PHYSICAL DRIVING FORCES AND MODELS OF CORONAL RESPONSES (Invited Paper)

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The reasons for nonstationary hydrodynamic flows in solar atmospere are reviewed. It is emphasized that a rapid local heating of corona or upper chromosphere can scarcely provide a very large mass of solar plasma ejections observed in corona and interplanetary space. We suggest that coronal transients and interplanetary ejections are produced by magnetic field evolving in solar atmosphere. Magnetic reconnection in current sheets can play essential role in this process. The suitable approximation of strong magnetic field is formulated. Some solutions of MHD equations in this approximation are demonstrated. Their applications to coronal conditions are discussed.

I. INTRODUCTION

During solar flares a significant fraction, if not the majority of the total energy, can release in the form of plasma motions in chromosphere and corona (see monograph by Svestka, 1976). For large flares these motions initiate in interplanetary space strong shocks with energy comparable with the total energy output. Webb et al. (1978a) have obtained accurate estimates of the mechanical energy budget in the 5 September 1973 flare, based on detailed observations of the different forms of motions. The mechanical energy appeared to be comparable with the change of magnetic energy in corona. This is not a surprising result because the magnetic energy in the corona is well known to be enough for all forms of flare energy release.

The other essentially nonsteady phenomena in solar atmosphere (e.g. coronal transients and flarelike events) differ from flares not only by scale (size, power) but also by their spatial structure and by the relative role of different energy release channels. Nevertheless, taking into account that magnetic field is the main energetic factor in active regions, we suggest that all those rapid phenomena have principally the same origin. Namely, we suppose that all those processes are due to magnetic forces.

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According to modern concepts (see Heyvaerts et al., 1977; Syrovatskii, 1977; Baum et al., 1978) the conversion of magnetic energy into other forms, particularly nonthermal, can be connected with the appearance and explosive distruction ('rupture') of current sheets (including the sheets which may appear in nonevolutionary force-free configurations, e.g. Burnes and Sturrock, 1972; Low and Nakagawa, 1972; Jockers, 1978; Bobrova and Syrovatskii, 1979). In order to interpret fast motions in the chromosphere and corona two different regimes of magnetic reconnection in current sheet are appropriate. First of them is the rupture of the laminar (nonturbulent) quasi-steady sheet as a result of some instability and excitation of anomalous resistivity. This regime provides the most powerful energy release and corresponds to the flash phase of flares. The second regime is the rapid reconnection in quasi-steady turbulent sheet. We think that this process can correspond to such coronal phenomena as rapid ascents of X-ray arches and transients.

We shall not discuss here the questions concerning the interpretation and simulation of large amplitude disturbances in the interplanetary space. These very important problems have been developed in detail by American scientists (for a review, see Dryer, 1979). It is more close to our consideration the magnetic field influence on the propagation of disturbances initiated at the coronal base (Steinolfson et al., 1978; Wu et al., 1978; Dryer et al., 1979). The channeling and blocking effects by the magnetic field were studied in those papers. Opposite to that passive role of the coronal magnetic field, we consider the last as a primary cause for plasma motions. For example, an emerging flux (like a magnetic piston) can cause the plasma to flow upward near the loop tops and downward near the loop legs (Nakagawa et al., 1976; Steinolfson et al., 1979). On the other hand, the rapid changes of magnetic field are necessary for explanation of fast motions. To our mind, those changes could result from the reconnection in current sheets. Before discussing these problems, let us consider briefly nonmagnetic mechanisms of plasma ejections.

2. THE MASS OF A PLASMA EJECTED BY FLARE HEATING OF SOLAR ATMOSPHERE

Energy release in the form of heating, particle acceleration etc. results finally in the local heating of the solar atmosphere to high temperatures (see Somov and Syrovatskii, 1976). Large gradient of pressure appears on the boundary of a heated region and can excite fast motions including shocks and plasma ejections.

Consider the at first sight most efficient process of ejections when the chromosphere is impulsively heated by flare electrons with energies $E_e \geqslant E_o \sim (10 - 30)$ keV. This process is really accompanied by the upward high-speed flow of heated plasma (Somov et al., 1977). Note here that the ejected mass is determined essentially by the position of the flare transition layer but not by the 'critical level' where the direct electron heating can be compensated by radiative cooling, as Brown (1973) affirms. Really, the true approach to the problem must include the heat conduction (Shmeleva and Syrovatskii, 1973) that can provide an effective transport of a flare energy even for very large energy fluxes (Somov, 1979a).

The depth ξ (in atoms cm^{-2}) above the flare transition layer depends mainly on the boundary energy flux F_o (erg $cm^{-2}s^{-1}$). For maximum values of F_o during the flash phase of flares, the depth ξ equals about $(1-3)\times10^{19}$ cm^{-2} (Somov et al., 1979). Hence the mass $\sim(2-5)\times10^{-5}$ g cm^{-2} can be ejected from unit area. Let S_1 be the area heated by energetic electrons during one 'elementary flare burst' (a single spike of hard X-rays). Bearing in mind some 'mean' impulsive flare, we suggest that S_1 is equal to the area of fast (5-10 s) small (1-3") chromospheric flash which coincides in time with hard X-ray burst (see Zirin and Tanaka, 1973). Then $S_1 \sim 5\times(10^{16}-10^{17})$ cm^2 and the ejected mass equals $M_1 \sim 10^{12}-3\times10^{13}$ g. Even if flare as whole ejects the total mass $M \sim 10\, M_1$, this is too small compared with the observed values $M_{obs} \sim 10^{16}-10^{17}$ g (e.g. Dryer et al., 1976a). Thus, the impulsive heating of the chromosphere by energetic electrons can not provide the observed mass. Note here that the fast heating by large thermal fluxes with the same F_o yields the same or somewhat less ejected mass (Sermulinia et al., 1979).

