rapidly (at a rate

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faster than the radiative and conductive cooling rate) up to a temperature of $\sim 10^7$ K. The resulting overpressure drives a mass flow upward along the loop at speeds of a few hundred km s^{-1} , which fills the flaring loops with a hot plasma in a process called “chromospheric evaporation,” giving rise to the gradual evolution of soft X-ray (SXR) emission. Observationally, this process should result in a derivative of the soft X-ray light curve closely matching the hard X-ray or microwave light curves, which is the essence of the Neupert effect.

Observational evidence of chromospheric evaporation has been documented in radio, EUV, and X-ray emission. As mentioned before, dense and cool materials from the chromosphere rise upward along the flaring loops to suppress radio emission at the decimetric wavelength. Aschwanden and Benz [1] analyzed 21 flares and detected a slowly drifting high-frequency cutoff at the frequency range between 1.1 and 3.0 GHz, with drift rates of $-41 \pm 32 \text{ MHz s}^{-1}$, the inferred speed is up to 360 km s^{-1} for the chromospheric evaporation. Using *Solar and Heliospheric Observatory/Coronal Diagnostic Spectrometer* observations at EUV, Czaykowska et al. [3] studied a two-ribbon flare in details and found that the fast upflows are localized at the outer sides of the flare ribbons, while the downflows is at the inside, which is consistent with the expectations from the chromospheric evaporation at the outside of the ribbons while the cooling down at the inside. Using the *RHESSI* data, Liu et al. [5] analyzed the spatial evolution of hard X-ray emission from an M1.7 flare on 2003 November 13. They found that the hard X-ray emission tends to rise above the footpoints and eventually merge into a single source at the same position of the loop-top source. Such the source movement is thought to be the X-ray evidence of the chromospheric evaporation. Ning et al. [8] analyzed an M1.1 flare on 2004 December 1 that showed evidence of chromospheric evaporation at both radio and hard X-ray emission. Using the joint observations of Hinode/EIS and RHESSI, Ning [7] found both EUV and X-ray evidences in the 2007 December 14 flare.

2 Observational Evidences

The evaporation starts the flare onset, and the mass move from the chromosphere upward to the corona. Therefore, the evidences are expectation to be observed at a wide layer from chromosphere to corona. In this section, we will review the evidences of chromospheric evaporation at radio, EUV lines and X-ray emissions respectively.

2.1 Radio Observations

Radio signatures of the evaporations were studied by [1]. The physical basis is that free-free absorption of plasma emission is strongly modified by the steep

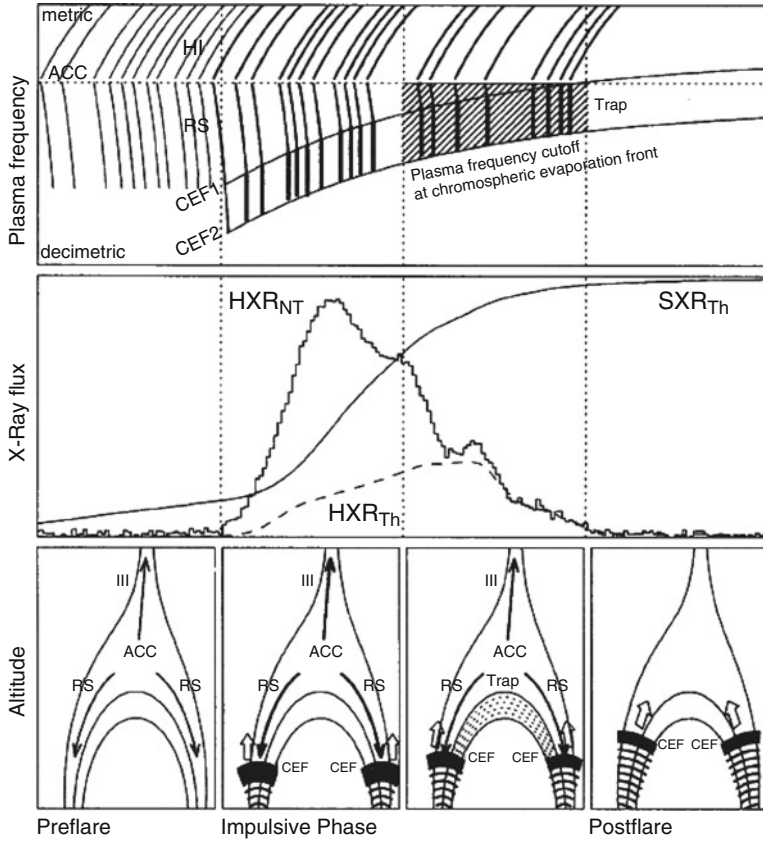


Fig. 1 Flare scenario illustrating the evolution of chromospheric evaporation (CEF), nonthermal hard X-ray emission (HXR_{NT}), thermal hard X-ray emission (HXR_{Th}), soft X-ray emission (SXR_{Th}), and radio emission (type III bursts, reverse-slope [RS] bursts, and plasma emission from a trap) (from [1])

density gradient and the large temperature increase in the upflowing flare plasma. The steplike density increase at the chromospheric evaporation front causes a local discontinuity in the plasma frequency, manifested as almost infinite drift rate in decimetric type III bursts. The large temperature increase of the upflowing plasma considerably reduces the local free-free opacity (due to the $T^{-3/2}$ dependence) and thus enhances the brightness of radio bursts emitted at the local plasma frequency near the chromospheric evaporation front, while a high-frequency cutoff is expected in the high-density regions behind the front, which can be used to infer the velocity of the upflowing plasma. Figure 1 shows the evolution the radio emissions of type III and reverse-slope (RS) bursts associated with HXR and SXR emissions. Figure 2

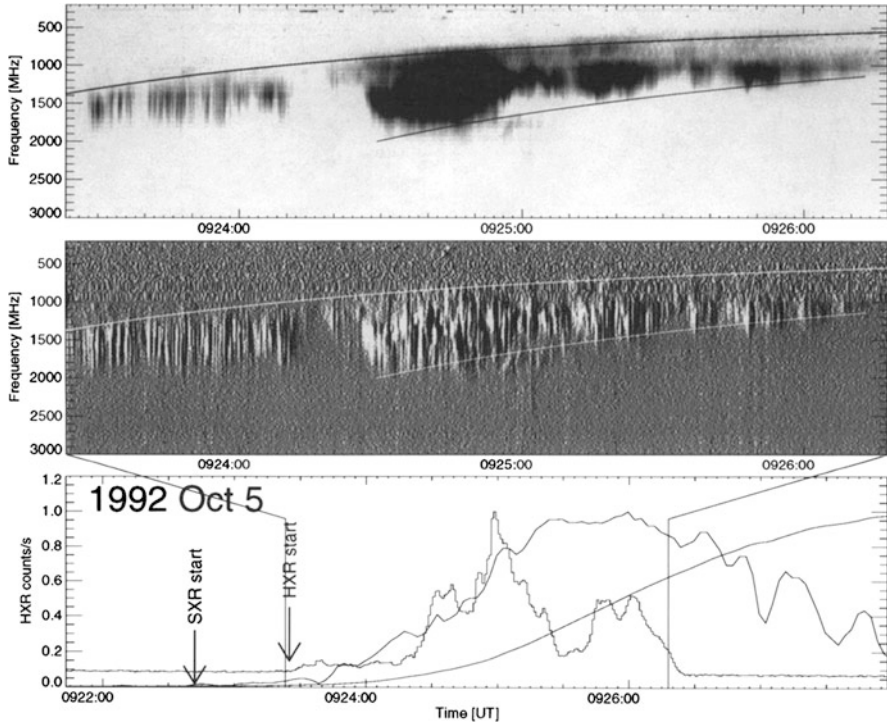


Fig. 2 *Top*: dynamic spectra of the 1992 October 5 flare. *Middle*: gradient-filtered of radio data. *Bottom*: BATSE HXR data (above 25 keV) and GOES SXR at 0.5–4 Å (from [1])

gives an example of flare on 1992 October 5. The Phoenix radio data are shown background-subtracted (top), and gradient-filtered (middle). Both the low- and high-frequency cutoffs show a negative drift rate, -14 MHz s^{-1} and -17 MHz s^{-1} respectively.

2.2 EUV Observations

The chromospheric upflow reaches speeds of a few hundred km s^{-1} from the footpoints upward to the top along the loop legs. Mass flow velocities appear as Doppler shifts in spectral measurements. For example, when the upflow reaches the corona, chromospheric evaporation leads to blueshifted upflow in the hot coronal lines [2, 6]. Due to momentum conservation, some chromospheric mass should move downward during the evaporation. This process is known as “chromospheric

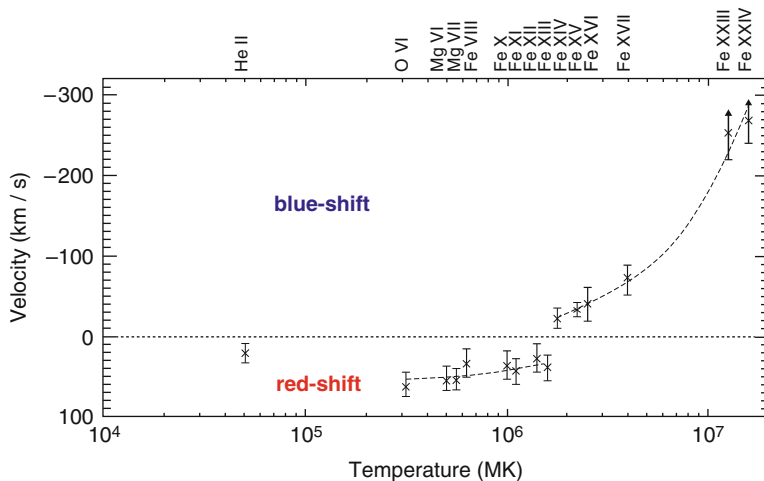


Fig. 3 Hinode/EIS observations showing the plasma velocity from a flare footpoint as a function of temperature for each emission lines. The *dashed lines* represent a weighted least-squares fit to the data points from 0.5 to 1.5 MK and 2.0 to 16 MK (from [6])

condensation.” In this case, the redshift is expected to be observed in the cool chromospheric lines. Such type of evaporation is expected to proceed “explosively.” The first unambiguous observations of explosive chromospheric evaporation were reported by Brosius and Phillips (see [4]). Figure 3 gives the mass velocity at a flare footpoint as a function of the temperature for the emission lines from the chromosphere to corona after the explosive evaporation. However, in some events, all the emission lines formed from chromosphere through transition region to corona appear blueshifted, such type of evaporation is to proceed “gently.” Brosius [2] observed the conversion from explosive to gentle chromospheric evaporation at one of the loop footpoints during an M1.5 flare, as shown in Fig. 4. As the flare evolves, the hot mass confined in lower and inner loops starts to cool down and the mass flows back to the chromosphere, while the electron beam driven evaporation occurs in higher and outer loops. This was reported by [3], who found that EUV lines show blueshifts at the outer edges of flare ribbons, and redshifts at the inner edges, as shown in Fig. 5.

2.3 X-ray Observations

HXR observations can, but do not necessarily, display signatures of evaporation. HXRs will be radiated by decelerating nonthermal electrons independent on whether

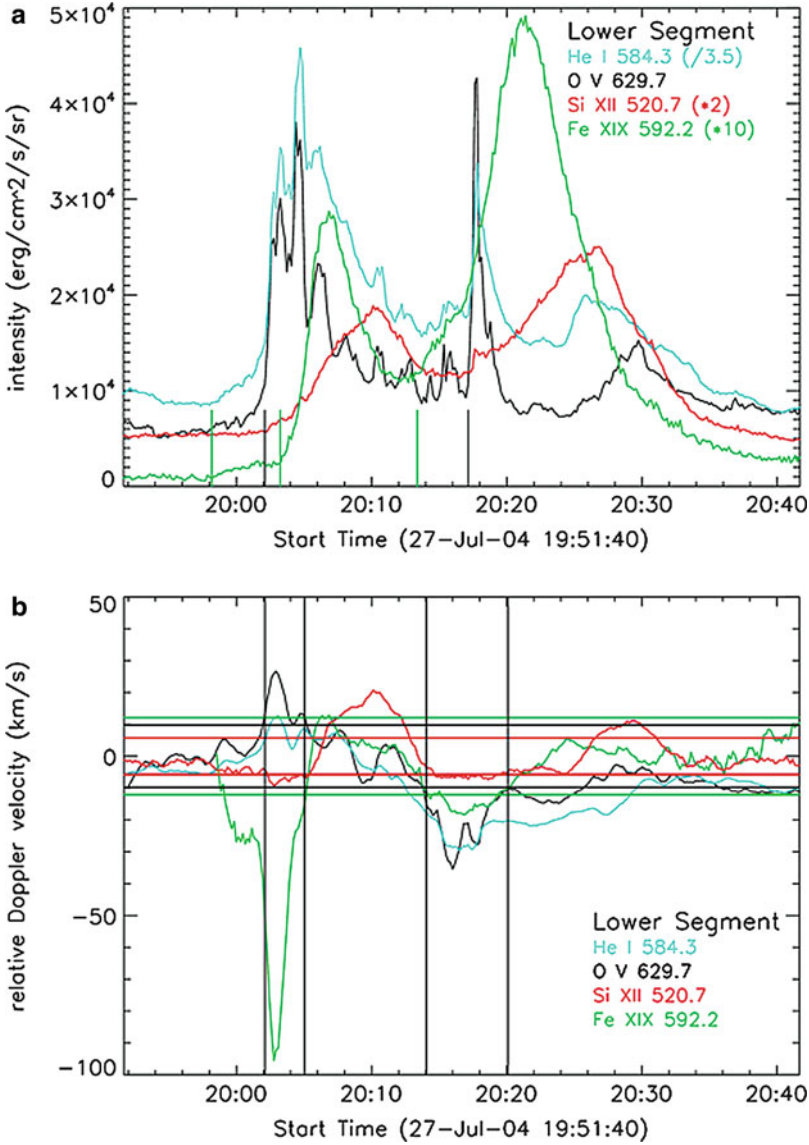


Fig. 4 Light curves (a) and Doppler velocities (b) measured in Fe XIX, Si XII, O V and He I for the SOHO/CDS slit segment. Vertical black lines indicate intervals of explosive (around 20:03 UT) and gentle (around 20:17 UT) chromospheric evaporation (from [2])

they produce chromospheric evaporation or not. It is possible that the nonthermal beams are stopped at progressively greater heights as evaporation fills the loops. From the observations, HXR footpoint sources are expected to be seen moving upward to the loop top along the two legs, as shown in Fig. 6. This is thought to

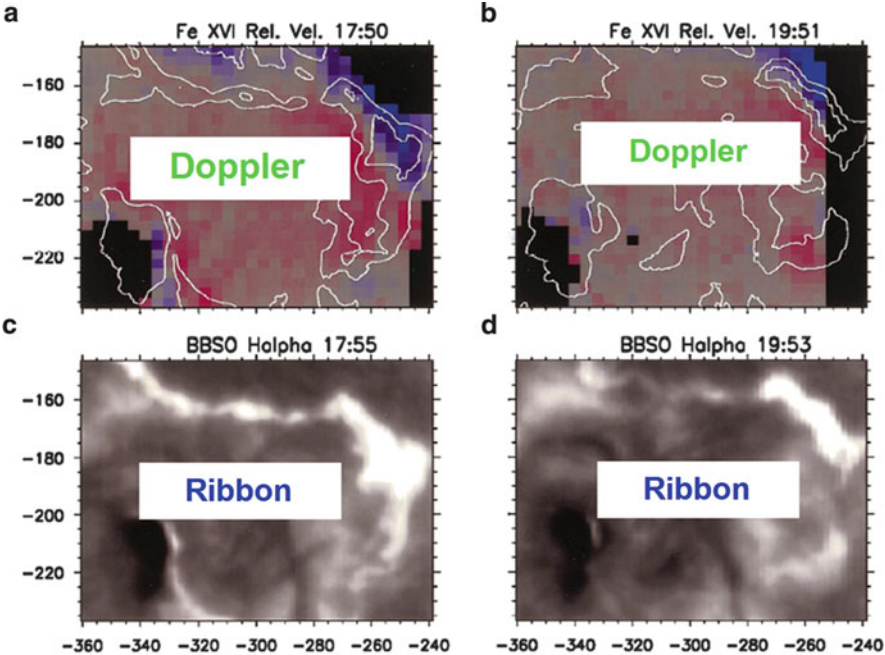


Fig. 5 (a) and (b): SOHO/CDS Fe XVI Dopplergrams with H α contours (white). (c) and (d): BBSO H α images (from [3])

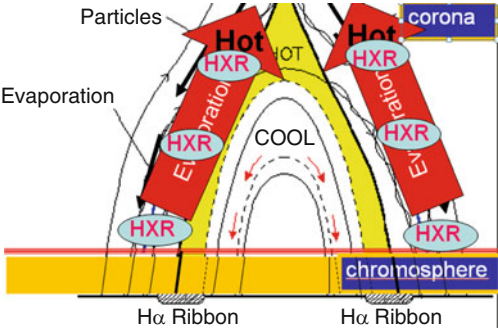


Fig. 6 Flare scenario illustrating the X-ray source movement from the footpoint to loop top caused by the evaporation

be the signature of chromospheric evaporation. *Reuven Ramaty High-Energy Solar Spectroscopic Imager* (RHESSI) observations display evidences of HXR source motions resulting from evaporation [5, 8, 9]. Considering projection effects, HXR emission tends to rise above the footpoints and eventually merge into a single source

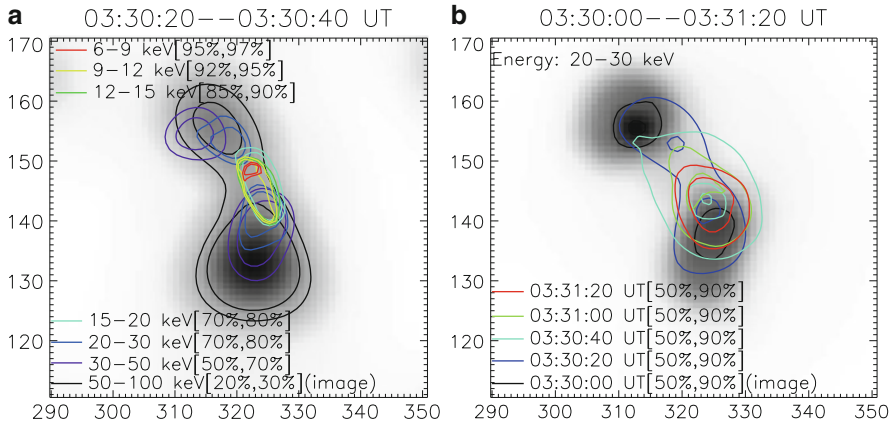


Fig. 7 (a) Energy-dependence X-ray source distribution at a given time of 03:30:20–03:30:40 UT and (b) time-dependence source distribution at a given energy band of 20–30 keV after the evaporation at the 2004 October 30 flare (from [9])

for limb events (e.g. [5]), while HXR emission tends to move closer and merge into a single source coincident with the looptop source for disk-center events [7–10]. Figure 7 gives the HXR source distribution dependence on the time and on the energy after the evaporation. The typical timescale of HXR source motion is a few tens of seconds, and the typical speed is about $200\text{--}400\text{ km s}^{-1}$. Figure 8 shows the source motion at 20–22 keV after the evaporation in detail [7]. A simulation result about the X-ray source motion after the evaporation can be found in paper by [11], as shown in Fig. 9. Observationally, such kind of HXR source motion caused by evaporation can only be seen at 10–30 keV. The plasma in the loops become denser than before due to evaporation. But its density is not high enough to stop the electrons with higher energies and, therefore, these electrons still penetrate deeper into the chromosphere. Observationally, HXR sources at higher energy (i.e. above 50 keV) do not show such motions, and stay at the same place during the flare. Figure 10 shows the joint observations between Hinode/EIS and RHESSI. The energy- and time-dependence distributions of X-ray sources are expected over the flare footpoint with a blue-shift.

3 Summary

Observational evidences of chromospheric evaporation have been documented at EUV, X-ray and radio emissions, respectively. In this paper, we briefly review the observations of this topic at the recent years. Usually, the chromospheric

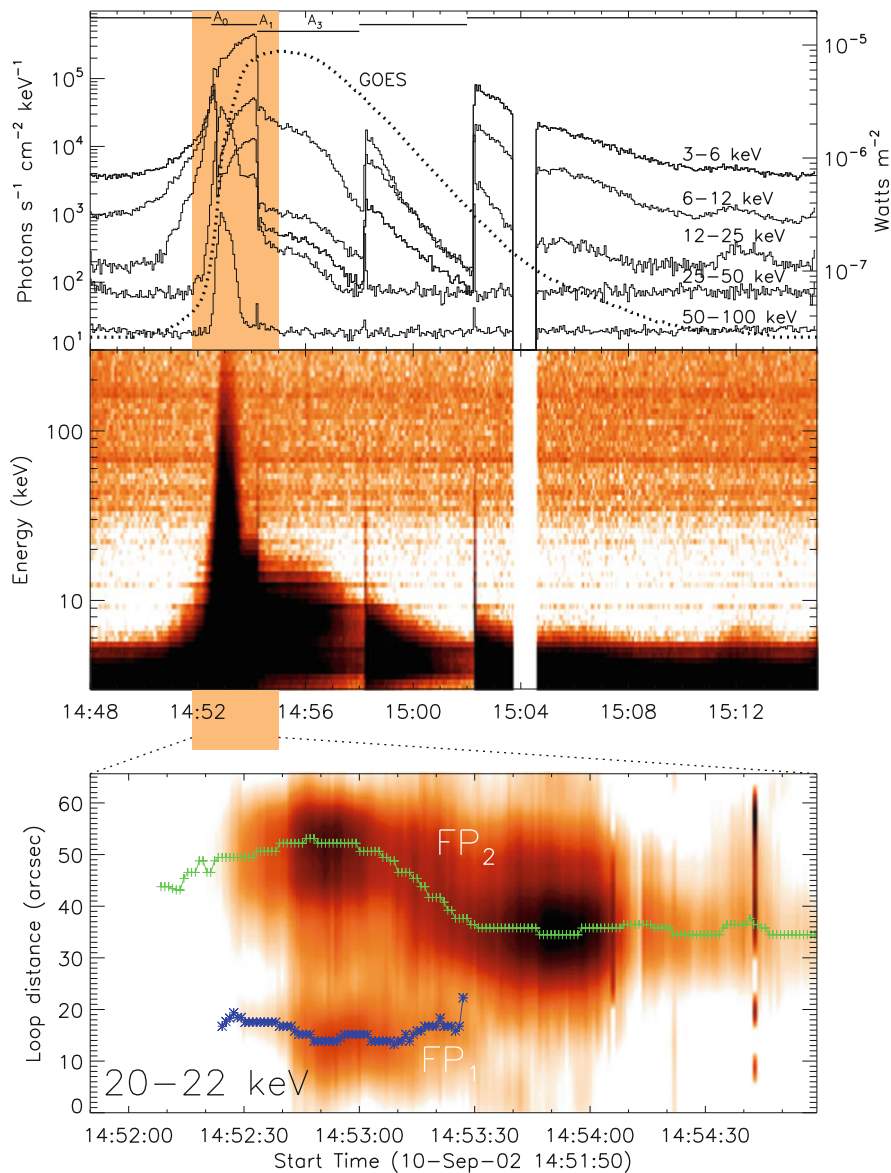


Fig. 8 *Top:* RHESSI (solid) and GOES 1–8 Å (dotted) light curves for the 10 September 2002 flare. *Middle:* RHESSI 3–300 keV X-ray emissions as a function of time, and the intensity jumps are due to RHESSI attenuator effects. *Bottom:* time evolution of the 20–22 keV brightness at both footpoints of FP₁ and FP₂ (from [7])

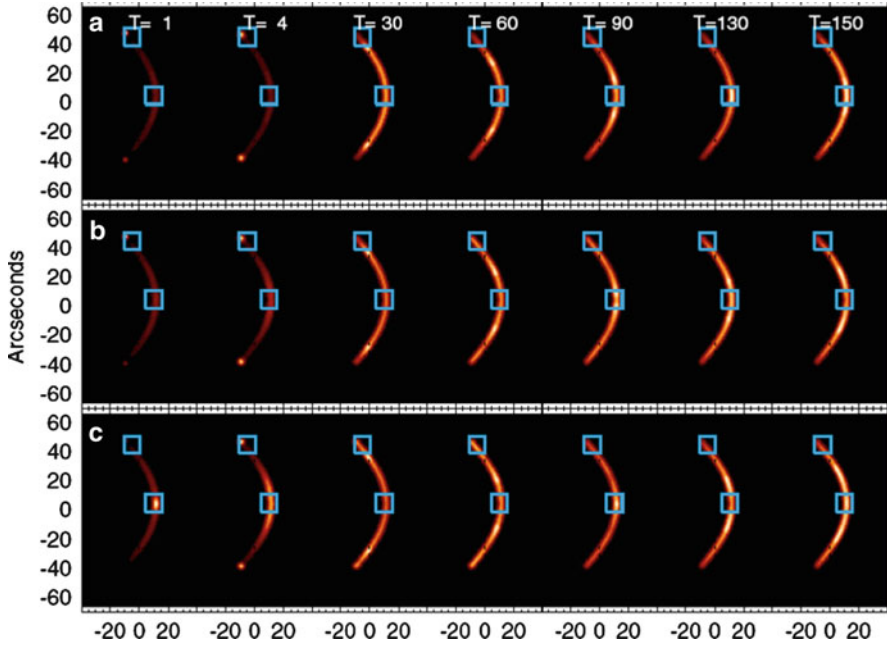


Fig. 9 Simulation of the time evolution of X-ray emissions along the flare loop for three kinds of beam pitch-angle distributions from *top* to *bottom* (from [11])

evaporation shows the mass upflows at the coronal layer, and the downflows at the chromosphere layer. This is the explosive evaporation, while the gentle case displays the mass upflows from chromosphere to corona. The chromospheric evaporation starts the HXR rising phase, and the typical timescale is about 30–80 s from the HXR observations [7]. However, it becomes as longer as about 2 min at the radio observations. The typical velocity is about 200–400 km s⁻¹ from both the EUV coronal (hot) line spectral observation and HXR source motions. However, the velocity is only several tens km s⁻¹ at the lower corona. HXR observations show source motion at not all energy band, but frequently seen at 10–30 keV. HXR emissions do not show the source motions at the energy band below 10 keV and above 30 keV. It is interesting to detect the evaporation velocity as a function of time from the observations.

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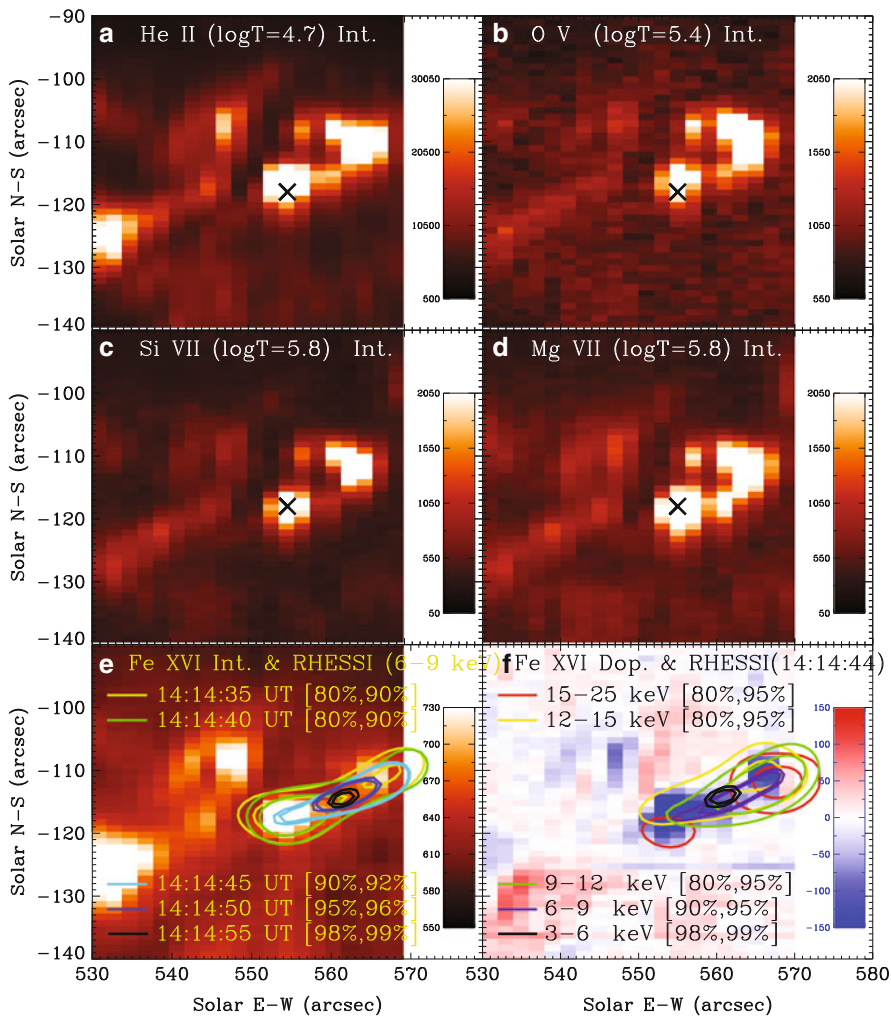


Fig. 10 Joint observation from Hionde/EIS and RHESSI of the 2007 December 14 solar flare. Intensity maps of He II (a), O V (b), Si VII (c), Mg II (d), Fe XVI (e), and Doppler velocity map (f) of Fe XVI. RHESSI X-ray contours are overplotted Fe XVI intensity map with the time-dependent source distribution at a given energy band of 6–9 keV, and Fe XVI velocity map with energy-dependent source distribution at a given time of 14:14:35 UT (from [9])

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Evolutionary of Discontinuous Plasma Flows in the Vicinity of Reconnecting Current Layers

L.S. Ledentsov and B.V. Somov

Abstract The question about the interpretation of numerical experiments on magnetic reconnection in solar flares is considered. A correspondence between the standard classification of magnetohydrodynamic discontinuities and the parameters characterizing the mass flux through a discontinuity and the magnetic field configuration has been established within a classical formulation of the problem on discontinuous magnetohydrodynamic flows. A pictorial graphical representation of the relationship between the angles of the magnetic field vector relative to the normal to the discontinuity plane on both its sides has also been found. The relations between the parameters of a two-dimensional discontinuous flow have the simplest form in a frame of reference where the magnetic field lines ($\mathbf{B}$) are parallel to the matter velocity ($\mathbf{u}$) the de Hoffmann–Teller frame. The question about the transformation of the magnetic field configuration when passing to a “laboratory” frame of reference where $(\mathbf{v} \cdot \mathbf{B}) \neq 0$, i.e., an electric field is present, is considered in this connection. The result is applied to the analytical solution of the problem on the magnetic field structure in the vicinity of a reconnecting current layer obtained previously by Bezrodnykh et al. The regions of nonevolutionary shocks are shown to appear near the endpoints of a current layer with reverse currents.

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1 Introduction

Magnetic reconnection plays a key role in the physics of many nonstationary phenomena in astrophysical and laboratory plasmas [5, 19]. Electric currents and fields are generated in a plasma during the interaction of magnetic fields, which slow down the reconnection process and accumulate the magnetic flux interaction energy called a free magnetic energy. Subsequently, this energy is released during a flare. In solar flares, magnetic reconnection takes place in high-temperature turbulent current layers [13, 15]. It rapidly converts the magnetic energy accumulated before a flare into the plasma particle energy during the flare. In this case, much of the flare energy is liberated in the form of fast, highly beamed plasma flows from the current layer—jets. The latter produce a complex pattern of magnetohydrodynamic (MHD) discontinuous flows outside the current layer that includes shocks of various types. An understanding of this pattern is needed, for example, to explain the observed properties of large eruptive flares, coronal mass ejections, and other geoeffective phenomena on the Sun.

Present-day numerical calculations of magnetic reconnection are performed in a wide variety of physical approximations and formulations of the problem. For example, Bezrodnykh et al. [2] consider a two-dimensional stationary reconnection model in a strong magnetic field that incorporates a thin Syrovatskii-type current layer [20] and four discontinuous MHD flows of finite length attached to its endpoints [9]. The solution of the problem found in analytical form provides the pattern of magnetic field lines near the reconnection region. The calculated field structure generally turns out to be fairly complex. A simple convenient interpretation of the changes in the magnetic field pattern on the discontinuity surface in a “laboratory” frame of reference, i.e., a frame of reference where the electric field inevitably associated with the magnetic reconnection process is nonzero, is needed for its explanation and comparison with the results of numerical simulations, for example, in the approximation of dissipative MHD [3, 4, 21].

The equations of ordinary hydrodynamics are known to have only two discontinuous solutions: a tangential discontinuity and a shock wave. In MHD, the presence of a magnetic field in a plasma leads to the existence of fast and slow Alfvénic shocks and other discontinuous solutions [1, 16, 18]. This introduces a considerable diversity into the possible types of flows in the vicinity of a reconnecting current layer. However, not all of the discontinuous flows are stable against the appearance of new discontinuities. In other words, not all of the discontinuous flows are structurally stable. If one discontinuity under an infinitely small perturbation in the medium immediately splits up into two or more other discontinuities localized on short time scales t in a small region of space, then it is called a nonevolutionary discontinuity. In particular, in contrast to the fast shock “reversing” the magnetic field lines (i.e., changing the sign of the tangential field component), the trans-Alfvénic shock (i.e., the shock for which the normal plasma inflow velocity is higher than the upstream Alfvén velocity, while the normal outflow velocity is lower than the downstream Alfvén velocity) is nonevolutionary [6, 12]. All of this creates

considerable difficulties in interpreting the results of numerical experiments on magnetic reconnection. Of course, in addition, the discontinuities can be destroyed with time in the process of dissipation, but, being within ideal MHD, we will not deal with these types of instability. The goal of this paper is to solve a comparatively simple kinematic problem. It is necessary to establish a correspondence between the well-known standard classification of two-dimensional discontinuous flows in an MHD medium [1, 12, 17] and the pictorial graphical representation of the relationship between the inclination angle of the magnetic field vector to the normal behind the discontinuity plane and the inclination angle to the normal ahead of the discontinuity under magnetic reconnection conditions, i.e., in the case where the electric field is nonzero.

The paper is organized as follows. The system of boundary conditions needed to solve the problem on the identification of the types of shocks on plane discontinuity surfaces is derived in the next section. Subsequently, a general formula that defines the relationship between the angles of the magnetic field vector relative to the normal to the discontinuity surface on both its sides is obtained. We consider the special cases of this formula corresponding to different types of MHD discontinuities; a correspondence between the standard classification of MHD discontinuities and the parameters characterizing the mass flux through a discontinuity (m) and the magnetic field configuration (a) was established. We found a pictorial graphical representation of the relationship between the angles mentioned above, which was applied to the analytical solution of the problem on the structure of a strong magnetic field in the vicinity of a reconnecting current layer [2].

2 Necessary Boundary MHD Equations

We will seek a solution of the formulated problem for an MHD discontinuity, i.e., a plasma region where the density, pressure, velocity, and magnetic field strength of the medium change abruptly at a distance comparable to the particle mean free path. The physical processes inside such a discontinuity are determined by kinematic phenomena in a gas. In contrast, for the hydrodynamic description, this discontinuity will have zero thickness and will occur on some discontinuity surface. In this case, the MHD equations should have discontinuous solutions and certain boundary conditions should be met on the discontinuity surface.

Within the local consideration of discontinuous MHD flows, it is convenient to restrict oneself to plane discontinuity surfaces and steady plasma flows, although both assumptions generally break down under real magnetic reconnection conditions. However, in this paper, we consider the relatively special question about the identification of discontinuities on plane surfaces attached to a reconnecting current layer in a frame of reference where an electric field associated with the reconnection process is present. As far as we know, there is no careful consideration of this question in the literature.

