Oblique Shocks in the Magnetic Reconnection Jet in Solar Flares

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Abstract

Strong radio emission of energetic electrons is observed in some solar flares. The origin of energetic electrons is, however, not fully known. In this letter we suggest that oblique shocks are created in reconnection jets in solar flares, and that energetic electrons are accelerated by shocks. We examine 2D MHD simulations of magnetic reconnection with high spatial resolution by assuming an anomalous resistivity model. As a result, magnetic reconnection is found to occur after a secondary tearing instability at the current sheet. We find that, during nonsteady Petschek reconnection, oblique shocks are created by an even mode of the Kelvin–Helmholtz-like instability in the reconnection jet when we assume an anomalous resistivity model. Furthermore, bursty, time-dependent reconnection ejects many plasmoids from the diffusion region, and creates shocks. We suggest that these shocks can be possible sites of particle acceleration in solar flares.

Key words: acceleration of particles — magnetohydrodynamics: MHD — Sun: corona — Sun: flares — turbulence

1. Introduction

In some solar flares, strong radio emission of energetic electrons is observed. In this letter, we describe 2D MHD simulations of magnetic reconnection with a high spatial resolution, while suggesting a possible origin of energetic electrons. In a previous letter (Tanuma & Shibata 2005; hereinafter Paper I), we suggested that internal shocks can be created by a secondary tearing instability, and that they are possible sites of particle acceleration. We also suggested that strong hard X-ray emission from energetic electrons accelerated there was observed with Yohkoh HXT (Masuda et al. 1994) in impulsive flares (see also a review by Miller et al. 1997). In this letter we suggest that oblique shocks are created in the reconnection jet (reconnection outflow) by an even mode of the Kelvin–Helmholtz-like instability, and that they are possible sites of electron energization.

The SXT and HXT on the Yohkoh satellite discovered evidence of magnetic reconnection in solar flares (Tsuneta et al. 1992; Masuda et al. 1994; see also a review by Shibata 1999). “Fast” magnetic reconnection (Petschek 1964) occurs due to anomalous resistivity (Ugai 1986; Shibata et al. 1992; Yokoyama & Shibata 1997; Tanuma et al. 2003). Tanuma et al. (1999, 2001) studied reconnection initiated by a shock wave. They found that the Petschek reconnection occurs because of a “secondary tearing instability” (Furth et al. 1963; see also Magara & Shibata 1999; Shibata & Tanuma 2001) and that fast shocks are created in a reconnection jet. Recently, Paper I found that weak internal shocks are created in the reconnection jet as the result of a secondary tearing instability and a bursty time-dependent reconnection in the diffusion region. In this letter, we suggest that energetic electrons could be accelerated if strong, oblique, internal fast shocks could be created in solar flares. To test this possibility, we investigated how the strong internal shocks are created in the reconnection jet by performing MHD simulations with a high spatial resolution and such parameters that are different from those of Paper I. It removes any numerical noise and resolves the diffusion region. These simulations led us to propose that multiple, oblique, fast shocks are created in solar flares, and that they are possible sites of particle acceleration. Our simulation method is described in the next section. In section 3, we present simulation results. Finally, we discuss the results.

2. Model of Numerical Simulations

We calculated the nonlinear, time-dependent, resistive, and compressible MHD equations of magnetic reconnection. A Harris-like current sheet was assumed, i.e., $B(x, z) = B_0 \tanh(z/l_{\text{init}})x$, $\rho_g(x, z) = \rho_{g0} + (B_0^2/8\pi)[1 - \tanh^2(z/l_{\text{init}})]$, and $\rho(x, z) = (y\rho_g/T) = \rho_0 + (y/T_0)(B_0^2/8\pi)[1 - \tanh^2(z/l_{\text{init}})]$ in the initial equilibrium conditions, where $B_0$, $\rho_0$, and $\rho_{g0}$ are dimensionless variables, and $x = (1, 0)$. The current-sheet half-thickness was $l_{\text{init}} = 1$ in the initial conditions. The ratio of gas to magnetic pressure was $\beta = 8\pi\rho_{g0}/B_0^2 = 0.2$ ($|z| \gg 1_{\text{init}}$). The sound velocity and temperature were $C_{z,0} \equiv (y\rho_g/\rho)^{1/2} = 1$ (uniform) and $T = T_0 = 1$ (uniform), respectively. The initial Alfvén velocity was $v_{a,\text{init}} = B_0/(4\pi\rho_0)^{1/2} \approx 2.45$ ($|z| \gg 1_{\text{init}}$). The model was the same as that of Paper I, but an anomalous resistivity model was used: $\eta = \eta_0$ for $v_d \leq v_c$, and $\eta = \eta_0 + \alpha(v_d/v_c - 1)^2$ for $v_d > v_c$ (Tanuma et al. 1999, 2001; Paper I), where $v_d(= J/\rho)$, $\rho$, $J$, and $v_c$ are the drift velocity, mass density, current density, and threshold of onset of the anomalous resistivity; $\eta_0$ (“background resistivity”) is 0.001, which is much larger than the “numerical resistivity” because of the quite fine grid size (Ugai 1999; Tanuma et al. 2001; Paper I). The other parameters are $\alpha = 10.0$ and $v_c = 100.0$.