The ejected mass is even smaller if the heating occurs not inside the chromosphere but in the corona. The numerical study of the dynamic response of a static corona on the finite-amplitude disturbances shows that the needed mass can be ejected upward only from beneath the boundary placed at the coronal base (Nakagawa et al., 1975). It means, for example, that for large flare spray the mass of order 10^{16} g is ejected from the chromosphere, but the physical reason for such ejection remains unspecified. In the numerical simulation for the fast-mode MHD shock propagation through the corona during the flare of 1973 September 5, the ejected mass is about 6×10^{16} g (Dryer and Maxwell, 1979). The mass of the same order is assumed for the model of interplanetary disturbances in the solar wind (Dryer et al., 1976b; Wu et al., 1976). All these approaches assume tacitly that the needed mass is ejected from the chromosphere just as a response to the flare energy release. Is really the large mass observed in the coronal transients and interplanetary disturbances taken directly from the chromosphere?

This question is especially critical one if we take into account that the most of transients ($\sim$ 60%) are not accompanied by observable chromospheric flares (Munro et al., 1979). This implies that the observed mass ought to be stored in the corona before a transient or raked up during it. Coronal transients are often associated with eruptions of prominences. For this reason, one can assume that the transients are caused by sudden heating and eruption of a previously cooler material, e.g. cold dense filament. Another possibility is that the needed mass is expelled from the low corona (Rust and Hildner, 1976). Then, the large area $S \sim 10^{21} - 10^{22}$ cm^2 is needed to gather the observed mass.

One more source of mass is conceivable. The current sheets could, in principle, accumulate coronal plasma. Magnetic field moves into the sheet and annihilates there whereas plasma carried with the field is accumulating in the sheet or in the neighbourhood. The current sheets are most effective in strong magnetic fields. For this reason, we consider in Section 2 the approximation of strong field.

3. THE MAGNETOHYDRODYNAMIC APPROXIMATION OF A STRONG FIELD

For a weak field, magnetic effects are only small corrections to hydrodynamic ones. Contrary to it, the strong field approximation (Syrovatskii, I966) implies that a magnetic force dominates over all other forces: gas pressure gradient, inertia etc. This means formally the following. The complete set of the ideal MHD equations can be written in the next dimensionless form (e.g. Somov and Syrovatskii, I974):

$$\varepsilon^2(\delta^{-I}\,\partial\vec{v}/\partial t + (\vec{v}\cdot\text{grad})\vec{v}) = -\gamma^2(\text{grad } p)/\rho - (\vec{B}\times\text{curl }\vec{B})/\rho, \quad (I)$$

$$\partial\vec{B}/\partial t = \delta \text{ curl } (\vec{v}\times\vec{B}), \quad (2)$$

$$\partial\rho/\partial t + \delta\,\text{div } \rho\vec{v} = 0, \quad (3)$$

$$\partial S/\partial t + \delta\,(\vec{v}\cdot\text{grad})S = 0, \quad (4)$$

$$\text{div }\vec{B} = 0, \quad (5)$$

$$p = p(\rho, S). \quad (6)$$

These equations contain three dimensionless parameters

$$\delta = VT/L, \qquad \varepsilon = V/V_A, \qquad \gamma^2 = p_o/\rho_o V_A^2, \quad (7)$$

where L, T, V, ρ_o, p_o and B_o are the scales of length, time, velocity, density, pressure and field strength, respectively. One can consider a magnetic field as a strong one if ["I" represents the integer "one" - Eds.]

$$\gamma^2 \ll I \text{ and } \varepsilon^2 \ll I. \quad (8)$$

It follows from Equation (I) that in zeroth order in small parameters (8) the magnetic field is a force-free one:

$$\vec{B}\times\text{curl }\vec{B} = 0, \quad (9)$$

or that is simply a potential field in the absence of currents.

In the first order, we neglect the gas pressure gradient in comparison with the inertia (Somov and Syrovatskii, I974):

$$\varepsilon^2(\delta^{-I}\,\partial\vec{v}/\partial t + (\vec{v}\cdot\text{grad})\vec{v}) = -(\vec{B}\times\text{curl }\vec{B})/\rho, \quad (I0)$$

which implies the approximation of a strong field and cold plasma:

$$\gamma^2 \ll \varepsilon^2 \ll I. \quad (II)$$

This approximation is especially appropriate for fast motions caused by changes of a strong field. Just this case is of interest for us.