From the viewpoint of an observer moving with the discontinuity plane in a homogeneous stationary plasma, the latter flows into the discontinuity on one side and flows out of it on the other side. In addition, neglecting the role of viscosity, thermal conductivity, and electrical resistivity, we will consider the behavior of the plasma in the approximation of ideal MHD. Thus, we will take the complete system of boundary conditions on the discontinuity plane [6, 18] in a Cartesian coordinate system with the x axis perpendicular to the discontinuity (y, z) plane as the initial one:

$$\{B_x\} = 0, \quad (1)$$

$$\{v_x B_y - v_y B_x\} = 0, \quad (2)$$

$$\{v_x B_z - v_z B_x\} = 0, \quad (3)$$

$$\{\rho v_x\} = 0, \quad (4)$$

$$\left\{ \rho v_x \left(\frac{v^2}{2} + w \right) + \frac{1}{4\pi} (B^2 v_x - (\mathbf{v} \cdot \mathbf{B}) B_x) \right\} = 0, \quad (5)$$

$$\left\{ p + \rho v_x^2 + \frac{B^2}{8\pi} \right\} = 0, \quad (6)$$

$$\left\{ \rho v_x v_y - \frac{1}{4\pi} B_x B_y \right\} = 0, \quad (7)$$

$$\left\{ \rho v_x v_z - \frac{1}{4\pi} B_x B_z \right\} = 0. \quad (8)$$

Here, the curly brackets denote the difference between the values of the quantity contained within the brackets on both sides of the discontinuity plane; for example, Eq. (1) means that the perpendicular magnetic field component is continuous:

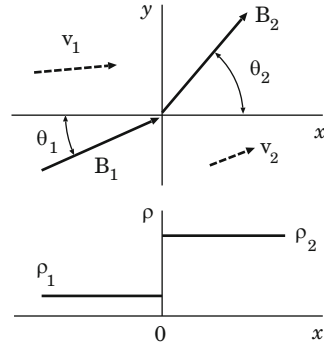
$$\{B_x\} = B_{x2} - B_{x1} = 0.$$

In what follows, the quantities marked by the subscripts “1” and “2” refer to the side from the discontinuity plane corresponding to the plasma inflow and outflow, respectively; below, for brevity, we will say to the left and to the right of the discontinuity, respectively, as is shown in Fig. 1.

Equations (2) and (3) are the ordinary electrodynamic continuity condition for the tangential electric field. The remaining equations express the continuity of the mass, energy, and momentum fluxes.

Of course, apart from these eight equations, it is implied that the equation of state for the plasma is given, for example, in the form of a dependence of the specific (per unit mass) thermal function w on density ρ and pressure p . As is well known, in contrast to the boundary conditions in ordinary hydrodynamics, the system of boundary conditions (1)–(8) does not break up into a set of mutually exclusive groups of equations and, hence, admits continuous transitions between the various

Fig. 1 Changes in magnetic field $\mathbf{B}$, velocity field $\mathbf{v}$, and plasma density ρ at the shock front $x = 0$. Since $E_z \neq 0$ in the coordinate system associated with the current layer, the velocity vectors are not parallel to the magnetic field vectors



types of discontinuous solutions as the plasma flow conditions change continuously. Such transitions occur through some discontinuities that simultaneously satisfy the boundary equations for two adjacent types of discontinuous flows, i.e., they can be attributed to both one type and the other one [17]. The presence of such transitions can be guessed if we pass from the discontinuous solutions to the limit of low-amplitude waves and keep track of their phase velocity diagrams (see, e.g., [14]). In this limit, fast and slow magnetoacoustic waves correspond to oblique shocks and an Alfvénic wave corresponds to a rotational discontinuity.

In the presence of transition solutions, the classification of discontinuities in MHD can only be relative. Indeed, a discontinuity of a given type can continuously pass into a discontinuity of another type as the plasma inflow and magnetic field parameters change gradually. As will be shown in the Sect. 5, the type of discontinuity can change when passing to a different point of the discontinuity surface. In any case, since a smooth transition is possible between discontinuities of various types, the local *external signatures* of the flow near the discontinuity plane are taken as a basis for their classification: the presence or absence of velocity v_x , and magnetic field B_x , components perpendicular to the plane (i.e., normal), the continuity or jump in density ρ . With respect to these signatures, the energy conservation law (5) is an additional condition: at the magnetic field strength, the velocity field, and the density jump found, Eq. (5) defines the jump in pressure p . Thus, bearing our objective of identifying the discontinuities near the reconnection region in mind, we can restrict our analysis to the remaining seven equations: (1)–(4) and (6)–(8). In addition, based on the expected applications of the results to two-dimensional magnetic reconnection models, we will restrict our analysis to plane discontinuous flows or, more specifically, flows in the (x, y) plane. For this purpose, we will rotate the coordinate system about the x axis in such a way that the velocity component $v_z = 0$. Substituting (1) into (8) will then yield the equation

$$\frac{B_x}{4\pi} \{B_z\} = 0.$$

It admits two different solutions.

1. At $B_x = 0$, from Eq. (7) using (4) we obtain

$$\rho v_x \{v_y\} = 0.$$

If $v_x = 0$, i.e., there is no plasma flow through the discontinuity, then $\{v_y\}$ and $\{B_y\}$ are arbitrary quantities, which corresponds to a tangential discontinuity. If, alternatively, $\{v_y\} = 0$, then, as we see from Eq. (3), the tangential component of the magnetic field vector undergoes a change when passing through the discontinuity due to the plasma flow. A perpendicular shock propagates.

2. At $\{B_z\} = 0$, from Eq. (3) we obtain

$$B_z \{v_x\} = 0.$$

Let us first consider the solution $\{v_x\} = 0$. Substituting this condition into Eq. (4) gives a solution with the new condition $\rho_1 = \rho_2$, which corresponds to an Alfvénic shock. The solution $B_z = 0$ leads us to a two-dimensional discontinuity pattern: the velocity and magnetic field vectors lie in the same plane orthogonal to the discontinuity plane.

Thus, four boundary equations remain in the two-dimensional case:

$$\{\rho v_x\} = 0, \quad \left\{ \rho v_x v_y - \frac{B_x B_y}{4\pi} \right\} = 0, \quad (9)$$

$$\{v_x B_y - v_y B_x\} = 0, \quad \left\{ \rho v_x^2 + p + \frac{B_x^2}{8\pi} \right\} = 0.$$

3 The Inclination of Magnetic Field Lines

Denoting the mean of two quantities by $\tilde{f} = (f_1 + f_2)/2$, we will write the system of Eqs. (9) in linear form with respect to the variables $\{v_x\}$, $\{v_y\}$, $\{r\}$ and $\{B_y\}$:

$$\begin{aligned} \{v_x\} - m \{r\} &= 0, & m \{v_y\} - \frac{B_x}{4\pi} \{B_y\} &= 0, \\ m \{v_x\} + \{p\} + \frac{\tilde{B}_y}{4\pi} \{B_y\} &= 0, \\ \tilde{B}_y \{v_x\} - B_x \{v_y\} + m \tilde{r} \{B_y\} &= 0. \end{aligned} \quad (10)$$

where we introduce new variables $r = 1/\rho$ and $m = \rho v_x$. Simultaneously solving the second and fourth equations of system (10) allows m^2 to be found for the case of an Alfvénic shock ($\{v_x\} = 0$) in a two-dimensional flow:

$$m^2 = \frac{B_x^2}{4\pi \tilde{r}}. \quad (11)$$

For nontrivial solutions of the linear system of Eqs. (10) to exist, the determinant composed of its coefficients must be equal to zero:

$$\begin{vmatrix} -1 & 0 & m & 0 \\ 0 & m & 0 & -B_x/4\pi \\ \tilde{B}_y - B_x & 0 & 0 & m\tilde{r} \\ m & 0 & \{p\}/\{r\} & \tilde{B}_y/4\pi \end{vmatrix} = 0$$

Let us expand the determinant:

$$\frac{\{p\}}{\{r\}} \left(\frac{B_x^2}{4\pi} - m r^2 \right) + m^2 \left(\frac{B_x^2}{4\pi} + \frac{\tilde{B}_y^2}{4\pi} - m r^2 \right) = 0.$$

The latter equation imposes constraints on the admissible values of the mass flux m :

$$m^2 = -\frac{\{p\}}{\{r\}} \frac{m^2 - B_x^2/4\pi\tilde{r}}{m^2 - (B_x^2 + \tilde{B}_y^2)/4\pi\tilde{r}}. \quad (12)$$

The quantity m^2 cannot be negative. Consequently, bearing in mind that, in view of Zemplén's theorem in MHD (see, e.g., [6]),

$$\{r\} = \frac{1}{\rho_2} - \frac{1}{\rho_1} = \frac{\rho_1 - \rho_2}{\rho_1 \rho_2} < 0,$$

two inequalities must hold: either

$$m^2 > \frac{B_x^2 + \tilde{B}_y^2}{4\pi\tilde{r}}, \quad (13)$$

or

$$m^2 < \frac{B_x^2}{4\pi\tilde{r}}. \quad (14)$$

As will be shown below, the former and latter cases correspond to fast and slow MHD shocks, respectively, while

$$m^2 = \frac{B_x^2}{4\pi\tilde{r}} \neq 0,$$

at which Eq. (12) cannot be satisfied corresponds, as we established above, to an Alfvénic shock. Solution of the system (10) allows to obtain the relationship between the tangential magnetic field components [7]

$$B_{y2} = \frac{2(B_x^2/4\pi - m^2\tilde{r}) + m^2\{r\}}{2(B_x^2/4\pi - m^2\tilde{r}) - m^2\{r\}} B_{y1}. \quad (15)$$

Let us divide both parts of (15) by B_x to obtain the relation that relates the angles between the magnetic field vector and the normal to the discontinuity surface on both its sides:

$$\tan \theta_2 = \frac{2 (B_x^2 / 4\pi - m^2 \tilde{r}) + m^2 \{r\}}{2 (B_x^2 / 4\pi - m^2 \tilde{r}) - m^2 \{r\}} \tan \theta_1. \quad (16)$$

4 Possible Relations Between the Angles

Based on Eq. (16), let us consider the function $\theta_2 = \arctan(a \cdot \tan \theta_1)$ at various admissible values of the coefficient

$$a = \frac{2 (B_x^2 / 4\pi - m^2 \tilde{r}) + m^2 \{r\}}{2 (B_x^2 / 4\pi - m^2 \tilde{r}) - m^2 \{r\}},$$

which shows by how many times the tangential magnetic field component will change when passing through the discontinuity.

4.1 Alfvenic Shock

In the special case of (11), Eq. (16) gives $a = -1$, $\theta_2 = -\theta_1$. In a plane Alfvenic shock, the tangential magnetic field component reverses its direction.

4.2 Slow Shock

In this case, inequality (14) is valid. Let

$$m^2 = \frac{B_x^2 - b^2}{4\pi \tilde{r}},$$

where the parameter b with the dimensions of a magnetic field will actually define the mass flux. From (16) we find

$$\tan \theta_2 = \frac{b^2 + 1/2 (B_x^2 - b^2) \{r\} / \tilde{r}}{b^2 - 1/2 (B_x^2 - b^2) \{r\} / \tilde{r}} \tan \theta_1.$$

Since

$$\frac{1}{2} \frac{\{r\}}{\tilde{r}} = \frac{r_2 - r_1}{r_2 + r_1} = -\frac{\rho_2 - \rho_1}{\rho_2 + \rho_1},$$

we have

$$a = \frac{b^2 2\rho_2 - B_x^2 (\rho_2 - \rho_1)}{b^2 2\rho_1 + B_x^2 (\rho_2 - \rho_1)}. \quad (17)$$

We will represent the derived coefficient as $(\alpha - \beta)/(\gamma + \beta)$, where

$$\alpha = b^2 2\rho_2, \quad \beta = B_x^2 (\rho_2 - \rho_1), \quad \gamma = b^2 2\rho_1.$$

In this case, the following inequalities hold: $\alpha > \gamma$ and $\alpha, \beta, \gamma > 0$. Consider the possible relations between the quantities α and β .

1. $\alpha > \beta$ ($a > 0$). For the coefficient a to be greater than 1, one of two inequalities, $\alpha > \gamma + 2\beta$ or $b^2 < B_x^2$, must hold, which is impossible, because $m^2 > 0$. Consequently, $0 < a < 1$.
2. $\alpha = \beta$ ($a = 0$). The value of $\theta_2 = 0$ corresponds to $a = 0$ at $\theta_1 \neq 0$, which means the disappearance of the tangential magnetic field component when passing through the discontinuity. Such a shock is called a switch-off wave. For it,

$$b^2 = B_x^2 \frac{\rho_2 - \rho_1}{2\rho_2}$$

and

$$m^2 = \frac{B_x^2}{4\pi} \rho_1.$$

3. $\alpha < \beta$ ($a < 0$). For the coefficient a to be less than -1 , one of two inequalities, $\alpha < -\gamma$ or $\rho_2 < \rho_1$, must hold. This condition cannot be met, because the density of the inflowing matter is less than that of the outflowing one in a shock. Consequently, $-1 < a < 0$. Thus, when passing through a slow shock, the tangential magnetic field component can change its direction, but it cannot increase in magnitude.

4.3 The Absence of a Mass Flux

Substituting $m = 0$ into Eq. (16) yields $a = 1$ and $\theta_2 = \theta_1$. Assuming that $m = 0$ in the second equation of system (10), we obtain the following: either zero B_x , which can correspond to both a tangential discontinuity and a perpendicular shock, or the absence of a change in B_y , which, as we see from the remaining three equations of system (10), is equivalent to the absence of any changes among v_x , p and v_y as well. However, in the latter case, $\{\rho\}$ is not necessarily equal to zero. When passing through the discontinuity, the temperature and density can then change. Such discontinuities are called contact ones.

4.4 Fast Shock

Let now in equality (13) be valid. Just as in the Sect. 4.2, let

$$m^2 = \frac{B_x^2 + \tilde{B}_y^2 + b^2}{4\pi \tilde{r}}.$$

After transformations similar to those made in the Sect. 4.2, we will obtain

$$a = \frac{(\tilde{B}_y^2 + b^2) 2\rho_2 + B_x^2 (\rho_2 - \rho_1)}{(\tilde{B}_y^2 + b^2) 2\rho_1 - B_x^2 (\rho_2 - \rho_1)}. \quad (18)$$

Let us again represent the coefficient a as $(\alpha + \beta)/(\gamma - \beta)$, with $\alpha > \gamma$ and $\alpha, \beta, \gamma > 0$. Consider the possible relations between γ and β .

1. $\gamma > \beta$ ($a > 1$). When passing through a fast shock, the tangential magnetic field component increases in magnitude. Indeed,

$$\tilde{B}_y^2 + b^2 > B_x^2 \frac{\rho_2 - \rho_1}{2\rho_1}.$$

2. $\gamma = \beta$ ($a \rightarrow \infty$). When $a \rightarrow \infty$ absence of a tangential magnetic field component upstream of the shock, its existence downstream of the shock is possible. This regime is called a switch-on wave. In this case,

$$\tilde{B}_y^2 + b^2 = B_x^2 \frac{\rho_2 - \rho_1}{2\rho_1}$$

and

$$m^2 = \frac{B_x^2}{4\pi} \rho_2.$$

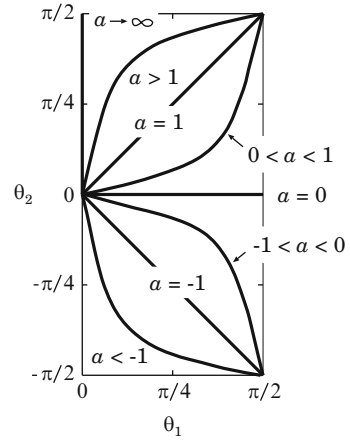
3. $\gamma < \beta$ ($a < -1$). We obtain a trans-Alfvenic discontinuity that increases the tangential magnetic field component and reverse its direction. For it,

$$\tilde{B}_y^2 + b^2 < B_x^2 \frac{\rho_2 - \rho_1}{2\rho_1}.$$

The listed cases are presented in the form of plots of the function $\theta_2 = \arctan(a \cdot \tan \theta_1)$ in Fig. 2 for $\theta_1 \in [0; \pi/2]$. The case of $\theta_1 \in [-\pi/2; 0]$ can be obtained by rotating the coordinate system about the x axis through the angle π . Consider the limiting values of θ_1 .

1. The field perpendicular to the discontinuity surface ($\theta_1 \rightarrow 0$). The angle θ_2 can take on the values $[0; \pi/2)$. The angle θ_2 is zero in any regime when $\theta_1 \rightarrow 0$. In

Fig. 2 Possible relations between the angles θ_1 and θ_2



contrast, the values of $\theta_2 \in (0; \pi/2)$ correspond to a switch-on wave at certain values of ρ_1 , ρ_2 and m^2 .

2. The field parallel to the discontinuity surface ($\theta_1 \rightarrow \pi/2$). Two values of the angle θ_2 are possible. The angle $\theta_2 = \pi/2$ can correspond to a tangential discontinuity perpendicular to the shock or to a contact discontinuity. The angle $\theta_2 = -\pi/2$ is the special case of an Alfvénic wave that reverse the direction of the tangential magnetic field component.

4.5 Passage to the “Laboratory” Frame of Reference

Let the primed coordinate system move relative to the “laboratory” one with a constant velocity $\mathbf{u}$ along the discontinuity plane in such a way that the directions of the z and z' axes coincide. Denote the field components $\mathbf{B}$ parallel and perpendicular to the relative velocity vector $\mathbf{u}$ by $B_{\parallel}$ and $B_{\perp}$, respectively. The Lorentz transformation formulas for the field can then be written

$$\mathbf{B}_{\parallel} = \mathbf{B}'_{\parallel},$$

$$\mathbf{B}_{\perp} = \left(\mathbf{B}'_{\perp} + \frac{1}{c} [\mathbf{u} \times \mathbf{E}'] \right) \left(1 - \frac{u^2}{c^2} \right)^{-1/2}.$$

Let the moving (primed) coordinate system be the de Hoffmann–Teller one, i.e., the coordinate system where $\mathbf{E}' = 0$. Taking the ratio of the parallel field component to the perpendicular one, we will then obtain a transformation of the angles in the moving (primed) and “laboratory” (unprimed) coordinate systems in the

$$\tan \theta' = \tan \theta \sqrt{1 - \frac{u^2}{c^2}}. \quad (19)$$

The transformation of the angles is the same on both sides of the discontinuity surface and, hence, cannot change the relation between the angles θ_1 and θ_2 . Equal angles remain equal; a larger angle will remain larger. Moreover, as we see from Eq. (19), the angles are just retained in the nonrelativistic limit. Thus, our graphical representation of the admissible relations between the angles on the discontinuity surface can be directly applied in the “laboratory” frame of reference.

5 Interpretation of Numerical Simulation Results

Bezrodnykh et al. [2] constructed a two-dimensional analytical model of stationary reconnection in a plasma with a strong magnetic field that included a thin current layer (CL) and four discontinuous MHD flows of finite length R attached to its endpoints. Figure 3 displays the configuration of discontinuities inside which the electric currents flow parallel to the z axis and the coordinate systems used in the calculations.

In specific astrophysical applications, in particular, in solar flares, the model of the so-called “superhot” turbulent current layer [15] should be used to determine the physical parameters of the central part of the reconnection region. The advantage of an analytical model is the possibility to investigate more general patterns that do not depend on the detailed assumptions of a physical reconnection model. Consider some of the properties of discontinuous flows in the vicinity of a current layer predicted by the analytical model of [2]. Let us turn to the inclination angles θ_1 and θ_2 of a magnetic field line to the inward normals (unit vectors directed away from the discontinuity surface on both its sides) at the point of intersection between this line and the discontinuity surface (see Fig. 3). The result of our work is a graphical representation of the possible relationships between the angles (Fig. 2). It can be used to determine the type of discontinuous MHD flows from the known magnetic field configuration in two-dimensional numerical simulations. Indeed, the plots of the function $\theta_2 = \arctan(a \cdot \tan \theta_1)$ in the region $\theta_1 \in (0; \pi/2)$ at various values of the coefficient a have no points of intersection inside the region. For this reason, a (and, hence, the type of discontinuity) is uniquely determined by the angles θ_1 and θ_2 . At $a > 1$, a fast shock takes place. At $a = 1$, there is no flow through the discontinuity surface. If $-1 < a < 1$, then a slow shock takes place. The value of $a = -1$ corresponds to an Alfvénic shock. Finally, $a < -1$ corresponds to a trans-Alfvénic wave.

Each value of the coefficient a is defined by three physical parameters: ρ_1 , ρ_2 and m^2 . Additional relations between these quantities can be established by choosing a certain model, while the interpretation of discontinuities at the limiting values of θ_1 is presented in the preceding section.

Figure 4 taken from [2] shows a gradual change in the angles of the magnetic field vector relative to the normal to the attached discontinuity surface at displacement l along this surface from the point $l = 0$ of attachment to the current layer and to its “free edge” $l = R$, $R = 1$, where the angles θ_1 and θ_2 are equal. In accordance with the graphical representation constructed in the previous section,

Fig. 3 The configuration of electric currents (*thick straight segments*) consists of a current layer (CL) and four segments (discontinuity surfaces) of finite length R attached to its endpoints; L is the current layer half-width

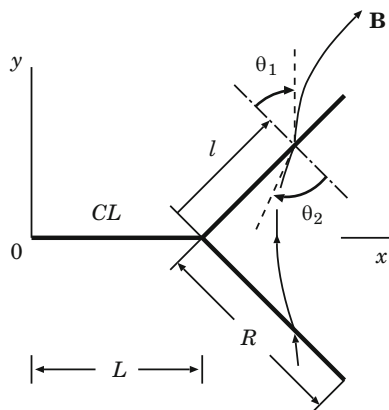
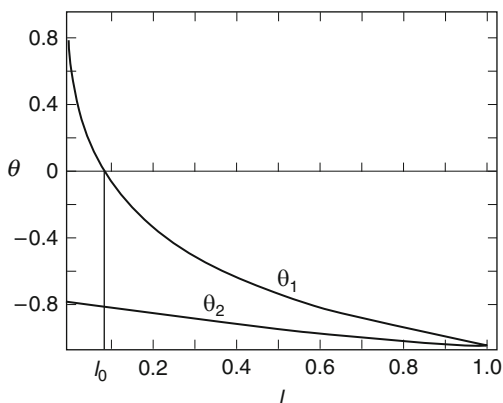


Fig. 4 Distribution of the angles θ_1 and θ_2 at the attached shock. The variable l is the distance measured from the current layer endpoint along the discontinuity surface; l_0 corresponds to the point at which the angle θ_1 changes its sign



the shock is trans-Alfvénic near the point of shock attachment to the current layer: the angles θ_1 and θ_2 are not equal and have opposite signs, with $-\theta_2 > \theta_1$. The situation changes when the angle θ_1 becomes zero at some value of $l = l_0$. At this point, the trans-Alfvénic shock turns into a switch-on one ($\theta_2 \neq 0$). Subsequently, the discontinuous flow passes into the regime of a fast shock. In the reconnection regime considered here, there is no transition from fast shocks (through a parallel shock) to slow ones. The presence of the latter is characteristic of Petschek's flow [10]. Recall, however, that Petschek's flow corresponds to the model problem on the reconnection of oppositely directed magnetic fields that are uniform at great distances from the reconnection region, formally at infinity. In the model of [2], the asymptotics of the magnetic fields at great distances is different—the field becomes hyperbolic, while the reconnection regime is probably typical of the case where the reconnection region (the magnetic field separator in the corona) is located not very high, at a comparatively small distance from the “magnetic obstacle,” i.e., the arcade of flare loops in the corona (see [15], Sect. 7.3). A similar situation arises in a nonstationary MHD model where the point of fast reconnection inside an

infinite current layer lies near a massive, slowly moving “magnetic island.” Trans-Alfvenic shocks are known to be nonevolutionary in both ideal and dissipative MHD [8, 11]. Moreover, they probably also remain nonevolutionary in a weakly collisional magnetized plasma in the vicinity of superhot turbulent current layers in solar flares. Therefore, we assume that the structure of the discontinuous flows near the endpoints of such a layer is complex. It may resemble the quasi-stationary pattern observed in numerical experiments within dissipative MHD (see, e.g., [21]). However, an essentially nonstationary pattern of discontinuous flows attributable to oscillatory disintegration of trans-Alfvenic shocks [8] is also possible. Although the latter possibility seems most likely, in general, this question requires further studies.

6 Conclusions

We established a correspondence between the standard classification of two-dimensional discontinuous flows in an MHD medium and the graphical representation (Fig. 2) of the relationship between the inclination angles of the magnetic field to the normal to the discontinuity surface on both its sides. This allows different areas of the MHD discontinuity surfaces attached to a reconnecting current layer to be identified with different types of MHD shocks. In particular, we found the regions of trans-Alfvenic shocks and switch-on waves, which are known to be nonevolutionary, near the endpoints of a current layer with reverse currents. The question about the pattern of discontinuous flows in the region of nonevolutionarity requires further studies.

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Analytical Models of Generalized Syrovatskii's Current Layer with MHD Shock Waves

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Abstract In the considered models the flow pattern near current layer with attached MHD shock waves is not prescribed but is determined from a self-consistent solution of the MHD problem in the approximation of a strong magnetic field. Generalized analytical solutions are found taking into account the possibility of a current layer rupture in the region of anomalous plasma resistivity. The global structure of the magnetic field in the reconnection region and its local properties near the current layer and attached discontinuities are studied. In the reconnection regime with reverse current in the current layer, the attached discontinuities occur to be trans-Alfvénic shock waves near the current layer edges. Two types of transition of nonevolutionary shocks into evolutionary ones along discontinuous flows are shown to be possible, depending on geometrical model parameters.

1 Introduction

1.1 Classical Models of Magnetic Reconnection

The effect of magnetic reconnection in current layers underlies many nonstationary phenomena in astrophysical and laboratory plasmas [1–3]. In connection with the

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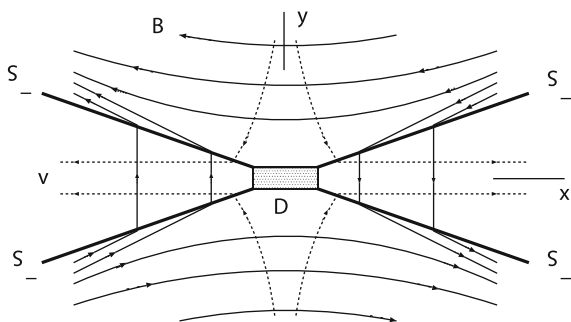
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Fig. 2 Petschek's flow consists of a small diffusion region D and four attached slow MHD shocks S_-



Here A is magnetic potential related to magnetic field B by formulas $A = \text{Im } \Psi$, $B_x - iB_y = \frac{d\Psi}{dz}$. Complexified magnetic potential Ψ is given by the relation

$$\Psi(z) = -\frac{ih}{2} \left[z \sqrt{z^2 - b^2} + (b^2 - 2\varepsilon^2) \ln(z + \sqrt{z^2 - b^2}) \right] + \text{const.}$$

Besides, in works [11, 12] distributions of plasma velocity v and density ρ have been found, which are a solution to the system (1).

Another classical reconnection model is called Petschek's flow [13] and is usually considered as an alternative to the Syrovatskii current layer. In the Petschek model, the reconnection of magnetic field lines takes place in a small diffusion region D , as is shown in Fig. 2, and the main conversion of magnetic energy into plasma thermal and kinetic energy takes place at four attached slow magnetohydrodynamic (MHD) shocks S_- of infinite length. The diffusion region differs significantly in its physical properties from a reconnecting neutral current layer [14]. First, the current density has a minimum at the center of the region D and a maximum at the center of the neutral layer. Second, as the plasma conductivity increases, the width of the region D decreases, while the width of the neutral layer grows.

1.2 New Analytical Models of Magnetic Reconnection

Since establishment of Syrovatskii's model, inspite of its successes, some reasons appeared for improvement this model. First of all, it concerns the existence of MHD shock-waves. Many arguments indicate that such shocks are attached to endpoints of Syrovatskii's current layer. This is confirmed through a numerical solution of the dissipative MHD equations [15, 16], see also [17–21]. Besides, MHD shocks appear in the above mentioned Petschek model [13] of magnetic reconnection.

Based on these arguments Somov and Markovskii [22] suggested a two-dimensional model that is a generalization of the models by Syrovatskii and Petschek. The magnetic field in this model like in Syrovatskii's and Petschek's ones

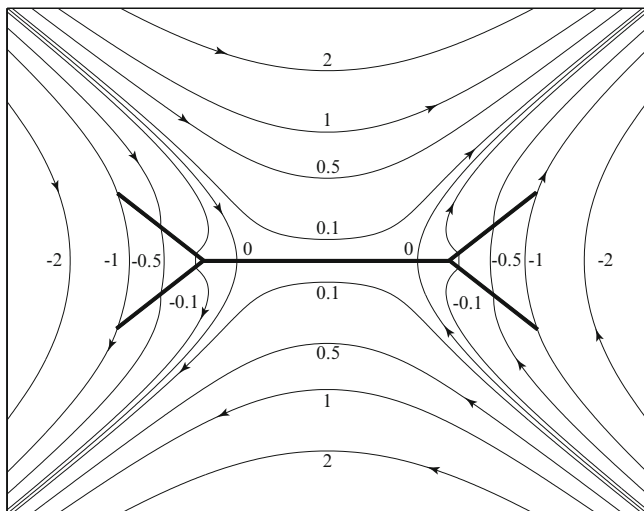


Fig. 3 Current structure (*thick straight-line segments*) and magnetic field lines (*thin curves with the field directions indicated by the arrows*) at $\alpha = 1/4$, $\beta = 1$, and $h = 1$. The field pattern is typical of the general case of physically significant solutions for the problem in the magnetic reconnection regime where reverse currents are present inside the current layer near its endpoints

is assumed to be potential in the exterior of the current configuration that includes Syrovatskii's current layer shown in Fig. 3 in the form of a horizontal cut of length $2b$ and four shocks attached to its endpoints at an angle $\pi\alpha$ as inclined cuts of finite length r . In this model it is supposed that normal magnetic field component at the shocks is prescribed to be β . The model also includes assumption that the magnetic field has linear growth at infinity with coefficient h . The type of shocks is not specified but should be found from a self-consistent solution of the problem.

The model described above is reduced to the Riemann–Hilbert problem [23] in an exterior of the system of cuts indicated in Fig. 3 by the straightline segments. The asymptotics of the solution of the problem found in [24] establishes that at a small “whiskers” length the correction to the field without whiskers, i.e., to the solution by Syrovatskii [7] without singularities, is of the order of $\sqrt{r/b}$. The complete solution of the problem and its interpretation are given in [25] and [26–28]. As an example, Fig. 3 represents the magnetic field pattern obtained from an analytical solution that corresponds to the parameters $\alpha = 1/4$, $\beta = 1$, and $h = 1$. This pattern is typical for the reconnection regime in which reverse currents are produced near the current layer endpoints (see [29]).

Another generalization of Syrovatskii's model is needed, because the current layer can disrupt into parallel current ribbons. This disrupt of a thin current layer can emerge from a tearing instability [30] or when a region of higher electrical resistivity, for example, anomalous resistivity due to the excitation of plasma turbulence, appears [31]. In [32] a simple analytical model has been suggested of a disrupting layer with an infinite width. The force of magnetic tensions proportional

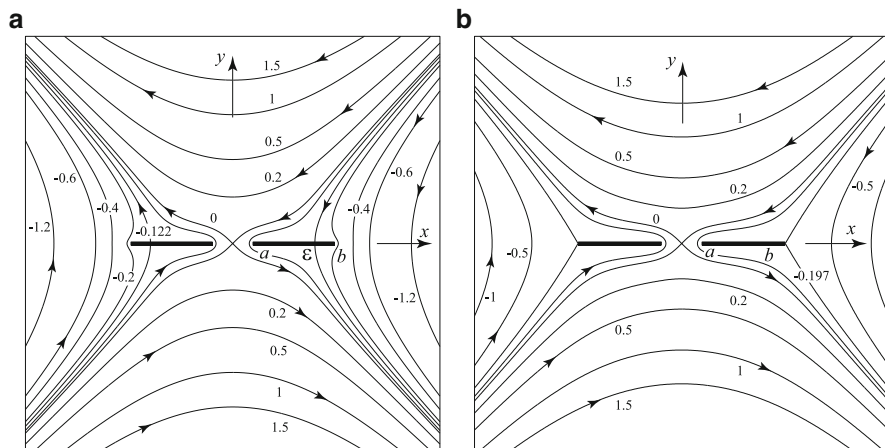


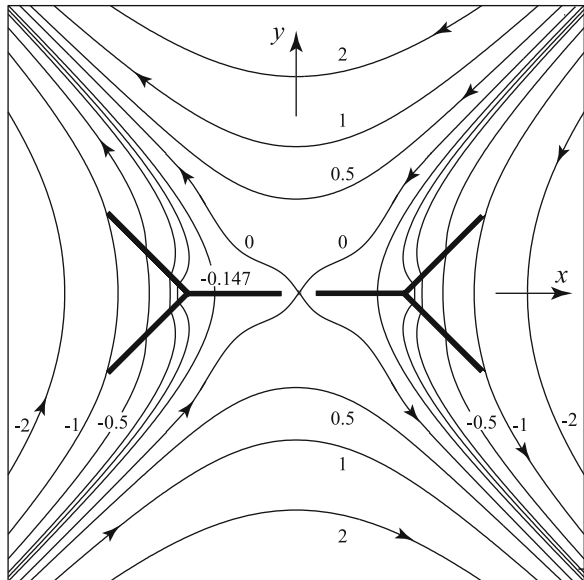
Fig. 4 Pattern of magnetic field lines in the vicinity of a current layer with a discontinuity but without attached shocks; **(a)** there are no reverse currents, **(b)** the total current in the layer is zero

to the magnitude of the discontinuity and tending to increase it acts on the edges of the discontinuity in the layer. A strong electric field capable of accelerating charged particles to high energies under astrophysical conditions (e.g., in solar flares) is induced inside the discontinuity [33].

To study the magnetic field structure in the exterior of a disrupting current layer with a finite width, in [28] two models have been suggested in which the current layer is depicted by two horizontal cuts in the plane. The first model does not contain MHD shocks (see Fig. 4). For the second model, the current configuration includes four attached MHD shocks shown in Fig. 5 in the form of cuts inclined at an angle $\pi\alpha$. At both cuts corresponding to the current layer, the magnetic field has no component normal to it, while at the cuts corresponding to the shocks, like in the model [22], the normal field component is specified by a constant β . The magnetic field has a linear growth at infinity and is limited in magnitude in the finite part of the plane, except for the current layer endpoints free from shocks, where it can have a power-law growth of the order of $-1/2$. This restriction of the mathematical formulation of the problem is not a consequence of any physical peculiarities. As in the problem of Syrovatskii's current layer, it stems only from the fact that the current layer thickness was formally taken to be zero. In the special case of Syrovatskii's current layer without reverse currents, this restriction is absent altogether, because the magnetic field vanishes at the current layer endpoints.

The present paper contains the following results. In Sect. 2 the general approach is outlined to finding the magnetic field in discussed models. In Sect. 3 we give solution to the model with continuous current layer with attached shocks. The Sect. 4 is devoted to finding magnetic field in the vicinity of a disrupting current layer when its endpoints are free from shocks or when two pairs of shocks are attached.

