A THEORETICAL MODEL OF A THINNING CURRENT SHEET IN THE LOW-\(\beta\) PLASMAS

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ABSTRACT

Magnetic reconnection is an important physical process in various explosive phenomena in the universe. In previous studies, it was found that fast reconnection takes place when the thickness of a current sheet becomes on the order of a microscopic length such as the ion Larmor radius or the ion inertial length. In this study, we investigated the pinching process of a current sheet by the Lorentz force in a low-\(\beta\) plasma using one-dimensional magnetohydrodynamics (MHD) simulations. It is known that there is an exact self-similar solution for this problem that neglects gas pressure. We compared the non-linear MHD dynamics with the analytic self-similar solution. From the MHD simulations, we found that with the gas pressure included the implosion process deviates from the analytic self-similar solution as \(t \rightarrow t_0\), where \(t_0\) is the explosion time when the thickness of a current sheet of the analytic solution becomes 0. We also found that a pair of MHD fast-mode shocks is generated and propagates after the formation of the pinched current sheet as \(t \rightarrow t_0\). On the basis of the Rankine–Hugoniot relations, we derived the scaling law of the physical quantities with respect to the initial plasma beta in the pinched current sheet. Our study could help us estimate the physical quantities in the pinched current sheet formed in a low-\(\beta\) plasma.

Key words: magnetohydrodynamics (MHD) – shock waves – stars: magnetic field

1. INTRODUCTION

It has been known that magnetic reconnection plays important roles in many energetic explosions such as solar flares (Zweibel & Yamada 2009; Shibata & Magara 2011). The knowledge about solar flares has been applied to other explosive phenomena, such as superflares on solar-type stars (Maehara et al. 2012) and soft gamma repeaters (Lyutikov 2006; Masada et al. 2010; Meng et al. 2014). In a classical standard model of solar flares, the magnetic energy is rapidly converted to the thermal and kinetic energies of plasmas through magnetic reconnection (Carmichael 1964; Sturrock 1966; Hirayama 1974; Kopp & Pneuman 1976; Shibata et al. 1995; Shibata & Magara 2011). Such rapid reconnection requires a much shorter current sheet width than its length (Yamada et al. 2010). Therefore, how thin current sheets are established is a central issue for understanding the origin of explosive phenomena.

Magnetohydrodynamics (MHD) processes to form a thin current sheet have been theoretically studied by many authors. In Forbes & Priest (1995) and Lin & Forbes (2000), they numerically and analytically investigated the formation of a current sheet below an erupting flux rope during a solar flare. There are many observational supports of the formation and the evolution of a current sheet between a cusped-shaped flaring loop and an erupting flux rope (Liu et al. 2010). The formation of a current sheet between an emerging flux and a preexisting magnetic field in the solar corona has been numerically studied (Forbes & Priest 1984; Shibata et al. 1992; Yokoyama & Shibata 1996). In McLaughlin et al. (2009), the collapse of a null point to a current sheet triggered by MHD fast-mode waves was investigated.

A long and thin current sheet can form many fine-scale current sheets inside it through some MHD instabilities. Such a current sheet can be fragmented to form magnetic islands or plasmoids through the tearing or plasmoid instabilities (Furth et al. 1963; Loureiro et al. 2007). The plasmoids are pulled against each other by a Lorentz force. When plasmoids collide with each other, new current sheets will be formed between them. This process can operate at multi-spatial scale, depending on the Lundquist number (Tajima et al. 1987; Shibata & Tanuma 2001; Bártá et al. 2011). Observational support from the direct imaging observation of the plasmoid formation in a solar flare was given by Takasao et al. (2012) for the first time. They also found that the plasma blobs (possibly plasmoids) in the current sheet collided or merged with each other before they were ejected from the current sheet. The pinching process of a current sheet by a Lorentz force like the last two processes above is the main subject of this paper.

In this study, we investigated the implosion process in which current sheets are pinched by the magnetic pressure gradient. In a previous study, the self-similar solution in the implosion process was analytically derived under the assumption that the gas pressure can be ignored (Imshennik & Syrovatski 1967; Bulanov & Olshanetskii 1984; Tajima et al. 1987). Forbes (1982) investigated the implosion process, including the gas pressure, and found the generation of two MHD fast-mode shocks as a result of a significant increase in the gas pressure at the center of the current sheet. However, in that simulation, the current sheet was incompletely pinched because of the low spatial resolution. To study the pinching process of the thin current sheet formed behind the shocks in detail, we performed a series of numerical simulations with a sufficient spatial resolution.