The parameter δ rates the local derivative $\partial/\partial t$ relative to the convective term $\vec{v}\cdot$grad. $\delta \gg I$ for quasi-steady flows, $\delta \ll I$ for small disturbances. Generally, $\delta = I$, then the equations are

$$\varepsilon^2\, d\vec{v}/dt = -(\vec{B}\times\text{curl }\vec{B})/\rho, \quad (I2)$$

$$\partial\vec{B}/\partial t = \text{curl } (\vec{v}\times\vec{B}), \tag{I3}$$

$$\partial\rho/\partial t + \text{div } \rho\vec{v} = 0. \tag{I4}$$

Expand the solution in the small parameter ε^2 (i.e. $f = f^{(o)} + \varepsilon^2 f^{(I)} + \ldots$). In zeroth order, magnetic field is determined by equation

$$\vec{B}^{(o)}\times \text{curl } \vec{B}^{(o)} = 0 \tag{I5}$$

with time-dependent boundary conditions. Following these conditions, the field changes and drives a plasma. Kinematics of this motion is determined by two equations. Equation (I2) means that the acceleration is perpendicular to magnetic field lines, namely

$$\vec{B}^{(o)}\cdot d\vec{v}^{(o)}/dt = 0. \tag{I6}$$

The second one coincides with the 'freezing-in' Equation (I3):

$$\partial\vec{B}^{(o)}/\partial t = \text{curl } (\vec{v}^{(o)}\times\vec{B}^{(o)}). \tag{I7}$$

The plasma density can be found from the continuity Equation (I4):

$$d(\ln \rho^{(o)})/dt = - \text{div } \vec{v}^{(o)}. \tag{I8}$$

Thus, the Equations (I5)-(I8) determine all unknown zeroth order variables. The same procedure can be continued to any order in the small parameter ε^2. For our aim, it is enough to consider only zeroth order, e.g. to neglect departures of the field from force-free (or potential) one.

To solve the problem we proceed as follows. At first, we find the force-free (or potential) field $\vec{B}^{(o)}(\vec{r},t)$ at any instant from the Equation (I5) and time-dependent boundary conditions. Then, the velocity $\vec{v}^{(o)}(\vec{r},t)$ follows from the Equations (I6) and (I7) and from the initial condition for velocity $\vec{v}^{(o)}(\vec{r},0)$ along the magnetic field. Finally, the continuity Equation (I8) and an initial plasma distribution $\rho^{(o)}(\vec{r},0)$ yield the density $\rho^{(o)}(\vec{r},t)$. Let us demonstrate this procedure by some examples.

4. HYDRODYNAMIC FLOWS NEAR A STEADY CURRENT SHEET

Let the steady current sheet be of the width 2b along the x axis (Figure I). The uniform electric field $\vec{E}_o$ is parallel to the z axis.

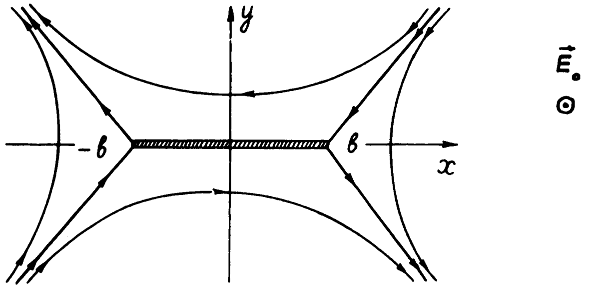

Figure I. The field lines near the steady current sheet.

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Two-dimentional motions near such a sheet can be described by the Equations (15)-(18) formulated in terms of the vector potential $\vec{A} = (0, 0, A(x,y))$:

$$\Delta A = 0, \qquad (19)$$

$$d\vec{v}/dt \times \text{grad } A = 0, \qquad (20)$$

$$dA/dt = 0, \qquad (21)$$

$$\partial\rho/\partial t + \text{div } \rho\vec{v} = 0. \qquad (22)$$

According to Equation (19) the 'potential' $A(x,y)$ is a harmonic function of a complex variable $z = x + iy$. For this reason, it is convenient to introduce the complex potential as the analytic function

$$F(z,t) = A(x,y,t) + iB(x,y,t). \qquad (23)$$

In the case under consideration the complex potential is of the form (Syrovatskii, 1971):

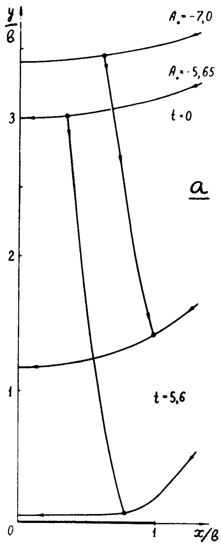
$$F(z,t) = \tfrac{1}{2}h_o z(z^2 - b^2)^{\frac{1}{2}} - \tfrac{1}{2}h_o b^2 \ln\frac{z + (z^2 - b^2)^{\frac{1}{2}}}{b} + A(t). \qquad (24)$$

The coefficient h_o is the magnetic gradient in the absence of current sheet (when $b = 0$). The term $A(t)$ is introduced formally to take into account the magnetic flux dissipation in the sheet.

The Equations (20)-(22) can be written in Lagrange variables as the set of four ordinary differential equations for the trajectories of 'fluid particles' $x(t)$ and $y(t)$ and the continuity equation for density

$$\rho(x,y,t)/\rho_o(x_o,y_o) = D(x_o,y_o)/D(x,y). \qquad (25)$$

Here $\rho_o(x_o,y_o)$ is the initial density distribution. The Jacobian on the right side can be found, for example, by the simultaneous calculations of three neighbouring trajectories.



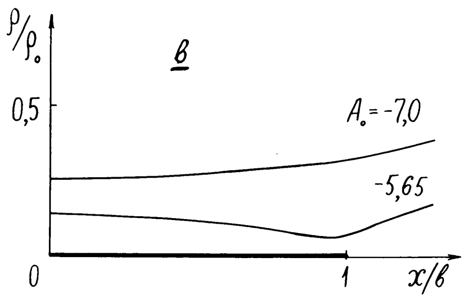


Figure 2. (a) Trajectories of fluid particles placed on two different field lines: $A(x,y,t) = A_o$ with $A_o = -5.65$ and $A_o = -7$. (b) The relative density distributions along the same field lines at the dimensionless time $t = 5.6$.