Fig. 5 Pattern of magnetic field lines in the vicinity of a disrupting current layer with attached shocks



2 General Approach to Finding the Magnetic Field

In the reconnection models we consider plane potential magnetic field $\mathbf{B} = (B_x, B_y, 0)$ in domain g , which is an exterior of current configuration depicted in the form of system of cuts in the complex $z = x + iy$ plane (Figs.3–5). This magnetic field is convenient to be written in complex form

$$B(z) = B_x(x, y) + iB_y(x, y).$$

The condition for the field component B_n normal to the line Γ depicting the current configuration mentioned in the Introduction is assumed to be prescribed in the models presented above. This field component is equal to zero at the current layer and to a constant β at the cuts in the complex plane corresponding to the shocks. It is easy to verify that B_n can be expressed in terms of $B(z)$ according to the formula

$$B_n = \operatorname{Re} [\nu(z) \overline{B(z)}], \quad (2)$$

where $\nu(z)$ is a complex unit normal, Re denotes the real part of the quantity in square brackets, and the overbar denotes complex conjugation.

The linear growth condition is set for the function $B(z)$ at infinity, which is reflected by the following asymptotics: $B(x, y) \sim ih\bar{z}$, $z \rightarrow \infty$, where h is a fixed real constant, the magnetic field gradient. This behavior of the field corresponds to the pattern of lines observed far from the hyperbolic null point.

To find the magnetic field B , it is convenient to use a complex conjugate function

$$\mathcal{F}(z) = u(x, y) + i v(x, y) = \overline{B}(z), \quad z \in g,$$

because it follows from the field potentiality that the function $\mathcal{F}(z)$ defined in this way is an analytic function of the complex variable z in domain g .

Substituting $\mathcal{F}$ for B in equality (2) and taking into account the above remark about the magnetic field component normal to Γ , we arrive to the Riemann–Hilbert problem for the analytic function $\mathcal{F}(z)$,

$$\operatorname{Re} [\nu(z)\mathcal{F}(z)] = c(z) \quad \text{on} \quad \Gamma, \quad (3)$$

where $c(z)$ is known. The equality $c(z) = 0$ holds at the current layer points, i.e., the boundary condition is homogeneous, and, if the model includes shocks, then the equality $c(z) = \beta$, where β is a fixed constant value of the magnetic field component B_n (model parameter), holds at the points of Γ depicting the shocks.

The following condition for a linear growth of the function $\mathcal{F}(z)$ at infinity follows from asymptotics (30) for the field B and definition (31) of $\mathcal{F}(z)$: $\mathcal{F}(z) \sim -ihz$, $z \rightarrow \infty$. It is assumed in the models being studied that the field B is symmetric, with its component B_x being even relative to the y axis and odd relative to the x axis and with the parity properties of its component B_y being opposite. These conditions can be written as $B(z) = \overline{B}(-\bar{z})$, $B(z) = -\overline{B}(\bar{z})$. It can be shown that when finding the magnetic field in the vicinity of a current layer without shocks, we can reduce problem (3) for the domain g to a similar problem in a canonical domain, the upper half-plane $H^+ = \{z: \operatorname{Im} z > 0\}$, by taking into account the mentioned above symmetry conditions. Thus, using formulas [23], the solution of the problem can be written directly via Cauchy-type integrals, which can be calculated in terms of elementary functions in the cases under consideration.

When finding the field in the models including shocks attached to the current layer, we can also simplify the original problem (3) by taking into account the field symmetry. In this case, however, it is reduced to a similar Riemann–Hilbert problem in the complex domain G —a quarter of the plane with a cut (see Fig. 6a). Therefore, to find the solution $\mathcal{F}(z)$, we will apply a conformal mapping $\zeta = \Phi(z)$ of the domain G onto the upper half-plane H^+ and pass to a similar problem in H^+ (Fig. 6a, b); so, we will obtain the solution $\mathcal{P}(\zeta)$ of the latter and then find the function $\mathcal{F}$ by substituting $\zeta = \Phi(z)$ into $\mathcal{P}(\zeta)$, i.e., we will write $\mathcal{F}$ as a superposition

$$\mathcal{F}(z) = \mathcal{P}[\Phi(z)]. \quad (4)$$

The mapping $\Phi(z)$ was constructed in [25]. Below, we provide the basic information about this mapping needed to analyze the models considered here.

Since the domain G is an (infinite) pentagon, the mapping $\Phi^{-1}(\zeta)$ inverse to $\Phi(z)$ can be written as the Christoffel–Schwarz integral [34]:

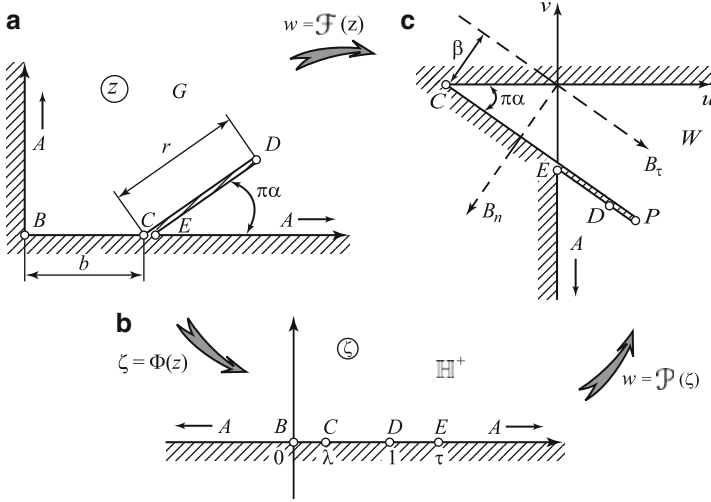


Fig. 6 Scheme for solving the Riemann–Hilbert boundary-value problem: (a) initial domain G (the first quadrant of the reconnection region) in the complex z plane, (b) upper half-plane, and (c) magnetic field hodograph domain

$$\Phi^{-1}(\zeta) = \mathcal{K} \int_0^\zeta t^{-1/2} (t - \lambda)^{-\alpha} (t - 1) (t - \tau)^{\alpha-1} dt. \quad (5)$$

where α is the inclination angle of the cut (CDE) divided by π (see Fig. 6a).

The points $\zeta = \infty$, $\zeta = 0$, and $\zeta = 1$ were chosen to be the preimages of vertices A , B , and D , respectively, while the preimages λ and τ of vertices C and E , along with the integrand factor $\mathcal{K}$, are to be found. These unknown quantities satisfy a system of nonlinear transcendental equations [34]. No analytical solution of such systems is known and the complex relation between the sought-for parameters and the geometrical parameters of the domain G can be established only numerically, using Newton's method. Note that the conformal mapping parameters λ , τ , and $\mathcal{K}$ appear in the derived analytical expressions for the function $\mathcal{P}$ from representation (4) for the conjugate field $\mathcal{F}$ (see Sects. 3, 4). Inversion of Φ^{-1} was made in [25].

To study the magnetic field, it is convenient to use the vector potential $\mathbf{A}$ in terms of which it is expressed according to the formula $\mathbf{B} = \text{curl } \mathbf{A}$. Since the field $\mathbf{B}$ under consideration is plane, only the third component of the vector $\mathbf{A}$ is nonzero, i.e., $\mathbf{A} = (0, 0, A)$. This component is reconstructed via the solution $\mathcal{F} = \overline{B}$ of the Riemann–Hilbert problem according to the formulas

$$A(x, y) = \text{Im} \Psi(z), \quad \Psi(z) = \int_0^z \mathcal{F}(t) dt.$$

We will call the function $\Psi(z)$ a complex potential of the field. It is easy to verify that the magnetic field B is directed tangentially to the isolines of the function A .

Therefore, we will represent it as a family of lines $A(x, y) = \text{const}$. In Figs. 1, 3–5, the values of this constant are given near the corresponding magnetic field lines.

The total current $\mathcal{J}$ through the current configuration is proportional to the circulation of the magnetic field B along an arbitrary closed contour enclosing this configuration. In terms of the analytic function $\mathcal{F}(z)$, we arrive to the following formula for the total current:

$$\mathcal{J} = -2\pi \operatorname{Im} [\operatorname{res} \mathcal{F}(\infty)].$$

Here, $\operatorname{res} \mathcal{F}(\infty)$ denotes the residue of the function $\mathcal{F}(z)$ at infinity, i.e., the coefficient of z^{-1} in its Laurent expansion taken with the minus sign.

3 Current Layer with Attached Shocks

According to what was said in Sect. 2, finding the magnetic field in the model of a current layer with attached shocks is reduced to the Riemann–Hilbert problem for the analytic function $\mathcal{F} = \overline{B}$. This function is given by formula (4), where Φ is the conformal mapping, which is inverse to integral (5) and function $\mathcal{P}$ is a solution of the corresponding boundary-value problem in H^+ , given by the integral

$$\mathcal{P}(\zeta) = -ih\mathcal{K} \int_{\lambda}^{\zeta} \frac{(t - \lambda)^{\alpha-1}}{(t - \tau)^{\alpha+1/2}} (t - p) dt - \frac{\beta}{\sin \pi\alpha}, \quad (6)$$

where

$$p = \frac{\beta}{h} \frac{\sqrt{\tau - \lambda}}{\pi^{3/2} \mathcal{K}} \Gamma(1 - \alpha) \Gamma\left(\alpha + \frac{1}{2}\right) + 2\alpha(\tau - \lambda) + \lambda.$$

Here, λ , τ , and $\mathcal{K}$ are parameters of the conformal mapping Φ^{-1} determined by the geometry of the domain G (see Fig. 6b), α is an inclination angle of the cut (CDE) divided by π (see Fig. 6a), β and h are, respectively, the normal field component at the shock and the linear field growth coefficient at infinity, and $\Gamma(s)$ is the gamma function [37].

The analytic function $w = \mathcal{F}(z)$ makes a conformal mapping of the initial domain G onto some domain $\mathcal{W}$ that, as follows from representation (6) for $\mathcal{P}$ as the Christoffel–Schwarz integral, is an infinite tetragon (Fig. 6c). We call $\mathcal{W}$ the magnetic field hodograph domain [34].

In [28] basing on analysis of the hodograph domain we have shown that, contrary to the expectations following from Petschek's model, the considered MHD discontinuities attached to Syrovatskii's current layer are not slow but trans-Alfvénic or intermediate shocks (see, e.g., [35], Chap. 7), i.e., the shocks for which the normal plasma inflow velocity is higher than the upstream Alfvén velocity, while the normal outflow velocity is lower than the downstream Alfvén velocity. Also it has been

found that two types of transitions from nonevolutionary shocks to evolutionary ones along discontinuous flows are possible, depending on the model parameters.

4 Disrupting Current Layer

A disrupting current layer is shown in the complex z plane in the form of two rectilinear cuts of equal length located on the real axis symmetrically relative to the coordinate origin, see Fig. 4a, b. Real positive numbers a and b , which define endpoints of the cuts, are parameters of the problem. Based on boundary conditions formulated in Sect. 1 and on approach to finding the magnetic field stated in Sect. 2, we constructed the conjugated magnetic field $\mathcal{F}(z) = \overline{B}(z)$ in the exterior of the cuts in the following form:

$$\mathcal{F}(z) = -ih \frac{z(z^2 - \varepsilon^2)}{\sqrt{(z^2 - a^2)(z^2 - b^2)}},$$

where number ε is a free real parameter that emerges when the problem is solved formally; from a physical point of view, it defines positions of the magnetic nulls. We will consider a situation where the field outside the current layer vanishes only at $z = 0$, which means the fulfilment of the inequality $a \leq \varepsilon \leq b$ for the parameter ε . The point $z = 0$ at the center of the reconnection region has a special status. We believe that the plasma density near this point can decrease in the reconnection process (see [11], Chap. 3, Sect. 2) to such low values that the reconnection becomes essentially collisionless and very fast. In other words, the plasma is not enough to produce a secondary current layer capable of suppressing the current layer disrupt.

Let us turn to the model with disrupting current layer and two pairs of shocks, which are attached to its endpoints. As it was stated in Sect. 2, the magnetic field in the model is given by the formula (4), where Φ is the conformal mapping, which is inverse to integral (5) and function $\mathcal{P}$ is the solution of the corresponding boundary-value problem in H^+ . Using the approach [36], we obtain the sought—for function $\mathcal{P}$ in form of the generalized Christoffel–Schwarz integral that is most convenient for its calculation and subsequent analysis of the magnetic field:

$$\mathcal{P}(\zeta) = -ih\mathcal{K} \int_0^\zeta t^{-1/2} (t-a)^{-3/2} (t-\lambda)^{\alpha-1} (t-\tau)^{-1/2-\alpha} P_3(t) dt.$$

Here, $P_3(\zeta)$ is a cubic polynomial with real coefficients the first of which (at ζ^3) is equal to unity. The expressions for the remaining coefficients can be found by the method from [36] in terms of the Appel function—a generalized hypergeometric function of two complex variables (see [37]). These expression are not given here, because they are cumbersome.

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Solar Convection and Self-Similar Atmosphere's Structures

A.A. Agapov, E.A. Bruevich, and I.K. Rozgacheva

Abstract We present a new model of large-scale multilayer convection in solar type stars. This model allows us to understand such self-similar structures observed at solar surface as granulation, supergranulation and giant cells. We study the slow-rotated hydrogen star without magnetic field with the spherically-symmetric convective zone. The photon's flux comes to the convective zone from the central thermonuclear zone of the star. The interaction of these photons with the fully ionized hydrogen plasma with $T > 10^5$ K is carried out by the Thompson scattering of photon flux on protons and electrons. Under these conditions plasma is optically thick relative to the Thompson scattering. We find the stationary solution of the convective zone structure. This solution describes the convective layers responsible to the formation of the structures on the star's surface.

1 Introduction

The systematic extreme ultraviolet and X-ray emission observations from Skylab station, Yohkoh, SoHO and Trace satellites give us the very interesting images of solar corona. The structures are similar to standard coronal loops that connected separate active regions together [1, 4], but their “foots” lean on the photosphere out of active regions. These regular structures cover the hole solar disk as the more large-scale chromospheric network. It's necessary note that photosphere and chromosphere have regular structures such as grains, supergrains and giant grains.

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The giant grains are discovered by the helioseismology's methods [1]. These giant grains have the regular structure, their sizes are about $3 \cdot 10^5$ km with regular plasma speeds of 100 m/s. The solar-like stars photospheres have the similar structures as Sun has: grains and supergrains. In this paper we present the simple model of the hydrogen star convective zone. The necessary condition of free convection (rises in plasma layers with thickness of only some times smaller then solar radius) is the Schwarzschild criterion—the specific entropy of plasma decreases with moving away from the star center. Such convection will develop when the temperature inside of small convective volume (convective cell) decreases slower than the temperature decreases in neighboring plasma [5]. If solar convection is laminar so such processes as granulation, chromosphere network and supergranulation may exist in the convective layers of different thickness. Therefore solar convective zone consists of the three layers at least. Let's consider the main assumptions of this model: The convective zone is the layer with the spherically symmetry distribution of plasma around the radiative transfer energy zone. In this layer the condition of real hydrostatic equilibrium is carried out. Also we consider a case when the layer under study consists of the ionized hydrogen plasma only (protons and electrons). This consideration allow us understand the mechanism of convection zone structures formation. In our model this layer is open system through which the energy flux moves upwards. So let's consider that the plasma conditions we can described as polytropic equation:

$$\frac{N}{N_0} = \left(\frac{T}{T_0} \right)^n, \quad (1)$$

where n is the polytropic index, N and T are the plasma concentration and temperature. These values are $N_0 \approx 5 \cdot 10^{27} \text{ m}^{-3}$ and $T_0 \approx 2 \cdot 10^6 \text{ K}$ at the bottom border of the layer. The layer thickness (the convective zone depth) is approximately $0.3 R_\odot$, where $R_\odot$ is the solar radius. The energy emission come to convective zone, the temperature $T_r \approx T_0$ near the bottom border of the convective zone. This flux is the reason of the development of laminar convection. The emission and ionized plasma interaction is carried out by photon scattering on electrons and protons in case where the photon energy don't exceed the value kT . The time of energy transmission from photons to plasma don't excess the value [3]:

$$t_0 = \frac{3m_p c}{8\sigma_T \varepsilon_r}, \quad (2)$$

if $kT_r \ll m_p c^2$, where k is the Boltzmann's constant, σ_T is the Thompson probability section of scattering, m_p is the electron mass, $\varepsilon_r = \frac{4\sigma_B}{c} T_r^4$, σ_B is the Stephan–Boltzmann constant. If $T_r = T_0$ then $t_0 \approx 0.1$ s.

The distance of free run for photons is equal to $\Delta \approx (\sigma_T N)^{-1}$. If $N = N_0$ then $\Delta \approx 3$ m.

In Δ^3 volume plasma and emission are in thermodynamic equilibrium almost because of the radiation is connected with matter. The thermal conductivity mechanism is made available by the next processes in our case. The plasma (heated

by radiation in value Δ^3) loses the energy by bremsstrahlung. The speed of these losses is $\varepsilon_{ep} = 1.6 \cdot 10^{-40} N^2 \sqrt{T} \text{ J m}^{-3} \text{ s}^{-1}$. The characteristic time of this process is equal to : $t_2 = 2.6 \cdot 10^{17} \sqrt{T} N^{-1} \text{ s}$, [3]. If $N = N_0$, $T = T_0$ we derive $t_2 \approx 10^{-6} \text{ s}$. At the other hand the bremsstrahlung heats up the electrons in the vicinity of the volume Δ^3 . The time taken for this process:

$$t_1 = \frac{3m_e c}{8\sigma_T E}, \quad (3)$$

where $E = 3/2 N k T m_e / m_p$ is the electron energy density. If $T_* \leq T \leq T_0$ and $N_* \leq N \leq N_0$, $N_* = 2.5 \cdot 10^{26} \text{ m}^{-3}$ we find $20 \text{ s} < t_1 < 160 \text{ s}$. The characteristic rate of thermal conductivity is equal to $v_\chi = \frac{\Delta}{t_1}$. So the thermal conductivity coefficient for the process described (for the same order of magnitude) is equal to:

$$\chi = v_\chi \lambda = \frac{\lambda \Delta}{t_1}, \quad (4)$$

where λ is the thickness of the shell warmed up.

The convective energy transfer is carried out thanks to macroscopic transports of the value Δ^3 . The temperature inside the volume Δ^3 is higher then the plasma temperature in the layers which are situated higher then the bottom border of the convective zone, see Sect. 1. Thus the Archimedean raising force acts on this value and gives him the acceleration $\frac{g \cdot \Delta T}{T}$, where g is the free fall acceleration on the bottom border of the convective zone, $\Delta T = T_0 - T$, $T < T_0$. The flotation process is retarded by viscosity. In our case the viscosity is the consequence of the Thompson scattering. The value Δ^3 is full of plasma and radiation. When this value moves the radiation is scattered by electrons of neighboring plasma. Thanks to the scattering the equalization of electron momentum takes place inside the volume Δ^3 and outside of one. This viscosity they called radiation viscosity. It characterized by the viscosity coefficient

$$\nu = \frac{1}{3} \frac{c}{\sigma_T N}, \quad (5)$$

If $N = N_0$ then $\nu \approx 6 \cdot 10^9 \text{ m}^2/\text{s}$. This value is similar to the value estimation taken from the analysis of observations. The floating is ended when the raising force is in equilibrium with viscosity forces. The characteristic time of convective floating is equal to

$$t_2 = \frac{\nu}{g \lambda \frac{\Delta T}{T}}, \quad (6)$$

where λ is correspond to characteristic scale of the convective layer (the mixing length). If $\lambda \approx 2 \cdot 10^8 \text{ m}$, $T = T_0$ and $\frac{\Delta T}{T} \approx 1$, $g = 2g_\odot$, where $g_\odot \approx 274 \text{ m/s}^2$ is the gravity force acceleration on the solar surface then we have $t_2 \approx 0.05 \text{ s}$. So at the bottom border of the convective zone the relation $t_1 > t_2$ is taken place. In this case the convective transfer is more effective then the heat conduction. Near the top border of the convective zone $N = 4 \cdot 10^{22} \text{ m}^{-3}$ and $\Delta_* = (\sigma_T N)^{-1} \approx$

10^3 km. In this case we can ignore the Thompson effect. The radiation of plasma propagates free up to solar photosphere. In the Sect. 2. we give the solutions of stationary convective zone structures in the hydrodynamics approximation with the heat conduction (4) and viscous (5) coefficients. These solutions have the solitary wave structure and describe the model of multi-layer convection. All the convective cells have the torus contour.

2 The Equations of the Stationary Convective Zone Structure

The set of simultaneous equations for the spherically symmetric stationary convective zone which rotates about z axis (because of the hydrodynamics approximation is correct) have the form:

$$(\mathbf{v}, \nabla)\mathbf{v} = \frac{\nabla(p + p_r)}{\rho} + \mathbf{g} + \nu \cdot \nabla \mathbf{v} - 2[\mathbf{v}, \boldsymbol{\omega}] , \quad (7)$$

is the motion equation, where $\boldsymbol{\omega}$ is the angular velocity of convective zone rotation, ρ is the plasma density, p is the plasma pressure, p_r is the pressure of radiation.

$$(\mathbf{v}, \nabla)T = \chi \cdot \Delta T , \quad (8)$$

is the heat conduction equation,

$$\nabla(p + p_r) + \rho \mathbf{g} = 0 , \quad (9)$$

is the hydrostatics equilibrium equation,

$$\frac{dp_r}{dr} = -\frac{\sigma_T N}{c} \frac{1}{4\pi r^2} L , \quad (10)$$

is the radiation transfer equation outside of the volume Δ^3 ,

$$\frac{dL}{dr} = 4\pi r^2 \varepsilon_{ep} , \quad (11)$$

is the bremsstrahlung of plasma equation inside the volume Δ^3 ,

$$dM = 4\pi r^2 \rho \cdot dr , \quad (12)$$

is the mass conservation equation. In the Eq. (8) we don't take the density of radiation ρ_r because of $\rho_r \ll \rho$ in solar-like stars.

The set of simultaneous equations (8)–(10) have used by A.S. Eddington in 1926 [2]. Let $p = NkT$, $N = N(r)$, $T = T(r)$ and state of plasma is described by the

polytropic equation (1). Let's take the variable $x = \frac{r}{\zeta}$, where $\zeta = \left(\frac{kT_0}{4\pi G m_p \rho_0} \right)^{1/2} \approx 13 \cdot 10^5 \text{ km} > R_\odot$. Then function $\tau = \frac{T}{T_0}$ with Eqs. (8)–(11) can be transformed to the following equation

$$(n+1) \frac{1}{x^2} (x^2 \tau_x)_x = -\tau^n + \alpha \tau^{2n+\frac{1}{2}}, \quad (13)$$

where index x means the differentiation with respect to x and $\alpha = \frac{\sigma T \varepsilon_p(N_0, T_0)}{4\pi G m_p \rho_0} \approx 10^{17}$.

Let velocity vector $\mathbf{v}$ has the $\{V, W, Z\}$ components in spherical coordinate system. The vector of angular velocity $\boldsymbol{\omega}$ has the following components: $\{\omega \cos \theta, -\omega \sin \theta, 0, 0\}$. Let $\theta \ll 1$. From the equations of the structure

$$(\mathbf{v}, \nabla) \mathbf{v} = v \cdot \Delta \mathbf{v} - 2[\mathbf{v}, \boldsymbol{\omega}]$$

$$(\mathbf{v}, \nabla) T = \chi \cdot \Delta T, \quad (14)$$

one can find the equation for $V(x)$ and $\tau(x)$:

$$V V_x = \frac{v}{\zeta} \cdot \frac{1}{x^2} \cdot (x^2 V_x)_x$$

$$V \tau_x = \frac{\chi}{\zeta} \cdot \frac{1}{x^2} \cdot (x^2 \tau_x)_x, \quad (15)$$

Equations (12) and (14) are simplified when we assume that the component of velocity V is decreased with the depth. This condition is in agreement with solar observations: the plasma spread out velocity in the photosphere decreases with the scale increasing from grains to giant grains. We choose the solution in the next form:

$$V = \frac{\sigma}{\zeta} v, \quad (16)$$

where σ is the free parameter.

This permits us to simplify the Eq. (14) and transform it to the following:

$$(n+1) \beta \tau_x = -\tau^{2n+1} (1 - \alpha \tau^{n+1/2}), \quad (17)$$

where $\beta = 7.5, \sigma = \frac{v_0}{\chi_0} \sigma$. Equation (16) has different solutions for the different values of n . Let's choose the value n (use the Schwarzschild criterion). According to this criterion the temperature inside the small element Δ^3 has to decrease with increasing of the distance from the star center slower then decreasing of plasma temperature occurs. The plasma is in the hydrostatic equilibrium and the radiation is absence.

Substitute $p_r = 0$, $p = NkT$ and $\rho = m_p \cdot N$ to (8). Use the polytropic equation (1), we find the relative change of the plasma temperature (radiation doesn't take into account) $|\frac{T_x}{T}|_0 = \frac{m_p g \xi}{(n+1)kT_0} \tau^{-1}$.

In the volume Δ^3 the temperature changes according to (16). Also let's take into account $\alpha \gg 1$ and $\tau \leq 1$. Then $\alpha \tau^{n+1/2} \gg 1$ and relative change of the temperature inside the volume Δ^3 is approximately equal to: $|\frac{T_x}{T}| \approx \frac{\alpha}{(n+1)\beta} \tau^{3n+1/2}$. In our case of the evolution of the convective instability $|\frac{T_x}{T}|_0 > |\frac{T_x}{T}|$ the number β is evaluate as: $\beta > \frac{\alpha}{223} \tau^{3n+3/2}$. In this case we have $\tau^n \gg (\alpha \sqrt{\tau})^{-1}$. Therefore $\beta > (223\alpha^2)^{-1}$.

At the other hand one can integrate the Eq. (16) because of ignoring the first member of the right part of the equation. Thus we obtain the next algebraic equation: $\tau^{-3n-1/2} - 1 = \frac{\alpha}{\beta} \frac{3n+1/2}{n+1} (x - x_0)$. Using this equation and the consideration that $\tau^n \gg (\alpha \sqrt{\tau})^{-1}$ we find that $\beta < \frac{3n+1/2}{n+1} \frac{x-x_0}{\alpha^2 \tau}$. The value of polytropic index n (1) is necessary to choose as to make up the next inequality:

$$\frac{1}{223\alpha^2} < \beta < \frac{3n+1/2}{n+1} \frac{x-x_0}{\alpha^2 \tau}, \quad (18)$$

For the solar-like star we have $\frac{1}{303} \leq \tau \leq 1$ and $10^{-4} \leq x - x_0 \leq \frac{2}{57}$. In this case the inequality (17) is realized for all n having the positive values. Let's choose $n = 3/2$. Then the accurate solution of the Eq. (16) for $\tau(x)$ one can find from the next algebraic equation:

$$\frac{2}{5\beta} (x - x_0) = \frac{1}{3} \left(1 - \frac{1}{\tau^3}\right) + \alpha \left(1 - \frac{1}{\tau}\right) - \frac{\alpha^{3/2}}{2} \left(\ln \frac{\alpha^{1/2} + 1}{\alpha^{1/2} - 1} - \ln \frac{\alpha^{1/2} \tau + 1}{\alpha^{1/2} \tau - 1}\right), \quad (19)$$

We have taken into account that $\tau(x_0) = 1$ here.

The solution for $\tau(x)$ has the solitary wave form. It's clear from the form of the Eq. (16). If $\alpha^{1/2} \tau \gg 1$ we have the asymptotic solution

$$\left(\frac{T_0}{T}\right)^5 \approx \frac{2\alpha}{\beta \xi} (r - r_0), \quad (20)$$

For $x \rightarrow x_0$ one can find that $\tau \approx e^{-\frac{2}{3\alpha\beta}(x-x_0)} \rightarrow 1$.

At last we can find the expression for the speed components W and Z . Then we examine the most simple case of the symmetric spreading out on the sphere surface when $W = Z$. Let's consider also that the angular velocity w we can take from the equation

$$V \frac{\partial W}{\partial x} + 2W \cdot \omega \cdot \cos \theta = \frac{v}{\xi^2} \cdot \frac{1}{x^2} \cdot \frac{\partial}{\partial x} \left(x^2 \frac{\partial W}{\partial x} \right), \quad (21)$$

As follows from the Eq. (20) the convective zone rotates differently. Thanks to the convection the redistribution of the rotatory moment inside the star takes place. This effect is accurately studied in [5]. Under conditions selected in our paper we can find

the equation for W from the first equation of the set of simultaneous equations (14). His form becomes simple enough

$$\frac{dW^2}{dl} = \frac{v}{\xi^2} \cdot \frac{1}{x} \cdot \frac{dW^2}{dl^2}, \quad (22)$$

if we change the θ and ϕ angular variables to l variable and $dl = \sqrt{d\theta^2 + \sin^2\theta d\phi^2}$. Among the multiple numbers of solutions of the Eq. (21) there is periodic solution. This periodic solution has the next form:

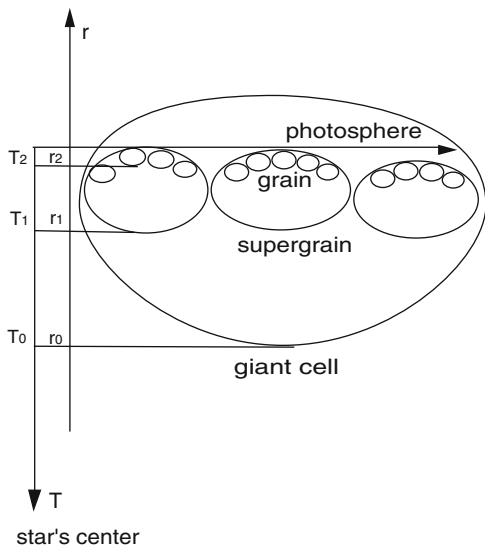
$$W = W_0 \cdot tg(W_0 \frac{\xi^2}{v} \cdot x \cdot (l - l_0)) \quad (23)$$

where W_0 is the speed peak value W , the point l_0 is situated at the radius x and is the start reading for l coordinate. On the surface of sphere with radius x_* plasma spreads out from l_0 point. So our model is symmetric there are many points $l_{0,i}$ on the surface of the sphere of radius x_* . The distance between the neighboring points is equal to $2\xi = x_0(l_{0,i} - l_{0,i-1}) = x_* \cdot \Delta l = \frac{\pi v}{W_0 \xi^2}$. Between these points there are two opposing plasma streams with velocities of opposite direction. These streams compensate each other at the distance ξ from the each points. So all the surface of the radius x_* breaks-down to the cells with diameters which are equal to ξ . All the number of these cells L we can calculate when we the surface square πx_*^2 divide by the cell square $\pi(\xi/2)^2$: $L = 4 \cdot (x_*/\xi)^2$. Then the velocity amplitude is equal to $W_0 = \pi v / 2\xi^2 x_* \sqrt{L}$. The kinetic energy density ε is proportional to W_0^2 . So the convective streams have the spectral energy distribution $\varepsilon \sim L \sim \varepsilon^{-2}$. The solutions of the convective zone structures (21) and (22) describes the stationary convection when all zone of the convective energy transfer consists of the layers with the different thickness. Every convective cell have the torus form. These solutions of this important problem are made for the first time. From the Eq. (19) follows the next conclusion: the convective zone differently rotates. Thanks to the convection the rotation moment redistribution inside the star is taken place. This effect is studied in detail in [5].

3 Summary and Conclusions

This model qualitatively describes the deep convective layers of the star under the supergrains layer. The plasma at these layers is the fully ionized. We don't study star's plasma at the highest convective under-photospheric layer where the turbulent processes are possible. In this turbulent layer there are necessary conditions for the generation of the long-scale magnetic field of the star. At the layers under this turbulent under-photospheric layer the convection is the stationary convection. Let's use the asymptotic solution (19) for the convective zone analysis. We have the convective zone consists of some layers with thickness of $\lambda_i, i = 0, 1, 2, \dots$. The

Fig. 1 The convective zone structure



temperature on the lower part of the layer's border is equal to T_i , on the top part is T_i^* . If we take the dependence of parameters from (15) and (19) on T_0 into account so we can find the relation between the velocity and temperature at the bottom and top borders of the neighboring layers:

$$V_{i-1}/V_i = (\lambda_{i-1}/\lambda_i)^2 (T_i/T_{i-1})^{3/2} (T_{i-1}^*/T_i^*)^5, \quad (24)$$

For the qualitative estimation let's substitute the characteristics of the convective layers associated with giant cells and supergrains into (23):

$$V_0 = 10 \text{ m/s}, T_0 \sim 2 \cdot 10^6 \text{ K}, \lambda_0 \sim 3 \cdot 10^5 \text{ km}$$

$$V_1 = 100 \text{ m/s}, T_1 \sim 10^6 \text{ K}, \lambda_1 \sim 3 \cdot 10^4 \text{ km}.$$

In this case we obtain that $T_{0*} \approx 0.4T_{1*}$ and the temperature on the top border of the layer λ_0 is smaller than the temperature on the top border of the layer $\lambda_1 < \lambda_0$. So we can see that λ_1 torus are situated into λ_0 torus. This qualitative analysis of the formulae (23) allows us to make the conclusion about the layers that are put one into another (see Fig. 1).

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SDO in Pulkovo Observatory

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Abstract We discuss effective applications of data obtained by both instruments of the Solar Dynamics Observatory: The Helioseismic & Magnetic Imager (HMI) and The Atmospheric Imaging Assembly (AIA). The purpose of this presentation is to show the most important problems of solar activity which are the main subjects in Pulkovo Observatory of the Russian Academy of Science. For these investigations uniform data sets of magnetic fields and coronal emissions in extreme ultra-violet bands are needed. Thus, we are planning to create SDO center in Pulkovo Observatory, which will help us in collaboration with existing SDO centers and provide more effective way of access to data for studies of the Sun.

1 Introduction

Topics of our interest are:

- Coronal helioseismology
- Coupling between the photosphere and white and EUV corona
- Dynamics of small-scale magnetic patterns in solar cycle
- Long-term sunspot oscillations
- Large-scale plasma motion in the Sun
- Developing of new mathematical methods of the image processing

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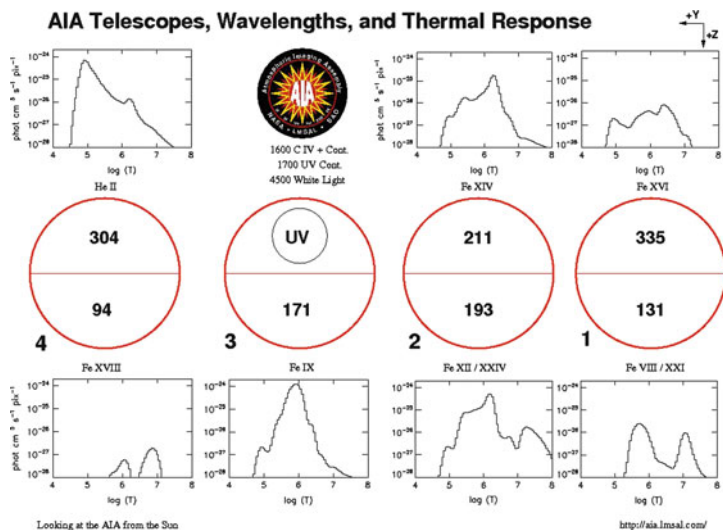
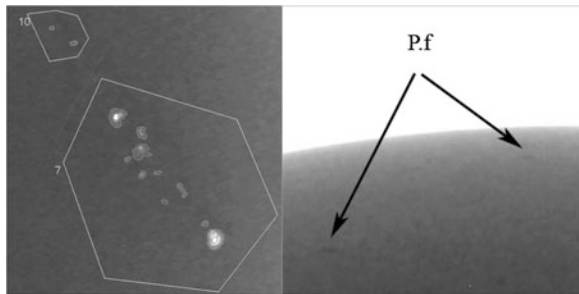


Fig. 1 A summary of the passbands of AIA, showing the filter responses of each of the EUV channels, and their arrangement on the instrument as viewed from the Sun

2 AIA and Coronal Helioseismology

Wave and oscillatory phenomena, which are an intrinsic attribute of the activity of solar and stellar coronae, is the subject of coronal seismology—a new and rapidly developing branch of astrophysics. Various approaches are used for the description of physical processes in flaring loops: kinetic, MHD, and electric circuit models are among them. Two main models are the most popular now in the coronal seismology. The first considers a loop as a resonator for MHD oscillations, whereas the second describes it in terms of an equivalent electric (RLC) circuit. Several detailed reviews are devoted to coronal seismology problems (see, e.g. [1, 2, 24, 25, 31]). Recent achievements in the solar coronal seismology are also referred in *Space Sci. Rev.*, vol. 149, No. 1–4 (2009). Nevertheless, some important aspects of the coronal seismology concerning diagnostics of physical processes and plasma parameters in solar and stellar flares require new approaches and further investigations. The Atmospheric Imager Assembly on a board of Solar Dynamics Observatory (SDO/AIA) produces full-disk multiwavelength extreme ultraviolet images of the Sun (see Fig. 1, from the SDO Data Analysis, by M. DeRosa and G. Slater).

