Oblique Shocks in the Reconnection Jet in Solar Flares

Syuniti Tanuma and Kazunari Shibata
Kwasan Observatory, Kyoto University, Kyoto, 607-8471, Japan

Abstract. The strong hard X-ray emission of energetic electrons is observed in some solar flares. The origin of energetic electrons is, however, not known fully. Then, we suggest that the internal shocks are created in the reconnection jet in solar flares, and that energetic electrons are accelerated by the shocks. In this paper, we examine 2D MHD simulations of magnetic reconnection with a high spatial resolution. As the results, the magnetic reconnection occurs after the secondary tearing instability at the current sheet. We find that, during the non-steady Petschek reconnection, the oblique strong shocks are created by the Kelvin-Helmholtz-like instability in the reconnection jet when we assume an anomalous resistivity model. The oblique shocks can be possible sites of the particle acceleration in the solar flares.

1. Introduction

The strong hard X-ray emission are observed in some solar flares. The origin of the energetic electrons, however, is not known fully. In this paper, then, we suggest that the internal shocks can be created in the reconnection jet and that the energetic electrons are created by these shocks. To examine its possibility, we perform 2D MHD simulations of the magnetic reconnection with a high spatial resolution.

2. Model of Numerical Simulations

We solve the nonlinear, time-dependent, resistive, compressible MHD equations. A rectangular computation box with 2D Cartesian coordinates in the x-z plane is assumed, where the x and z axes are in the horizontal and vertical directions, respectively. We assume an ideal gas, i.e., \[ p_g = (\gamma - 1)\epsilon, \] where \( \gamma \) is the specific heat ratio (=5/3). The velocity and magnetic field are \( \mathbf{v} = (v_x, v_z) \) and \( \mathbf{B} = (B_x, B_z) \). Gravity is neglected.

In the initial equilibrium conditions, a Harris-like current sheet is assumed, i.e., \( \mathbf{B}(x, z) = B_0 \tanh(z/l_{\text{init}}) \mathbf{x}, \) \( p_g(x, z) = p_{g0} + (B_0^2/8\pi)[1 - \tanh^2(z/l_{\text{init}})], \) \( \rho(x, z) = (\gamma p_g/T) = \rho_0 + (\gamma/T_0)(B_0^2/8\pi)[1 - \tanh^2(z/l_{\text{init}})], \) where \( B_0, p_{g0}, \) and \( \rho_0, \) are dimensionless variables, and \( \mathbf{x} = (1, 0). \) We assume in the initial conditions that the current-sheet half-thickness is \( l_{\text{init}} = 1. \) The ratio of gas to magnetic pressure is \( \beta = 8\pi p_{g0}/B_0^2 = 0.2 \) (\(|z| \gg l_{\text{init}}\)). The initial gas pressure is \( p_{g0} \) outside the current sheet, and \( p_{g0} + B_0^2/8\pi = p_{g0}(1 + 1/\beta) \) inside the current sheet. The total pressure is uniform. We assume \( p_{g0} = 1/\gamma = 0.6 \) and \( \rho_0 = 1. \) The sound velocity and temperature are \( C_s \equiv (\gamma p_g/\rho)^{1/2} = 1 \) (uniform) and
\( T = T_0 = 1 \) (uniform). The initial Alfvén velocity is \( v_A^{\text{init}} = B_0/(4\pi\rho_0)^{1/2} \simeq 2.45 \) \(|z| \gg z_0^{\text{init}}\). We assume an anomalous resistivity model as follows: \( \eta = \eta_0 \) 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. 2001; Tanuma & Shibata 2005, 2006), where \( v_d(\equiv J/\rho) \), \( \rho \), \( J \), and \( v_c \) are the dimensionless, relative ion-electron drift velocity, mass density, current density, and threshold above which the anomalous resistivity sets in. We also assume that the resistivity does not exceed \( \eta_{\text{max}} = 1 \). In this paper, we assume a “background resistivity” \( \eta_0 = 0.001 \), which is sufficiently larger than the “numerical resistivity” because of the grid size (Tanuma et al. 2001; Tanuma & Shibata 2005, 2006). The other parameters are \( \alpha = 10.0 \) and \( v_c = 100.0 \).