The results of numerical solution (Somov and Syrovatskii, I974) show that for considered supersonic ($\gamma^2 < \varepsilon^2$) flows near the steady current sheet there is no stationary solution for density. The last decreases monotonously on the field lines flowed to the sheet (Figure 2). As a final result, the current sheet could be surrounded by a very low density plasma if the electric field $\vec{E}_0$ acts long enough. This conclusion is important for the problem of stabilization and rupture of current sheets in laboratory and cosmic plasmas (e.g. Bulanov et al., I977; Bulanov and Sasorov, I978).

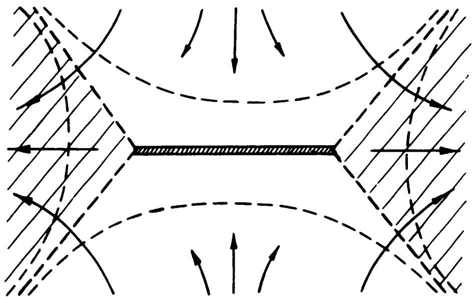


Figure 3. The regions of contraction (two shaded sections) and of depletion (the rest) near the steady current sheet. The motion trajectories of plasma are shown schematically by arrows.

On the field lines which have undergone magnetic reconnection in the sheet a plasma is raked-up and compressed (Figure 3). Together with a plasma flowed out from the sheet, this compressed plasma moves rapidly in opposite directions along the x axis. Fast plasma ejections are met in the laboratory and numerical simulation of the current sheet development and rupture (Kirii et al., I979; Podgornii and Syrovatskii, I979). Ivanov and Platov (I977) consider a similar process (the raking-up near a neutral line, i.e. b=0) as a model for loop prominence formation.

5. THE RUPTURE OF A CURRENT SHEET AND THE FORCES ACTING ON ITS EDGES

Magnetic field is weak or fully absent inside a current sheet. However the external strong field can penetrate in one or more points (gaps) when and where the sheet ruptures. In this case, magnetic forces begin to act on the edges of appearing gap and will enlarge it (Somov and Syrovatskii, I975).

In the simplest case one can imagine that before the rupture the current sheet is a uniform current in the y = 0 plane (Figure 4a) flowing parallel to the z axis and creating the magnetic field jump from B_0 at $y < 0$ to $- B_0$ at $y > 0$. Let us assume that the gap is formed parallel to the z axis as a result of some instability of the sheet. 2a(t) is the width of this gap (Figure 4b). In the approximation of ideal conductivity, the field lines at the surface of the current sheet remain the same as before rupture. The magnetic field is potential one outside currents and, if a new current does not arise along X-type neutral line at the gap (see more general case in Somov and Syrovatskii, I975), the complex potential is of the particularly simple form

$$F(z,t) = iB_0(z^2 - a^2(t))^{\frac{1}{2}}. \qquad (26)$$

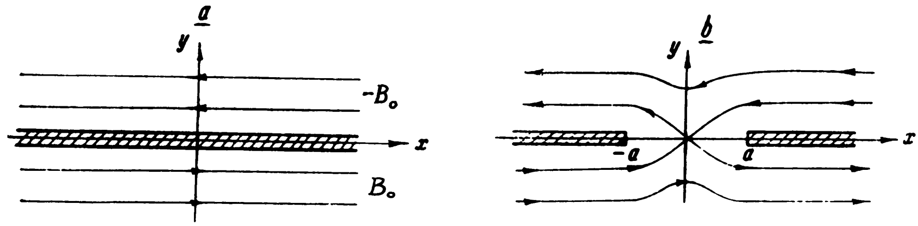


Figure 4. Field lines near the primary sheet (a) and ruptured one (b).

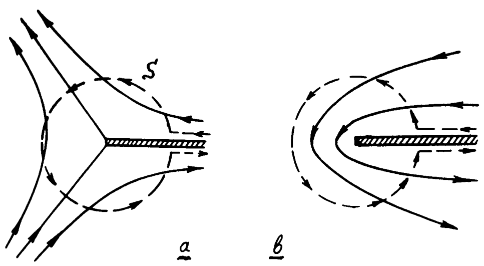


Figure 5. Field structure near free edge of the sheet (a) and the edge of the gap.

Let us derive a formula for the force acting on the edge of gap. The simplest method to determine this force is to calculate the Maxwell stresses at the surface S (Figure 5) bounding an edge:

$$\vec{F} = \frac{1}{4\pi} \oint_S \left[-\vec{B}(\vec{B}\cdot\vec{n}) + \tfrac{1}{2}B^2\vec{n}\right] dS. \quad (27)$$

A real current sheet always has a finite width. But according to (27) there is no magnetic force on the free edge ($F \sim z^{3/2}$, Figure 5a) of the sheet. On the other hand, at the edge of the gap ($F \sim z^{1/2}$, Figure 5b) the magnetic force acts along the current sheet (the x axis):

$$F_x = \pm \tfrac{1}{8} B_o^2 a \quad \text{at} \quad x = \pm a. \quad (28)$$

This result can be understood easily in another way, namely by calculation of the free magnetic energy for the system of currents

$$W = - \tfrac{1}{8} B_o^2 a^2. \quad (29)$$

Since W equals the work done by the magnetic field on the expansion of the gap, the force acting on the gap edge is

$$F_x = \partial(\tfrac{1}{2}W)/\partial a = \pm \tfrac{1}{2} B_o^2 a \quad \text{at} \quad x = \pm a. \quad (30)$$

Thus, the edges of the gap are subjected to magnetic forces which are proportional to the width of the gap and tend to enlarge it.