Fig. 2 *Left panel:* An example of the automatic detection of observation using SDO/HMI umbra and penumbra boundaries 2010.11.28. *Right panel:* An example of polar faculae



3 Automatic Detection of Sunspots and Faculae in White Light in Cycle 24 Using HMI/AIA Data

The stability of automatic detection methods of solar activity can be tested by means of available long-term series of observational data. The white-light solar data observed by SOHO/MDI [27] are the most solid set of data for checking algorithms and software tools for extraction of solar active regions, allowing selection of both sunspots and solar faculae. The homogeneous HMI/AIA data of the solar atmosphere in white light allow extension historical sunspot data. Using SDO data we propose to solve the following problems. The extension of the sunspot data sets added new parameters such as a vector boundary of sunspots, sunspot's core and faculae, and it can be easily compared with magnetic and coronal structures in other wavelength bands (Fig. 2, left panel). In the present time, due to changing of methods of polar faculae registration, collecting of long-term data in the “old system” in observatories Noricura, Greenwich, Kislovodsk etc. is finished. SDO helps to continue these data sets and investigate cycle 24 and hopefully cycle 25 (Fig. 2, right panel). The stability of automatic detection methods of solar activity can be tested by means of available long-term series of observational data. For this study we use the set of full-disk (level 1.8) calibrated synoptic daily continuum and line-of-sight magnetogram observations. The data almost continuously covers the time period from 1996 until 2011 with the cadence of 4 continuum and 15 magnetogram observations per day. In automatic mode, a total of 31,988 sunspots were located. We give the results of automatic extraction of the sunspots, sunspot umbra and faculae in white light and the comparison of these parameters with the results of manual extraction in accordance with the observation at Kislovodsk Solar Station.

Correlation analysis with manual processing gives correlation coefficient $R = 0.99$. Figure 3 represents monthly averages of sunspot areas over cycle 23. In the automatic mode, a total of 31,988 sunspots were located. Correlation analysis with manually processed data give the following relationships: $A_{spot}^{MDI} = -20(\pm 15) + 0.95(\pm 0.013)A_{spot}^{Kisl}$, $R = 0.987$ [29].

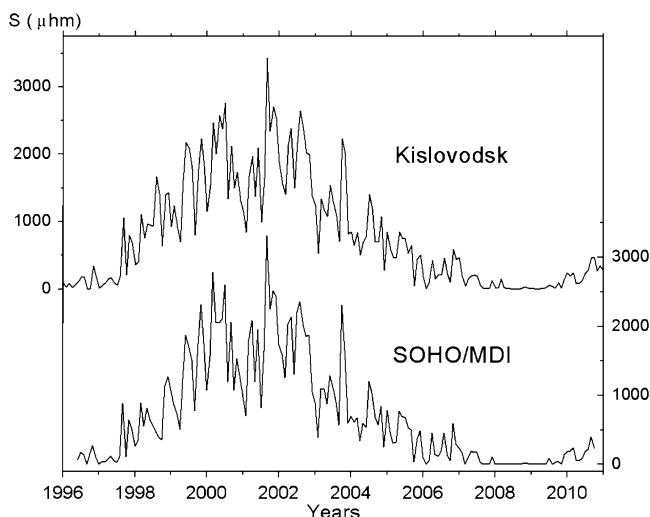


Fig. 3 Monthly averages of the sunspot areas from automatic processing of SOHO/MDI observational data (*lower graph*) compared with the results of manual processing of the sunspots' area from Kislovodsk Solar Station data (*upper graph*). Areas presented in millionths of the solar hemisphere

4 Dynamics of the Small-Scale Magnetic Elements

SOHO/MDI data have revealed, in addition to the rotation of solar plasma and magnetic field, a random motion of magnetic elements.

High latitude regions have a longitudinal speed of about 100 m/s both along and against the solar rotation. It arises, probably, from supergranulation [4, 5]. SDO/HMI data shows dynamics of small-scale magnetic elements in details. Figure 4 displays the evolution of the quiet sun during 12 h on May 5, 2010 in the line-of-sight component of the magnetic field. The 5×5 pixels filtering are applied. Size of one pixel is 0.001 in sine latitude and 0.1° in longitude. Figure 4 shows the dynamics of the three magnetic elements *b* during the 10 h by circles. And, by the end of this time these elements merge into one magnetic element. So, we observe the transport of the magnetic elements of the same polarity due to plasma motion. For example, the complicated dynamics of the magnetic elements could be shown by analyzing features of the magnetic elements *I* and *II*: area, magnetic flux, rotation and meridional displacements. The area of the magnetic element *I* increases with the decrease of the absolute value of the line-of-sight component of the strength of the magnetic field during the first 2 h. It is a process of the magnetic diffusion. After that, the area decreases and the magnetic flux of the same polarity increases. It tells about the emerge of the magnetic flux inside the same magnetic element. But, then, the diffusion of the magnetic flux returns. Should be mentioned that the element *I* rotates about $1^\circ/\text{day}$ slowly than Carrington rotation rate. But, the magnetic element *II* rotates slightly faster than Carrington

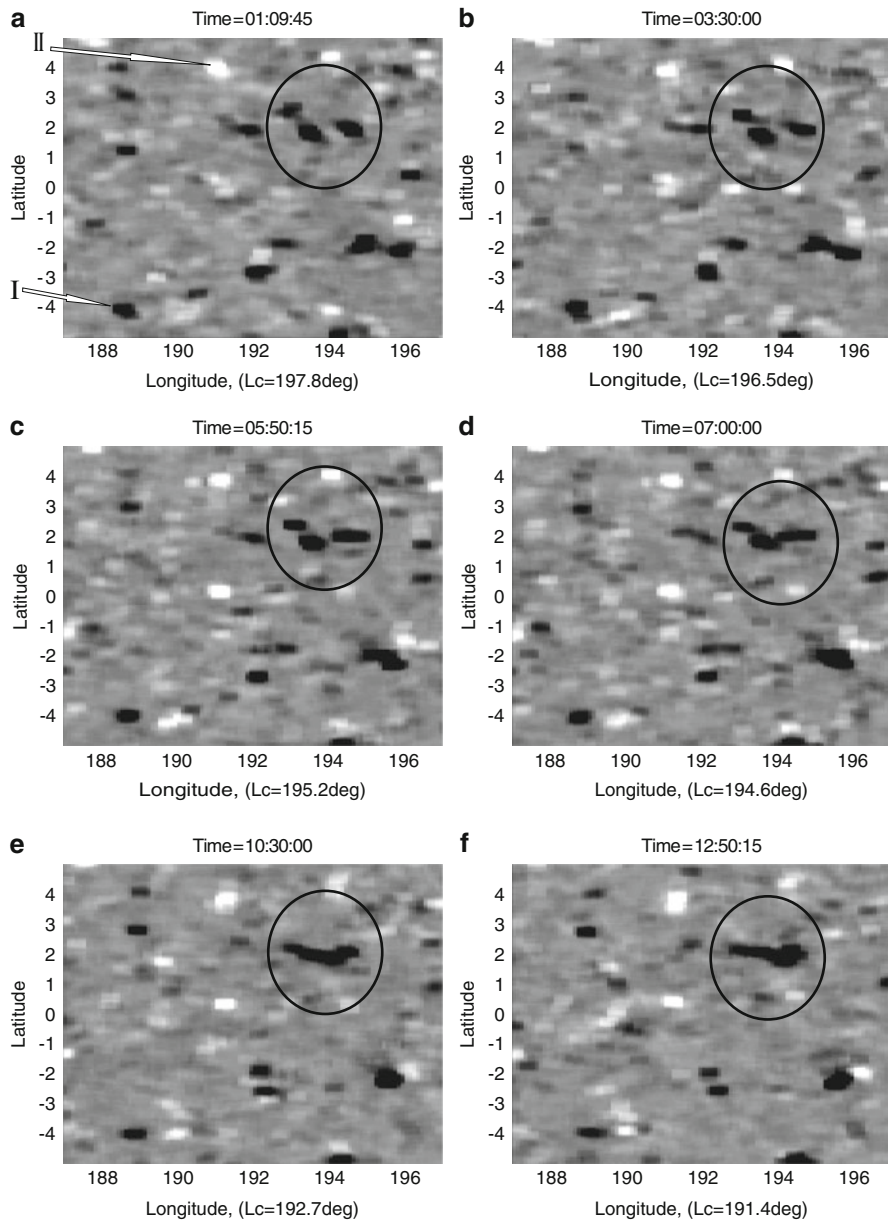


Fig. 4 Examples of consecution of the magnetic SDO/HMI frames in Carrington system, $\pm 5^\circ$ from the central meridian and the equator. *Circles* show the considered magnetic elements going into one element due to supergranulation

rotation rate ($\sim 0.6^\circ/\text{day}$). The both magnetic elements display displacements toward the equator. Therefore, the analyzing of the individual magnetic elements helps to identify the emerge and the dissipation of the magnetic flux together with the dynamics of the magnetic elements.

5 The Topology of the Small-Scale Magnetic Fields from HMI Data

Topology and geometry of the small-scale magnetic field pattern are analyzed in frame of the theory of random fields [18]. The excursion set of sufficiently regular random field is almost always a finite union of convex sets [3, 30]. Then it is possible to regard magnetograms as stratified manifolds where each stratum is a binary set that consists of clusters or “islands” formed by pixels that above of given level. The complexity of each stratum can be described by the Minkowski Functionals [19, 26]. We apply only two functionals: Euler characteristic and perimeter. The Euler characteristic measures the topological complexity of the magnetic field as an alternating sum of the number of singular points. Otherwise, perimeter makes the follow physical sense. Let us consider a statistical topography of the random field. Its total variation is a sum of modules of the field gradients in the open region. Geometrically, the total variation can be measured by the sum of equidistant isolines. So, for two hills of the same height the total isoline length is less for steeper peak. Thus, variations of any physical field are expressed by co-area formula. This formula tells that total perimeter of “islands” measures a portion of the level set to total variation of the field (Chan and Shen, 2005). We are going to apply this technique to HMI data for the following two problems. The first problem is to find an appropriate model of the magnetic field of quite Sun. The second task is related to the description of the active region dynamics using the morphological functionals.

6 Long-Term Oscillations of Sunspots and Their Magnetic Field

Several decades ago long-term sunspot oscillations (with periods of about tenth-hundreds minutes) were detected by ground-base observational data ([10]—microwave emission; [6, 7, 16]—spectral and white-light data). Since that time, there are a lot of discussions in respect to both about a reality of the long-term sunspot oscillations and a probable nature of this phenomenon. There are several spatial types of these oscillations: torsional (θ -mode), radial (r -mode), latitudinal (φ -mode) and longitudinal (λ -mode) [22].

The torsional and radial modes are examined by Nagovitsyn and Nagovitsyna [20, 21]. The evidences of latitudinal and longitudinal periodic motions are discussed

by Nagovitsyn and Vyalshin [9,23]. A response of long-term oscillation phenomena in upper solar atmosphere is considered by Gelfreikh et al. [15]. The radial-vertical mode of long-period oscillations of sunspot are investigated by Efremov et al. [11–14]. Here, using the HMI data, Efremov, Parfinenko & Solov'ev present an example of independence of magnetic field oscillations in widely separated sunspots. Figure 5 shows the results of analysis of magnetic data performed with the HMI (SDO) during the 80 h (3.3 days, from 4 to 7 June 2011). Two sunspots in different active complexes in one hemisphere are selected for the simultaneous studying. These sunspots are enough spaced. They have the same polarity and approximately equal values of magnetic strength: 1,880 G (S1) and 1,800 G (S2). But sunspot S1 has a little higher value. This fact, as we try to show below, may be a key in determining of status of the periodic signal. Is it a real or artifact? The location of the selected spots on the solar disk at time of 18:00, 05 June 2011 is shown in the upper left panel of Fig. 5. Since the spots are widely separated in space, the start of observational series is chosen when the spot S2 is at the central meridian. The end of the time-series corresponds to the time when the spot S1 reaches the central meridian. Temporal variations of the magnetic field at the centers of these sunspots (extreme counts) are shown in the upper right panel. The crossing of trends arises due to geometrical effect of projection of the line-of-sight component of the magnetic field strength. The position of the crossing corresponds to the chosen observation period. The middle panel of Fig. 5 shows two wavelets of the oscillatory part $h_1(t)$ and $h_2(t)$ (trends caused by projection effect are removed) of magnetic field in sunspots S1 (left) and S2 (right), correspondingly. The horizontal line drawn through the maximum of global wavelet for $h_1(t)$ (spot S1) indicates the level of periods which corresponds, in this case, to the global period of 17–18 h. Let note that the level of the global maximum for $h_2(t)$ (it is marked in the Fig. 4 by small line between the wavelets), is slightly lower than for $h_1(t)$ and corresponds to the period of 19–20 h. The form of the wavelet for h_2 (left middle panel) shows a small systematic difference in pe-riods which is maintained throughout the observational session. We know from our earlier study that the period (P) of long-term sunspot oscillations depends clearly on magnetic field strength (H) of sunspot: $P(H)$. This observational dependence is such that for the fields less than 2,700 gauss, the period of sunspot oscillations decreases with growing of magnetic field [14]. Taken into account that the magnetic field in sunspot S1 is about 80 G greater than in S2, we conclude that this small difference in the magnitudes of the field corresponds well to the dependence $P(H)$ mentioned above. Besides, this result agrees well with the theoretical model of “shallow” sunspot [28]. Finally, the time-series of $h(t)$ and $h_2(t)$ are presented together in the lower part of Fig. 5. In general, throughout the investigated interval of the time (80 h) any noticeable synchronization of two series of signals is missing. This is in the range of 17–20 h; they are reliable and have sharply defined peaks in power spectrum indicating the regularity of the periodic component in the original signal. During the small time-interval (only near the end of data set), some visible synchronization of two time-series in phase can be noted. This, of course, is not an evidence of a physical connection between the two oscillating systems. This is a simple mathematical consequence of the fact that the periods of these two

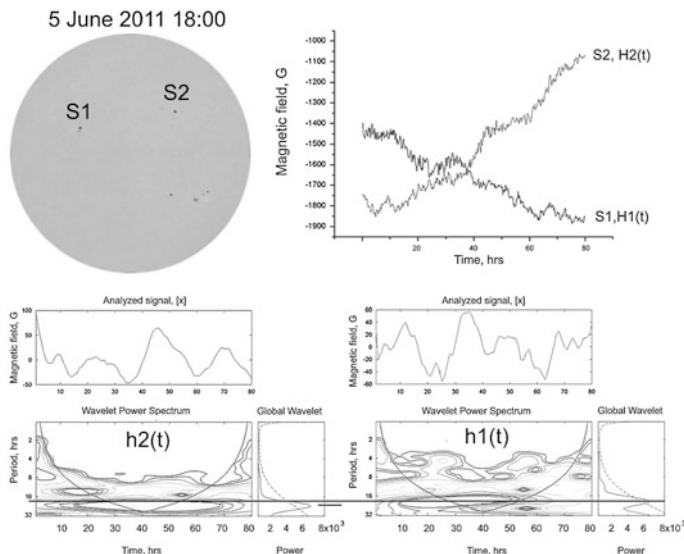


Fig. 5 Simultaneous observations of the magnetic field time-variations in two widely separated sunspots of approximately equal magnetic strength: 1,800 G (S2) and 1,880 G (S1)

oscillations are slightly different from each other (about 2 h). Because this small difference in the global periods, a partial time-synchronization will appear inevitably over a small time span on the large observation interval of 80 h as some weak visible effect. Thus, our example clearly shows that magnetic fields of spatially separated sunspots which belong to the different active regions oscillate independently. So, it should be for the two oscillating systems with similar parameters, but without any physical connection between them. As we see, the analysis of oscillations of these two sunspots do not reveal the artifacts with the periods of 12 and 24 h, which could be caused by the known orbital effects of SDO. But our other investigations of long-term sunspot oscillations show that the false harmonics usually arise in the power spectra of sunspots when the magnetic field strength in the sunspots is higher than 2,000 G.

7 Studying of Large-Scale Motions in the Solar Photosphere

Tracing of motions of magnetic structures on solar magnetograms can be used for study of global flows in solar photosphere. For example, Hathaway and Rightmire [17] tracked the motions of small-scale magnetic flux concentrations on MDI full disk images to reveal large-scale meridional flows and their evolution in solar cycle. The SDO HMI magnetograms can also be used for that purpose. They have higher space and time resolutions than MDI images, and, therefore, allow incorporating in the study both smaller tracers and lower speeds of the flows. In particular, it

would be interesting to investigate dependence of the motions upon the area of the magnetic elements and their heliospheric coordinates. At present the length of the HMI time series is as short as months. But in future, with accumulating of longer series, they also can be used to study the evolution of the large-scale flows and its dependence on the phase of the 11-year solar cycle.

8 Conclusion

In this paper we have presented our topic of interest for future investigations using the observation of the Solar Dynamics Observatory. Here, we display only several slides among the numerous unsolved problems of solar physics. Uniform data of the Doppler velocity, magnetic field and EUV coronal emission of the Solar Dynamic Observatory permit to investigate and solve these problems.

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Variations of Microwave Emission and MDI Topology in the Active Region NOAA 10030 Before and During the Power Flare Series

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Abstract The pre-flaring activity of the X3.0 (GOES class) flare that occurred in active region (AR) NOAA 10030 on 2002 July 15 is analyzed. We concentrate on the microwave (MW) emission and MDI topology in AR during some days before the X3 flare. Daily solar multi-wave observations with the RATAN-600 radio telescope in the range of 1.8–5.0 cm have been used. The following features in MW emission of AR have been detected in pre-flare period: (1) We have a new evidence of the MW “peculiar” source which has appeared 1 day before the X3 flare as a result of formation of δ -configuration in the central part of sunspot group. Such “peculiar” sources have been observed earlier with RATAN-600 in many active regions as precursors of high flare activity. (2) The most intensive local MW sources registered in the regions between the main sunspots coincide with the places of greatest values of the emission measure calculated by two lines of UV emission (SOHO/EIT). (3) The MW emission associated with the leading sunspot in AR considerably decreased (disappeared) 2 days before the X3.0 flare probably due to the changing of height of the emitting region. A topological method for detecting in AR a new emergence of magnetic flux using successive SOHO/MDI images of the solar disk is proposed. It was found that a number of disconnected components increases directly

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before the series of power flares (in particular, before the event on 2002 July 15) or accompanies this process.

1 Introduction

Dynamic regime of flaring active regions encloses the large range of a scale solar atmosphere which includes subphotospheric layers accessible for helioseismology, the photospheric layer and a chromosphere observed in various sites of a spectrum and the low corona visible in radio range from short cm- to decimeter waves. Observations of various tracers in all these layers allowed come to conclusion that preparation and realization of flares are connected with a change of complexity of magnetic structures in all layers of the solar atmosphere. It was shown by Künzel [15] that development of photospheric situation marked by the formation of δ -configuration phenomenon is one of the most impressive precursors of high flare activity. The complex topology of the magnetic field in strong flares has been studied in [10]. However, any uniform scenario of a complex of the processes leading to a flare is absent till now.

In this paper we tried to restore the general physical context previous and accompanying flare activity of the AR NOAA 10030 (July, 2002). Particularly, we focused on the pre-flare stage of the most powerful flare X3.0 (GOES class) occurred in this AR on 2002 July 15 when it was located close to the center of the solar disk (N19°, W01°). This eruptive event (X3.0 flare accompanied by coronal mass ejections (CMEs) has been investigated from different points of view in papers [4, 6, 9, 16, 19].

The favourable conditions for study this region (the situation near the equator and elongation along the EW-direction) provided to keep trace of the microwave emission (MW) of the AR some days before the X3.0 flare and during the flaring series and compare with the magnetic structure development in it.

We used in this investigation regular (daily) groundbase multi-wavelength radio observations in microwaves made with *the Large reflector radio telescope of the Russian Academy of Sciences* (RATAN-600) with very high sensitivity in flux measurements (several mJy) due to the huge effective area of the antenna system. The topology of magnetic fields variations of the AR at a photosphere level was studied on the base of 96-min magnetic data of *the Michelson Doppler Imager* (MDI) on board *the Solar and Heliospheric Observatory* (SOHO).

The topological complexity of magnetic fields is described by the alternating sum of the number of singular points of the observed line-of-sight scalar component of the magnetic field, the so-called “Euler characteristic” [11]. In order to reduce the measurement error for the intensity of the magnetograms, we estimate “Euler characteristic” for a set of excursions of the field above a specified level. In addition, an original method for the evaluation of new magnetic elements, which are possibly related with an emerging flux [12] was considered.

The methods and results of observations are presented in Sect. 2. The description of flaring activity of AR NOAA 10030 in July 2002 and information on charged particle fluxes produced by flares in this region and detected by satellites GOES, CORONAS-F are given in Sect. 2.1. The microwave manifestations of pre-flaring activity of the X3.0 flare of July 15 based on RATAN-600 solar observations are presented in Sect. 2.2. A topological method for detecting the magnetic fields variations of the AR at a photosphere level using 96-min SOHO/MDI data and its application for AR 10030 are given in Sect. 2.3. In Sect. 3 discussion and conclusions are presented.

2 Methods and Results of Observations

2.1 *Flaring Activity of AR NOAA 10030 and Charged Particle Fluxes*

According to GOES satellite data, the AR 10030 has generated 46 flares of soft X-ray (SXR) emission during 10 days from 10 to 19 July 2002: six flares belonged to M-class and two flares—to X-class. The first X3.0 flare has occurred on July 15 (the beginning at 19:59 UT, the peak at 20:08 UT and the end 20:14 UT). The second X1.8 flare occurred on July 18 (the peak at 07:44 UT).

Unfortunately, the solar observatory CORONAS-F [14] was in the shade of the Earth during both X-flares, and there was no possibility to measure their hard X-ray (HXR) and gamma-emission which have been detected by SONG-instrument [17] on board CORONAS-F [14] during three M-class flares occurred on July 11, 17 and 18. Information on these three flares is presented in Table 1.

During flares M8.5 (July 17) and M2.2 (July 18) gamma-emission has been detected up to channel of 1.6–4.8 MeV whereas during the flare M5.8 (July 11) maximal energy of the photons detected by SONG-instrument was no more than 550 KeV. Values of power index of HXR-spectra presented in Table 1 demonstrate that the neutral solar flare emission became harder during evolution of target AR.

One can see Fig. 1 that the strong enhancement of solar proton flux has been detected on July 16 at a geostationary orbit by GOES-10 and at low altitude orbit at 350 km in polar caps by CORONAS-F only after X3 flare. The intensive peak of particles with the low energies (up to 10 MeV) accelerated on CME bow shock is clearly seen on July 17 by both satellites. This peak is significantly weaker for particles with the energies more than 30 MeV and it is practically absent for particles with the energies more than 50 MeV. One can see also in Fig. 1 a weaker additional enhancement of solar proton flux in all energy range (up to 50 MeV) on July 19. Possibly it can be caused by both M2.2 and X1.8 flares, which were observed on July 18. The same times, there was no significant additional increasing of solar proton flux from flare M8.5 (July 17). The X3 flare on July 15 produced the most powerful flux of solar protons at the period under consideration. One may propose

Table 1 SXR, HXR and gamma-rays in three solar flares in AR NOAA 10030

N	Data dd/mm/yy	UT of flares	SXR	Flare coordinates	UT of flares	Emax	Power index of HXR spectra
		GOES SXR, hh:mm	class GOES		SONG SXR, hh:mm	SONG, Mev	
1	11/07/02	14:44–14:51–14:57	M5.8	N21E58	14:47–14:49	0.17–0.55	3.3
2	17/07/02	06:58–07:13–07:19	M8.5	N22W17	07:10–07:14	1.6–4.8	3.1
3	18/07/02	03:32–03:37–03:40	M2.2	N22W27	03:33–03:34	1.6–4.8	3.2

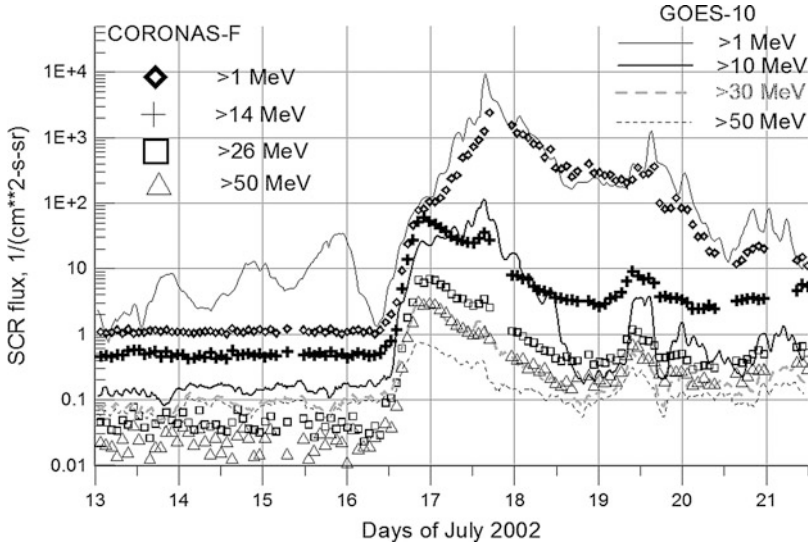


Fig. 1 Time profiles of solar protons observed at geostationary orbit (GOES-10) and at polar caps at 350 km (CORONAS-F)

that the conditions for more efficient acceleration of the charged particles have been created during the evolution of the AR 10030, when it moved across the solar disk. Gamma-emission of the last flares (M8.5 and M2.2) detected by SONG-instrument indicates that the charged particles have been accelerated in these flares to high energies.

2.2 Microwave Observations with the RATAN-600 Radio Telescope

In July 2002 solar observations were carried out daily at the radio telescope RATAN-600 [13] at about 11 UT. South sector of main mirror and flat periscope mirror were used in these observations. The knife-like diagram pattern of the radio telescope (FWHM) in this case was determined by the following relations: $\Theta_{horizontal}(\text{arcsec}) = 8.5 \times \lambda \text{ (cm)}$, $\Theta_{vertical}(\text{arcmin}) = 6.5 \times \lambda \text{ (cm)}$. Right

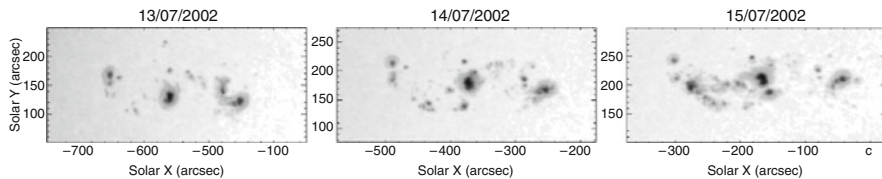


Fig. 2 Fragments of SOHO/MDI images (intensity) on July 13, 14 and 15, 2002 at the nearest moments to RATAN-600 observations

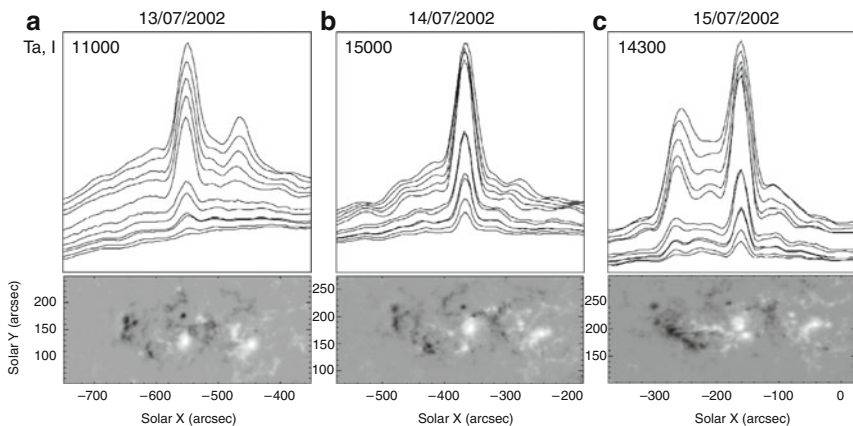


Fig. 3 (a)–(c) Fragments of one-dimensional solar scans (RATAN-600) on July 13, 14, 15, 2002 and SOHO/MDI magnetograms at the nearest moments to radio observations. Scans are ascending (upward) from the short wavelengths 1.8 cm (16.4 GHz) operating range

and left circularly polarized components (RCP and LCP) were recorded at 31 wavelengths simultaneously within the wavelength range of 1.8–5.0 cm while the Sun crossed the fixed antenna diagram. The total intensity (Stokes parameter “I”) and the circularly polarized component (Stokes parameter “V”) are calculated as $I = RCP + LCP$ and $V = RCP - LCP$.

Here we pay attention to the evolution of MW emission of AR 10030 during some days before the eruptive event on 2002 July 15. The AR 10030 was in a complicated magnetic morphology. Its development during 3 days (July 13–15) one can see in Fig. 2a–c. The fragments of solar SOHO/MDI images (intensity) of the AR at the nearest moments to RATAN-600 observations with orientation changed to fit radio observations are shown.

Figures 3a–c (*top panels*) show fragments of one-dimensional RATAN-600 solar scans (Stokes “I”) associated with the AR 10030 at several wavelength in the range of 1.8–5.0 cm. In Fig. 3a–c (*bottom panels*) one can see the fragments of solar SOHO/MDI images (magnetogram) of the AR. The orientation of the solar image was changed to fit the RATAN-600 observations. The co-alignment between the scan and the magnetogram image is assumed to be of 5 arcsecs.

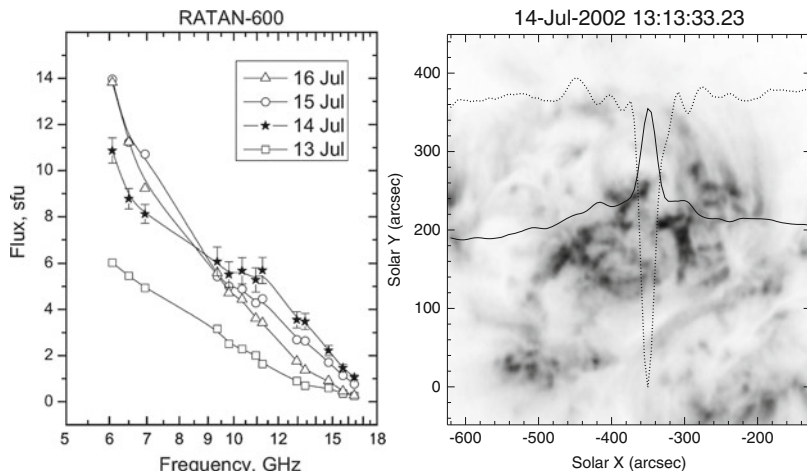


Fig. 4 (*Left panel*) Total flux MW spectra of the brightest radio source associated with the central part of the AR. (*Right panel*) Fragment of the RATAN-600 solar scan at $\lambda = 2.11$ cm overlaid on the AR image (negative) in *EM* obtained by SOHO/EIT data (Fe IX-X 171 Å–Fe XII 195 Å), (*SolarSoft*, data base of spectral lines *Chianti*). *Solid line*—Stocks “I”, *dotted line*—Stocks “V”

The fragments of scans show a most bright radio source associated with the central part of the AR. It dominated in the MW emission of the AR during 3 days before the X3.0 flare. According to [6, 9, 16] this region was initial site for the eruption. Then, the rotation of main (central) spot of sunspot group was detected [4].

The evolution of the total flux spectra of the most intensive radio source associated with the central part of the AR is presented in Fig. 4 (*left panel*). The spectra are constructed as the upper envelope of spectra obtained by using all wavelengths in the operating range. The error in the flux measurements (10–15%) is determined by the accuracy of both the separation of radio source on one-dimensional solar radio scan and the calibration technique. For absolute calibration we used the *Moon* and *Crab Nebula* observations made with the RATAN-600 radio telescope. The total fluxes in microwaves measured at different solar stations (*Nobeyama*, *Pentintonton*, *Learmonth*) were also taken into account.

It is well seen that the intensity of the central radio source increased considerably at highest frequencies with time at least 1 day before the X3 flare (at about 30% in antenna temperature T_a (K)). We note that there were not any powerful flares during 13–15 July in the AR under consideration.

Figure 4 (*left panel*) shows the microwave spectrum to be monotonically decreasing to higher frequencies on the July 13 (the empty square curve). This evidences the gyro-resonance emission [8, 24] of radio source associated with the central spot seen in Fig. 2a. The spectrum on July 14 shows significant changes—the emission at highest frequencies increased considerably and the local peak at frequencies 10–11 GHz has appeared (the black star curve). Figure 2 show a new sunspots and small pores nearby observed in the central part of the AR on July 14.

A main (central) sunspot in a positive magnetic field formed a δ -configuration with a small negative spot [16]. It may be proposed that the changes in the total flux spectra on July 14 are due to an appearance of a new “peculiar” microwave source associated with new formed δ -configuration. Such “peculiar” sources have been registered many times with RATAN-600 radio telescope in flaring active regions 1–2 days before the powerful flares (e.g. [2, 3, 20]).

The spectrum of July 14 contains an additional contribution from the source associated with the emergence of the new small spot of opposite polarity within the penumbra of the main spot (as shown in Fig. 1c in [16]). So, the spectra of July 14 may be explained as sum of the gyro-resonance emission of the sunspot-associated radio source and “peculiar” microwave source. The last one is characterized by high spectral index at short cm-wavelengths and by maximum of emission at 3–4 cm [20].

We paid attention to other features on the radio scans well visible at short microwaves—a small local peaks on both sides of the intensive radio source associated with central sunspot (see, Fig. 3 (*top panels*)). For a detailed identification of these peaks we used SOHO/EIT Data. In Fig. 4 (*right panel*) the fragment of scans (Stocks “I, V”) at wavelength of 2.11 cm (14.25 GHz) are compared with the emission measure (EM) distribution obtained for the pair of Fe lines in order to identify bright dense interspot parts in the active region that can contribute to MW emission.

The method used to obtain the EM -distribution is the following: the emission measure is radiant emittance of medium (plasma in case of solar atmosphere) and it depends on electron density:

$$EM = \int_{R_0}^R N_e dr = \bar{N}^2 \Delta R [cm^{-5}], \quad (1)$$

where N_e —electron density, R_0 and R —borders of a emitting layer, $\Delta R = R - R_0$ —a thickness of a emitting layer. EM is proportional to superficial brightness of a source with $\tau \gg 1$ in a coronal loops.

The EM can be defined on a relation of temperature responses of two near UV lines [5]. Observations of SOHO/EIT are performing in four UV spectral lines. The temperature and an EM can be defined from pair of lines Fe IX–X (171 Å)–Fe XII (195 Å) and Fe XII (195 Å)–Fe XV (284 Å), using *SolarSoft* and a database of spectral lines *Chianti*. The temperature response of spectral lines $R_i(T)$ whence can be obtained [22]. In Fig. 4 (*right panel*) the distribution of EM obtained from Fe IX–X (171 Å)–Fe XII (195 Å) lines is shown by halftones (negative).