In Section 2, the self-similar solution proposed by Tajima et al. (1987) is introduced. In Section 3, we show the basic equations and initial conditions of the numerical simulations. In Section 4, we show the results of the numerical simulations and the power-law scalings of the physical quantities in the pinched current sheets. In Section 5, we analytically derive the power-law relation shown in Section 4. Finally, we summarize our conclusions in Section 5.

2. SELF-SIMILAR SOLUTION

The situation that we consider is similar to that of Tajima et al. (1987), where a current sheet is pinched by a Lorentz
force (see Figure 1). They derived a self-similar solution using the two-fluid equations and Maxwell’s equations. Here we derive the same self-similar solution from 1D MHD equations to review the basic physics and assumptions of the solution. We start from the basic 1D MHD equations:

\[
\frac{\partial}{\partial t} \rho = - \frac{\partial}{\partial x} (\rho v),
\]
\[
\frac{\partial}{\partial t} (\rho v) = - \frac{\partial}{\partial x} \left( \rho + \rho v^2 \right) - \frac{j B}{c},
\]
\[
\frac{\partial}{\partial t} B = - \frac{\partial}{\partial x} \left( v B \right),
\]
\[
\frac{\partial}{\partial x} B = \frac{\partial}{\partial x} B.
\]

Now we introduce dimensionless physical quantities as below:

\[
x^* \equiv x/L,
\]
\[
t^* \equiv t/\tau_A,
\]
\[
\rho^*(t^*) \equiv \rho(t)/\rho_0,
\]
\[
p^*(x^*, t^*) \equiv p(x, t)/p_0,
\]
\[
B^*(x^*, t^*) \equiv B(x, t)/B_0,
\]
\[
v^*(x^*, t^*) \equiv v(x, t)/V_A,
\]

where \(L\) is the initial current sheet thickness, \(V_A\) is the initial Alfvén speed, \(\tau_A \equiv L/V_A\) is an Alfvén timescale, and the initial values of the physical quantities are denoted by 0. The dimensionless equations can be written as

\[
\frac{\partial}{\partial t^*} \rho^* = - \frac{\partial}{\partial x^*} (\rho^* v^*),
\]
\[
\frac{\partial}{\partial t^*} (\rho^* v^*) = - \beta_0 \frac{\partial}{\partial x^*} p^* - \frac{\partial}{\partial x^*} \left( \rho^* v^2 \right) - \frac{\partial}{\partial x^*} \left( B^2 / 2 \right),
\]
\[
\frac{\partial}{\partial x^*} B^* = (v^* B^*),
\]

where \(\beta_0 = p_0 / (B_0^2 / 2)\) is the initial plasma beta.

We introduce a scale factor \(a(t)\) as follows

\[
v^* = \frac{\dot{a}(t^*) x^*}{a(t^*)},
\]

where a dot represents the time derivative. An ansatz is that the velocity is linear in \(x^*\). For simplicity, we assume that the density in the current sheet, is spatially uniform. From Equations (5) and (8), we obtain

\[
\rho^* = \frac{1}{a(t^*)}.
\]

We also assume that \(B^*\) is linear in \(x^*:\ B^*(x^*, t^*) = B^*(t^*) x^*\). Then from Equations (7) and (8), we obtain

\[
B^* = \frac{1}{a(t^*)^2}.
\]

We take the following form of the scaling factor

\[
a(t^*) \propto \left( \frac{t_0 - t}{t_0} \right)^k = \left( 1 - \frac{t}{t_0} \right)^k,
\]

where \(t_0\) is the explosion time when the thickness of the current sheet goes to 0. Neglecting the gas pressure term in Equation (6), we obtain the index \(k\) and explosion time \(t_0\) from Equations (6) and (11)

\[
a(t^*) \propto \left( 1 - \frac{t}{t_0} \right)^{2/3},
\]
\[
t_0 = \frac{\sqrt{2}}{3}.
\]

From Equation (12), we obtain the following expressions:

\[
v^*(x^*, t^*) \propto \left( 1 - \frac{t}{t_0} \right)^{-1},
\]
\[
\rho^*(t^*) \propto \left( 1 - \frac{t}{t_0} \right)^{-2/3}.
\]
3. BASIC EQUATIONS AND INITIAL CONDITIONS OF NUMERICAL SIMULATIONS