We normalize the velocity, length, and time by the sound velocity \( (C_s) \), initial current-sheet thickness \( (H) \), and \( H/C_s \), respectively. The units of normalization are \( C_s \sim 150 \text{ km s}^{-1} \), \( H \sim 3000 \text{ km} \), and \( \tau \equiv H/C_s \sim 20 \text{ s} \). The units of temperature, density, gas pressure, and magnetic field strength are \( T_0 \sim 2 \times 10^6 \text{ K} \), \( n_0 \sim 10^9 \text{ cm}^{-3} \), \( p_0 \sim 10^{-1} \text{ erg cm}^{-3} \), and \( B_0 \sim 2 \text{ G} \), respectively. The grid number is \( (N_x, N_z) = (13000, 1300) \), and the grid size is \( (\Delta x, \Delta z) = (0.013, 0.013) \) (uniform). We assume that the top \( (z = +8.45) \) and bottom \( (z = -8.45) \) surfaces are symmetric boundaries, and that the right \( (x = +84.5) \) and left \( (x = -84.5) \) ones are periodic. The simulation box size is \( (L_x, L_z) = (169.0, 16.9) \). The magnetic Reynolds number is \( \text{Re}_m^{\text{init}} \equiv v_A^{\text{init}}L_x/\eta_0 \sim 414505 \sim 4 \times 10^5 \). We use a 2-step modified Lax-Wendroff method. The resistivity is initially enhanced \( (\eta = 1.0) \) for a short time \( (t < 4.0) \) in the central region of the current sheet (Tanuma & Shibata 2005, 2006).
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3. Results

The current sheet is unstable because of the tearing instability (Tanuma et al. 2001; Tanuma & Shibata 2005, 2006). The magnetic dissipation time, Alfvén time, and tearing instability time scale are $\tau_{\text{dis}}^{\text{init}} = l_{\text{init}}^2/\eta_0 \sim 100$, $\tau_A^{\text{init}} = l_{\text{init}}/v_A^{\text{init}} \sim 0.4$, and $\tau_t^{\text{init}} = (\tau_{\text{dis}}^{\text{init}}/\tau_A^{\text{init}})^{1/2} \sim 6$, respectively. The tearing instability is initiated in the current sheet by the initial perturbation. The current sheet becomes gradually thinner in the nonlinear phase of the tearing instability ($t \sim 7 - 18$). Its length is comparable to the most unstable wavelength of the tearing instability, i.e., $\lambda_t \sim 5.6 \text{Re}_{\text{m,z}}^{1/4}$ (Tanuma et al. 2001; Tanuma & Shibata 2005, 2006), where $\text{Re}_{\text{m,z}} \equiv v_A^{\text{init}} l_{\text{init}}^{\text{init}}/\eta_0 \sim 500$, $\lambda_t \sim 20$ in this simulation.

The current-sheet thickness in this phase is $l_t \sim (\lambda_t/2)\text{Re}_{\text{m,t}}^{-1/2}$, i.e., the current sheet corresponds to a Sweet-Parker sheet. The magnetic Reynolds number is $\text{Re}_{\text{m,t}} = (\lambda_t/2)v_A^{\text{init}}/\eta_0$ in this phase. Furthermore, $\text{Re}_{\text{m,t}} \sim 2500$ and $l_t \sim 0.2$. These values explain our results well.