So, in Fig. 4 (*right panel*) one can see a small local peaks (Stocks “T”, solid line) on both side of the brightest polarized radio source associated with the central sunspot in the AR. These peaks are associated with darker regions on the negative image which show the location of coronal loops parts with high dense and high temperature 1–2 day before the powerful X3 flare. Note, that the most dark region on the EM -image (negative) coincides with the place where X3 flare occurred on July 15 [16]. In this place the EM reached values of $EM = 4.47 \times 10^{28} cm^{-5}$ on

July 14 at 07:33 UT and $EM = 2.44 \times 10^{28} \text{cm}^{-5}$ on July 15 at 13:33 UT. It is close to the maximum values of EM over all solar disk during this time of 5.0×10^{28} and $2.5 \times 10^{28} \text{cm}^{-5}$, respectively.

So, the favorable location of AR along the E–W direction provided separation of the interspot microwave radio sources associated with the most dense and hot parts of the loops in the pre-flaring period.

We pay attention to other unusual features of MW emission of the AR observed during the pre-flaring stage (Fig. 3 (*top panels*)): the MW emission associated with the leading sunspot is directly observable at longest microwaves in the operating range (1.8–5.0 cm (16.4–6 GHz)) on July 13, but at the same time the radio emission associated with the follower part of the sunspot group is not detectable. On July 14 the intensity of radio emission associated with the central part of the AR increased significantly (at about 30%) at all wavelengths in comparing with the previous day. On July 15 (10 h before the X3 flare) the MW emission associated with the follower part of the sunspot group became to be detectable in the operating range, but the MW emission associated with the leading spots is not detectable at all wavelengths (Fig. 3a–c *top panels*). Probably, these microwave features are explained by the reconstruction of the magnetic field configuration in the AR before the X3 flare and may be connected with change of heights of the emitting regions [7].

2.3 Magnetic Fields Geometry and Topology Using MDI Images

We used the topological method for detecting the new emergence of magnetic flux [12]. This method uses the number of pixels in the image that can be distinguished from a specified value to a prescribed threshold (the number of disconnected components). Precisely, we assume that two pixels i and j from the field Ω of magnetogram are equivalent, $p_\varepsilon(i) \approx p_\varepsilon(j)$, if their values $p_\varepsilon(*)$ do not differ to within a specified threshold ε ; that is, $|p_\varepsilon(i) - p_\varepsilon(j)| < \varepsilon$. Denoted as $\mu_{iso}(\Omega)$ the number of pixels in the area that *could not be distinguished* from the central pixel $c(\Omega)$, to within ε :

$$\mu_{iso}(\Omega) = \sharp\{i \in \Omega : p_\varepsilon(i) \approx p_\varepsilon(c(\Omega))\}, \quad (2)$$

where the symbol $\sharp$ is an integer number of pixels. Since $\mu_{iso}(\Omega)$ is equal to the number of components that cannot be distinguished to within ε , the value of ε determines the accuracy with which we can distinguish pixels in the area, or the resolution. Using topological ideas of ε -connectivity [18], we assume that the number of ε -connected components, $C(\varepsilon)$, is equal to the number of pixels that cannot be distinguished to within ε , or $C(\varepsilon) = \mu_{iso}(\Omega)$. The number of *disconnected* or *distinguishable* components will then be the number of pixels that can be distinguished to within ε : $D(\varepsilon) = N - C(\varepsilon)$, where N is the total number

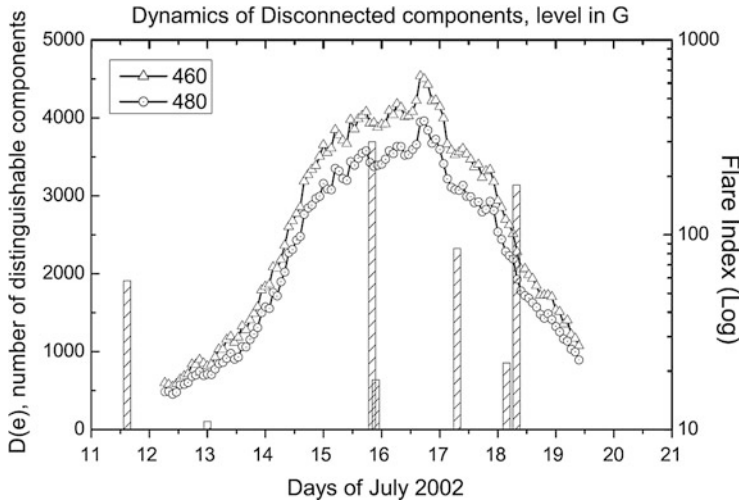


Fig. 5 Dynamics of disconnected components (level in Gauss). Vertical grey bars are all values of flare activity index

of pixels in the area Ω . Choose for the first magnetogram in sequence such a value ε that $D(\varepsilon) \approx 0$. Then increasing of $D(\varepsilon)$ with time could be interpreted as an emergence of new magnetic flux. We apply this approach to the series of active regions and found that flare activity of an AR is closely related to the time evolution of $D(\varepsilon)$. An increase in $D(\varepsilon)$ precedes or accompanies a series of M and X flares. Figure 5 shows the evolution of $D(\varepsilon)$ for the AR 10030. The level of threshold corresponds to $B_z = 460$ G and $B_z = 480$ G. At this figure we can see gradually increasing the number of $D(\varepsilon)$ before the first big flare, after which growth continues and falls is along with the series of next flares.

Another topological characteristic which can be used for describing the evolution of magnetic field is “Euler characteristic” χ which could be defines for so called excursion set of field [11]. We may consider that line-of-sight component of the Solar magnetic field represents the height $X(\mathbf{t}) = B_z(\mathbf{t})$, of the random magnetic surface above or below some “zero plane”, that passes through magnetic topography by the points $\mathbf{t} \in (t_1, t_2)$. We do not know the exact form $X(\mathbf{t})$ and allow it to be two-dimension random field in some sense. The most interesting thing is statistics of extremes $X(\mathbf{t})$ i.e. $\max X(\mathbf{t})$ or “hills” and $\min X(\mathbf{t})$ or “valley” of random field in some domain of the magnetograms [1]. In practical application $X(\mathbf{t})$ is given by numerical matrix of line-of-sight component values in bounded region Ω of square lattice pixels $\mathbf{t} \in \mathbb{Z} \otimes \mathbb{Z}$.

Topography of the random field $X(\mathbf{t})$ can be so complex, that the set of their levels $A_u(X(\mathbf{t})) = \{\mathbf{t} | X(\mathbf{t}) = u\}$ on the height u may not contain connected elements, such as isolines. However excursion set $A_u(X(\mathbf{t})) \geq \{\mathbf{t} | X(\mathbf{t}) = u\}$ of sufficiently regular field over level u is almost always definite union of convex sets [1]. The number of connected components of excursion set A_u is called “Euler characteristic”. The

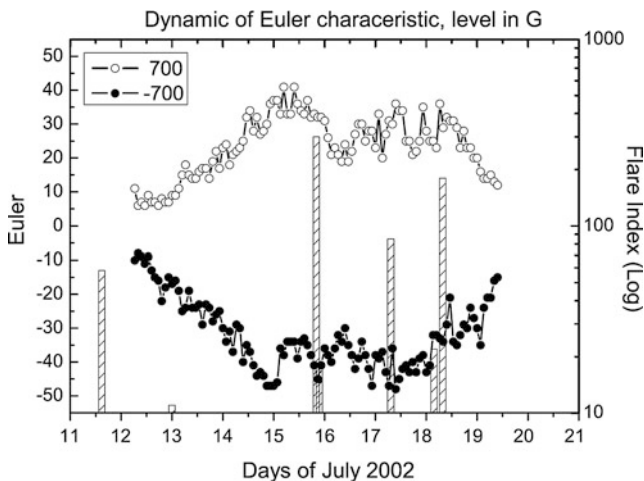


Fig. 6 Evolution of “Euler characteristics” $\chi(u)$ for the interval of observed flare activity of the AR 10030 for levels $u \geq 700$ G and $u \leq -700$ G. Increasing $\chi(u)$ for these levels is accompanied by rise of flare activity and successive local depressions on a base of common weak increasing. Vertical grey bars are powerful values of flare activity index

pixels of A_u are marked as black and the remaining pixels as white. Thus we obtain a binary image. Each matrix of magnetogram consists of collection of binary images indexed by the level of excursions. Then we calculate Euler characteristic for each level u with the help of formula

$$\chi(u) = \#(\text{islands}) - \#(\text{holes}), \quad (3)$$

where islands are produced by pixels belongs to excursion set.

There is another definition the Euler characteristic which comes from Morse theory [1]. Let us denote as n_0, n_1, n_2 , the number of maxima, saddles and minima of the field, where $\nabla X(\mathbf{t}) = 0$ on the set $\mathbf{t} \in M$. Then $\chi(M) = n_0 - n_1 + n_2$. So, the Euler characteristic measures topological complexity for $M \equiv A_u$ of magnetic field for points set $X(\mathbf{t}) \geq \mathbf{u}$. For numerical calculation we use successive images of SOHO/MDI magnetograms with time step 96 min for the AR 10030. Size of each fragment is 200×200 pixels. Euler characteristic was calculated for each fragment and set of levels u . In Fig. 6 the curves show the evolution of the Euler characteristic for levels $u = -700$ G and $u = 700$ G and the columns of the Flare productivity in logarithmic scale. Increasing of Euler characteristic by absolute value before the flare on July 16 could be interpreted as an increasing the number of field extremes (maximums and minimums) caused by elements of new emergence magnetic flux. From the other side the same effect could be caused by destruction of saddle points.

3 Discussion and Conclusion

The AR NOAA 10030 under consideration had high flare productivity—46 flares during period of July 10–19, 2002. Six flares of class M (GOES) and two flares of class X (GOES) have been detected in this AR. In this paper we paid attention to the features revealed before and during the series of flares with the powerful flare X3 occurred on 2002 July 15 in the center of solar disk. This flare has generated strong enhancement of solar proton flux detected on July 16 both at geostationary orbit by GOES-10 and at low altitude of 350 km in polar caps by CORONAS-F. It means that conditions for effective acceleration of the charged particles have been created in this AR.

The daily solar multi-wave radio observations in microwaves made with the RATAN-600 radio telescope revealed some features in radio emission of the AR 3 days before the X3 flare. The most remarkable of them is the change of the spectrum of radio emission of the most intensive radio source associated with the main spot in central part of the sunspot group. This fact may be explained as the appearance of the “peculiar” MW radio source as a result of new flux emergence of opposite polarity inside the pre-flare magnetic configuration—the main (central) sunspot in a positive magnetic field formed δ -configuration with a small negative spot which has detected on July 14. The formation of δ -configuration phenomenon is known as one of impressive precursors of high flare activity (e.g., [10, 15]). The “peculiar” MW sources have been observed earlier with the RATAN-600 and Large Pulkovo radio telescope in many flaring active regions as predictors of powerful flares (e.g., [2, 3, 20]).

Favorable orientation of the AR along the E–W direction provided the detailed identification of the sources of enhanced radio emission in the interspot regions during the pre-flare stage—they coincide with the dense and hot parts of the loops in the pre-flaring period. We came to this conclusion as a result of identification the fragments of one-dimensional solar scans (RATAN-600) with images of solar disk given in *EM* (emission measure) which were obtained by the relation of temperature responses of two near UV lines [21].

We paid attention to the unusual “decreasing” of the MW emission of the leading part of sunspot group 2 days before the X3.0 flare. It was observed simultaneously with “increasing” of the MW emission of the source associated with the follower part of the group. One may propose that it is due to the change of the height of the emitting region above the leading spot in the group. And other option, we can assume the existence of the restructuring process in the group which can occur with the acceleration of a small number of electrons. So, reduced MW emission can come from smaller number electrons coming into a footpoint, and larger number of protons which do not produce gyro-synchrotron emission (this different topology of electron and proton acceleration was discussed in recent papers [22, 23]).

We have applied a topological technique to trace regimes accompanying flare activity in the AR. We used two approaches: (1) we examined the change in the number of pixels distinguishable up to a prescribe threshold, it has been chosen so

that on the first magnetogram before the flare all the pixels in the selected window were indistinguishable. So, the appearance of distinct pixels in the evolution of the AR can be interpreted as the emergence of new magnetic flux. This increase one can see in the curve of maximum near July 17, 2002. (2) We estimate the complexity of the field—"Euler characteristic" (χ), it provides to measure imbalance in the number of extremes of the field and saddles (for a Gaussian stationary field imbalance is not). We see an increase in the absolute value of χ for the two selected levels (-700 and 700 G) with the increase in the flare productivity, i.e. the second approach allows to trace a flaring phase of the AR 10030.

We have revealed and analyzed several observed features in MW emission of high flaring AR with a complex magnetic structure. We concentrated on the MW emission and MDI topology in the AR during some days before the X3 flare and had a new evidence of the MW "peculiar" source which has appeared 1 day before the X3 flare as a result of formation of δ -configuration in the central part of sunspot group. This fact may be used for prediction of powerful flares as well as the results of application of the topological method to detecting a new emergence of magnetic flux in the AR. In this paper we investigated only one AR 10030, but it is necessary to obtain the statistical data about results of studying and comparing the flaring active regions with the same and different structure and magnetic topology. It would be a good basis for future flares investigations.

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Scenario of Evolution of the Epoch of Minimum at the Final Stage of Cycle 23

N.A. Lotova and V.N. Obridko

Abstract The paper is devoted to the study of 11-year cycles of solar activity with an emphasis on the peculiar features observed at the final stage of Cycle 23. The study is based on information on the solar wind stream structure and its relation to the solar wind sources in the corona.

1 Cycles of Solar Activity

Traditionally, the principal characteristic used to analyze the evolution of solar activity is the number of sunspots (Wolf numbers, R_z). In the variety of magnetic fields observed in the Sun, sunspots constitute the low-latitude, small-scale component [1]. In the recent years, much attention has been given to another, more universal characteristic—the intensity of the global coronal magnetic field $I_{Br}(t)$ [2]. The time variations of the two aforementioned parameters, $R_z(t)$ and $I_{Br}(t)$, are compared in Fig. 1.

Figure 1 shows that the time limits of the cycles and their epochs according to these two parameters differ. It has been established recently that the time variation of the general intensity of the global coronal magnetic field $I_{Br}(t)$ is a factor of similar or even greater importance than the Wolf numbers $R_z(t)$. So, new methods have appeared that allow us to study the solar periodicity using information on the solar wind, which is an extension of coronal magnetic fields to interplanetary space. The study is based on the radio occultation method, which was first developed and

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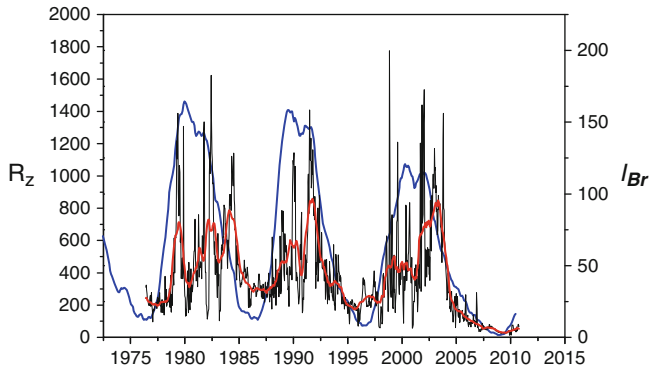


Fig. 1 Cycles of solar activity in the $R_z(t)$ and $I_{Br}(t)$ indices

applied by Vitkevitch [3, 4]. Later on, it was used for complex investigations of the solar wind along with spacecraft observations and provided fundamental results [5].

2 Radio Occultation Method and Solar Wind Studies Today

A new modification of the radio occultation method has been proposed lately for the study of the solar wind transonic transition region where the supersonic solar wind streams are formed [6]. It is used at relatively small distances to the Sun, $R \sim (2.5\text{--}70) R_s$, for simultaneous studies of the radio wave scattering angle $2\theta(R)$ in the strong scatter mode (quasar sources, wavelength $\lambda \sim 3$ m) and the scintillation index $m(R)$ in the weak scatter mode (maser sources of the water vapor line, wavelength $\lambda = 1.35$ cm). Observations are carried out at the Pushchino Radioastronomic Observatory of the Russian Academy of Sciences with the use of large radio telescopes DCR-1000 and RT-22. The plots of radial dependence of the scattering angle $2\theta(R)$ and the scintillation index $m(R)$ based on observation data are illustrated in Fig. 2.

The radial dependence of the scattering characteristics $2\theta(R)$ and $m(R)$ makes it possible to localize the boundaries of the solar wind transonic transition region in interplanetary space: the inner boundary R_{in} (the solar wind sound point) and the outer boundary R_{out} . Two modifications of the radio occultation method used simultaneously increase the number of the occulted sources involved and allow wide-range sounding of the solar wind. Regular experiments on multi-source scintillation measurements of circumsolar plasma at the distances $R \sim 70 R_s$ provided us with radio maps of the solar wind structure for 1987–2010 based on the annual data on the position of the solar wind transonic region boundaries—the inner boundary $R_{in}(\varphi)$ (the heliolatitude contour of the solar wind sonic point) and the outer boundary $R_{out}(\varphi)$ —in Cycles 22 and 23 (Fig. 3).

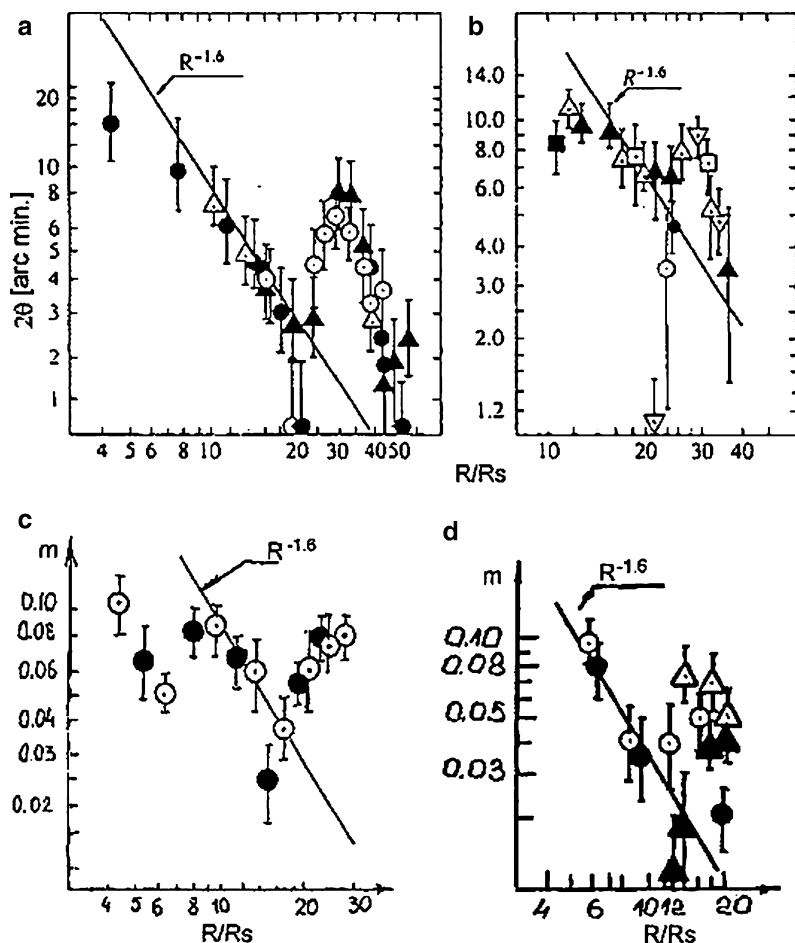


Fig. 2 Examples of radial dependence of the radio wave scattering angle $2\theta(R)$ and scintillation index $m(R)$: (a) observations of the sources 3C133 (O), 3C154(Δ), 3C162(□), and 3C172 (▽), $\lambda = 2.9$ m, June 1991; (b) observations of the sources 3C215 (Δ) and 3C225 (O), $\lambda = 2.7$ m, August 2001; (c) observations of the source IRC-20431, $\lambda = 1.35$ cm, December 2001; and (d) observations of the sources 0-IRC-20431 (Δ) and W31(2), $\lambda = 1.35$ cm, December 1997

Radio maps visualize the stream structure of the solar wind and reveal the typical features of the spatial latitude profile of the transonic region in various epochs of the activity cycle.

In the course of Cycle 23, significant differences could be seen in the evolution of the epoch of minimum 2006–2009 (Fig. 3) compared to the similar epoch in Cycle 22, e.g., a strongly increased lateral dimension in a narrow equatorial latitude zone in 2008 and an abnormal development of the transonic region in a wide range of heliolatitudes in 2009–2010.

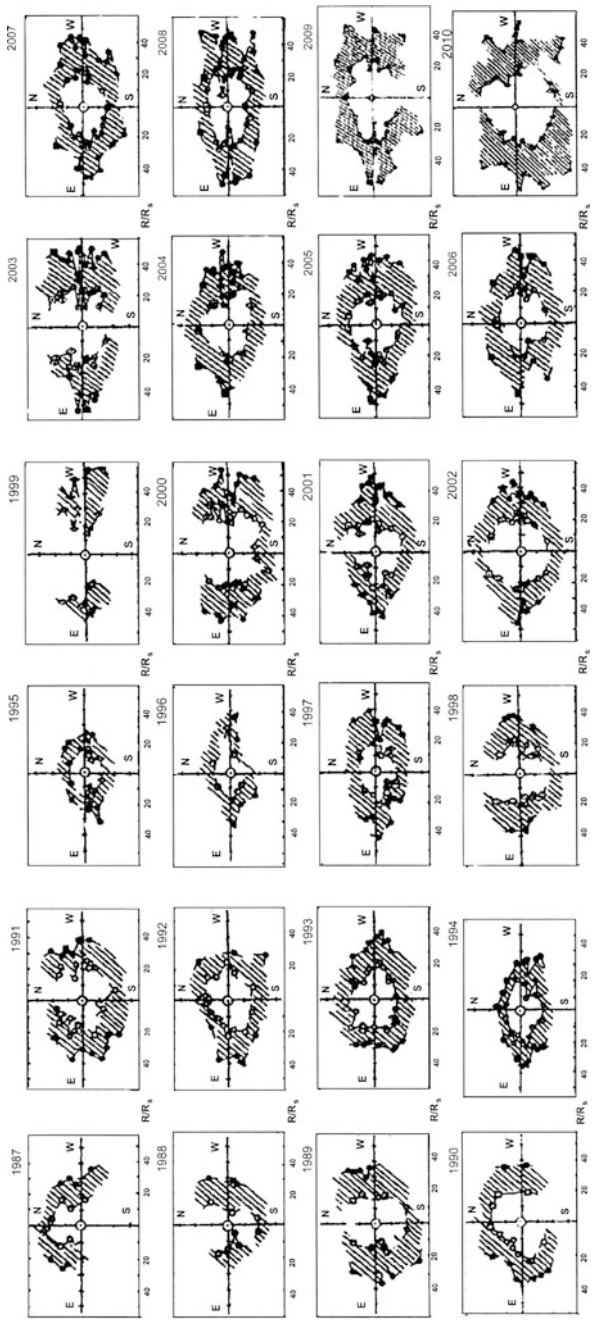


Fig. 3 Radio maps of the solar wind stream structure for 1987–2010

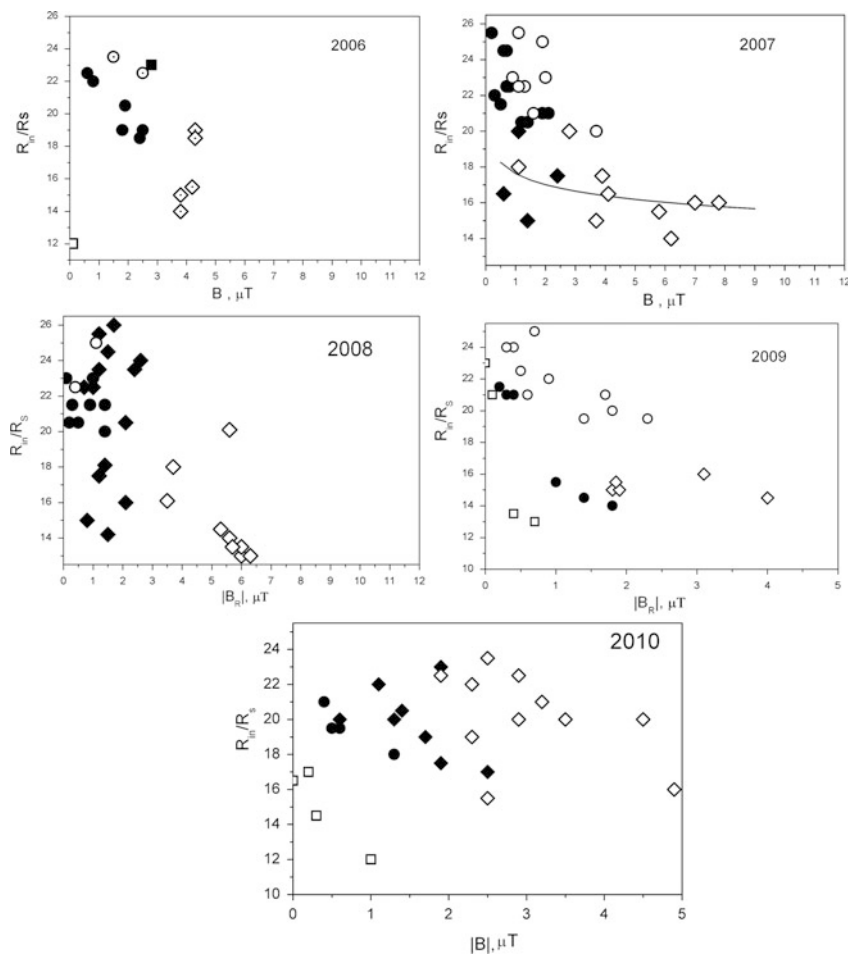


Fig. 4 Correlation diagrams of the sonic point R_{in} as a function of the magnetic field intensity $R_{in} = F(|B_R|)$ for the specific stream types in the period 2006–2010

3 Correlation Analysis of the Types and Structure of the Solar Wind Streams

The alternative method for analyzing the structure and types of the solar wind streams is based on the study of the correlation dependence $R_{in} = F(|B_R|)$, i.e. the dependence of the position of the solar wind sonic point R_{in} on the magnetic field intensity at the conjugate source point in the solar corona at $R = 2.5 R_s$ [6]. The $R_{in} = F(|B_R|)$ correlation diagrams break up into two branches—the types of streams of the solar wind (Fig. 4). Table 1 presents the results of the correlation analysis—the stream types and their sources for Cycle 23.

Table 1 Stream types and their sources

N ₀	Stream type	Magnetic field intensity $ B_R $	Magnetic field structure	White-light corona structure	Symbols
1	Fast stream	Strong field	Open field lines	Large CH and polar CH	◇
2	Fast stream	Medium and weak field	Open field lines	Equatorial CH	◆
3	Fast stream	Weak field	Open structure in weak field	Side lobe of streamer	△
4	Very fast stream	Very weak field	Open field lines	Vicinity of the zero line, probably, current sheet discontinuity	□
5	Slow stream	Medium and weak field	Closed field lines, loops, or mixed magnetic structure	Streamers	●
6	Slow stream	Weak field	Mixed structure	Weak diffuse glow: vicinity of streamer	○
7	Slow stream	Weak and medium field: descending branch on the correlation diagram	Periphery of equatorial CH	Zone between streamer and dark region	▼
8	Slow stream	Weak field: ascending branch on the correlation diagram	Mixture of loop structures of different scales	Zone between streamer and CH	▲
9	Slow stream	Medium field	Extended magnetic loops	Loop of the neutral line from CH to equator	■
10	Slow stream	Medium and strong field	Mixed	Zone between streamer and CH	▽
11	Slow stream	Weak field in a narrow band of $ B_R $	Open field lines	Small equatorial coronal holes (ECH)	◆

Figure 4 represents the correlation diagrams for a specific period—the epoch of minimum of Cycle 23 (2006–2009).

Figure 4 shows that the final period of Cycle 23 was characterized by numerous flux components with a large parameter R_{in} , i.e., slow components, dominating the solar wind [7].

4 Scenario of Evolution of the Epoch of Minimum of Cycle 23

According to the radio maps represented in Fig. 3, the epoch of minimum of Cycle 23 started in 2006 when the solar wind transonic region contracted crosswise and began approaching the Sun. A similar process of evolution of the transition region was observed at the minimum of Cycle 22 in 1995–1997. In Cycle 23, however, this evolution was not completed, but was changed suddenly by the reverse process: in 2007, the transonic region began to move away from the Sun. In 2008, the evolution of the transition region took an absolutely odd form: its transverse size increased significantly in a narrow equatorial zone (Fig. 3). Such a radical change of the heliolatitude shape of the transition region was accompanied by the appearance of a formerly unknown, slow component of the solar wind in the $R_{in} = F(|B_R|)$ correlation diagram (Fig. 2, Table 2). Thus, a new type of the low-speed solar wind streams (Fig. 2) arose in 2008 under the conditions of extraordinarily low Wolf numbers $R_z \sim 0$ (Table 2).

This increased the role of the slow wind in the equatorial region and was manifested in the shape of the transition region (Fig. 3, radio map for 2008).

In 2009, the solar activity began to increase very slowly to judge from the Wolf numbers, but the total number of sources of the low-speed streams still remained rather large (Table 2) and, as a result, the activity of the solar wind was low. This is seen on the corresponding radio map (Fig. 3), where the solar wind transition region remains wide and is located far from the Sun.

Table 2 illustrates time variations in the parameters of solar activity in the photosphere, solar corona, and solar wind. The tabulated data show that in 2007–2008, the Wolf numbers came close to zero and lost their value as indicator of evolution of the solar activity. On the other hand, the role of the solar wind parameters that reflect physical processes in the Sun increased significantly.

The particularities of evolution at the final stage of Cycle 23 resulted in an extension of the epoch of minimum (2006–2009) and of the cycle as a whole.

5 Conclusion

The new methods of investigation of the solar wind have expanded our knowledge of the evolution of solar activity. The experimental data obtained allowed us to reconstruct the scenario of abnormal evolution of processes in the solar wind at the final stage of Cycle 23.

Table 2 Time variations in the different parameters

N	Year	Rz Annual mean	Stream sources $IB_R [\mu T]$	Number of components in the diagram $R_{in} = F(B_R)$	Stream sources in the diagram $R_{in} = F(B_R)$	Stream types in the diagram $R_{in} = F(B_R)$
1.	2006	15.2	150	3	Streamers Vicinity of streamers Large PCH	Slow Slow Fast
2.	2007	7.5	80	4	Streamers Vicinity of streamers Large PCH Large ECH	Slow Slow Fast Fast
3.	2008	2.9 December: 1.7	50	6	Periphery of HCS Streamers Vicinity of streamers Large PCH Large ECH Small ECH	Slow Slow Slow Fast Fast Slow
4.	2009	3.1	20	4	Streamers Vicinity of streamers Large PCH	Slow and fast Slow Fast
5.	2010	16.5 (preliminary)	50	4	Periphery of HCS Periphery of HCS Streamers ECH Large Large PCH	Slow and fast Slow and fast Slow Slow Slow

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Solar Magnetic Fields as a Clue for the Mystery of the Permanent Solar Wind and the Solar Corona

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Abstract We analyze, generalize, and interpret the data for the permanent solar wind (PSW) velocities measured on board of Ulysses (SWOOPS). A finding of a principal importance extracted from Ulysses' observations is a discovery of the clear-cut inverse coupling between the SW velocities and the solar magnetic fields (SMF) (*the stronger close MF, the slower SW, and vice versa*), which points to the solar wind plasma deceleration by the SMFs below the source surface. Taking SMF into consideration leads to the alternative paradigm of the SW: *flow deceleration instead of the acceleration*. In such a case, both the SW and solar corona are converted into products of the interaction of an initial high-velocity plasma outflow ejected from the photosphere by solar magnetic fields. The latter not only divide initial fluxes into fast and slow parts, but also create and heat the corona through capture and stoppage of plasma in magnetic traps and the subsequent plasma heating. Observational arguments are presented in favor of the suggested idea.

1 Instead of the Introduction

The permanent solar plasma outflow named *the solar wind* (SW) was discovered in circumterrestrial space in 1958. Over more than 30 years this event was studied only on the base of the SW parameters measured by near-Earth space-crafts (SC) within a very thin heliospheric layer of $\pm 7^\circ$ near the ecliptic. The largest part of the heliosphere beyond this extremely narrow latitude belt was inaccessible for near-Earth SC. Obviously, theories of the SW origin based on such poor observations, in particular, on the flow acceleration data, are unlikely to be reliable. It was clear that

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an extra-ecliptic spacecraft mission was necessary to probe the heliosphere from the ecliptic up to the solar poles. Early in 2002 such a spacecraft named Ulysses was launched to an extra-ecliptic nearly polar orbit and began its climb to the south pole of the heliosphere. The Ulysses' unprecedented ~ 18.5 -year mission involved nearly three orbits, covered two solar minima and one maximum, and ended in June 2009, because of the lack of the energy resources. Measurements of the SW parameters were carried out within the Ulysses/SWOOPS project [1]. Here, the data on the permanent SW velocities are analyzed and interpreted versus heliographic latitudes, heliocentric distances, and solar magnetic fields regardless to any previous hypotheses.

In this study we also use IPS-SW velocities measured at EISCAT station in the north of Finland [4] and at VLBA-USA stations [11], solar full-disc magnetograms (NSO Kitt Peak and SOHO/MDI), and XUV images of the solar corona (Yohkoh and EIT/SOHO) taken from the Internet archive.

2 Ulysses' Permanent SW Velocity Data and General Results

The Ulysses' total data on the SW velocity taken from [5] are reproduced in Fig. 1 as a function of heliographic latitude, heliocentric distances, and solar activity phases (three diagrams in polar coordinates for each of the three orbits of the spacecraft around the Sun).

We can see that SW velocity patterns $V(\varphi, r)$ in the epochs of minimum and maximum solar activity differ noticeably. In the low activity phases (panels "a" and "c") the heliosphere is sharply divided into two parts with a narrow boundary between them. Inside the streamer belt of the heliosphere, the slow SW (< 500 km/s) prevails. Outside the streamer belt up to polar heliographic latitudes, only the fast permanent SW with the velocity of 700–800 km/s was observed. In the high-activity epoch, the slow-speed SW dominates over the whole heliosphere, with admixture of fast streams (500–800 km/s) coming from coronal holes (CH) of the active Sun ("b"). In the epochs of both quiet and active Sun, the slow SW is bound only with closed magnetic field structures of the streamer belt or the with active regions, whereas the high-velocity SW flows (500–800 km/s) are always associated either to the open magnetic fields of coronal holes in the epoch of the active Sun or to the weakest background and open polar magnetic fields of the low-active Sun.

Among Ulysses/SWOOPS measurements, the following results seem to be of a crucial importance due to their direct connection to the origin of the permanent solar wind:

1. The discovery of the regular high-velocity (750 ± 50 km/s) SW flows in the high-latitude heliosphere beyond the streamer belt. This result refutes the statement that coronal holes are the sole sources of the fast SW; this outflow is an event of only the low-activity epoch.
2. The discovery of the direct dependence of the SW velocity on both the phases of the solar activity cycle and the solar magnetic fields.

