TWO-DIMENSIONAL MHD SIMULATION OF CHROMOSPHERIC EVAPORATION DRIVEN BY MAGNETIC RECONNECTION IN SOLAR FLARES

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ABSTRACT

Two-dimensional MHD simulation of chromospheric evaporation associated with a solar flare is performed. The thermal energy driving the evaporation is first supplied from the coronal magnetic field through the magnetic reconnection mechanism and is transported into the chromosphere by heat conduction. Nonlinear anisotropic heat conduction is taken into account. The results show that temperature distribution is similar to the cusp-like structure of long-duration-event (LDE) flares observed by the soft X-ray telescope (SXT) aboard Yohkoh satellite. We also reproduced the observation of the radio free-free emission.

Key words: Sun: flares — plasmas — MHD — conduction

1. INTRODUCTION

Chromospheric evaporation (Neupert 1968; Hirayama 1974) is an important process, which may be associated with energetic phenomena in the solar corona, such as flares, jets, and transient brightenings. The primary energy release mechanism for these phenomena is now believed to be the magnetic reconnection (e.g., Petschek 1964) high in the corona (e.g., Tsuneta 1996; Shibata 1996; Shimizu et al. 1992; Yokoyama & Shibata 1994, 1995). When the released energy in the corona through magnetic reconnection mechanism is transported to the upper chromosphere (Fig. 1), the dense plasma in the chromosphere is suddenly heated-up and expands because of this energy input. The induced pressure-gradient force drives the plasma to go up to the corona along the field lines. This upflow process is called the chromospheric evaporation.

There are several pieces of observational evidence for this process. One of them is the blue-shifted component of the soft X-ray spectral lines emitted from flare plasmas. The first observations were reported by Feldman et al. (1980). Recently, Yohkoh Bragg Crystal Spectrometer (BCS) shows a very good performance in observing this line-profile and the results were also reported by some authors (e.g., Culhane et al. 1992; Doschek et al. 1992; Ding et al. 1996). As for the imaging observation, Yohkoh Soft X-ray Telescope (SXT) is the first instrument which enabled us to see the moving plasma in the corona (Hudson et al. 1994; Savin 1996). Rather indirect but important evidence for the evaporation is the fact that the observed soft X-ray loops are filled with denser plasma than the surrounding corona. It is quite natural to consider this is the consequence of the evaporated plasma.

Figure 1. Schematic illustration of a solar flare. Thick solid lines show magnetic field.

The theoretical study on the evaporation has been done by means of numerical simulations. Many were performed as one-dimensional hydrodynamic simulations along a rigid magnetic loop without assuming particular energy release mechanism (e.g., Nagai 1980; Peres and Reale 1993; Fisher et al. 1985; Mariska et al. 1989; Gan et al. 1992). These simulations were quite developed in the sense that they take into account various physical processes, such as optically-thin radiation, Spitzer-type nonlinear heat conduction, viscosity, electron beam and so on. They were quite successful; they can reproduce various observed spectral line profiles in X-ray and ultraviolet ranges. Recently, Hori et al. (1997) developed a 'pseudo two-dimensional' model of flare loops based on the one-dimensional simulation. They considered a system of multiple loops each of which is geometrically adjacent, so that they are apparently two-dimensional, but thermally isolated, so that it can be dealt with by hydrodynamic one-dimensional calculation. The flare energy is injected successively from the innermost to the outermost loop, by which an energy release as well as magnetic field line closure by magnetic reconnection process is implicitly assumed. On the other hand, the two-dimensional

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simulation of flares have been developed, though they do not include the effect of evaporation. (Forbes & Priest 1983; Magara et al. 1996; Ugan 1996; Forbes & Malherbe 1991). Forbes, Malherbe, & Priest (1989) constructed a semi-analytical model of flare loops, by considering jump conditions across various discontinuities in the two-dimensional model of a flare. Although each of these fields is advanced well, there was still no self-consistent simulation of chromospheric convection including both the energy release by magnetic reconnection and thermal and dynamical processes of the evaporation.