We normalized the velocity, length, and time by the sound
velocity \((C_{x,0} \sim 150\ \text{km}\ \text{s}^{-1})\), initial current-sheet thickness \((I_{\text{init}} \sim 3000\ \text{km})\), and \(I_{\text{init}}/C_{x,0} \sim 20\ \text{s}\), respectively. The units of temperature, density, gas pressure, and magnetic field strength were \(T_0 \sim 2 \times 10^6\ \text{K}, n_0 \sim 10^4\ \text{cm}^{-3}, p_{g0} \sim 10^{-1}\ \text{erg}\ \text{cm}^{-3}\), and \(B_0 \sim 2\ \text{G}\), respectively. The grid number was \((N_x, N_z) = (13000, 13000)\), and the grid size was \((\Delta x, \Delta z) = (0.013, 0.013)\) (uniform). The magnetic Reynolds number was \(R_{\text{m,t}} = \eta^{-1}_A \left(\Delta x N_z/\eta_0\right) = 4 \times 10^5\). We used a 2-step modified Lax–Wendroff method. The resistivity was initially enhanced \((\eta = 1.0)\) for a short time \((t < 4.0)\) in the central region of the current sheet (see also Ugai 1986; Paper I).

3. Results

The current sheet is unstable to the tearing instability (Furth et al. 1963). The magnetic dissipation time, Alfven time, and tearing instability time scale are \(t_{\text{dis}} \equiv I_{\text{init}}^2/\eta_0 \sim 1000\), \(\tau_A = I_{\text{init}}/v_A \sim 0.4\), and \(t_{\text{t}} = (t_{\text{dis}} \tau_A)^{1/2} \sim 20\), respectively. Figure 1 displays the time variation of the spatial distribution of the gas pressure in the reconnection region. Figure 1a shows that the current sheet becomes thin in a nonlinear phase of the tearing instability \((t \sim 7–17)\), the length of which is the most unstable wavelength of the tearing instability. It is \(\lambda_t \sim 5.6 R_{\text{m,t}}^{1/2} I_{\text{init}}^{1/2}\) (Magara & Shibata 1999; Tanuma et al. 2001; Paper I) where \(R_{\text{m,t}} \equiv v_A^{-1}(I_{\text{init}}/\eta_0) \sim 5000\); \(\lambda_t \sim 47\) in this simulation. The current-sheet thickness in this phase is \(l_t \sim (\lambda_t/2) R_{\text{m,t}}^{-1/2}\). The current sheet becomes a Sweet (1958)–Parker (1957) one (Tanuma et al. 1999, 2001; Paper I). The magnetic Reynolds number is \(R_{\text{m,t}} = (\lambda_t/2) v_A^{-1} I_{\text{init}}/\eta_0\) in this phase. Our results are consistent with these values: \(R_{\text{m,t}} \sim 5 \times 10^4\) and \(l_t \sim 0.1\).

Figure 1b shows that the sheet becomes unstable to the tearing instability again at \(t \sim 20\) (“secondary tearing instability”). Many quite small plasmoids appear due to the secondary tearing instability.