We can draw two conclutions from this analysis. First of them is a condition of sheet rupture. If the sheet has a finite thickness d and if the gas pressure inside the sheet is balanced by external magnetic pressure, then the sheet can be ruptured in the case of $d < a$. This conclusion follows from a comparison of the magnetic force in (30) with the opposing gas pressure force $pd = (B_o^2/8\pi)d$. The second conclusion is the exponential growth of small perturbations in the initial stage of the rupture.

At nonlinear stage of the rupture, in order to write an equation of motion for the edge, it is necessary to take into account the mass of a plasma involved in motion and also thermal effects (Bulanov and Sasorov, 1978). The gas mass can be estimated as

$$M(a) = n_s m_i a d \quad (g\ cm^{-1}) \tag{31}$$

where n_s is plasma concentration inside the sheet. Thermal effects being disregarded,

$$d(Mda/dt)/dt = F_x(a) = \tfrac{1}{8}B_o^2 a. \tag{32}$$

From this it follows that the edge moves with the constant acceleration

$$\ddot{a} = 4\pi V_A^2/3d \quad \text{where} \quad V_A = B_o/(4\pi n_s m_i)^{\frac{1}{2}} \tag{33}$$

is the characteristic value of Alfven velocity.

To take into account thermal effects it is necessary to solve hydrodynamic equations $y = 0, |x| \leq a$. The boundary conditions for pressure at the point $x = a$ describes the action of the magnetic force:

$$p(a) = F_x(a)/d. \tag{34}$$

The solution of this self-similar problem (Bulanov and Sasorov, 1978) does not change the conclusion about the character of the current sheet rupture affected by magnetic force. It is essential here that after the rupture the plasma motion velocity becomes larger than the velocity (33).

6. PLASMA ACCUMULATION INSIDE CURRENT SHEET

Let us estimate the plasma mass inside a quasi-steady laminar current sheet. Following Syrovatskii (1976), we use the momentum and continuity equations and Ohm's law from which there follow the order of magnitude relations (see Figure 6):

$$n_o V_d b = n_s V_x a, \quad B_s^2/8\pi = 2n_s kT = \tfrac{1}{2}n_s m_i V_x^2,$$
$$cB_s/4\pi a = \sigma E_o. \tag{35}$$

Here n_o is plasma concentration near the sheet, a and b are half-thickness and half-width of the sheet, $V_d = cE_o/B_s$ is the drift velocity near the sheet, E_o is the electric field along the zeroth line of primary field. From (35) it follows that the velocity of plasma outflow from the laminar sheet equals the Alfven velocity inside the sheet (cf. Sweet, 1958)

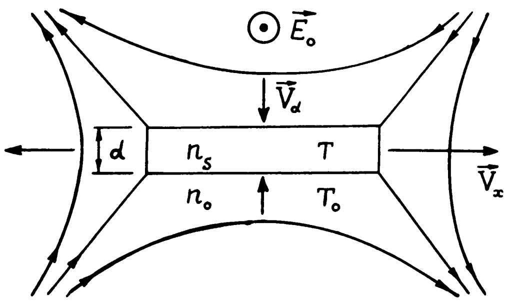


Figure 6. Steady model of the sheet.

$$V_x = B_s/(4\pi n_s m_i)^{\frac{1}{2}}, \tag{36}$$

and the other quantities are expressed through the external parameters n_o, E_o and h_o and through the sheet temperature T (Syrovatskii, 1976). Specifically, in the current sheet being of the length l the plasma mass is

$$M = 4abl m_i \approx 5 \times 10^{-21}(E_o^5 n_o^4/h_o^7)^{1/3} T^{1/6} l \quad (g). \tag{37}$$

Let, for example, $n_o = 5 \times 10^8$ cm^{-3}, $E_o = 4 \times 10^{-4}$ CGSE, $h_o = 5 \times 10^{-7}$ G/cm and $T = 8 \times 10^4$ K (cf. Syrovatskii, 1976). Then $M \approx 141$g. If $l = 10^{10}$ cm, the current sheet mass has to be $M \sim 10^{11}$g. This value is small compared to the observed ejected mass (Section 2).

Thus, the mass of plasma accumulated inside the current sheet is insignificant. However, another point is essential. Firstly, an emerging strong magnetic field can lift plasma from the chromosphere and low corona. Some part of the ascendant plasma flows into the current sheet and is ejected from it with near Alfven velocity. This ejected plasma moves together with and along the magnetic field lines which have undergone the reconnection inside the sheet. Partially those fast motions of plasma are directed upward through corona. Secondly, during a fast reconnection the above-mentioned raking-up mechanism can also give rise to upward fast ejection. The total mass involved in both processes can be estimated roughly if one takes into account that the emerging and reconnecting magnetic field pushes out a plasma practically from the whole coronal volume placed below the current sheet. For the sheet at the height $h \sim 3 \times 10^{10}$ cm, this mass is of about $m_i n h^3 \sim 10^{16}$ g if $n \sim 3 \times 10^8$ cm^{-3} is the coronal plasma concentration.