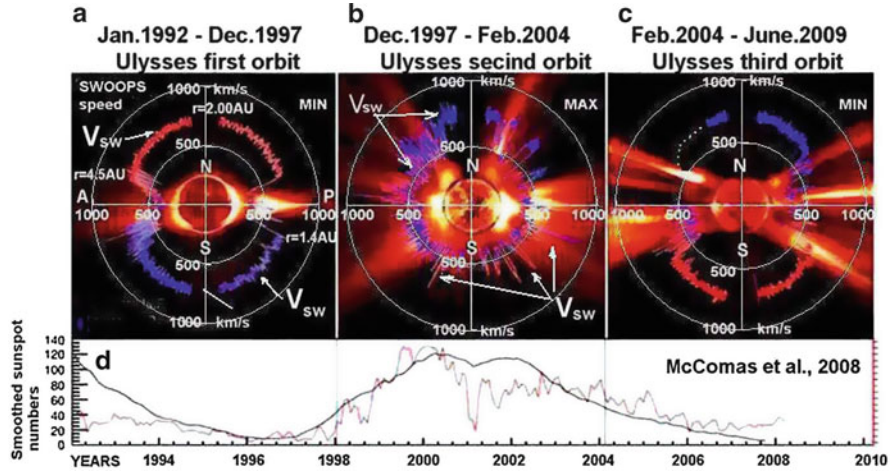


Fig. 1 Three diagrams of SW velocity in polar coordinates $V_{SW}(\phi, r)$ are combined with 195A FeXII corona image (EIT/SOHO) inside and white-light corona (LASCO-2) outside the solar disc for minimum, maximum, and minimum activity (panels “a”, “b”, and “c”, respectively). On each frame, *Aphelion* (A) (at $r \approx 5$ AU) is on the *left* and *Perihelion* (P) (at $r \approx 1.24$ AU) on the *right*. Ulysses orbited the Sun in counterclockwise direction (A $\rightarrow$ S $\rightarrow$ P $\rightarrow$ N $\rightarrow$ A). *Horizontal white straight lines* indicate the heliographic equator, the angles with it denote heliographic latitudes ($\phi = 0^\circ \pm 90^\circ$), the two concentric *white circles* show the SW velocity scale: 500 and 1000 km/s. The letters N and S denote the north and south solar poles. On the diagrams, SW velocities are presented with thin uninterrupted *white lines*. The *white spots* directly outside the solar limb are coronal streamers. Below, on panel “d”, a solar activity graph is plotted in the sunspot numbers

Until now, the problem of the connection between solar magnetic fields and the SW velocity was not considered, mainly due to the lack of the extra-ecliptic observations. Unlike the high-latitude SW, the low-latitude wind (only accessible for near-Earth spacecrafts) does not show remarkable velocity variations over an activity cycle. Only Woo and Habbal [12], [13] analyzed the relationship between the SW velocity and coronal structures using the Ulysses/SWOOPS data. In particular, they found that the velocity variations of 750 ± 50 km/s in the high-latitude SW of the quiet Sun epoch are not the Alfvén wave features but reflect the structures in the underlying background solar corona. In contrast, we discuss the link between the SW velocity and the magnetic field. The results are generalized below in Table 1.

The most interesting feature following from Table 1 is the *clear-cut inverse relationship* between the SW velocity and solar magnetic fields:

*the weaker close SMFs (<100 G), the higher SW speed (>500 km/s)
and vice versa*

the stronger close SMFs (>100 G), the lower SW speed (<500 km/s).

Table 1 The link between the SW velocity and the magnetic field

SW velocity (km/s)	Activity phases and magnetic field types associated
700–800	Only the phenomena of low-activity epochs: they are bound with open magnetic fields on the solar poles and with the background magnetic fields (BGMF) between the streamer belt and polar regions connected to the magnetic network [2]. These magnetic fields of different types display the same property: they are equally transparent for the <i>fast</i> flows
500–800	A fast SW outflow with such velocities is always associated with coronal holes in any activity phase
< 500	Slow SW is always coupled with closed magnetic fields of active regions and the streamer belt

This indicates the deceleration of the flow in the solar magnetic fields, which is only possible for a direct contact of the initial high-velocity outflow with magnetic structures deeply beneath the source surface.

Thus, we conclude that fast and slow SW streams originate from the interaction of the initial high-velocity (~ 900 km/s) outflow with the solar magnetic fields, which divide SW into the fast and slow fluxes [9].

Flow deceleration diagnosed near the solar magnetic field footpoints confirms that the high-speed plasma outflow near the photosphere surface is a real phenomenon rather than a hypothesis.

3 The Main Parameters of the Initial SW: V_0 and N_{p0}

A procedure of the estimation of the initial values of the velocity V_0 and proton density N_{p0} in the SW, based on Ulysses/SWOOPS measurements, is given in detail in [9]. As it was shown, in low-activity epochs the high-latitude fast SW goes out practically without deceleration in the solar magnetic fields, but under the influence of the solar gravitation. A fast SW motion may be considered as radial movement of protons from the Sun through its gravitational field, with the initial velocity V_0 . In this case a particle velocity $V(r)$ at any distance r from the solar centre is described by the following formula:

$$V(r) = \left[V_0^2 - GM_{Sun} \left(\frac{1}{r_{Sun}} - \frac{1}{r} \right) \right]^{1/2} \quad (1)$$

where G is the gravitation constant, M_{Sun} is the solar mass, r is a heliocentric distance, V_0 is an initial SW velocity at $r = R_{Sun}$. The formula (1) allows us to calculate either the SW velocity at any distance r in the heliosphere (if the initial velocity V_0 is known), or $V_0(r)$ (if $V(r)$ is known for a fixed r). Taking (according to Ulysses' data) $V_{SW} = 800$ km/s, we obtain $V_0 \approx 900$ km/s. The initial SW proton

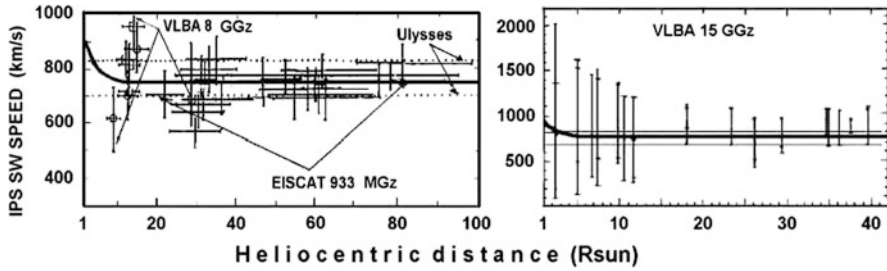


Fig. 2 IPS–SW velocities– $V_{SW}(r)$ are presented from [4] on the *left* and [11] on the *right*, measured in the high-latitude ($\varphi \geq \pm 60^\circ$) neighbouring ($r < 100 R_{Sun}$) heliosphere during the Ulysses’ first pass from the south to the north solar pole (1994–1995). *Two dotted lines in the left graph and two thin straight lines in the right* show the interval of Ulysses in-situ SW velocities in the outer high-latitude heliosphere

density $D_{po} = 2.8 \times 10^5 \text{ cm}^{-3}$ was evaluated in [9] using the conservation law for the SW total proton flux and the requirement of the compensation for the corona energy losses. Thus, the general parameters of the initial SW outflow are

$$D_{po} \approx 3 \times 10^5 \text{ cm}^{-3} \quad \text{and} \quad V_o \approx 900 \text{ km/s}$$

A good agreement of the IPS velocities, within uncertainty, with the Ulysses’ data shows that, in the same way as in the outer heliosphere, the high-latitude SW velocities in the inner heliosphere are concentrated in the interval 700–800 km/s. The velocity model for the high-latitude SW– $V_{sw}(r)$, calculated using the formula (1) for $V = 750 \text{ km/s}$ at $r = 750 \text{ km/s}$ and presented on both graphs in Fig. 2 (thick black solid lines), is consistent with SW velocity measurements and provides an evidence supporting our concept.

4 The Formation of the Solar Corona and Solar Wind

It is well known that, unlike those of the solar wind, the parameters (the density and temperature) of the main solar corona (SC) reveal a direct connection with the magnetic field strength:

the stronger are close magnetic fields, the brighter and hotter the coronal structures

and vice versa,

the weaker are close magnetic fields, the weaker and “colder” the features of the corona.

This indicates that SMFs take part in the formation of the corona and in its heating through capture and stoppage of plasma flows in magnetic traps. Generalizing observations of SW and the corona, we come to the following scheme:

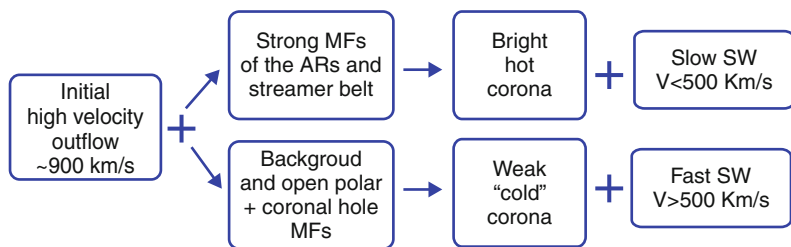


Fig. 3 This scheme demonstrates the process of formation of the solar corona and SW through deceleration of the initial fast outflow in the solar magnetic fields

Strong SMFs $\rightarrow$ bright hot corona + slow solar wind.

Weak SMFs $\rightarrow$ weak “cold” corona + fast solar wind.

Inserting the initial high-velocity outflow into this scheme, we obtain the united process of the formation of the solar wind–solar corona (Fig. 3).

Thereby we have come to the conclusion that the solar corona and the permanent fast and slow-speed solar wind are closely related to each other and represent products of the interaction between the initial fast outflow and the solar magnetic fields, in the form of plasma deceleration rather than acceleration.

The following question arises: what is the source of the initial outflow?

The answer: only the solar photosphere, a gigantic reservoir of plasma and energy, may generate such a super-high-speed initial plasma outflow.

Argumentation.

1. Only the photosphere is able to provide such a stable SW velocity $\sim 700 - 800$ km/s for two successive minima of activity in the high-latitude heliosphere (Fig. 4). Such a stability of SW velocity is unlikely to be due to random explosive events in the background magnetic field.
2. Another argument follows from the analysis of the element abundance in SW, coronal, and photospheric plasma carried out in [3]. It is known that in hot coronal structures connected with strong magnetic fields, as well as in the slow SW, the abundance of chemical elements with a low (< 10 eV) first ionization potential (FIP)—Al, Mg, Ca, Si and Fe—is by a factor of 4–5 higher than that in the photosphere (FIP-effect). On the contrary, the abundances of the same elements in the high-velocity SW and in the photosphere plasma were found to be close to each other [3]. This indicates that fast SW with the velocity 700–800 km/s originates directly from the initial high-speed outflow without being trapped by solar magnetic fields. Slow SW and the solar corona arise from the same initial plasma ejections trapped by the magnetic fields of the Sun.

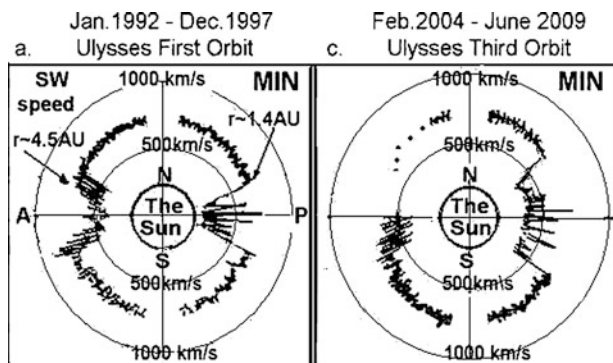


Fig. 4 The SW velocity $V_{SW}(\varphi)$ in the minima of 22^d and 23^d cycles. Despite the facts that the time lapse between panels “a” and “c” is equal to the length of the activity cycle, and that on the *right panel* the streamer belt is twice as thick as that on the *left panel*, the SW velocities beyond the streamer belt are exactly the same: 700–800 km/s

5 A Possible Non-thermal Mechanism of Generation of the Initial High-Speed Outflow Within the Solar Photosphere

Within the classic theory of stellar atmospheres, steady high-velocity eruptions of photosphere matter remain unexplained. The photosphere temperature is too low and the plasma density is too high to produce the observed high-velocity outflow in terms of thermodynamics. Similar phenomena, when extremely high velocities of a plasma stream cannot be described thermodynamically, are known as “‘Strange’ Fermi Processes”. Milovanov and Zeleni [6] showed that such high magnetic plasma super-streams can arise in the fractional medium under any force field effects as a power-law non-thermal tail. Such a situation is close to that considered here. According to Mogilevsky [8], with respect to some parameters, photospheric plasma represents a fractal medium. This means that the plasma is constructed of self-similar 3D elements, forming chains, with empty channels between them [7]. In our case, a power field as an outward radial draught (the difference between the pressure of the plasma inside and outside the photosphere) orients the channels in the radial direction with the plasma moving out with acceleration. Reaching the photosphere boundary, these flows are ejected in the form of the initial high-velocity outflow.

The initial high-velocity outflow may be considered as a solar plasma loss, which is required to maintain the energy balance in the solar interior. This means that the solar wind is not a surface event, but should be referred to global problems of the Sun as a star [10].

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Two Types of Coronal Bright Points in the 24-th Cycle of Solar Activity

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and Isroil Sattarov

Abstract We applied an automatic program for identification of coronal bright points (CBPs) to the data obtained by SOHO/EIT observations taken at the wavelength 195 Å, in the time interval from the end of the 23rd to the early 24th solar cycle. We studied the total number of CBPs and its variations at the beginning of the given cycle of solar activity, so that the development of the solar activity could be predicted with the use of CBPs. For a primary reference point for the 24th solar cycle, we took the emergence of a high-latitude sunspot with the reversed polarity, which appeared in January, 2008. We show that the observed number of CBPs reaches the highest point around the minimum of the solar activity, which in turn may result from the effect of visibility. The minimum solar activity at this time provides the opportunity to register the number of CBPs with the highest accuracy, with its uniform latitudinal distribution. We also study the properties of CBPs in a new 24th cycle of solar activity. It is shown that variations in the cyclic curve of the number of coronal bright points associated with variations in the solar activity, for the latitudes of the quiet Sun to be anticorrelation characteristic changes in the number CBPs to the solar activity, and the observational data are for the regions of active formations on the Sun almost identical on character on the equatorial latitude,

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but this have lightly expressed character in high-latitude zone. To explain the cyclic curves of variation in the number of coronal bright points in connection with the solar cycle in different latitudinal zones, we suggest a hypothesis of the existence of two types of coronal bright points: those associated with the quiet corona and those related to active formations.

1 Introduction

Bright X-ray points (XRTs, or coronal bright points, CBPs)—see their identification in [6, 7, 9] and Sattarov et al. 2005—are small-scale magnetic structures with a temperature about 2–4 million K and the average lifetime between 8 h and 2 days [1–3]. XRTs were first discovered in 1969 in soft X-ray beams. XRTs are too small (less than 60 arcsec in diameter) to have been discovered with the previous generation of X-ray telescopes. Later, the coronal bright points were also identified with the extreme ultraviolet telescope (EIT/SOHO) and called EIT bright points or EUV bright spots, according to the wavelength at which they were observed. Bright points are recorded in the photosphere, in the corona, and in the transition zone between the chromosphere and corona [4].

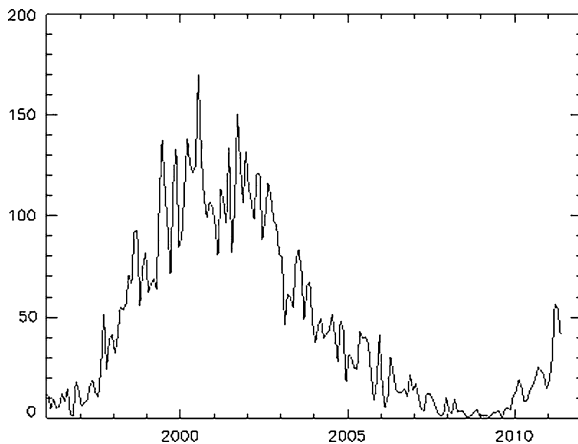
Bright points have central cores of approximately 10,000 km in diameter and generally occur over the areas of opposite magnetic polarity in the photosphere, when the regions of opposite polarity meet and destroy each other, releasing energy that heats the gas above the photosphere to 1–2 million K [5]. These small, point-like formations can also occur when a newly emerging magnetic field interacts with the existing magnetic field in the corona, again with the release of magnetic energy, which heats the gas. Being short-lived, transient objects, they are distributed almost evenly over all latitudes [2], and are observed in the equator, in active and quiet regions, and in coronal holes.

Despite the fact that the bright points have been studied both theoretically and observationally, numerous questions related to the formation of these bright spots in the lower corona (or rather in the transition zone between the chromosphere and corona) still remain unanswered. For example, it is still unclear how these structures emit energy and what is their role in the formation of the solar activity and solar radiation, whether they have a pronounced magnetic field, whether the effect of visibility is the only mechanism responsible for the anticorrelation of CBPs figures and sunspots, how the transients evolve, what is their connection with the corona heating and solar wind, etc.

The Sun, as a magneto-active star, has a strong magnetic field which, on average and on a large scale, is described as a magnetic dipole. The axis of this dipole changes its direction to opposite approximately every 11 years, which corresponds to the 11-year cycle of the solar activity; in turn, it is reflected in the sunspot cycle measured by the Wolf numbers.

Based on the data from SOHO, we constructed a curve of solar activity for the time interval from 1996 to 2011 (see Fig. 1). In the course of observations in

Fig. 1 Solar activity from 1996 to early 2011



Tashkent, the first appearance of a high-latitude spot was recorded on November 3, 2008 in the northern hemisphere, which indicates the beginning of the 24th solar cycle. The diagram (Fig. 1) shows that the second observed minimum of the 23rd solar cycle occurred at the end of 2008–2009, which yields the length of the cycle of about 13 years. The dynamics of the observed variations in the cycle and the comparison with the cyclic changes in the curves of CBPs for different latitudes can indicate not only the degree of the dependence of the CBPs phase on the phase of the solar activity cycle, but also will make it possible to hypothetically forecast the subsequent likelihood in the development of the further cycle.

2 Observations and Data Analysis

In this study, we use a series of data (an average of four images per day) obtained from the extreme ultraviolet telescope EIT installed on board of SOHO mission. The data were taken at the wavelength $195 \text{ \AA}$ from the 23rd to the beginning of the 24th solar cycle. The data were obtained using an automatic program for identification of coronal bright points (CBPs); the cyclic curves for different latitudes are constructed on the basis of these data. We studied the distribution of coronal bright points on the solar disk for different latitudes, for the quiet Sun and for active regions on the Sun, in order to identify patterns of variations in the 23rd—the beginning of 24th cycles of solar activity and to predict the nature of the cycle development.

We plotted the curves of the cyclic changes in the number of CBPs for the quiet Sun (QS) and for active regions of the Sun (AS) at different latitudes, to detect a connection with the cycle of solar activity (SA), processing the data derived from the images at the wavelength of 19.5 nm (EIT/SOHO) within the time interval 1996–2010. The hypothesis of the existence of two types of CBPs is discussed in more detail in the studies of Davis et al. [1] and Sattarov et al. (2010) [8].

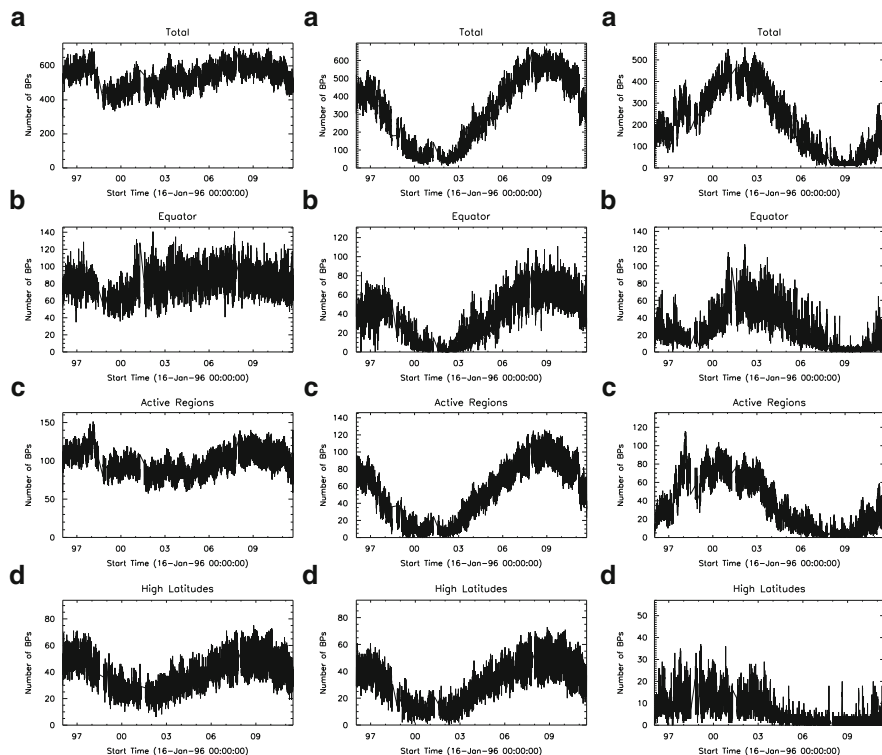


Fig. 2 Cyclic variations in the *curves* for the width of CBPs

The diagrams in Fig. 2 show the cyclic curves of the variation in the number of CBPs: a curve of the total number of CBPs and of the numbers in three latitude zones, at the equator between $+5^\circ$ and -5° , in the zone of active formations ($\pm |25^\circ - 35^\circ|$), and at high latitudes ($\pm |45^\circ - 55^\circ|$). For the sake of clarity, the curves are compared with the total variation in the CBPs number, for the quiet and active Sun, respectively. It is clearly seen that the variation in the number of CBPs at different latitudes has a different character depending on the variation in the phase of the solar activity. Thus, in the zone of active regions and at the equator, variations in the total number of CBPs (Fig. 2, column GN) do not show a significant link with the solar activity cycle, whereas at the high latitude zone the anticorrelation between the number of CBPs and the phase of solar activity is clearly seen. This result will be discussed below. When the data for the total width of the Quiet Sun (QS) and for active regions in the Sun (AS) are separated, another pattern is seen: in this case, correlation is observed only at high latitudes.

The analysis of the cyclic curves showed the decline in the total cyclic curve for the CBPs (Fig. 2 GN-a-d) in 1998 and 1999. This decline is noted at the equator and visible for the active regions. The solar activity grew steadily to its maximum in the early 2000–2001, without sharp peaks. There was a sharp drop in the solar activity

in the 23rd cycle in 2001 followed by an increase in 2002 (the difference was about 2 years).

For the cyclic curve of the total number of CBPs in high-latitude areas (Fig. 2 GN, d) a reverse relationship with the cycle of solar activity is observed; the number of CBPs was in anticorrelation with variations in the solar activity.

It should be noted that the number of CBPs varies with the solar activity, and the cyclic curve of CBPs not only displays a distinct two-humped pattern, but also shows anticorrelation between the number of CBPs and the cycle of solar activity at all latitudes of the quiet Sun. The tendency of the cyclic curve of CBPs on the quiet Sun to form the double-hump shape is seen for the equator, in active regions, and at high latitudes. For high-latitude zones of solar activity, this trend is not typical. Also the AS and QS's cyclic curves more clearly display the solar activity minimum in 2009, while on the curve of solar activity (Fig. 1) the minimum value can be traced from the end of 2008 to the middle 2009. Also can be more clearly observe of the solar activity minimum in 2009 on the AS and QS's cyclic curves of the number CBPs, while on the curve of solar activity (Fig. 1) the period of solar activity minimum can be traced with the end of 2008 to the middle 2009.

The cyclic curve of the CBPs on the active Sun is more or less consistent with the solar activity phases (see Fig. 2 GN, a), with the only difference in the graph of the total number CBPs in the cyclic curve AS (Fig. 2, AS, a), which shows a small but rather sharp jump in the number of CBPs from mid-1997 to early 1998; the SA curve during this period displays a more uniform growth. At the equator, the maximum number of CBPs is seen only in 2001–2002, and an increase in the CBPs number begins only in 2000, while the growth of the 23rd solar activity cycle starts at the beginning of 1998 and reaches its peak in 2000 (with the delay of about 2 years). The equator is characterized by a broad profile of the cyclic curve of CBPs for both the active and quiet Sun.

3 Conclusion

The following conclusions can be made from the study:

- Variations in the number of CBPs during a solar cycle cannot be explained only by the effect of visibility for the equatorial and high latitudes.
- The number of CBPs at different latitudes varies differently, depending on the phase of solar activity.
- To explain the cyclic curve of variations in the number of coronal bright points in connection with the solar cycle in different latitude zones, we suggest the hypothesis of the existence of two types of coronal bright points: those connected to the quiet corona and to active formations.
- The difference between the numbers of coronal bright points in the years of the minimum and maximum of the solar activity for the same latitude is different for the quiet and active Sun, and one can trace the following relationship: the

Quiet Sun displays an inverse relationship, with the double-humped shape of the distribution, and with the number of CBPs in anticorrelation with the cycle of the solar activity; in the Active Sun's regions, the variations of the number of CBPs almost correspond to those of the solar activity cycle.

- Regarding the SA curve, we can forecast the development of the next peak in the late 2012–early 2013. A more detailed analysis and conclusions, with more data obtained for 2 years to provide further CBP analysis will confirm our hypothesis.
- Our suggested determination of the solar activity using cyclic curves and two types of CBPs (AS and QS) describes more clearly the phase of the cycle in the corona.

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The Self-Similar Shrinkage of Force-Free Magnetic Flux Ropes in a Passive Medium of Finite Conductivity

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Abstract The twisted magnetic flux tube (magnetic rope) is a typical and the most important element of solar activity. Normally, the magnetic flux ropes in solar atmosphere are surrounded by a quasi-potential magnetic field, which provides the pressure balance in the cross-section. We have found a new exact MHD-solution for dissipative evolution of a thin magnetic flux rope in passive resistive plasma, with the growing gas density in the rope. The magnetic field inside the flux rope is assumed to be force-free: it is a set of concentric cylindrical shells (envelopes) filled with a twisted magnetic field [$B_z = B_0 J_0(\alpha r)$, $B_\phi = B_0 J_1(\alpha r)$, Lundquist, Phys. Rev. **83**, 307–311 (1951)]. There is no dissipation in a potential ambient field outside the rope, but inside it, where the current density can be sufficiently high, the magnetic energy is continuously converted into heat. The Joule dissipation lowers the magnetic pressure inside the flux rope, thereby balancing the pressure of the ambient field; this results in radial and longitudinal contraction of the magnetic rope with the rate defined by the plasma conductivity and the characteristic spatial scale of the magnetic field inside the flux rope. Formally, the structure shrinks to zero within a finite time interval (the dissipative magnetic collapse). The compression time can be relatively small, within a few hours, for a flux rope with a radius of about 300 km, if the magnetic helicity initially trapped in the flux rope (the helicity is proportional to the number of magnetic shells in the rope) is sufficiently large. This magnetic system is open along its axis of symmetry and along the separatrix surfaces, where $J_1(\alpha r) = 0$. On the rope's axis and on these surfaces, the magnetic field is strictly longitudinal, so the plasma will be ejected from the compressing tube along the axis and separatrix surfaces outwards in both directions (jets!) at a rate substantially exceeding that of the diffusion. The obtained solution can be applied to the mechanisms of coronal heating and flare energy release.

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1 Introduction

Twisted magnetic flux ropes are one of the most important and interesting elements of solar activity complexes. Due to the twisting of the field, these magnetic structures can store a large amount of free magnetic energy (non-potential energy of electric currents), and therefore they are able to provide energetically both the heating of the corona and the flare energy release. These magnetic elements are also able to transfer the magnetic energy from the photosphere and sub-photospheric layers to the corona by torsional MHD waves, and provide heating of the corona by dissipation of the waves on discontinuities. According to Parker's idea [7], numerous thin, unresolved magnetic flux tubes can provide heating of the corona due to mutual contact interactions, thereby creating a number of fine current sheets, in which the magnetic energy is converted into heat by reconnections of magnetic field lines. These fine magnetic elements could also be used to model large-scale solar flares, if it were somehow possible to force a large number of thin magnetic flux ropes to interact simultaneously, i.e. to create instantaneously a great number of small-scale current sheets [2, 16, 17].

The twisting of the magnetic field in a flux rope provides another important property of these magnetic structures: when a longitudinal magnetic field tends to reduce the length of a flux tube, the mutual repulsion of the rings of the azimuthal magnetic field, strung on the common magnetic axis, by contrast, stretches the tube along its length. It is the twisting of the magnetic rope (i.e. the azimuthal component of the field) that provides the fast coronal mass ejections (CME) into the interplanetary space against the gravity. No other effective driver for CME can be assumed.

The above arguments show that the theoretical study of the properties of the magnetic flux ropes remains a challenging problem of solar physics.

Magnetic elements in the solar atmosphere are prevented from fast transverse expansion primarily not by the difference in the gas pressure between the element and the surrounding medium but by the pressure of the ambient magnetic field, which is non-twisted and quasi-potential. Due to the potentiality of the external field, no dissipative processes in the ambient medium occur, but they do so inside the magnetic element, where the current density can be sufficient and the energy of the electric currents continuously dissipates via Joule heating. This dissipation leads to resistive time evolution of the system on timescales appreciably longer than the Alfvén time scale. The problem of resistive diffusion of a force-free field in the one-dimensional case (a plane-current sheet) was analyzed by [3].

In the current study, we have obtained new analytical MHD solutions describing the self-similar compression of twisted force-free magnetic flux tubes [4] immersed in resistive plasma. The compression of the rope is a result of the work done by the ambient medium, in which there is no dissipation and the magnetic pressure of the quasi-potential field outside the rope remains constant. Formally, the obtained solutions yield a decrease to zero in the transverse cross-section of the magnetic flux tube over a reasonable finite time. Of course, in reality, the flux rope does

not collapse totally, since, according to the physical sense of the problem, these solutions cease to be valid at some moment: as the element is compressed, the density and pressure of the gas within the element increase, the force-free approximation becomes unacceptable, and the self-similar compression ceases. However, the transverse size of the magnetic flux tube can decrease substantially over this time. In a resistive medium, the dissipation process will inevitably continue over new small space scales, but its description has to be appreciably more complicated and is unlikely to be self-similar.

The effect of dissipative contraction of force-free magnetic structures balanced by a quasi-potential ambient field should be considered as an important general mechanism of the formation of fine magnetic elements with force-free internal fields in conductive cosmic plasma.

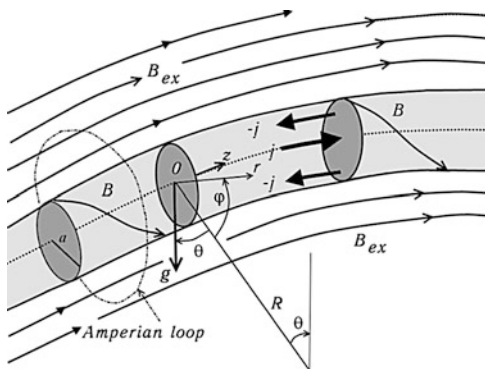
Before turning to the central problem of this study, we will recall the definition of a magnetic flux rope and present some formulae characterizing its equilibrium and the forces acting along its radius of curvature.

2 Definition of a Magnetic Flux Rope and the Zero Value of the Total (Net) Electric Current Along the Rope

We define a magnetic flux rope as a weakly curved magnetic flux tube twisted around its central longitudinal axis, with a circular cross section of the radius a and a large radius of curvature $R \gg a$ (Fig. 1). Due to the latter condition, the toroidal deviations of the magnetic field structure in the rope from the pure cylindrical symmetry can be neglected, and the magnetic field inside the flux rope can be represented as $\mathbf{B}\{0, B_\varphi(r), B_z(r)\}$. We use a local cylindrical coordinate system r, φ, z with its z -axis directed along a tangent to the central axis of the slightly curved tube (Fig. 1).

The twisting of the field in the flux tube can be produced by vortex motions of highly conducting plasma in the convective zone under the photosphere. Each

Fig. 1 The geometry of a twisted magnetic flux tube. The magnetic field displays the helical structure with two components: $B_\varphi(r)$ and $B_z(r)$. $\mathbf{B}_{ex}$ is an ambient quasi-potential field. The black thick arrows indicate the current density j . The net current through the cross-section πa^2 is equal to zero; $\mathbf{g}$ is the gravitational acceleration



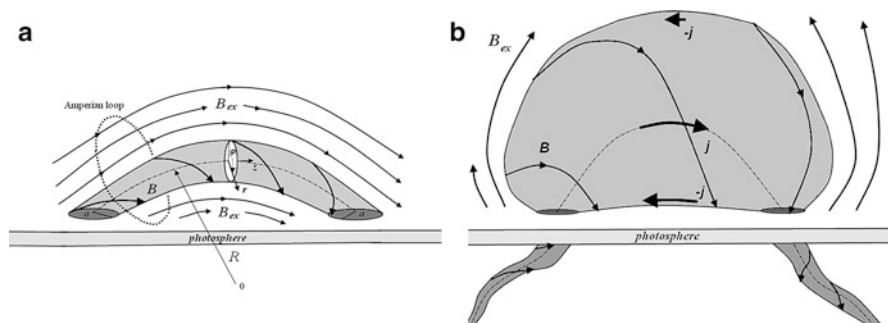


Fig. 2 (a) A low-lying magnetic loop (slightly curved magnetic flux rope) in the strong ambient magnetic field, quasi-parallel to the rope; (b) An expanded magnetic flux loop in a relatively weak ambient field of the solar corona. The arising new flux does not have time to penetrate into the plasma elements of the environment, but it can push apart mechanically the elements of the ambient medium at a great distance due to the excess of magnetic pressure in the arising loop. The electric currents in the loop are shown by *thick black arrows*: the net current is zero, the opposite currents near the photosphere exceed substantially those in the corona

plasma vortex has a well-defined finite size, so that the azimuthal field $B_\varphi(r)$ is to be limited by some finite cross-section πa^2 . The restriction of the B_φ – field in the radial direction means that this component can still differ from zero at the inner side of the tube, i.e. $B_\varphi|_{r=a-0} \neq 0$, but it must be zero at the outer side of the tube, at $r = a + 0$: $B_\varphi|_{r=a+0} = 0$. This condition can be considered as a formal definition of the radius of the transverse cross-section of the rope – a . As the azimuthal field $B_\varphi(r)$ falls with the distance from the rope axis more rapidly than r^{-1} , the total electrical current through the transverse cross-section of the flux rope, which can be calculated in the MHD-approximation using Ampère's law, $\text{rot} \mathbf{B} = 4\pi c^{-1} \mathbf{j}$, turns out to be zero:

$$I = \int_0^\infty j_z 2\pi r dr = \frac{c}{4\pi} \int_0^\infty \frac{1}{r} \frac{\partial}{\partial r} (r B_\varphi) 2\pi r dr = \frac{c}{2} r B_\varphi(r) \Big|_0^\infty = 0. \quad (1)$$

This fact can be illustrated more clearly when Ampère's law is presented in its integral form:

$$\oint_{(l)} \mathbf{B} d\mathbf{l} = \frac{4\pi}{c} \int_{(S)} \mathbf{j} d\mathbf{S} \equiv \frac{4\pi}{c} I. \quad (2)$$

Here, $d\mathbf{l}$ is a linear element of an arbitrary closed contour (the Amperian loop) and I is the total (net) current enclosed by this loop. If this loop is located in a plane normal to the rope axis and lies beyond the radius of the rope, where $B_\varphi(r) = 0$ (Figs. 1 and 2), it is obvious that the left-hand side of (1), representing the circulation of the magnetic field around the Amperian loop, vanishes. Consequently, the total current enclosed by the loop is also zero: $I = 0$.

Based on these arguments, we can draw the following general conclusion as applied to the Sun: for any magnetic structure appearing at the solar surface over

a relatively short time (hours to months) we can identify the spatial scale for which the total current enclosed by a closed loop of this size vanishes. To satisfy this condition, it is sufficient to place an Amperian loop along sufficiently distant points on the solar surface, such that the field of the newly formed object (a magnetic filament, a sunspot, faculae, or an entire active region) does obviously not reach these locations. The time for the diffusion of the magnetic field into the surrounding medium is $\tau = \tilde{n}^{-2} 4\pi\sigma\delta^2$, where σ is the conductivity of the plasma and δ is the skin depth. If the Amperian loop is sufficiently large ($l \gg \delta$), the diffusion time for the spatial scale of l will be huge, and the magnetic field of the new object cannot penetrate the ambient plasmas at a great distance to reach the points of the large closed loop l . Thereby, in these points the field of the new active structure will certainly remain nonexistent, its circulation around this loop will also be zero, and the total net current enclosed by the Amperian loop will vanish (Figs. 1 and 2a).