Figure 1c displays that three large islands with a size of \(\sim 2\) are created at \(t \sim 24\). They are ejected along the current sheet at \(t \sim 28\) (figure 1d). The bow shocks in front of the plasmoids also propagate along the current sheet. After plasmoid ejection, the current sheet becomes very thin, and the density decreases in the current sheet. The drift velocity \((v_d)\), reconnection rate \((v_{in}/v_A)\), and inflow velocity \((v_{in})\) toward the diffusion region increase quickly. Figure 2 shows the time variation of the drift velocity \((v_d)\), density, current density, and reconnection rate \((|\eta J_z|)\). An anomalous resistivity sets in because the drift velocity \((v_c)\) increases above the threshold \((v_c = 100)\) at \(t \sim 20.5\). The drift velocity, however, remains at around \(v_d \sim 100\) at \(t \sim 20.5–22.5\). At \(t \sim 22.5\), because it increases greatly, at a strong anomalous resistivity is excited. Then, a small diffusion region with a strong resistivity appears, and nonsteady Petschek (1964)-like (fast) reconnection starts. Many shocks are created inside the reconnection jet by bursty, time-dependent reconnection in this phase (Paper I).

Figure 1e shows that the reconnection jet starts to oscillate near to the diffusion region in a late phase \((t \sim 32)\) of fast reconnection. The reconnection rate, defined by \(|\eta J_z|\), is very high in this model (figure 2d). The even mode of the Kelvin–Helmholtz-like instability is excited (Biskamp 2000).

Then, the reconnection jet starts to oscillate in the current sheet. Figure 3a shows the 2D spatial distribution of the gas pressure at the central region at \(t \sim 35\). Many fast shocks are created in the reconnection outflow. The drift velocity reaches \(v_d \sim 340\) at \(t \sim 36\). The electric resistivity reaches the maximum value in this phase. The reconnection rate \((v_{in}/v_A)\) and the inflow velocity \((v_{in})\) also increase. During nonsteady fast reconnection, many plasmoids are created and ejected along the current sheet. Figure 1c shows, for example, that an island is ejected in the left direction from the diffusion region.

Figure 1f shows that, by the bursty, time-dependent reconnection and the even mode of the Kelvin–Helmholtz instability, many strong shocks are created in the reconnection jet. The velocity of strong shocks is \(\sim 0\), once created outside of the diffusion region, although they move inside and near to the diffusion region. Another type of shocks is also building up behind the plasmoids when the fast jet flow collides with the slower moving plasmoids. Figure 4 displays profiles of some variables in \(z=0.0\) at \(t = 35.0\). The profiles of the variables show that many pressure and density jumps are created in the jet (figures 4a and 4b). Figure 4c displays the profiles of the jet \((v_z)\), local sound \((C_s)\), and local fast wave \((\sqrt{(C_s^2 + v_d^2)})\) velocities. It also shows that the velocity of the jet becomes supersonic, i.e., the shocks are strong. Figure 4d displays the profile of \(v_z\). It shows that the reconnection jet oscillates in the current sheet. The inflow velocity \((v_{in})\) toward the diffusion region and the reconnection rate \((v_{in}/v_A)\) continue the oscillation. As a result, the reconnection jet becomes turbulent (Matthaeus & Lamkin 1985, 1986; Lazarian & Vishniac 1999; Fan et al. 2004) in the latest phase (figures 1d and 3b).

We examine the dependence of an anomalous resistivity model on the results. They are shown in figure 2. We assumed \(\eta_0 = 0.005\) and \(v_c = 20\) in the previous letter (Paper I). As a result, the even mode of the Kelvin–Helmholtz-like instability does not occur. In this letter, however, the diffusion region is shown to become smaller than that in the previous letter. On the other hand, the resistivity and reconnection rate (defined by \(|\eta J_z|\)) increase to be much higher than those given in the previous letter. Thus, the instability occurs because the asymmetric pattern can appear easily.

4. Discussion

The even mode of the Kelvin–Helmholtz-(like) instability has been observed in reconnection sites of some previous simulations (Biskamp et al. 1998; Arzner & Scholer 2001; Tanuma 2000; Tanuma & Shibata 2002, 2004). In this letter, however, we reveal that it occurs in the reconnection jet examining an anomalous resistivity model. Note that the difference between this letter and the previous one (Paper I) is an anomalous resistivity model: \(\eta_0 = 0.005\) and \(v_c = 20\) in the previous letter, while \(\eta_0 = 0.001\) and \(v_c = 100\) in this letter. In this letter, then, an asymmetric structure easily appears in the diffusion region, so that the even mode of the Kelvin–Helmholtz-like instability occurs, while it is not mentioned in the previous letter. As shown in figure 2d, in this model the reconnection rate \((|\eta J_z|)\) is 4-times larger than that of our previous model (Paper I). Furthermore, in this model the
Fig. 1. 2D spatial distribution of gas pressure in the reconnection region. The normalization units of time, length, and velocity are 20 s, 3000 km, and 150 km s$^{-1}$, respectively. (a) The current sheet becomes thin in a nonlinear phase of the tearing instability. (b) Many small plasmoids are created by the secondary tearing instability. (c) Anomalous resistivity sets in at $t \sim 22$, so that Petschek-like reconnection starts. Then, the reconnection rate $|\eta J_z|$ increases drastically. Large plasmoids are created. (d) They are ejected along the current sheet. Bow shocks also propagate with the plasmoids. (e) The even mode of the Kelvin–Helmholtz instability occurs, which also creates standing shocks. (f) The oscillation continues in a late phase of reconnection. It also creates fast shocks which are almost standing. The reconnection jet enters into a turbulent state.