In the study described in this paper, we performed two-dimensional MHD simulation of chromospheric evaporation associated with a solar flare. In this simulation, the thermal energy driving the evaporation is first supplied from the coronal magnetic field through the magnetic reconnection mechanism and is transported into the chromosphere by heat conduction. Nonlinear anisotropic heat conduction is taken into account (Yokoyama & Shibata 1997).

2. MODELS

The initial condition is in magnetohydrostatic equilibrium with antiparallel magnetic fields in the xz-plane, between which there is a current sheet in \(|z| < 0.4\delta\). In the current sheet, the direction of magnetic field in yz-plane changes continuously from negative \(z\) to the positive \(z\) with keeping its magnitude. The gas pressure is assumed to be uniform initially. The plasma beta, the ratio of the gas pressure to the magnetic pressure is taken to be \(\beta = 0.2\). A low-temperature and dense region, whose density is \(10^5\) times that of the other area, is located near the z-axis (\(0 < z < \delta\)). This represents a very idealized model of a chromospheric plasma. The magnetic field lines are well line-tied to the x-axis because of the large inertia of the heavy plasma. The units of length, velocity, and time in the simulations are \(\delta\), \(C_\alpha\), and \(\tau \equiv \delta/C_{\alpha}\), respectively, where \(\delta\) and \(C_\alpha\) are the thickness of the initial current sheet and the sound speed of the initial condition. We assume that the units of the temperature, the density, and the length are, respectively, \(T = 2 \times 10^6\) K, \(n_{\text{normal}} = 10^9\) cm\(^{-3}\), and \(\delta = 3 \times 10^8\) cm. In dimensional unit, the sound velocity is \(C_\alpha \approx 170\) km s\(^{-1}\), the Alfvén velocity is \(V_A \approx 410\) km s\(^{-1}\), and the heat conduction coefficient is \(\kappa \equiv \kappa_0 T^{3/2} \approx 5.7 \times 10^{14}\) erg cm\(^{-1}\) K\(^{-1}\) s\(^{-1}\) (where \(\kappa_0 \approx 10^{-26}\) erg cm\(^{-1}\) K\(^{-1}\) s\(^{-1}\) is the constant coefficient), and so the sound-crossing time \(\tau_s \equiv \delta/C_\alpha \approx 18\) s, the Alfvén-crossing time \(\tau_A \equiv \delta/V_A \approx 7.4\) s, and the heat conduction time \(\tau_{\kappa} \equiv \delta^2/(\kappa_{\text{normal}} k_B T/\kappa) \approx 2.2\) s, where \(k_B\) is the Boltzmann constant. Thus, the conduction time is comparable to the Alfvén-crossing time.

In order to initiate the magnetic reconnection, we assumed an anomalous resistivity model. The functional form is \(\eta = \kappa_0 T^{3/2}\) for \(\eta \geq \eta_c\), where \(\alpha\) is the nondimensional resistivity parameter, \(\eta_c \equiv J/\rho\) the nondimensional (relative ion-electron) drift velocity, \(\eta_c\) the threshold above which anomalous resistivity sets in, and \(J = (J_x^2 + J_y^2 + J_z^2)^{1/2}\) the total current density. [In dimensional form \(\eta_c \equiv J/(en)\), where \(e\) is the elementary charge and \(n\) is particle number density.] We also assumed that there is an upper limit, \(\eta_{\text{max}}\), for the resistivity. In this study, we fixed \(\eta_c = 100\), \(\alpha = 0.1\), and \(\eta_{\text{max}} = 0.3\). For the initial perturbation, we imposed a localized resistivity near the point \((x, z) = (0, 20)\)

\[
\eta = \eta_0 \left[ 2 \left( r/r_\eta \right)^3 - 3 \left( r/r_\eta \right)^2 + 1 \right] \quad \text{for} \quad r \leq r_\eta,
\]

and 0 for \(r > r_\eta\), where \(r \equiv \sqrt{x^2 + (z - 20)^2}\) is the distance, \(r_\eta = 0.86\) is the size of the region. The magnetic Reynolds number is \(Rm \equiv V_A \delta/\eta_0 \approx 25\).

