EFFECT OF THE ELECTROMAGNETIC FIELD OF AN ERUPTIVE PROMINENCE ON CORONAL AND CHROMOSPHERIC PLASMAS (CME AND H\(\alpha\) EMISSION FORMATION)

B. P. Filippov

The frozen-in chromospheric and coronal plasma motions during an eruption of a filament with a magnetic field configuration described by the inverse polarity model are considered. At the initial stage of the filament motion the magnetic field compresses the chromospheric gas within two strips located symmetrically about the inversion line. The compression is accompanied by plasma heating and emission enhancement in the line H\(\alpha\). The distance between the strips increases with filament altitude above the photosphere. This mechanism is sufficient to describe the dynamics of H\(\alpha\) emission kernels in two-ribbon flares. In the corona region in which the magnetic pressure of the filament field is greater than the gas pressure, plasma rarefaction and cavity formation occur. Near the boundary \(\beta = 1\) the plasma is decelerated and its density increases, which corresponds to the formation of an outer shell of the CME.

Representing the prominence eruption as an evolution of the current system, which lost its stability, we can interpret many features of coronal and chromospheric changes as natural consequences of the frozen-in plasma motion. The filament motion with the current induces an electric field \(\vec{E}\), which generates the plasma drift in the surrounding space with velocity

\[
\vec{v}_d = \frac{\vec{E} \times \vec{B}}{B^2}.
\]

The inhomogeneity of the velocity field \(\vec{v}_d\) gives rise to the compression and rarefaction regions according to the continuity equation

\[
\frac{\partial \rho}{\partial t} + \text{div} \rho \vec{v} = 0.
\]

Let us consider the 2-D motion of the forward linear current along the z axis. Allowing for the boundary condition on the photosphere, the magnetic field of the current \(I\) is described by the vector potential

\[
A_z = \frac{I}{c} \ln \left[ \frac{x^2 + (y + h)^2}{x^2 + (y - h)^2} \right],
\]

where \(h\) is the altitude of the current above the photospheric surface \(y = 0\). Ascending with velocity \(v_0\), the current induces the electric field

\[
\vec{E} = -\frac{1}{c} \frac{\partial A_z}{\partial t}.
\]

On the oy axis above the filament the plasma velocity increases linearly with distance from the latter

\[
v_y = v_0 \frac{y}{h},
\]

and

\[
\text{div} \vec{v} \bigg|_{z=0} = \frac{v_0}{h} \left( 1 + \frac{2y^2}{y^2 - h^2} \right).
\]
is positive everywhere above the filament and decreases tending to \(3v_0/h\) at a large distance. Therefore, one can expect corona rarefaction above the filament [1]. Of course, these relationships are fulfilled only in the region in which the gas pressure forces are small compared with the magnetic forces. Approximately, as the boundary of this region we can take the surface on which

\[
\beta \equiv \frac{8\pi nkT}{B^2} = 1.
\]

Since outside the region \(\beta < 1\) the plasma remains unperturbed, the coronal material decelerates and its density increases near the surface \(\beta = 1\).

In the majority of cases, the material ejections in the solar corona, which are also called coronal mass ejections (CMEs), have the following spatial structure: a bright kernel that has the same location as the eruptive prominence is surrounded by a dark cavity whose density is smaller than that of the surrounding corona. The cavity is edged by a bright loop with increased density. The velocity of motion of all of the three ejection parts is different and increases with distance from the kernel to the outer loop [2]. Evidently, the frozen-in motion of the coronal plasma in the ascending prominence field is in good correspondence with the CME formation: the rarefaction above the filament results in cavity formation, and the increasing density, by its location and motion character, simulates the external front part of the CME.

To obtain more detailed information on the character of material motion in the corona during the filament eruption, we must consider the solution of the entire system of MHD equations. This system of equations was solved numerically in the approximation of ideal conductivity for a 2-D geometry. The motion of the forward current \(I\) was assumed to be given. Figure 1 shows the density distribution in the root as the filament with current \(I = 1.5 \times 10^{11}\) A ascends with velocity \(v_0 = 500\) km/sec. Parameters that determine an external dipole field are chosen to fulfill the conditions that determine the equilibrium loss by the current. Directly above the filament a cavity is formed immediately, and an increased-density region is formed above and on the side of the latter. The lines of force of the current field that ascend together with the filament “draw apart” the coronal plasma. A dense shell is formed around a “bubble” with \(\beta < 1\), from which the plasma is “swept out.” The shell velocity is greater than the filament ascent velocity.

The MHD calculations in [3] support the conclusion on the near-boundary formation of the equality between the pressures of the plasma and magnetic field of the filament current \(\beta = 1\) of the dense shell that travels faster than the eruptive filament. In the region where the magnetic pressure prevails (\(\beta < 1\)) the coronal material rarefaction and the cavity formation occur. The maximum influence of the filament current field on the chromosphere is at the initial eruption stage when the filament is not far from the chromosphere.

In the case of equilibrium in the model of inverse polarity, at a certain distance from the filament, the structure of the field in the chromosphere is as if there were two saddle points under the photosphere. First, the points diverge as the current ascends, then they converge when the current reaches altitudes that are comparable to the scale of the external field, and, finally, one of the points also ascends along the line \(z = 0\). These two late stages are not interesting for the analysis, since the influence of the current field on the chromosphere becomes insignificant.

Let us consider the density variations near the saddle point located on the photospheric surface at \(z = s\) as the current ascends. The expression for \(\text{div } \vec{v}\), in which we allow for \(x \approx s \gg h\) and denote \(t = z - s\), has the form

\[
\text{div } \vec{v} = \frac{4Iv_0}{c^3} \frac{t(t^2 - 5y^2)}{s^2(t^2 + h^2)^2}.
\]

Obviously, above the photosphere there are two compression regions and two rarefaction regions, which are separated by the lines \(t = 0, y = \pm t/\sqrt{5}\) (Fig. 2). The upper regions originate from the specific features of the magnetic field: the plasma accelerates as it approaches the singular point and decelerates as it moves away from the latter. The lower regions emerge from the specific features of the electric field \(E_z\), which changes its sign on the line \(y = 0\), creating the divergent and convergent flows. To estimate the compression and, hence, the degree of heating of the chromospheric plasma (since compression brings about a temperature increase: the adiabatic compression of a monatomic gas implies \(\Delta T \sim \Delta \rho^{2/3}\)), we must integrate with respect to the time and obtain the density variation over a finite interval.