Figure 1 shows the time variation of the spatial distribution of the gas pressure in the reconnection region. The current sheet becomes very long so that it becomes unstable to the tearing instability again at $t \sim 17$ (“secondary tearing instability”). Many small plasmoids are created along the diffusion region by the secondary tearing instability in the thin current sheet just after the Sweet-Parker sheet is created. The top panel of figure 1 shows that many small islands are coalesced to become three large islands with the size of $\sim 2$ at $t \sim 24$. The 2nd panel of figure 1 shows that three islands are propagating along the current sheet at $t \sim 28$. The bow shocks in front of the plasmoids are also propagating along the current sheet. After the plasmoid-ejection, the current sheet becomes much thinner, and density in the current sheet decreases. At the same time, the drift velocity ($v_d$), reconnection rate ($v_{ni}/v_A$), inflow velocity ($v_{ni}$) toward the diffusion region increase drastically. The drift velocity ($v_d$) reaches the threshold ($v_c = 100$) at $t \sim 21.5$ so that the anomalous resistivity sets in. The drift velocity, however, remains around $v_d \sim 100 - 110$ at $t \sim 20.5 - 22.5$. At $t \sim 22.5$, the drift velocity increases a great deal above the threshold. The strong anomalous resistivity is excited so that non-steady Petschek-like (fast) reconnection starts. Then, a small diffusion region with a strong electric resistivity appears, and is
accompanied by slow shocks. Due to the bursty, time-dependent reconnection, many high pressure regions are created inside the reconnection jet (Tanuma & Shibata 2005).

The 3rd panel of figure 1 shows that the reconnection jet starts to oscillate near the diffusion region in the latest phase \((t \sim 32)\) of fast reconnection. Figure 2 shows the time variation of the reconnection rate defined by \(|\eta J_y|\). It is very high in this model so that Kelvin-Helmholtz(-like) instability (odd-mode) is excited. Then, the reconnection jet starts to oscillate in the current sheet. Many fast shocks are created in the reconnection outflow. The drift velocity reaches \(v_d \sim 350\) at \(t \sim 36\). The electric resistivity reaches the maximum value in this phase. The reconnection rate \((v_{in}/v_A)\) and inflow velocity also increases. During the non-steady fast reconnection, many plasmoids are created and ejected along the current sheet. The 3rd panel of figure 1 shows, for example, that a small island is ejected in the left direction from the diffusion region.

The bottom panel of figure 1 shows that the Kelvin-Helmholtz instability continues at the central region (Tanuma & Shibata 2006). By the bursty, time-dependent reconnection and Kelvin-Helmholtz instability, many strong shocks are created in the reconnection jet. They are almost standing shocks, once created outside the diffusion region, although they move in and near the diffusion region. Many pressure and density jumps are created in the jet. The reconnection jet becomes supersonic. They are strong shocks. The inflow velocity \((v_{in})\) toward the diffusion region and reconnection rate \((v_{in}/v_A)\) continue the oscillation. As a result, the reconnection jet becomes in a turbulence in the latest phase (the bottom panel of figure 1).

4. Discussion

Kelvin-Helmholtz(-like) instability occurs in the reconnection jet so that many strong, oblique shocks are created in this paper \((\eta_0 = 0.001\) and \(v_c = 100\); Tanuma & Shibata 2006). On the other hand, in another model with different parameters, the instability does not occur and only weak fast shocks are created \((\eta_0 = 0.005\) and \(v_c = 20\); Tanuma & Shibata 2005). In the typical mode, the reconnection rate shown in figure 2 is very higher than that in the other model.

If many strong, oblique shocks are created by the Kelvin-Helmholtz(-like) instability in the reconnection jet, they can be possible sites of acceleration of the energetic electrons (Tanuma & Shibata 2005, 2006). Time-dependent bursty reconnection can also explain the X-ray observation of non-steady plasmoid ejections in the downflows toward the magnetic loop. Furthermore, the oscillation of jet may also explain the oscillations in “downflow” toward the magnetic loop.

Acknowledgments. The authors thank the anonymous referee for careful reading and fruitful comments.

References