Fig. 2. Time variation of (a) maximum drift velocity ($v_d$), (b) the density at the maximum $v_d$, (c) the current density at the maximum $v_d$, and (d) reconnection rate $|\eta J_z|$. (a) The solid line and dotted line show the result of this letter ($\eta_0 = 0.001$, $v_c = 100$) and that of the previous letter ($\eta_0 = 0.005$, $v_c = 20$), respectively. In this letter, the drift velocity is much larger than that in the previous letter. (b) The density is almost the same between the two results. (c) In this letter, the current density is larger than that in the previous letter. (d) In this letter, because the reconnection rate is much ($\sim 4$ times) higher than that in the previous letter, the oscillation of reconnection jet is excited and oblique shocks are created along the current sheet.

Fig. 3. Snapshots of 2D spatial distribution of gas pressure in the reconnection region (a) at $t \sim 35$ and (b) at $t \sim 40$. (a) Some oblique shocks (high pressure regions) are created near to the diffusion region in the current sheet because the reconnection jet oscillates due to the even mode of the Kelvin–Helmholtz instability. (b) In a later phase, the reconnection jet becomes turbulent.
Fig. 4. Profiles of variables in $z \sim 0.0$ at $t \sim 35$ (a late phase). (a), (b) The plots of the gas pressure and the density show that many strong shocks (jumps) are created. (c) The profiles of $v_x$, the local sound velocity ($C_s$), and the magnetosonic velocity [$v_{ms} = \sqrt{\left(v_A^2 + C_s^2\right)}$] are displayed by the solid line, dash-dotted line, and dotted line, respectively. We plot both $C_s$ and $-C_s$, and both $v_{ms}$ and $-v_{ms}$, $v_A$ and $C_s$ are nearly identical because the magnetic field is very weak at $z=0$. The jet is usually supersonic, while it becomes subsonic at the shocks. The jumps are strong fast shocks. (d) The profile of $v_z$ represents that the reconnection jet actually oscillates along the current sheet due to the even mode of the Kelvin–Helmholtz instability.

As shown by our simulations, many strong, fast, oblique shocks are created by the Kelvin–Helmholtz instability in the reconnection region. These shocks are possible sites for the particle acceleration in solar flares (Tsuneta & Naito 1998; see also Paper I). We also reveal that another type of shocks, i.e., a propagating shock, is created by the nonsteady reconnection and plasmoid ejection. On the other hand, many radio bursts are observed at solar flares. The frequency of radio emission is calculated by $f \propto \sqrt{\frac{l}{m}}$, i.e., depending on the height of shocks if the emission mechanism is the fundamental plasma emission. The origin, mechanism, and location of a dm-wave burst (negative drift) are still unknown. The time-dependent bursty reconnection, however, can explain them by time-dependent plasmoid ejections. We conclude that these nonsteady plasmoid ejection and bow shock propagation can be possible candidates to explain the slowly drifting structures (Karlický 2003) and narrow-band dm-spikes (Bártá & Karlický 2001) in the upward jet. The time scale for the tearing instability in the diffusion region is $\sim 30 \left(l/10^5 \text{ cm}\right)^{3/2} \left(v_A/10^8 \text{ cm s}^{-1}\right)^{-1/2} \left(T/10^6 \text{ K}\right)^{3/4}$ s if we assume Spitzer conductivity, and that the diffusion region thickness is $l$. This explains the time scales associated with the above phenomena (Karlický 2003). The duration of drift is determined by the time scale of the propagation of plasmoids; for example, it is $30$ s if a plasmoid propagates for $30000$ km at velocity of $\sim 1000$ km s$^{-1}$. Furthermore, time-dependent bursty reconnection can also explain the X-ray observation of nonsteady plasmoid ejections in downflows toward the magnetic loop (McKenzie & Hudson 1999; Asai et al. 2004).

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