Thus, the depth of penetration of the new field into the plasma elements of the environment is very small compared to the size of the object, but, of course, a strong magnetic field rising from the photosphere into rarified upper layers of the solar atmosphere can mechanically push apart the environment with a weak magnetic field frozen into it, which is an external field in relation to the new arising magnetic flux. In this case, the geometrical boundary separating the two magnetic fluxes (new and previous) can be pushed up into the solar corona at a very large distance from the ascent of a new field, but, evidently, the electrodynamics condition $I = 0$ (which certainly holds in the photosphere and below, in the convection zone) can never be violated by such mechanical expansion of a magnetic flux tube in the corona (Fig. 2b).

3 The Lateral Equilibrium of a Flux Rope and the Force Acting Along the Radius of the Rope's Curvature

Due to the inequality $R \gg a$, one can analyze separately the condition for transverse equilibrium of a slightly curved magnetic flux rope and the dynamics of the same flux tube along its length L (or along its radius of the curvature R). The equilibrium over the cross section of a relatively thin magnetic rope is established within a short time, while the establishment of the equilibrium along L (or R) may take substantially longer or be nonexistent in the dynamical case.

Therefore, let us consider, firstly, the transverse equilibrium of a magnetic flux rope, assuming that stratification of the gas due to the gravity can be neglected inside the tube. That is, we assume that, like the magnetic field, the pressure P and density of the gas ρ in the tube are independent of the rotation angle φ .

In MHD-approximation the equation of force-balance has the form

$$-\nabla P + (4\pi)^{-1}[\text{rot } \mathbf{B} \times \mathbf{B}] + \rho \mathbf{g} = 0. \quad (3)$$

The radial component of the equation is (Fig.1)

$$\frac{\partial P}{\partial r} + \frac{B_z}{4\pi} \frac{\partial B_z}{\partial r} + \frac{B_\varphi}{4\pi r} \frac{\partial}{\partial r}(r B_\varphi) = \rho g \cos \theta \cos \phi. \quad (4)$$

To derive the condition for transverse equilibrium of the rope in its most general form, in terms of the values averaged over the cross section (the angular brackets $\langle \rangle$ are used here to denote these values), we multiply (4) by r and take the integral over the angle φ from 0 to 2π and over the radial distance r from 0 to a . Due to the integration over φ , the term with the force of gravity $\rho g \cos \theta \cos \varphi$ disappears. The integration by parts over r yields:

$$8\pi P(a+0) - 8\pi \langle P \rangle + B_z^2(a+0) - \langle B_z^2 \rangle + r^2 B_\phi^2|_0^{a+0} = 0. \quad (5)$$

The latter term on the left-hand side vanishes, since by definition of the rope $B_\varphi|_{r=a+0} = 0$. Denoting quantities corresponding to the external medium with the subscript “*ex*,” we will find that in terms of the averaged values the condition for transverse equilibrium of the rope with arbitrary distributions of the fields $B_z(r)$ and $B_\varphi(r)$ has the same simple form as that in the case of an untwisted cylindrical tube [6, 7, 9–14]:

$$8\pi \langle P \rangle + \langle B_z^2 \rangle = B_{z,ex}^2 + 8\pi P_{ex}. \quad (6)$$

It is important to emphasize that the azimuthal magnetic field is absent in this averaged balance of the total pressure, independent of the strength of the field, the uniformity or non-uniformity of its twist inside the flux rope volume, and the type of the field structure inside the rope (force-free or not).

Using (6), we can calculate the force acting along the radius of the curvature of the rope [9–11, 14], but here we should consider several cases, depending on the strength and structure of the external magnetic field.

1. The external field is weak or absent, $B_{ex} \approx 0$;

$$F_R(1) = \frac{1}{R} \left(\frac{\langle B_\varphi^2 \rangle}{8\pi} - \frac{\langle B_z^2 \rangle}{4\pi} \right) - \langle \rho \rangle g \cos \theta = -\frac{1}{R} \frac{\langle B_z^2 \rangle}{8\pi} (2 - \kappa) - \langle \rho \rangle g \cos \theta \quad (7)$$

Equation (7) contains the total twisting of the field in the rope defined by the relation $\kappa = \langle B_\varphi^2 \rangle / \langle B_z^2 \rangle$. In this case, the magnetic part of the force F_R vanishes when $\kappa_0 = 2$. The physical sense of this result is very simple: only for this degree of total twisting the tension of the B_z —field, which is equal to $-\langle B_z^2 \rangle (4\pi)^{-1} \pi a^2$ and which tends to shrink the rope along its length, is precisely balanced by the mutual repulsion (pressure) $\langle B_\varphi^2 \rangle (8\pi)^{-1} \pi a^2$ produced by the “rings” of the azimuthal field strung on a common magnetic axis of the rope. This pressure acts to extend the length of the rope [6, 7, 9–11, 14].

2. The external magnetic field differs from zero, $B_{ex} \neq 0$, and is quasi-parallel to the rope axis, so that in the surroundings of the curved rope it displays approximately the same radius of curvature as that of the magnetic loop.

In this situation, the external curved quasi-longitudinal field acts on a given magnetic flux tube (loop) not only along its minor radius r , due to the compressing transverse pressure $B_{z,ex}^2 (8\pi)^{-1}$, but also along R , due to the magnetic tension of the curved field lines of the B_{ex} -field. Thus, the term $\frac{1}{R} \frac{B_{z,ex}^2}{4\pi}$ will add to F_R , which, using (6), can be written as follows [14] :

$$\begin{aligned} F_R(2) &= \frac{1}{R} \left(\frac{\langle B_\phi^2 \rangle}{8\pi} - \frac{\langle B_z^2 \rangle}{4\pi} + \frac{B_{z,ex}^2}{4\pi} \right) - \langle \rho \rangle g \cos \theta \\ &= \frac{1}{R} \left(\frac{\langle B_\phi^2 \rangle}{8\pi} + 2(\langle P \rangle - P_{ex}) \right) - \langle \rho \rangle g \cos \theta. \end{aligned} \quad (8)$$

The first term in (8) is positive if the gas pressure in the loop exceeds the external pressure. Therefore, for any arbitrarily small (!) twisting of the field in the rope, the loop can be prevented from rising only by gravitation.

3. The external field has the same longitudinal component as that in the (2) case, and, besides, displays an additional component transverse to the rope axis $B_{\perp,ex}$:

$$F_R(3) = \frac{1}{R} \left(\frac{\langle B_\phi^2 \rangle}{8\pi} - \frac{\langle B_z^2 \rangle}{4\pi} + \frac{B_{z,ex}^2}{4\pi} - \frac{R}{a} \frac{B_{\perp,ex}^2}{8\pi} \right) - \langle \rho \rangle g \cos \theta \quad (9)$$

Here, the term with the transverse external field contains a large factor R/a . Thus, even a relatively weak transverse field can confine a strongly twisted magnetic loop. If the two terms in (9) describing the action of the external field on the magnetic flux rope cancel out each other, we return to the case (1), when the sign and magnitude of the magnetic part of the force $F_R(3)$ depend only on the total twisting of the field into the flux rope: $\kappa = \langle B_\phi^2 \rangle / \langle B_z^2 \rangle$.

4 Relaxation of the Magnetic Field in a Flux Rope to a Linear Force-Free State

In sub-photospheric layers of the Sun, the plasma parameter $\beta \equiv 8\pi P B^{-2} \gg 1$ and the behavior of the magnetic field are primarily determined by plasma flows (usually, not only $\beta \gg 1$, but also $\rho V^2/2 \gg B^2/8\pi$ i.e., $V^2 \gg V_A^2$, where V is the macroscopic velocity of the flows and $V_A = B(4\pi\rho)^{-1/2}$ is the Alfvén velocity). However, the situation changes dramatically when a fairly strong magnetic field (several hundred Gauss or more) enters the photosphere and higher layers.

Here, in the MHD-equation of the momentum for a plasma element,

$$\rho \left(\frac{\partial \mathbf{V}}{\partial t} + (\mathbf{V} \nabla) \mathbf{V} \right) = (4\pi)^{-1} [\text{rot } \mathbf{B} \times \mathbf{B}] - \nabla P + \rho \mathbf{g}, \quad (10)$$

the non-magnetic forces on the right-hand side become negligible, so that the magnetic system begins to move with a velocity comparable to the Alfven velocity towards the equilibrium, in which the magnetic force is also close to zero: $[\text{rot } \mathbf{B} \times \mathbf{B}] \rightarrow 0$. This is a force-free state described by the equations

$$\text{rot } \mathbf{B} = \alpha \mathbf{B}, \quad (\nabla \alpha \cdot \mathbf{B}) = 0. \quad (11)$$

In (11), α is some pseudo-scalar function of the coordinates. When $\alpha = \text{const}$, the magnetic field is a *linear* force-free field. Such a field possesses a remarkable property: if its helicity, which specifies the degree of coupling (knotting) of the field lines, is preserved in a given volume of plasma,

$$K = \int_V \mathbf{A} \times \mathbf{B} dV, \quad (12)$$

then the distribution (11) with $\alpha = \text{const}$ corresponds to the minimum magnetic energy of the system [18]. (In (12), $\mathbf{A}$ is the magnetic vector-potential). This means that a magnetic field, as it becomes to be an independent system (free from the influence of gas pressure and gravity) in a region with a small plasma parameter, should relax into a linear force-free state. However, such relaxation can be possible only due to a variation of topological characteristics of the system, i.e. due to reconnections of magnetic field lines in numerous small-scale current sheets. This process is dissipative, and it lowers the total magnetic energy of the system, but the helicity of the field (12) varies appreciably slower than the magnetic energy.

Therefore, the helicity plays the role of a topological invariant: $K \approx \text{const}$ [15]. The solution for the variational problem, for the conditional energy extremum for a given invariant K leads to the linear force-free field $\text{rot } \mathbf{B} = \text{const } \mathbf{B}$ [15]. For a cylindrically symmetrical flux rope with two components of the field $\mathbf{B}\{0, B_\phi(r), B_z(r)\}$, this distribution has the form [4]:

$$B_z(r) = B_0 J_0(\alpha r), \quad B_\phi(r) = B_0 J_1(\alpha r) \quad (13)$$

Under the photosphere, $\beta \gg 1$, and the structure of the magnetic field is determined by plasma motions. Near the photosphere, $\beta \approx 1$, and the vertical gradient of the magnetic buoyancy is sufficiently large to split a thick tube into a number of thin flux ropes. In upper layers, $\beta \ll 1$, and the inner structure of the rope tends to (13) due to topological relaxation. This relaxed part of the rope is compressed slowly under the external magnetic pressure due to Joule dissipation in the rope. White thick arrows indicate the rates of the compression.

We would like to emphasize that the phenomenon of relaxation of a magneto-plasma system to the force-free state (13) has been confirmed in experiments with “diffusion pinches,” when at the periphery of a plasma pinch a longitudinal

field of the opposite sign aroused spontaneously during the relaxation processes, i.e. function $J_0(\alpha r)$ passed through its first root [15].

At the boundary between the magnetic flux rope and the external medium, the following pressure balance is to be maintained:

$$B_z^2(\alpha a) + B_\phi^2(\alpha a) = B_{z,ex}^2 + 8\pi(P_{ex} - P_{in}(a)). \quad (14)$$

This is exactly in agreement with the general condition (6), since for the field (13)

$$\langle B_z^2(\alpha r) \rangle = \frac{2\pi}{\pi a^2} B_0^2 \int_0^a J_0^2(\alpha r) r dr = B_0^2 (J_0^2(\alpha a) + J_1^2(\alpha a)) = B_z^2(\alpha a) + B_\phi^2(\alpha a). \quad (15)$$

As the coordinate of the boundary surface of the magnetic flux rope αa , we choose one of the zeros of the function $J_1(x_n) = 0$, i.e. $(\alpha a)_n = x_n$, where x_n is one of the terms of the infinite sequence

$$x_n = 3.8317; 7.0166; 10.1735; 13.3237; 16.4706; 19.6158; 22.7601; 25.9037 \dots \quad (16)$$

In this case, the requirement $B_\phi|_{r=a+0} = 0$ is automatically satisfied, and there is no need to introduce the concept of a surface current at the interface between the flux rope and the external longitudinal field, which would be unavoidable if the coordinate of the rope boundary did not coincide with a zero of $J_1(\alpha r)$. Depending on the choice of the value of x_n , we will have a rope in which the longitudinal field changes its sign with the distance from the axis once, twice, three times, or more.

It is easy to verify that, for a force-free magnetic flux rope with the boundary $(\alpha a)_n = x_n$, $\langle B_z^2(\alpha r) \rangle = \langle B_\phi^2(\alpha r) \rangle = B_0^2 J_0^2(\alpha a)$, i.e. the total twisting of such a flux rope does not depend on our choice of the numerical value of the boundary coordinate $(\alpha a)_n$; it is always constant and equal to unity: $\kappa = 1$, ($\kappa < \kappa_0 = 2$). In accordance with (3)–(5), this means that, in a medium with a weak external field or an external field whose transverse component compensates the repulsive action of the curvature of the ambient field, such a magnetic flux rope will tend to reduce its length.

5 The Resistive Evolution of a Force-Free Magnetic Flux Rope in a Potential Ambient Field

Let us consider a magnetic flux rope which came to equilibrium with the surrounding external field after the relaxation to the state (13) in layers above the photosphere. We assume that the length of the relaxed part of the rope is only slightly longer than its cross-section radius, i.e., the rope has not yet emerged far into the corona. This means that we are considering low-lying magnetic loops with fairly strong magnetic fields, making it possible to neglect the pressure differences $P_{ex} - \langle P_{in} \rangle$ and $P_{ex} - P_{in}(a)$ in the formulae (6) and (14) and to use the force-free approximation (Fig.2a).

In resistive plasma, the parameters of the force-free solution (13) will slowly vary with the time: $B_0 = B_0(t)$ and $\alpha = \alpha(t)$. Thus, here we consider the problem of relatively slow dissipative evolution of the relaxed part of the flux rope filled with the field (13) in a passive external medium. The system of MHD equations, whose solution we seek, has the following form. The equation of the momentum of the medium in the force-free approximation is

$$\rho \left(\frac{\partial \mathbf{V}}{\partial t} + (\mathbf{V} \nabla) \mathbf{V} \right) = (4\pi)^{-1} [\text{rot } \mathbf{B} \times \mathbf{B}]. \quad (17)$$

The induction equation is

$$\frac{\partial \mathbf{B}}{\partial t} = -\text{rot}(v_m \text{rot } \mathbf{B}) + \text{rot}[\mathbf{V} \times \mathbf{B}], \quad (18)$$

where $v_m = c^2(4\pi\sigma)^{-1}$ is the magnetic viscosity, σ is the conductivity of the plasma. Finally, the continuity equation is

$$\frac{\partial \rho}{\partial t} + \text{div}(\rho \mathbf{V}) = 0. \quad (19)$$

We will seek a solution for the velocity field in the form $\mathbf{V} = \{V_r(r, t), 0, V_z(z, t)\}$, where

$$V_r(r, t) = -\gamma(t)r, V_z(r, t) = -\gamma(t)z. \quad (20)$$

and the function $\gamma(t)$, which has the dimension of inverse time, must be determined. Substituting (20) into (17), we find that the inertial term in (17) vanishes, so that $[\text{rot } \mathbf{B} \times \mathbf{B}] = 0$, i.e. the force-free field distribution with the variables $B_0(t)$ and $\alpha(t)$ will be maintained inside the rope over the time if $\gamma(t)$ obeys the condition

$$\frac{d\gamma}{dt} = \gamma^2. \quad (21)$$

This equation has the simple solution

$$\gamma(t) = \frac{\gamma_0}{1 - \gamma_0 t}, \quad (22)$$

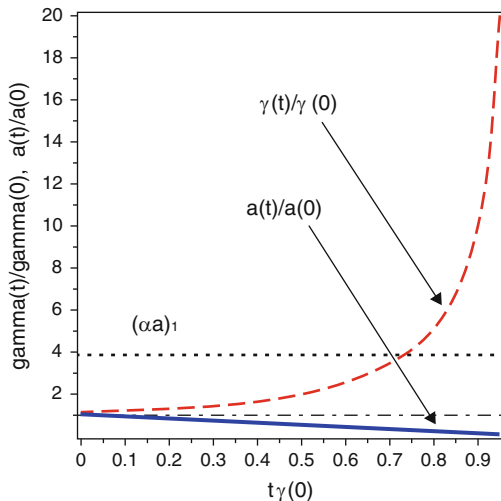
where $\gamma_0 = \gamma(0) = \text{const}$ is an unknown constant. As we see, over a finite time $t \rightarrow t_0 = \gamma_0^{-1}$ the function $\gamma(t)$ reaches infinitely large values.

Equation (18) written in terms of its components takes the form

$$\frac{\partial B_\phi}{\partial t} = -v_m \alpha^2 B_\phi - \frac{\partial}{\partial r}(V_r B_\phi) - \frac{\partial}{\partial z}(V_z B_\phi), \quad (23)$$

$$\frac{\partial B_z}{\partial t} = -v_m \alpha^2 B_z - \frac{1}{r} \frac{\partial}{\partial r}(r V_r B_z). \quad (24)$$

Fig. 3 The time functions $\gamma(t)$, $a(t)$ are given in dimensionless units. The dimensionless radius of the sphere is indicated by the dotted line for the case $(\alpha a)_1 = \text{const} = 3.8317$. The time on the horizontal axis is expressed in the units $\tau_c \equiv \gamma_0^{-1}$, and varies within the range 0–0.95. When $t\gamma_0 \rightarrow 1$, the function $\gamma(t)/\gamma(0)$ tends to infinity, and $a(t) \rightarrow 0$



Substituting the distribution (13) with the time-dependent variables $B_0(t)$ and $\alpha(t)$, together with the components of the plasma velocity in the form (20), we obtain from (23), as well as from (24), the same system of two ordinary differential equations

$$\gamma = \frac{1}{\alpha} \frac{d\alpha}{dt}, \quad (25)$$

and

$$\frac{\partial \ln B_0(t)}{\partial t} = -\frac{c^2}{4\pi\sigma} \alpha^2 + 2\gamma, \quad (26)$$

which define the dependences $\alpha(t)$ and $B_0(t)$. It follows from (22) and (25) that

$$\alpha(t) = \frac{\alpha_0}{1 - \gamma_0 t}. \quad (27)$$

Since the rope boundary is specified as $(\alpha a)_n = x_n = \text{const}$, it is obvious that $a(t) = a_0(1 - \gamma_0 t)$ and $V_r(a) = -\gamma(t)a(t) = -\gamma_0 a_0 = \text{const}$, i.e. the transverse radial compression of the rope occurs uniformly, at a constant rate (Fig.3).

The continuity equation (19) takes the form

$$\frac{\partial \rho}{\partial t} + \frac{\rho}{r} \frac{\partial}{\partial r} (r V_r) + \rho \frac{\partial}{\partial z} V_z = 0 \quad (28)$$

and, upon substitution of (20), yields the solution for the density inside the rope $\frac{\partial \ln \rho(t)}{\partial t} = 3\gamma$, whence we obtain

$$\rho(t) = \frac{\rho_0}{(1 - t/t_0)^3}. \quad (29)$$

According to the requirement for the normal component of the velocity to be continuous at the boundary, we will assume that outside the rope,

$$V_r(r, t) = -\gamma(t)a^2/r, \quad V_z(r, t) = 0. \quad (30)$$

Then, from the induction and continuity equations,

$$\frac{\partial B_{z,ex}}{\partial t} = -\frac{1}{r} \frac{\partial}{\partial r}(r V_r B_{z,ex}), \quad \frac{\partial \rho_{ex}}{\partial t} + \frac{\rho_{ex}}{r} \frac{\partial}{\partial r}(r V_r) = 0 \quad (31)$$

a simple external solution follows:

$$B_{\phi,ex} = 0, \quad B_{z,ex} = \text{const} \neq 0, \quad \rho_{ex} = \text{const}.$$

The continuity condition for the total pressure should be satisfied at the movable boundary of the rope ($\alpha a = x_n$):

$$B_{z,ex}^2 = B_0^2(t) J_0^2(x_{1,n}) + 8\pi(P(a) - P_{ex}). \quad (32)$$

If we neglect the contribution of the gas pressure and take $B_{z,ex} = \text{const}$, we will find that $B_0 = \text{const}$ as well. Then, it follows from (26) that

$$\frac{c^2}{4\pi\sigma} \alpha^2 = 2\gamma. \quad (33)$$

Taking into account (22) and (27), this condition is satisfied if the conductivity grows, as the plasma in the compressed rope is heated:

$$\sigma(t) = \sigma_0 \frac{\alpha(t)}{\alpha_0}. \quad (34)$$

Using (33), we estimate the collapse time t_0 in this case as

$$t_0 = \gamma_0^{-1} = \frac{8\pi\sigma_0 a_0^2}{c^2 x_n^2}. \quad (35)$$

If we do not resort to the hypothesis (34), and assume that the conductivity of the plasma remains constant with the time, then (26) should be used to calculate the dependence $B_0(t)$, and (32) to calculate the time behavior of the pressure difference $8\pi[P(a(t)) - P_{ex}]$. However, we believe that the dependence (34) better corresponds to the nature of the process under consideration.

Let us now present a numerical estimate for the compression time t_0 . Under the assumption of the fine structure of the magnetic field in the solar chromosphere we take the characteristic scale a_0 (the radius of the rope cross-section after its relaxation to a linear, force-free field) to be of 300 km (3×10^7 cm), and the

conductivity of the solar plasma in the region of the temperature minimum and in the transition region of the order of 10^{11} s^{-1} [5, 8]; $\sigma_0 = 3 \times 10^{11} \text{ s}^{-1}$. Then,

$$t_0 = \gamma_0^{-1} = \frac{75 \times 10^{11} 10^{15}}{10^{21} x_n^2} = \frac{7.5 \times 10^6}{x_n^2} \text{ s}. \quad (36)$$

The choice of the root of the function $J_1(\alpha a)$ from the sequence (16) strongly influences numerical estimates for t_0 . As it was noted above, the numerical value of this root is determined by the magnitude of the helicity K of the field in the rope, which is unknown. If this helicity is high, then, possibly, $x_n^2 \gg 1$. For example, if we take the fifth term in the sequence (16), $x_n = 16.472$, we obtain $t_0 = 2.8 \times 10^4 \text{ s}$, less than 8 h. This comes about in spite of the fact that, in accordance with (34), we have allowed a possible sharp growth in the plasma conductivity with the time! It is remarkable that the formula (35) for the time of compression contains the plasma conductivity only for the initial state of the plasma, i.e. σ_0 , the smallest of its values. This “choice” of the minimum value of plasma conductivity is well consistent with the general principle of maximum entropy production in any dissipative system; therefore, we can assume that the regime (34) is the most likely of all. The rate of the radial compression is $V_r(a) = -\gamma_0 a_0 = -\frac{c^2 x_n^2}{8\pi \sigma_0 a_0} = -\frac{a_0}{t_0}$. Numerically it yields $-3 \times 10^7 / 3 \times 10^4 \approx -10^3 \text{ cm/s}$; this is substantially lower than the Alfvén velocity, which is approximately equal to 10^8 cm/s in layers above the photosphere (for a field of 300 G and a plasma number density $\approx 10^{12} \text{ particles/cm}^3$).

6 The Open Structure of a Magnetic Flux Rope, Plasma Jets Along the Axis, and Separatrix Surfaces

Up to now, we have been discussing the uniform compression of the magnetic flux rope without taking into account the fact that the system has an axis of symmetry and separatrix surfaces, in which the magnetic field is strictly longitudinal. It means that the system is open along these directions. In the diffusion equation (18), the term $\text{rot}[\mathbf{V} \times \mathbf{B}]$ contains, in fact, only the velocity transverse to the magnetic field; therefore, the dissipative evolution of the magnetic field in a passive medium cannot determine the plasma flow directed strictly along the field. On the axis of symmetry of the system and on separatrix surfaces, the magnetic force has no component along the longitudinal direction, regardless of the field strength; therefore, to calculate the plasma motion along these directions we should take into account the gradient of the gas pressure in the equation of momentum (17), no matter how small the plasma parameter $\beta \equiv 8\pi P / B^2$ is. (If we denote the direction along the axis of the sphere by the letter z , then the equation of momentum in this direction will have the form $\rho(\partial V_z / \partial t + V_z \partial V_z / \partial z) = -\partial P / \partial z$). In the limit, the effect discussed here can be described by specifying the velocity field in the form:

$$\mathbf{V}(r, t) = -\gamma(t)r\mathbf{e}_r + [-\gamma(t)z + f(t, z)\delta(J_1(\alpha r))]\cdot\mathbf{e}_z + 0\cdot\mathbf{e}_\phi. \quad (37)$$

Here, $\delta(J_1(\alpha r))$ is the delta-function of the argument $J_1(\alpha r)$, and $f(r, t)$ is a positive function at $z > 0$ and negative at $z < 0$. In this case, all the above results remain valid, as in the formula (23) the longitudinal component of the velocity V_z is used in combination only with the factor $J_1(\alpha r)$. According to the well-known property of the delta function ($x\delta(x) = 0$), addition of the new singular term with the delta function $\delta(J_1(\alpha r))$ to the expression (37) does not affect anything except the axis of symmetry and separatrix surfaces. When using (37) instead of (20), we obtain a qualitatively correct description of the phenomenon: the plasma is ejected from the uniformly contracting flux rope precisely along the axis of its symmetry and the separatrix surfaces (see Fig.4), but the velocity of these jets will remain indefinitely large and the plasma density at the axis will tend to zero. In order to estimate realistically the ejection velocity V_{jet} , maintaining its connection with the physical parameters of the problem, we will consider the case of a steady gas outflow from a magnetic flux rope in a very thin cylinder of a small radius Δ and a length of $2L$, whose axis of symmetry coincides with that of the sphere. We will write the condition of the mass flux balance: the amount of gas entering the cylinder due to uniform compression of the rope per a unit of time through the lateral surface $2\pi\Delta 2L\rho(t)V_n$ (where V_n is the component of the compression rate normal to the surface of the cylinder) should be equal to the mass of gas flowing through both ends of the cylinder: $2\pi\Delta^2\rho_0 V_{jet}$. (We assume that on the axis, at the output from the rope $z = \pm L$, the initial plasma density ρ_0 is maintained). The radial velocity V_n is obviously equal to $V_n = \gamma\Delta$. Accounting for the mass balance, we obtain:

$$2\gamma L\rho(t) = \rho_0 V_{jet}. \quad (38)$$

In this expression, the unknown infinitely small radius of the virtual cylinder Δ is eliminated, and using (34) we obtain exponentially rapid increase in the ejection velocity with the time:

$$V_{jet} = 2V_z(L)\exp\left(\frac{3\gamma_0 t}{1 - \gamma_0 t}\right). \quad (39)$$

7 The Further Evolution of a Magnetic Flux Rope

According to the obtained solution, in addition to the radial compression, longitudinal flows of plasma also arise in the dissipative magnetic flux rope (Figs. 4 and 5). The velocities of the flows grow linearly with z , $V_z(r, t) = -\gamma(t)z$. At the edges of a magnetic loop of a length $2L$, the rate of the longitudinal constriction of the rope is $V_z(L, t) = -\frac{\gamma_0 a_0 L}{2a_0(1-t/t_0)}$. It will inevitably reach the Alfven velocity, so that the obtained solution is no longer valid, as $t \rightarrow t_0$.

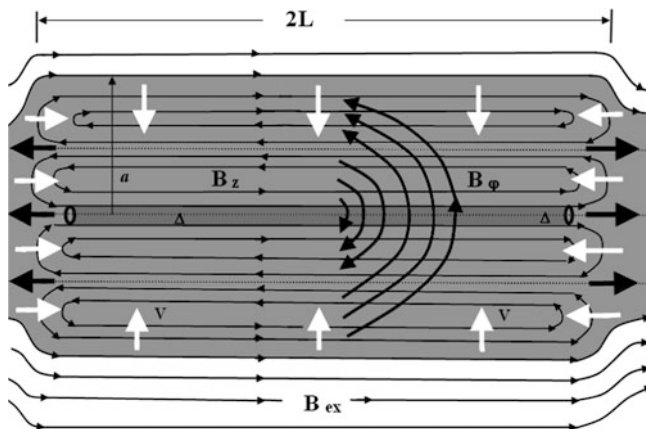


Fig. 4 The inner structure of the relaxed part of a flux rope (the curvature is neglected) described by (13), for the case $(\alpha a)_2 = x_2 = 7.0166$, when B_z changes its sign twice along the radius a , while B_ϕ only once. B_ϕ -components are shown in the central part of the figure by *thick solid lines*. The rates of dissipative contraction of the rope are shown by *thick white arrows*, and *thick black arrows* mark the plasma flows $\mathbf{V}_{jet}$ along the axis of symmetry and along the separatrix surface, where $J_1(\alpha r) = 0$ (their positions are indicated by *thin dot-and-dashed lines*). The cylinder of a small radius Δ , concentric to the rope axis, is painted in a *darker color* (the model of this virtual cylinder is used in Sect. 6 to calculate the rate of the plasma outflow along the rope axis)

Taking into account this circumstance and the discussion in Sect. 3 concerning the possible dynamics of twisted magnetic-field tubes, we can imagine two fundamentally different scenarios for the behavior of a magnetic flux rope with a force-free internal field, depending on the structure of the external magnetic field.

Firstly, if the external field is unable to confine the magnetic flux tube near a photospheric layer, the twisted flux rope will be ejected into the corona, greatly increasing its length. However, due to the very large difference (roughly five orders of magnitude!) between the diffusion and the Alfvén velocities, the velocity V_z remains appreciably lower than V_A for a substantial time. This means that the solution found above remains valid, and the transverse compression of the magnetic flux rope continues, ultimately leading to the appearance of fine-structure magnetic elements in the corona.

In the second scenario, in the presence of a transverse component of the ambient magnetic field, the magnetic flux rope can be confined in relatively low chromospheric layers. As the twisting of the field in the force-free rope is below its equilibrium value ($\kappa = 1$), it is subject to the dynamical (not only dissipative!) longitudinal shrinkage. In this case, the relaxed part of the magnetic flux rope can take the shape of a magnetic sphere with a force-free internal structure in the form of magnetic toroids inscribed in spherical layers ([1], Fig. 6).

On the symmetry axis of this sphere, the magnetic field is strictly radial and sign-alternating along the radius. The external magnetic field around the sphere is

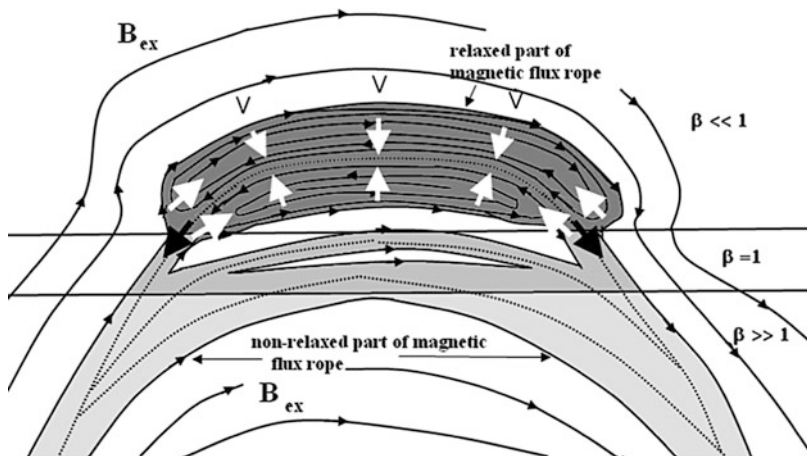


Fig. 5 Here, $J_0(\alpha r)$ and $J_1(\alpha r)$ are Bessel functions of the first kind of the zeroth and first order, respectively, and α is a constant with the dimension of inverse length. The product αa is the dimensionless coordinate of the lateral surface of the magnetic flux rope. Its numerical value is determined by the initial helicity of the magnetic field K (i.e., the total coupling of the fluxes of the longitudinal $B_z(r)$ and azimuthal $B_\varphi(r)$ fields) in the relaxing part of the magnetic flux tube

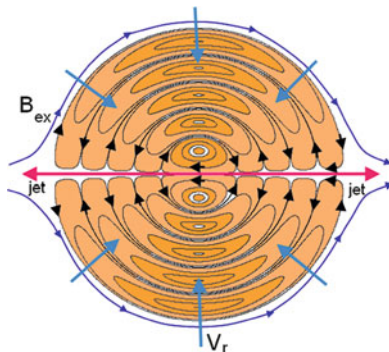


Fig. 6 The meridional cross-section of a force-free magnetic sphere for the case $\alpha R = 17.22$ (five magnetic toroids inscribed in a sphere). On the axis of symmetry, the magnetic field is radial and sign-alternating along the radius. The external potential field B_{ex} provides a pressure balance on the boundary. The azimuthal field B_φ , whose *ring-shaped* lines encompass the axis of symmetry, is not shown in the figure. This field changes its sign at the spherical surfaces that separate one magnetic toroid from another. The *radial* arrows indicate uniform compression of the magnetic sphere

potential. It is remarkable that an exact analytical solution similar to that presented above can be obtained for such a force-free magnetic sphere, which describes its self-similar uniform dissipative compression (Fig. 6). The resistive evolution of a force-free spheroid will be analyzed in a special paper.

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Solar Activity Indices in the Cycles 21–23

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Abstract A stable cyclicity of correlation coefficients K_{corr} for some solar activity indices versus $F_{10.7}$ was found after monthly averaged values analysis. These indices are: Wolf numbers, 10.7 cm radio flux $F_{10.7}$, 0.1–0.8 nm background, the total solar irradiance (TSI), Mg II UV-index (280 nm core to wing ratio) and monthly flare count/10. The correlation coefficients of the linear regression of these solar indices versus $F_{10.7}$ were analyzed for every year in solar cycles 21, 22 and unusual cycle 23. We found out that the values of yearly determined correlation coefficients K_{corr} for solar activity indices versus $F_{10.7}$ show the cyclic variations with stable period closed to half length of 11-year cycle.

1 Introduction

We have studied monthly averaged values of six global solar activity indices in cycles 21, 22 and 23. Most of these observed data we used in our paper were published in Solar-Geophysical Data bulletin.

The unusual cycle 23 was examined carefully: the rise, decline, minimum and maximum phases were studied separately. This study made it possible to determine that correlation coefficients of linear regression for some solar activity indices versus $F_{10.7}$ have the different values for different cycle's phases. In [7] there were analyzed solar cycles 18–20 and pointed out that correlation for spot numbers versus $F_{10.7}$ does not show the close linear connection during all the activity cycle. To achieve to

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best agreement in approximation of spot numbers values by $F_{10.7}$ observations it was proposed to approximate the dependence $W-F_{10.7}$ by two linear regressions [7]: for the low solar activity and for the high activity ($F_{10.7}$ more than 150). In our paper we found out that the linear correlation was violated not only for maximums of solar activity cycles but for minimums of the cycles too.

2 Global Activity Indices

Then we have to say a few words about solar indices studied in this paper.

The Wolf numbers is a very popular, widely used solar activity index: the series of Wolf numbers observations continue more than 200 years.

The solar radio microwave flux at wavelengths 10.7 cm $F_{10.7}$ has also the longest running series of observations started in 1947 in Ottawa, Canada and maintained to this day at Penticton site in British Columbia. This radio emission comes from high part of the chromosphere and low part of the corona. $F_{10.7}$ radio flux has two different sources: thermal bremsstrahlung and gyro-radiation. These mechanisms give rise to enhanced radiation when the temperature, density and magnetic fields are enhanced. So $F_{10.7}$ is a good measure of general solar activity.