In our opinion, many observations give an evidence for the hypothesis that current sheets lie at the tops of loops or arcades of loops. In particular, X-ray and EUV observations of flares and flarelike events (for a review, see Somov, 1978) together with observations and calculations of magnetic field confirm that many active processes in the corona are conditioned by emergence of a new magnetic flux. The front boundary of the flux is observed as an arcade of interrelated loops whose density is two-three orders of magnitude larger than that in the surrounding corona. Observed temperature distributions and plasma motions agree qualitatively with the assumption that at the tops of loops there are current sheets where the magnetic reconnection occurs. Optical observations (e.g. Mahmudov et al., 1979; Ostapenko, 1979) also confirm that at the tops of the loops plasma condensations are formed from which plasma flows out with the initial acceleration (due to magnetic forces in the current sheet model).

7. APPEARANCE OF A CURRENT SHEET IN A PLASMA MOVING IN DIPOLAR FIELD

In the presence in plasma of singular lines of magnetic field (zeroth lines $\vec{B} = 0$, $\vec{E} \neq 0$ in the simplest case) a continuous flow becomes impossible because the freezing-in Equation (2) is violated at these lines (Syrovatskii, 1978). As has been shown by Syrovatskii (1971), current sheets appear where the singular lines would take place in the absence of plasma (Figure 7). The general solution including the reversed currents at the sheet edges is shown in the Figure 7b. In this case, there exists magnetic force acting on the current sheet edges (cf. Figure 5b). Besides, magnetic field strength becomes infinite value here in the approximation of infinitely thin sheet. For the quasi-steady sheet the structure of magnetic field near the edges is shown in Figure 5a.

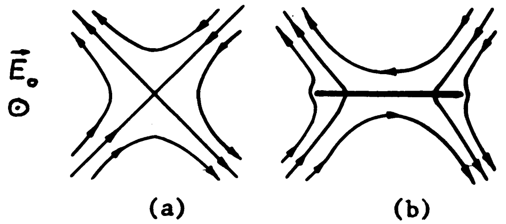


Figure 7. Formation of current sheet in the place of a singular zeroth point (line along the z axes). (a) Field lines near the primary zeroth point, (b) the potential field outside the sheet.

We consider below an idealized model when a current sheet appears as a result of capture and extension of a dipolar magnetic field by a plasma flow (Somov and Syrovatskii, I972). Let the two-dimensional dipole be placed in the base of a semicylindrical region (Figure 8) inside of which the conditions (II) are satisfied. We assume that the magnetic flux is conserved on the boundary R of the region expanding in accordance with a specified law $R = R(t)$. For the corona, such a boundary is the simplest model of transition region from the low corona where the strong field conditions are satisfied to the upper corona in which the kinetic energy of solar wind dominates.

In zeroth order in small parameter ε^2, the 'potential' A(x,y,t) is defined by the Laplace equation (I9) with boundary conditions

$$A(x,y,t) = A(r,\varphi,t) = \begin{cases} 0 & \text{if } y = 0 \\ (\alpha m/r)\sin\varphi & \text{if } r = R(t) \end{cases} \qquad (38)$$

and singularity of the dipolar type at $r = 0$. Here r and φ are polar coordinates, α is a fraction of the magnetic flux penetrating through the boundary at the initial time (α = I in Figure 8a), $R_o = R(0)$. The solution in terms of the complex potential is obvious:

$$F(z,t) = im/z - im(\alpha R - R_o)z/(R^2R_o). \qquad (39)$$

Pattern of magnetic field lines is shown in Figure 8b. If the radius R(t) of the boundary increases to a value larger than $2R_o/\alpha$, then a zeroth point of X-type appears inside the region. This point coordinate is

$$z_o = iR(R_o/(\alpha R - R_o))^{\frac{1}{2}}. \qquad (40)$$

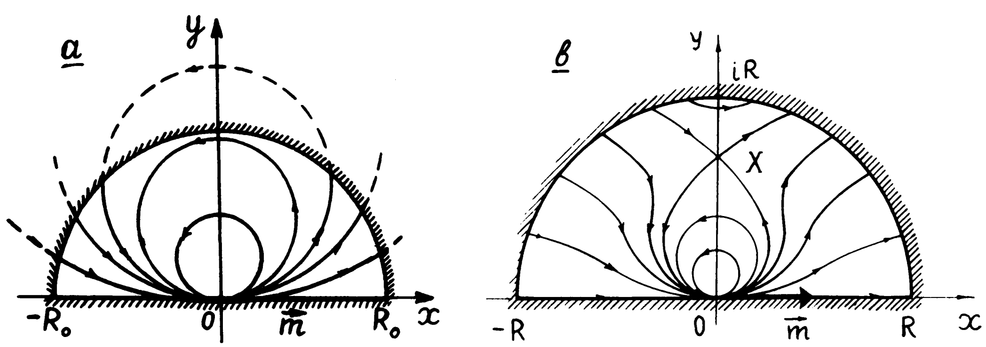


Figure 8. The magnetic field of two-dimensional dipole m with field lines frozen into the boundary R(t). (a) The dipolar field has penetrated through boundary at t = 0. (b) Magnetic field in the absence of plasma.