We solved the two-dimensional \((\partial/\partial y = 0\) but \(\eta_\gamma, B_\gamma \neq 0\), nonlinear, time-dependent, resistive, compressible MHD equations. Ohmic heating and heat conduction are taken into account for heat loss and gain. The conduction coefficient is the Spitzer-type one which is proportional to \(T^{3/2}\) (e.g., Priest 1982). We also assumed that it is anisotropic, working only to the direction along the magnetic field line so that \(\kappa_\parallel = \kappa_0 T^{3/2}\) and \(\kappa_\perp = 0\) in the simulations, where \(\kappa_\parallel\) and \(\kappa_\perp\) are, respectively, the conductivity along and across the magnetic field. (Note that, if we take into account the numerical diffusivity, the conductivity across the field line is not zero. The effective value would be estimated as \(\kappa_\perp/k_\parallel \ll 0.01\). The specific heat ratio is taken to be \(\gamma = 5/3\). In the numerical procedures, the modified Lax-Wendroff method is adopted for the MHD part of the calculations, and the red-and-black over-relaxation method is adopted for the heat conduction part (e.g., Hirsch 1988). We assumed the symmetry and only the first quadrant \((0 < x < 17\) and \(0 < z < 125\) were calculated with nonuniform grids, the number of which is \(340 \times 440\). The minimum grid size is \(\Delta x = 0.01\) and \(\Delta z = 0.08\).

3. RESULTS

Figure 3 shows the results. Because of the enhanced resistivity around \(z = 20\), magnetic reconnection starts at this point. The reconnected field lines together with the frozen-in plasma are ejected from this
X-point to the positive and negative z-directions due to the tension force of the reconnected field lines. To complement this outflow, an inflow takes place from ±z-directions of the current sheet. At the boundary between this inflow and the outflow a slow-mode MHD shock is formed (Petschek 1964). Since there are two sets of inflow, from the positive and the negative z-directions, a pair of the shocks is formed, emanating from the neutral point. At these shocks, the plasma is heated up by shock heating. A pair of hot plasma jet is thus ejected from the neutral point. At the same time, a heat conduction front propagates from the hot region between the pair of shocks. It propagates rapidly (t = 10 and t = 20 of Fig. 3) because the temperature becomes lower as the temperature becomes higher. The conduction of heat is only in the direction along the magnetic field line. The outer edge of the conduction front, therefore, traces the magnetic field lines extending from the X-point. This temperature distribution (t = 25 of Fig.3) is very similar to the cusp-like structure of solar flare loops, which are observed by the soft X-ray telescope of Yohkoh satellite (Tsuneta et al. 1992). In the density distribution, a growing plasma mound can be seen (from t = 20 to t = 25 of Fig.3). This is the direct consequence of the so-called chromospheric evaporation. The chromospheric plasma is heated up and expands suddenly due to the penetration of the heat conduction front. The induced pressure-gradient drives a back-flow toward the corona. This flow carries up the dense plasma into the corona.

In this simulation, we found two interesting new features. One is a pair of high pressure humps on the evaporated plasma mound (around s = 9 at t = 30 of Fig. 4 indicated by a thick arrow; also indicated by a red arrow in Fig. 5). They are formed because of the collision between the evaporation flow from the chromosphere and the reconnection flow from the neutral point. The density enhancements in these humps are about three times as the environment density. Thus, it would be observable with e.g. the soft X-ray telescope aboard Yohkoh. The other new finding is the high density blob above the flaring loops (around (x, z) = (0, 12) at t = 25 of Fig. 3). This blob is formed by compression at the fast-mode MHD shock at the collision site of the reconnection jet with the flare loops. This blob has been already found as a high pressure region by Magara et al. (1996) by the adiabatic MHD simulation (see also Uga 1996). In our case this is seen as a dense blob. This is because our fast-mode MHD shock is isothermal shock while it was adiabatic one in Magara et al. The density jump is much larger for the isothermal shock than for the adiabatic one. This high density blob may have some relation with the observed hard X-ray source above the flare loops (Masuda et al. 1994).
4. DISCUSSION

The present results shown in this paper is that of the first two-dimensional calculation of coronal evaporation. The evolution from the rise phase to the early part of the decay phase of a solar flare should be reproduced in this simulation. In order to check this, we observe 'the computation results under the response of the real X-ray instruments, such as the soft X-ray telescope on board the Yohkoh satellite. It is compared with the real flares on the sun.