Mg II index derived from daily solar observations of the core-to-wing ratio of the Mg II doublet at 279.9 nm provides a good measure of the solar UV variability and can be used as a reliable proxy to model extreme UV (EUV) variability during the solar cycle [5, 6]. The Mg II observation data were obtained from several satellite's (NOAA, ENVISAT) instruments. NOAA started in 1978 (during the 21st, 22nd and the first part of the 23rd solar activity cycles), ENVISAT was launched on 2002 (last part of the 23th solar activity cycle). We used both the NOAA and ENVISAT Mg II index observed data.

Data of GOES observations of the X-ray 0.1–0.8 nm background were taken from Solar Geophysical Data bulletin. This permanent 0.1–0.8 nm background monitoring of solar disk at the 0.1–0.8 nm range is a good indicator of solar corona activity without flares.

Data of the TSI SOHO observations for the 23rd cycle and earlier TSI observations from 1985 to 2000 (when TSI was observed by Earth Radiation Budget Satellite—EBRS) were also taken from Solar-Geophysical Data bulletin.

3 Global Activity Indices in the Cycle 23

It is well known that the last cycle 23 was unusual among most of solar activity indices. Figure 1 demonstrates that for all activity indices in the 23rd solar activity cycle one can see two maximums separated one from another on 1.5 year approximately. We see that similar double-peak structure exists in cycle 22 but in the cycle 21 the double-peak structure is not evident. We see that there are

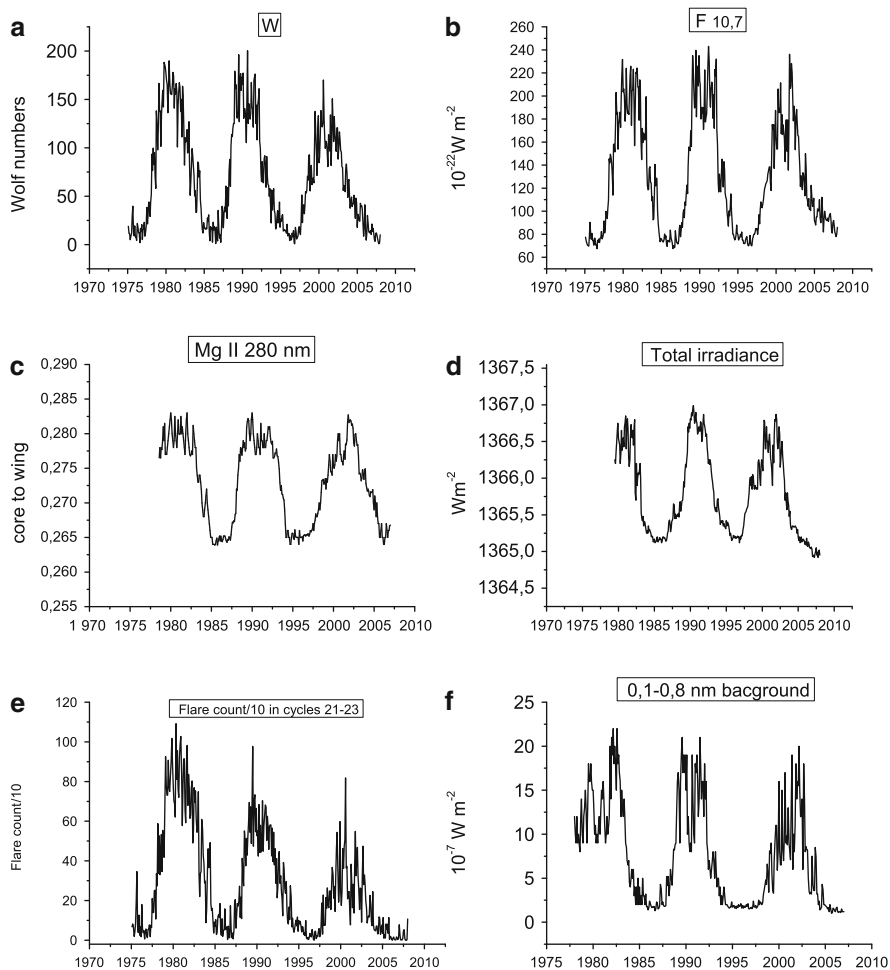


Fig. 1 The monthly averages for (a) Wolf numbers, (b) $F_{10.7}$ (high chromosphere and low corona), (c) Mg II core to wing ratio (chromosphere), (d) TSI (photosphere), (e) flare count/10 and (f) 0.1–0.8 nm (corona) in the cycles 21, 22 and 23

displacements in both maximum occurrence time of all these indices in the 23rd solar cycle. In this unusual cycle 23 the monthly averages values for Wolf numbers during 8 months exceeded 113 and most of sunspot groups were less in size, their magnetic fields were less composite and characterized by the greater lifetime near 2nd maximum in comparison with values near the 1st maximum. In the cycle 23 the $F_{10.7}$ radio flux and the TSI have the lowest values from 2007 to 2009 (the beginning of the cycle 24) all over of these indices observation period.

We assume that the probable reason of such double-peak structures is a modulation of the 11-year fluxes variations by both of the quasi-biennial and 5.5 year

Table 1 Correlation coefficients for activity indices versus $F_{10.7}$ at the rise, decline and maximum phases of the cycle 23

Correlation between activity indices	Rise phase of the cycle 23	Decline phase of the cycle 23	Maximum of the cycle 23	All over the cycle
$W-F_{10.7}$	0.919	0.961	0.742	0.939
$Mg\ II-F_{10.7}$	0.963	0.964	0.757	0.879
$TSI-F_{10.7}$	0.879	0.949	0.743	0.920
$0.1-0.8\ nm-F_{10.7}$	0.899	0.814	0.773	0.812
$Flare\ count/10-F_{10.7}$	0.905	0.895	0.785	0.890

cyclicly. The different time of 1st and 2nd maximums appearance may be caused by the difference in fluxes formation conditions (for our indices) at different atmosphere's altitudes of the Sun.

Figure 1 also shows that for all solar indices in the cycle 23 the relative depth of the cavity between two maximums is about 10–15%.

When studied five activity indices in 23rd solar activity cycle we separated out rise phase (from October 1997 to November 1997), cycle maximum phase (from November 1997 to July 2002) and decline phase (from Jul 2002 to Jan 2006). The results of correlation study of solar activity indices versus $F_{10.7}$ in cycle 23 are presented in Table 1. We see that the maximum values of correlation coefficients K_{corr} reached for the rise and decline phases of the cycles. According to our calculations the highest values of correlation coefficients K_{corr} we see in connection between W and $F_{10.7}$. Correlation coefficients K_{corr} between 0.1 and 0.8 nm background flux versus $F_{10.7}$ are minimal of all correlation coefficients determined here.

The cyclic variation of fluxes in different spectral ranges and lines at the 11-year time scale are widely spread phenomenon for F, G and K stars (not only for the Sun). The chromospheric activity indices (radiative fluxes at the centers of the H and K emission lines of Ca II—396.8 and 393.4 nm respectively) for solar-type stars were studied during HK project by [2] at Mount Wilson observational program during 45 years, from 1965 to the present time. Authors of the HK project supposed that all the solar-type stars with well determined cyclic activity about 25% of the time remain in the Maunder minimum conditions. Some scientists proposed that the solar activity in future cycle 24 will be very low similar to activity during the Maunder minimum period.

4 Correlation Study of Solar Activity Indices in the Cycles 21–23

We analyzed the interconnection between activity indices W , $Mg\ II$ core to wing ratio, TSI and numbers of flares versus $F_{10.7}$ for the 21st, 22nd and 23th solar cycles. At Fig. 1 we can see the correlation between solar indices and radio flux $F_{10.7}$ in the

Table 2 Coefficients of linear regression (intercept— a_{ind}^{cycle} and slope— b_{ind}^{cycle}) for solar activity indices versus $F_{10.7}$

Solar activity indices	The cycle 21	The cycle 22	The cycle 23	Averaged standard error σ
W^{cycle}	a = -66.4 b = 1.11	a = -56.43 b = 1.012	a = -67.27 b = 1.001	$\bar{\sigma}_a = 3.0$ $\bar{\sigma}_b = 0.02$
$F_{Mg II}^{cycle}$	a = 0.257 b = 1.20E-04	a = 0.256 b = 1.09E-04	a = 0.257 b = 1.20E-04	$\bar{\sigma}_a = 5.25E-04$ $\bar{\sigma}_b = 4.0E-06$
F_{TSI}^{cycle}	a = 1,364.312 b = 0.011	a = 1,364.044 b = 0.013	— —	$\bar{\sigma}_a = 0.058$ $\bar{\sigma}_b = 4.4E-04$
$F_{flare\ count/10}^{cycle}$	a = -35.38 b = 0.58	a = -21.33 b = 0.39	a = -25.88 b = 0.33	$\bar{\sigma}_a = 2.05$ $\bar{\sigma}_b = 0.014$

cycles 21–23. The interconnection between activity indices corresponds to the linear regression equation:

$$F_{ind}^{cycle} = a_{ind}^{cycle} + b_{ind}^{cycle} \cdot F_{10.7}. \quad (1)$$

where F_{ind}^{cycle} is the activity index flux,

a_{ind}^{cycle} is the intercept of linear regression,

b_{ind}^{cycle} is the slope of linear regression.

In Table 2 we present coefficients of linear regression. These results from Table 2 show that coefficients of linear regression (a_{ind}^{cycle} and b_{ind}^{cycle}) for solar activity indices versus $F_{10.7}$ differ for different activity cycles (21–23). In Table 2 we also present standard errors σ of intercept and slope values averaged among the cycles 21–23 (because σ_a and σ_b differ negligible for these cycles). One can also see at Fig. 1 that the difference of regression coefficients of the linear regression for solar activity indices versus $F_{10.7}$ is the most significant for cycle 23.

Many authors pointed that there was very high level of flared activity in cycle 21 and very low level of flared activity in cycle 23. We see (Table 2) that the difference of $a_{flare\ count/10}^{23}$ and $b_{flare\ count/10}^{23}$ values from similar values, determined for cycles 21 and 22, is more significant among all the different cycle's regression coefficients. Figure 2 also demonstrates that the flared activity in the 23rd cycle almost twice weaker (flare count/10 versus $F_{10.7}$) in comparison to 21st cycle.

We have to point out that close interconnection between radiation fluxes characterized the energy release from different atmosphere's layers is the widespread phenomenon among the stars of late-type spectral classes. It was shown in [3] for solar-type stars of F, G, K and M spectral classes that the summary areas of spots and values of X-ray fluxes increase gradually from the sun and HK project stars with the low spotted discs to the highly spotted K and M-stars for which in [1] the zonal model of the spots distributed at the star's disks was constructed.

We've also calculated yearly averaged values K_{corr} of linear regression for solar activity indices versus $F_{10.7}$ for cycles 21, 22 and 23. Yearly averaged values K_{corr} were determined for each year (we use the monthly averaged values for 12 months of the year with half-monthly time interval, 24 values for every year in all).

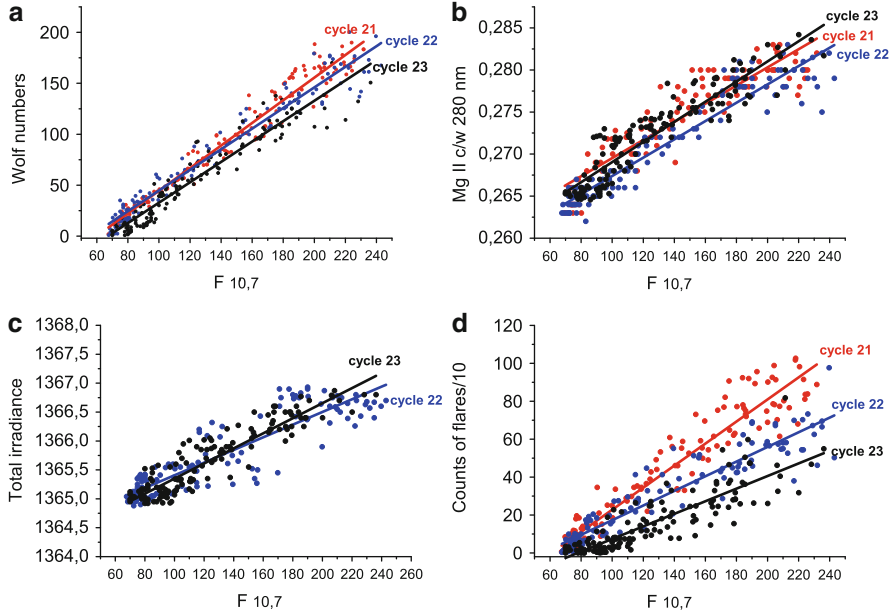


Fig. 2 Correlation between monthly averages for solar indices versus $F_{10.7}$ radio flux in cycles 21–23. (a) Wolf numbers, (b) Mg II core to wing ratio, (c) TSI and (d) flare count/10

Figure 3 demonstrates the results of our correlation calculations of these solar activity indices versus $F_{10.7}$ (K_{corr} variations during the cycles 21–23). We can see that all the K_{corr} have the maximum values at the rise and decline phases. It's easy to estimate the value of period of K_{corr} cyclic variations as 5.5 years approximately. We assumed that this new cyclicity (characterized with period's value equal to half length of 11-year cycle) is important for the successful forecasts of the solar activity indices fluxes.

The cyclic behavior of K_{corr} can be explained by following assumption: we imagine that some activity index consists of:

$$F_{ind}(t) = F_{ind}^{background}(t) + \Delta F_{ind}^{AR}(t). \quad (2)$$

where $F_{ind}^{background}(t)$ is the background flux rising continuously with increasing of solar activity and $\Delta F_{ind}^{AR}(t)$ is the additional flux from the active regions.

We consider that:

$$F_{ind}^{background}(t) = a_1 + b_1 \cdot F_{10.7}^{background}(t). \quad (3)$$

$$\Delta F_{ind}^{AR}(t) = a_2 + b_2 \cdot \Delta F_{10.7}^{AR}(t). \quad (4)$$

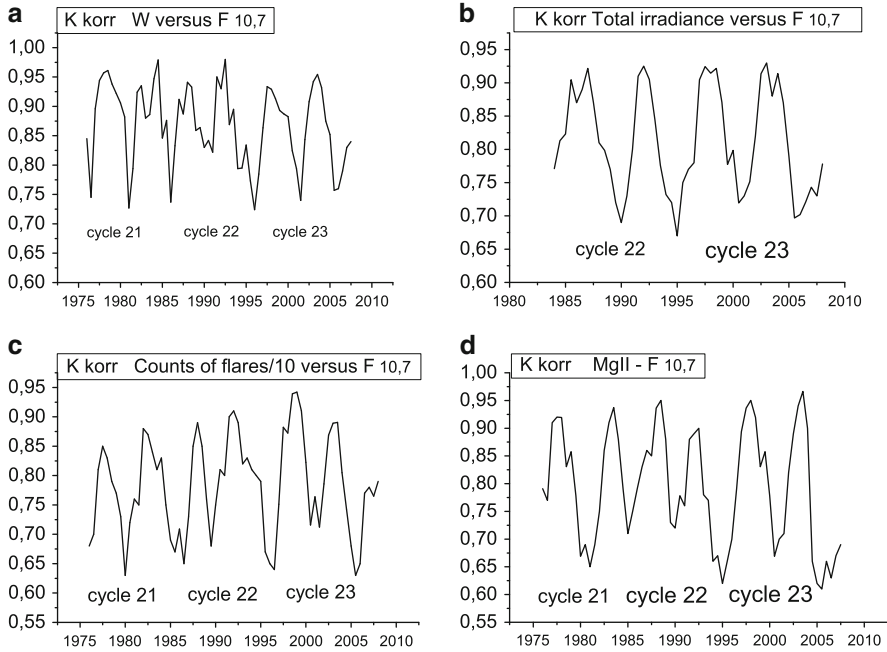


Fig. 3 Yearly calculated correlation coefficients of linear regression K_{corr} for (a) Wolf numbers, (b) TSI, (c) flare count/10 and (d) Mg II UV-index versus $F_{10.7}$ in solar cycles 21, 22 and 23

The coefficients a_1 and b_1 vary from a_2 and b_2 in different power for our different activity indices. For Wolf numbers this difference is small, but for 0.1–0.8 nm background and number of flares the difference between a_1 , b_1 and a_2 , b_2 is more significant than for W and TSI.

During the rise and decline cycle's phases the dependence $F_{ind}(t)$ versus $F_{10.7}(t)$ is approximately linear because coefficients a and b from Table 2 (which described all the cycle) are close to a_1 and b_1 (see Fig. 2) and relative addition flux from active regions $\Delta F_{ind}^{AR}(t)$ is neglect with respect to $F_{ind}^{background}(t)$. So additional flux from active regions cannot destroy a balance in the close linear correlation between $F_{ind}(t)$ and $F_{10.7}(t)$ and respective values of K_{corr} reach their maximum from all the cycle.

During the minimum of activity cycle both values $F_{ind}^{background}(t)$ and $\Delta F_{ind}^{AR}(t)$ are small, but additional flux from active regions is not neglect in relation to background flux that has the minimum values from all the activity cycle. Therefore values a and b from Table 2 cannot describe the linear regression to a considerable degree in cycle's minimum and values of K_{corr} reach their minimum values in the cycle.

During the maximum of activity cycle $\Delta F_{ind}^{AR}(t)$ often exceeds $F_{ind}^{background}(t)$ so disbalance in linear regression between activity indices increases and values of K_{corr} reach their minimum values in the activity cycle too.

5 Conclusions

For a long time the scientists were interested in the simulation of processes in the earth's ionosphere and upper atmosphere. It's known that the solar radiance at 30.4 nm is very significant for determination of the Earth high thermosphere levels heating. In [5, 6] it was showed that the for solar 30.4 nm radiance fluxes forecasts (very important for Earth thermosphere's heating predictions) there were more prefer to use Mg II 280 nm observed data unlike usual $F_{10.7}$ and Wolf number observations. In [4] for these purposes have developed a two-component model of solar emission in the EUV range (10–105 nm). It was shown that the intensity ratios of individual spectral lines depend only on the solar activity level.

In this paper we found out the cyclic behavior of yearly values of correlation coefficients K_{corr} of linear regression for W, TSI, Mg II 280 nm and flare count/10 versus $F_{10.7}$ during solar activity cycles 21, 22 and 23 (see Fig. 3). Since we show that yearly values of K_{corr} have the maximum values at the rise and decline phases so the linear connection between indices is more strong in these cases. It means that the forecasts of solar indices, based on $F_{10.7}$ observations in our case, will be more successful at rise and decline cycle's phases.

We also determined that the yearly values of K_{corr} are characterized by cyclic variations with the period that is equal to half length of period at 11-year time scale (5.5 years about). Our study of linear regression between solar indices and $F_{10.7}$ confirms the fact that at minimum and at maximum cycle's phases the nonlinear state of interconnection between solar activity indices (characterized the energy release from different layers of solar atmosphere) increases.

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Coronal Mass Ejections on the Sun and Their Relationship with Flares and Magnetic Helicity

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Abstract A brief review of results of observations of coronal mass ejections (CMEs), obtained aboard SOHO and STEREO during last decades, is presented. CMEs velocity, acceleration, mass and angular width are considered in connection with the flares arising simultaneously with the CMEs from common solar regions. Statistically properties of associated CMEs-flares are related. So, higher the flux of the associated flare is, higher the CME mass, angular width and velocity are. The magnetic helicity of an active region (AR) seems to play an important role whether a confined or a CME-associated flare will be produced. The dynamical evolution of CMEs shows that the total energy is kept relatively constant and the energy released in radio and other forms of radiations is not significant. Slow CMEs could propagate through heliosphere with a deflection of $2 - 30^\circ$ both in latitudinal and azimuthal direction obeying an interaction with the surrounding magnetic field. Faster CMEs tend to propagate radially. Data from published scientific papers and Internet have been used.

1 Introduction

Since the first observations of coronal mass ejections with OSO-7 the occurrence of CMEs in association with flares and eruptive filaments is well known and continuously investigated [1–12]. CMEs origins are due to magnetic reconnection in low corona and accompanied by a global reconstruction of magnetic field in corona.

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Temporally and spatially correlated CMEs and solar flares, arising from common ARs, are called associated CMEs/flares.

The results of white-light observations with the LASCO (Large Angle Spectroscopic Coronagraph) aboard SOHO (Solar and Heliospheric Observatory) and coronagraphs COR1, COR2 and Heliospheric imagers HI-1 and HI-2 aboard STEREO (Solar Terrestrial Relations Observatory) A and B are used. Two coronagraphs C1 and C2 LASCO with fields of view (FOVs) $2.2 - 6.0 R_s$ and $4.0 - 30.0 R_s$ correspondingly give images at the distances up to $\sim 30 R_s$. From STEREO a CME could be followed up to Earth and farther. STEREO/COR1 has FOV $1.5 - 4.0 R_s$ and COR2 has FOV $2.0 - 15 R_s$, Heliospheric imager HI1 is pointed 13.2° away from the Sun showing a 20° field near the ecliptic plane ($\sim 15 - 90 R_s$, HI2 is pointed 53.4° from the Sun and has a 70° field centered around the ecliptic plane ($\sim 90 - 300 R_s$). Statistical relations between different properties of CMEs and associated flares are considered. Analysis of correlations is useful in understanding of space weather and for prediction of CMEs and Solar Energetic Particles (SEP) events occurrence.

2 CMEs Characteristics in the Nearest Heliosphere

Observations aboard SOHO gives a rich information in the nearest heliosphere up to distances of $20-30 R_s$. CMEs properties vary in large range inside the LASCO FOV. At the initiation phase the CME velocities are mainly low ($10 - 100 \text{ km s}^{-1}$) increasing during acceleration phase and reaching their peaks. The CME accelerations range from several m s^{-2} to several dozens or hundreds m s^{-2} and more. So the CME on 1997 November 6 had a very large acceleration $a \sim 7,300 \text{ m s}^{-2}$ and peak velocity $V \sim 2,150 \text{ km s}^{-1}$ [6]. Unsuccessful CMEs could after deceleration return to the solar surface. The ranges of CME properties are presented in Table 1. In the first column the considered properties are listed, in the second one their ranges are given composed by us on the results from [6, 12–14]. In the third, fourth and fifth columns the averages according to data from [9, 13, 14] are presented. Periods of observations: 1996–2003 [9]; 1996–2008 for all CMEs and CMEs with angular width $W > 30^\circ$ [13]; 1996–2006, V for all flare-associated CMEs, a for M class and stronger, W for B-X class flare-associated CMEs correspondingly [14]. We must remember that the values of V , a and W are suffered from the projection effect. Daily rate (DR) means numbers of CMEs per day. DR ranges from solar minimum to solar maximum.

3 Relations Between CMEs and Flares Properties: CMEs Velocities with LASCO SOHO and STEREO

Detailed analysis based on comprehensive series of observations was fulfilled by Aarnio et al. [14] who used the LASCO CME catalog and GOES (Geostationary Operational Environmental Satellite) flare database. During 1996–2006 13,862

Table 1 CME properties

Characteristics	Range	Average value		
Velocity V km s ⁻¹	$20 < V < 3,000$	489	466–470	495 ± 8
Acceleration a m s ⁻²	$-218 < a < 7,300$	1.6×10^{15}	41–60	-3.7 ± 1.3
Energy E erg	$10^{27} < E < 10^{33}$	47		42–80
Mass M g	$10^{12} < M < 10^{16}$			
Angular width W°	$5 < W < 360$			
Daily rate DR	$0.5 < DR < 6$			

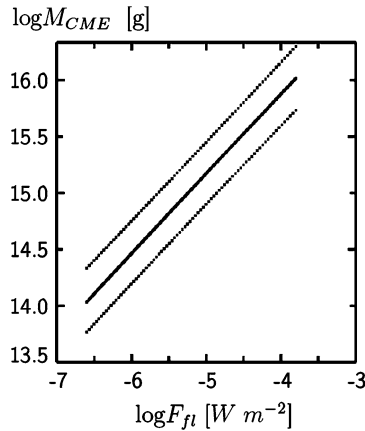


Fig. 1 Statistical relation between log MCME, where the CME mass M is in g, and log Ffl, where the flare flux is in $W\ m^{-2}$, based on observations aboard SOHO and constructed according to results by [14]. Average line with its dispersion is shown by *solid* and *dashed* lines

identified CMEs and 22,674 flares were registered. Basing on their own considerations and consulting results of other investigators [10], Aarnio et al. [14] deduced criteria for CME-flare temporal separation to be $10^m - 80^m$ and position angular difference $\pm 45^\circ$, i.e. the events have to occur in the same quadrant on the solar disc. Then using these criteria 826 associated CMEs/flares pairs have been chosen from 12,050 X-ray flares with known positions and 6,733 CMEs with well measured mass. Position angles (PAs) of flares are computed from their heliographic coordinates and for CMEs the central position angles (CPAs) are taken.

Statistically CME mass increases with flare flux. Figure 1 demonstrates the statistical linear dependence deduced in [14] including halo CMEs and CMEs with purely defined mass: $\log M_{CME} = (18.5 \pm 0.27) \times 0.70 \log F_{fl}$, where the mass M of an CME is in g and the flare flux F is in Watt m^{-2} (see also Figs. 15 and 16 [14]). The results agree with the results presented in [7], [9], [10]. CME width increases with CME mass, as it is shown in Fig. 13 [14], and with the associated flare energy as it can be seen from Table 2 composed by us on the basis of the data presented in Fig. 12 [14]. CME acceleration decreases with flare flux and the slow CMEs accelerate and the fast ones decelerate what is demonstrated in Figs. 8 and 11 [14] correspondingly.

Table 2 CME widths for different flare classes

Flare class	CME width, °
B	42 ± 14
C	53 ± 0.9
M	63 ± 1.8
X	80 ± 10

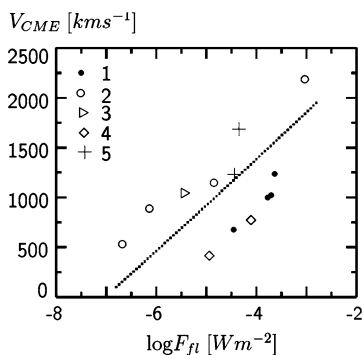


Fig. 2 Relation between observed velocities V in km s^{-1} and fluxes of associated flares F in W m^{-2} : 1—four CMEs on 1997.11.24–25 AR NOAA 9286, Table 1 [11]; 2—four CMEs on 1997.15.16, 1997.01.23, 1998.06.11 and 1997.11.06 in the order of increasing velocities V , Table 1 [6]; 3—CME on 2000.10.25 [12]; 4—two CMEs on 2003.11.20 AR NOAA 10501 [15]; 5—two CMEs on 2003.10.18, AR NOAA 10484 [16]

CMEs accompanied by flares have higher velocities (average $V = 495 \pm 8 \text{ km s}^{-1}$) than CMEs not associated with flares (average $V = 422 \pm 3 \text{ km s}^{-1}$) [14]. Relation between CME velocity and flare flux is presented in Fig. 2 composed by us using the data from [6], [11], [12], [15], [16]. Statistically CME velocity increases with the flux of the associated flare. Average dashed curve is shown.

The velocities appear to increase with height in the COR1 and COR2 LASCO SOHO FOV in different range for fast and slow CMEs as it has been shown by Vourlidis et al. [4] who investigated 11 flux-rope CMEs. For fast CMEs $V = 600 - 900 \text{ km s}^{-1}$ at $\sim 15 R_s$, for slow CME $V = 200 - 250 \text{ km s}^{-1}$ at $\sim 16 - 20 R_s$ as it can be seen from Fig. 3 composed by us from data of [4]. The total energy and mass are kept relatively constant. Usual CMEs achieve the escape velocity at $\sim 8 - 10 R_s$. For fast CMEs the kinetic energy is greater than the potential and magnetic energies. For slow CMEs the potential energy is greater than the kinetic and magnetic energies [4].

STEREO data permit to follow CMEs up to $100 - 150 R_s$. In [17] Wang et al. propagation and expansion of the flux-rope slow gradually accelerated CME observed on 2007 October 8 were investigated up to distances of $\sim 70 R_s$, using the HI1 and HI2 STEREO B images. The CME showed a constant velocity 21.3 km s^{-1} phase and then a constant acceleration (7.6 m s^{-2}) phase [18]. At the distances of $50 - 70 R_s$ the velocity of the leading edge of the CME reached 270 km s^{-1} and expansion velocity was possibly $\sim 90 - 100 \text{ km s}^{-1}$ (Fig. 4).

Fig. 3 Evolution of velocity V with height (in R_s) for fast and slow CMEs:

1—1997.02.23,
2—1998.06.02,
3—1997.10.30, LASCO
SOHO. *Smoothed curves*,
constructed on the results by
[4], are shown

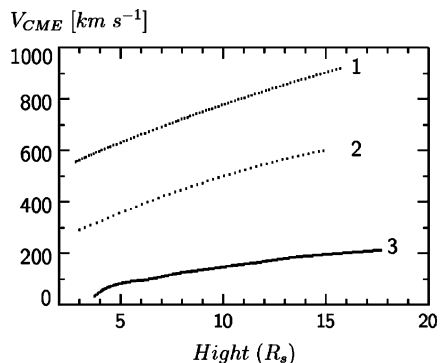
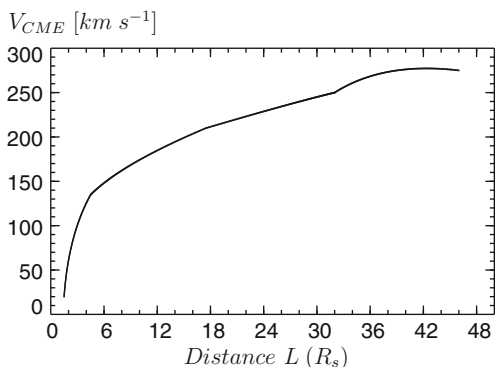


Fig. 4 Evolution of velocity V of the slow gradually accelerated CME on 2007.10.08, propagating through heliosphere, STEREO B. The distance, L in R_s , of the CME flux-rope center is measured from the surface of the Sun



4 Magnetic Helicity

The magnetic helicity in an active region (AR) is known to be inhomogeneous and changeable quickly with time. The flare activity depends on the magnetic helicity and its changes. Nindos and Andrews [19] analyzed 78 events with big flares observed during the period of 1996–1999. They showed that statistically preflare helicity of the ARs (and overlying corona) producing big flares not accompanied by CMES is smaller than the helicity of the ARs producing CME-associated big flares. They found that $\alpha = 0.018 \pm 0.010 \text{ Mm}^{-1}$ for the first type ARs when flares were not accompanied with CMEs and $\alpha = 0.035 \pm 0.018 \text{ Mm}^{-1}$ for the second type ARs in that associated CMEs-flares arouse (see Fig. 1 [19]).

5 Latitudinal and Longitudinal Deflection of CMEs

At solar minimum a CME at early stage may deflect for $20\text{--}30^\circ$ from high latitudes toward the lower ones and beyond $5\text{--}6 R_s$ propagates at almost constant position angle, as it was for the case of the slow gradually accelerated helical CME on

8 October 2007 [17], [18]. Similar slow deflected CME on 8 November 2008 was studied by Kilpua et al. [20]. The slow CMEs have difficulties to overcome straining forces of overlying magnetic field, they obey polar magnetic field of the Sun, i.e. CMEs tend to propagate from regions with high density of magnetic energy to sites of lower magnetic energy density near heliographic current layer. The slow CME on 8 October 2007, observed aboard the STEREO B, continuously deflected toward the ecliptic plane from the position angle $PA \sim 306^\circ$ to $PA \sim 276^\circ$ and beyond $5.5 R_s$ in the COR2 FOV it propagated almost radially (see Fig.2 [18]).

Earlier the CMEs deflection in the meridian plane was investigated for example during 1972, 1974 in [21]. The average deflection of $\sim 2.2^\circ$ toward the ecliptic plane was found. On the basis of observations with the LASCO, EIT, MDI SOHO and base-ground $H\alpha$ images during 1996–2002, 124 structured flux-rope CMEs with known information on the associated source regions (SRs) were analyzed in [22], [23]. The region where the CME originated is named SR. The SRs regions were identified by pre- and post-eruptive events, such as prominences, expanding loops, dimming. The PA of an AR (or a flare) is calculated from its heliographic coordinates. Spatial and temporal coincidence between the CME and its SR is necessary.

Comparing the PAs of the CMEs and SRs they found that during 1996–1998 (near solar minimum) the central positions angles CPAs of the structured CMEs deflected for $\sim 20^\circ$ to lower latitudes toward the solar equator. At times of the high activity (1999–2002) the deviations fluctuated towards the solar poles or equator without a systematical trend. Yashiro et al. [10] investigated spatial relation between associated flares and CMSs comparing their PAs. For 1996–2005 LASCO SOHO observations they found 496 flare-CME pairs considering limb events. They concluded that differences between flare PAs and CME central PAs were of $\sim 17^\circ$.

The simultaneous observations in white-light from different view-points made with the wide-angle imagers HI-1 and HI-2 aboard STEREO A and B give a possibility to track a CME remotely from the Sun up to the Earth and to derive the CME shape and direction of its propagation through the heliosphere. Distinct differences in the CME morphology observed simultaneously from the COR2-A and COR2-B are seen as a result of projection through different lines of sight from two STEREO space crafts. An optically thin structure of a CME is seen differently from different points of view. The examples of observations are shown in Fig.1 [24]. The observations by STEREO, A and B, have shown that the direction of propagation of CMEs (or their pieces) might deflect from the direction Sun–Earth, monotonically or with some temporal fluctuations, toward the east (in some cases) or the west (in other cases) for $5 - 30^\circ$ up to the distances of $100 - 150 R_s$ [25–27]. So the 2008 August 30 CME deflected toward the east for $\sim 30^\circ$ and crossed the Sun–Earth line.

Different techniques are proposed to reconstruct the true direction of propagation, 3D velocities and configuration. The results obtained with different methods agree enough well and azimuthal angles differ for about 10° . The 3D velocity might be greater than the projected velocity for $100 - 200 \text{ km s}^{-1}$.

6 CMEs Accompanying Phenomena

CMEs propagation through the solar corona is accompanied by shocks, accelerating particles and radio bursts. The disturbances in the heliosphere following the most strong CMEs might arrive the Earth and cause storms. The CMEs accompanied with the metric type II radio bursts have a velocity of $\sim 600 \text{ km s}^{-1}$, and the CMEs associated with long waves radio emission (up to kilometric wavelengths) are the fastest ones with average speed of about $1,500 \text{ km s}^{-1}$. The radio phenomena occurring in conjunction with flares and eruptive filaments are widely studied [5], [28–32].

7 Summary

A brief review of results of observations of coronal mass ejection and flares during last decades from SOHO, STEREO and GOES is presented. Relations between different properties of CMEs and associated flares are considered. Aarnio et al. [14] examining a numerous data set have found that the CMEs, associated with flares, statistically have higher velocities ($495 \pm 8 \text{ km s}^{-1}$) than the CMEs not associated with flares ($422 \pm 3 \text{ km s}^{-1}$), CME mass and width increases with flare flux, i.e. the stronger the associated flare is, the more massive CME occurs, and CME acceleration decreases with flare flux. The total energy and mass of CME preserve relatively constant. For fast CMEs the kinetic energy is greater than the potential and magnetic energies. For slow CMEs the potential energy is greater than the kinetic and magnetic energies [4]. Observations from STEREO A and B have shown that a CME might strongly expand. So for the CME on 8 October 2007 the radius of the CME became $\sim 20 R_s$ when its leading edge reached the distance of $\sim 70 R_s$ from the Sun [17]. Statistically preflare helicity of the ARs (and overlying corona) producing big flares accompanied by CMEs is high [19]. The simultaneous observations with the wide-angle imagers HI-1 and HI-2 aboard STEREO A and B give a possibility to track a CME to the distances of almost $100 - 150 R_s$ and to derive the CME shape and direction of its propagation through the heliosphere.

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