One can see that electric field at the X-type point is nonzero. Therefore, that point is singular in the sense of the freezing-in condition and should be eliminated from the region where the analytic function F(z,t) is defined. This can be done by the cut on the complex plane along the y axis from the boundary to a certain height h (Figure 9).

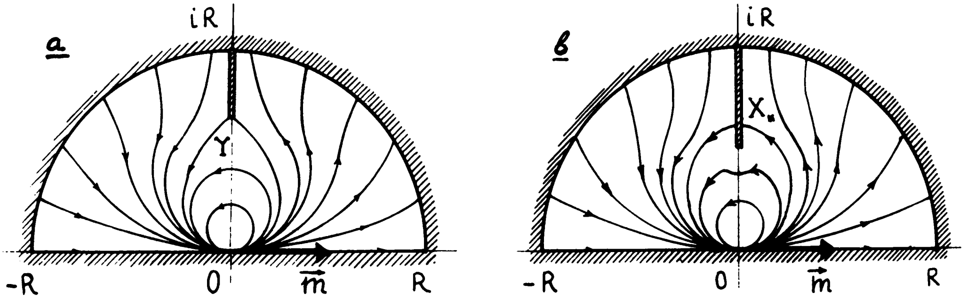


Figure 9. Field lines corresponding to the solution with the current sheet. (a) The length of the cut is minimal, (b) the cut length is larger than the minimal one (there is a reversed current at the sheet edge).

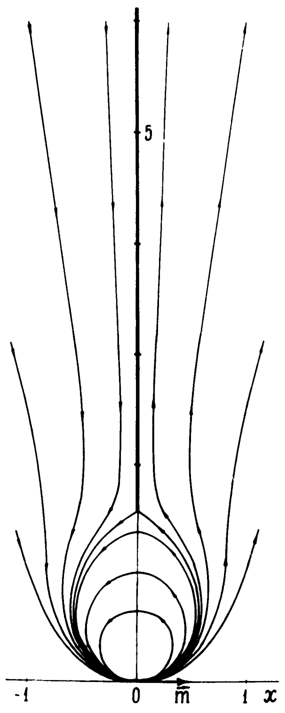


Figure 10. Magnetic field corresponding to the asymptotic solution at $\alpha R/R_0 \to \infty$.

The introduction of the cut takes into account that at the time when the zeroth point appears ($R = 2R_0/\alpha$) a current sheet begins to develop. The width of this sheet increases with increasing R(t). The general solution of the problem with the sheet (Somov and Syrovatskii, 1972) is shown schematically in Figure 9. If $\alpha R/R_0 \to \infty$, the solution rapidly reaches its asymptotic form which is shown in Figure 10 for the case when the reversed current is absent. Near the dipole, the field has its usual structure. At large distances, the field lines tend to become radial straight lines.

The model considered has the advantage that it can be fully elaborated conserving the fundamental physical essence of the phenomenon. As a result, the conditions under which neutral layers (coronal streamers) are produced when a plasma flows in a dipolar field become more clear. Three such basic conditions can be indicated:

(1) A sufficiently high conductivity of the plasma, so that the magnetic field can be regarded as frozen into the plasma (outside the sheet itself where reconnection takes place).

(2) The existence of a boundary or a sufficiently narrow transition layer between the region where magnetic forces dominate (magnetic cavity) and the region where the plasma kinetic energy dominates (solar wind).

(3) A penetration of the magnetic field from

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the magnetic cavity into the region of wind, i.e. the 'capture' of the field by the solar wind.

8. PLASMA MOTIONS INITIATED BY RAPID CHANGES OF MAGNETIC FIELD

Current sheets screen magnetic fluxes from different sources. For example, if a new emergent flux differs in direction from magnetic field of an active region, then the new and old fluxes are separated by the surface of screening current (Figure IIa). The magnetic reconnection in this current sheet changes the topology of field: the fluxes are redistributed between different sources (Figure IIb). This can illustrate the slow (quasi-steady) evolution of magnetic fields in an active region.

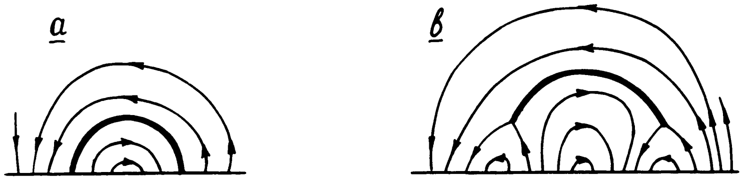

Figure II. Current sheet between the new and old fluxes.

Another situation is more interesting in connection with nonsteady phenomena like flares, transients etc. During the slow evolution described above some critical state can be reached. Starting from it the continuous deformation of magnetic field does not lead to a new equilibrium, and fast dynamic phase of evolution begins. For the case under consideration the critical state corresponds to an approximate equality of effective dipolar moments for the new emerging and old magnetic fields. The sudden transition to the eruptive phase is accompanied by the fast change of magnetic topology.

To illustrate this affirmation let us consider the potential field for four magnetic regions of interchanging polarities on the photosphere. If the normal component is uniform in each region it is easy to calculate the magnetic field in the whole space over the photosphere (Syrovatskii, I979). For symmetrical case the neutral (zeroth) point of magnetic field appears on the symmetry axis (Figure I2a). The neutral point height is determined by the effective moments of the weak (background) field (S_o,N_o) and of the evolving internal bipolar group (N,S). When the magnetic moment of the internal group increases, the neutral point ascends with increasing velocity upto an infinite height if the effective moments of the emerging and background fields become nearly equal. In this time the 'closed' magnetic configuration turns into the 'open' one (Figure I2b). Note that the background field can be very weak provided its total flux (or, more exactly, the effective magnetic moment) is comparable with the emerging flux. It is important that the ascent velocity of the neutral point increases infinitely near the critical value of the emerging magnetic moment.