Figure 5 is thus derived image of the 'observation' of the simulation results after taking into account of the response of the thin Aluminium filter of the Yohkoh soft X-ray telescope (Tsuneta et al. 1991). This X-ray distribution is very similar to the cusp-like structure of loops of a long-duration-event (LDE) flare, which are observed by the soft X-ray telescope of Yohkoh satellite (e.g., Tsuneta et al. 1992; Tsuneta 1996; Forbes & Acton 1996). It is noted that the cusp-like structure of a flaring loop is found in various coronal phenomena in different scales, i.e., microflares in an active region (Yoshida et al. 1996), large-scale arcade reformation (Hiei et al. 1993).

At the same time, however, the density distribution of the simulation results suggest that the plasma in the cusp region is not the evaporated plasma. In relation with this fact, an interesting suggestion was reported by Flanaoka (1994). They found that there is no cusp-like structure in 17 GHz radio observation of an LDE-flare loop. Only a dense loop at the inner edge of the X-ray loop can be observed. The radio emission observed in 17 GHz is due to the thermal free-free radiation from optically thin plasma. It is expressed as $S \propto n^2T^{-1/2}$. This shows that the plasma in cusp-shape structure of the LDE-flare loop is hotter and less-denser than the inner edge of the same loop. Our simulation results reproduced the radio observation. The hot plasma distribution outside the dense loop may correspond to the observation of the hot plasma ridge outside the intense soft X-ray loop of the impulsive flare (Tsuneta 1997).

Through the very simplified model of a flaring loop, Fisher & Hawley (1990) estimated the apex temperature $T_A$ of strongly evaporating case. The saturated upper limit of this temperature is given by

$$T_A = \left( \frac{2.3}{\kappa_0} \right)^{2/7} \frac{Q L^2}{\kappa_0}$$

where $Q$ is the volumetric heating rate, $L$ is the apex to footpoint loop length, and $\kappa_0$ is the Spitzer coefficient of heat conduction. In our present reconnection model, heating rate is given as

$$Q = \frac{B^2 V_A}{4\pi L}$$

If we adopt the value which is used in the simulation, that is $B = 6$ G, $V_A = 410$ km s$^{-1}$, and $L = 3 \times 10^4$ km, where $L$ is the loop length estimated from the simulation results (see Fig. 5), $T_A$ becomes $5.3 \times 10^6$ K. This value is roughly consistent with the simulation results (see the upper-right panel of Fig. 4).

We finally give a comment on the observation of the reconnection inflow region (Fig. 6). In the reconnection process shown in this paper, an inward flow is required from the ambient to the reconnection site. Hudson & Kahn (1996) claimed that there is no evidence of this inflow found by the imaging observations by the Yohkoh SXT and that this is the strong

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Figure 6. Temporal evolution of one-dimensional plot in the inflow region near the magnetic neutral X-point along z = 20. The shown quantities are density (upper panel) and inflow velocity (lower panel). As time proceeds, density near the neutral point (x = 0) decreases.

evidence for rejection of the magnetic reconnection model of solar flares. According to the results we obtained, however, this could not be true. The temporal evolution of X-ray images in our simulations show that there is no apparent feature of inflow to the reconnection region. On the contrary, we found an expansion wave outward of the reconnection site (Fig.6). This is due to the successive suction of the plasma into the neutral point. This expansion wave has very low emission measure which is at least, four orders of magnitude smaller than the brightest feature in the solar flare. In conclusion, we suggest that the non-finding of the inflow feature can not be the strong evidence of the rejection of the reconnection model.

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