In real conditions, in a plasma of high conductivity, a current sheet is formed in the place where the neutral point should be in the absence of the plasma (Figure 7). The reconnection in the current sheet provides the transition from a 'closed' magnetic field to an 'open' one. The reconnection rate determines the real velocity for the 'neutral point' ascent. It is known that maximum rate can be near (about a few tenthes) to the Alfven velocity in the undisturbed plasma. For the corona it corresponds to values of order thousand km/s. The coronal reconnection is especially effective if the current sheet is turbulent after a flare. Under such a rapid reconnection the real picture of magnetic field lines differs only a little from the ideal picture of potential field with the neutral point (Figure I2).

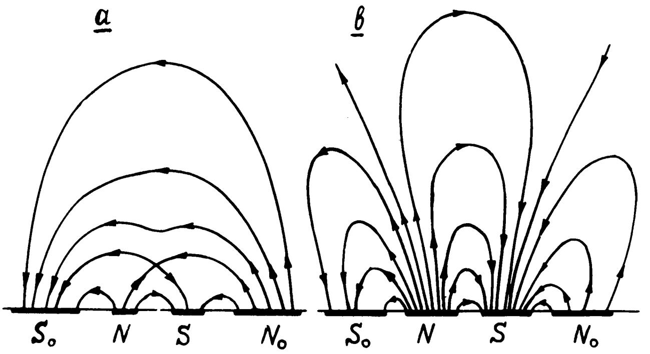

Figure I2. Potential field of four magnetic regions. (a) The close configuration with the neutral point. (b) The critical instant when the magnetic field opens.

Thus, the emergence of a new magnetic field can result in very fast coronal motions accompanied by the reconnection and appearence of regions with an 'open' field (coronal holes). To our mind, it is very likely that such changes of strong magnetic field can cause coronal transients.

9. CONCLUSIONS

For the interpretation of fast motions in the corona (under condition of relatively slow photospheric motions) the concept of current sheet is of principial weight. Current sheets can store magnetic energy during a long slow evolution. This pre-eruption phase can be ended with a sudden explosion (a fast dynamic phase) by transition into turbulent state. The explosion is accompanied by direct ejections from a current sheet and by fast plasma motions in the surrounding chromosphere and corona due to rebuilding of total magnetic structure. One example of such a motion is considered by Somov and Syrovatskii (I979).

On the other hand, under certain conditions, the turbulent current sheet over ascending magnetic flux can penetrate rapidly (with Alfven velocity of order one) into the upper corona and rebuild the whole mag-

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netic field from a 'closed' topology to an open one. These processes create the fast plasma motions and the direct conversion of magnetic energy into heat and radiation.

Note here that Webb et al. (1978b) found a tendency for transients to occur on the borders of growing coronal holes. Rust (1978) presented 'the observational evidence for transient field openings' in the transient coronal holes. According to the model considered above (contrary to the model by Pneuman, 1979) this process is result from the magnetic reconnection in the ascending current sheet (the 'neutral point') on the tops of rising loops.

For more detailed investigation of disturbances in the upper corona and interplanetary space and for comparison of them with definite active processes on the Sun it is necessary, to our mind, to develop further the nonsteady MHD theory for radiative plasma flows in the low corona and upper chromosphere. What should future theory take into account (see also Somov, 1979b)?

First of all, this is an inhomogeneity of magnetic field and an existence of the singular (neutral in the simplest case) lines. Models with homogeneous or simple dipolar field do not have those lines and are not satisfactory for this reason. The singular lines appear, for example, if besides the emerging 'dipolar' magnetic flux there is an 'old' background field in the corona. This field, even if it is very weak, occupies large areas and can have a sufficient magnetic flux to interact with local emerging flux efficiently.

The second important factor is the solar wind. It has a strong influence on the magnetic field in upper corona. Development of solar wind theory (especially near the Alfven surface) is necessary to interprete coronal streamers and holes. In the internal corona the magnetic field structure is determined mainly by photospheric sources, but this is not true for the upper corona. The approximation of a strong field is not applicable here. The investigation of a transition region between the internal corona with a strong magnetic field and the upper corona with a weak field trailed by solar wind is an important problem to be solved.

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DISCUSSION

Moore: What is your opinion of the magnetic field configuration proposed by Jerry Pneuman as being appropriate for magnetic field reconnection in a filament eruption?

Somov: Contrary to the model by Dr. Pneuman, the opening of magnetic field in our model is result from magnetic reconnection in the ascending current sheet.

Nakagawa: (Comment) Those models presented represent schematic illustrations at best. I do not believe these figures or models, unless realistic physically self-consistent computations are carried out and numerical results are obtained to compare with observations.

Callebaut: Have you thought of or tried to introduce some of the features of the Uchida-Sakurai model? I think in particular of the interchange in stability which allows the reconnecting surface and volume to be increased very much.

Somov: I think that the interchange instability considered in the Uchida-Sakurai model is very essential for the case when gas pressure gradient is not disregarded. But this instability can hardly be effective in a strong magnetic field in the internal (low) corona, especially if a small field component exists along a current sheet.