All the physical quantities are functions of \( x \) and \( t \), and a magnetic field is considered only in the \( y \)-direction \( (B_y) \). The basic equations are as follows:

\[
\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x}(\rho v_x) = 0,
\]

\[
\frac{\partial}{\partial t} (\rho v_x) + \frac{\partial}{\partial x} \left( \rho v_x^2 + p - \frac{B_y^2}{8\pi} \right) = 0,
\]

\[
\frac{\partial B_y}{\partial t} + \frac{\partial}{\partial x} (v_y B_x - B_y v_x) = 0,
\]

\[
\frac{\partial e}{\partial t} + \frac{\partial}{\partial x} \left( e + p + \frac{B_y^2}{8\pi} v_x \right) = 0,
\]

where the total energy density \( e \) is defined as

\[
e \equiv \frac{p}{\gamma - 1} + \frac{\rho v_x^2}{2} + \frac{B_y^2}{8\pi}.
\]

We adopted the self-similar solution of Tajima et al. (1987) as the initial condition of the numerical simulations. The initial condition is not an MHD equilibrium, and after the simulation’s start, the current sheet is pinched by a Lorentz force as shown in Figure 1. The time interval and the last time are, respectively, \( \Delta t = 0.9 \) and \( t_{ad} = 1.5 \) in the units of the Alfvén time.

Note that these scaling laws are the same as those originally obtained by Imshennik & Syrovatski (1967).

If the plasma adiabatically evolves, we obtain the gas pressure in the following form:

\[
p^* (t^*) \propto \left( 1 - \frac{t^*}{t_0} \right)^{-4/3},
\]

\[
B_y^* (x^*, t^*) \propto \left( 1 - \frac{t^*}{t_0} \right)^{-10/9}.
\]

This indicates that the gas pressure at the center should explode as \( t \to t_0 \). Therefore, the pinching by a Lorentz force must stop, which cannot be described by the linear theory. To study the evolution near and after the explosion time, we numerically investigated the implosion of a current sheet using 1D MHD simulations.

Figure 2. Time evolution of the implosion process of \( p \) and \( B \), with \( \beta_0 = 10^{-2} \). The time interval and the last time are, respectively, \( \Delta t = 0.9 \) and \( t_{ad} = 1.5 \) in the units of the Alfvén time.
4.2. $\beta_0$ Dependence of Physical Quantities the in the Current Sheet

We investigate the $\beta_0$ dependence of the physical quantities in the current sheet. The range of $\beta_0$ is $-3 \leq \beta_0 \leq -1$. The $\beta_0$ dependence of the density, pressure, and magnetic field strength just behind the shock is shown in Figure 4. Figure 4 implies the following power-law relations:

\begin{align}
\frac{\rho}{\rho_0} &\propto \beta_0^{-0.58}, \\
\frac{p}{p_0} &\propto \beta_0^{-0.97}, \\
\frac{B_y}{B_{0y}} &\sim \text{const.}
\end{align}

(25) (26) (27)

We also investigated the $\beta_0$ dependence of the thickness of a current sheet, $L_{\text{min}}$. We define $L_{\text{min}}$ as follows

\begin{align}
L(t) = \frac{B_{\text{max}}(t)}{J_{\text{max}}(t)},
\end{align}

(28)

where $B_{\text{max}}(t)$ and $J_{\text{max}}(t)$ are, respectively,

\begin{align}
B_{\text{max}} &\equiv \max \left( B_y(x, t) \right), \\
J_{\text{max}} &\equiv \max \left( J_x(x, t) \right).
\end{align}

As a result of the numerical simulations, we found a scaling law of $L_{\text{min}}$ (see Figure 5) as follows:

\begin{align}
L_{\text{min}} &\propto \beta_0^{1.177}.
\end{align}

(29)

To study the formation of the pinched current sheet in detail, we tracked selected Lagrangian particles that are initially located in the initial current sheet. Figure 6 shows the trajectories of the Lagrangian particles on the time–distance diagrams of $v_x$, $\log_{10} p$, and $B_y$. From Figure 6, we found that the pinched current sheet is formed from a part of the initial current sheet, not from the whole of it.

5. THE ANALYTICAL DISCUSSION ABOUT THE $\beta_0$ DEPENDENCE OF PHYSICAL QUANTITIES

In this section, we aim to derive the $\beta_0$ dependence of the physical quantities in the pinched current sheet. As shown in Section 4, we confirmed the pinching process by the Lorentz force and the propagation of the MHD fast-mode shocks, as shown in Forbes (1982). Behind the fast-mode shock, the physical quantities are determined by the Rankine–Hugoniot
equations
\[
\frac{v_{x2}}{v_{x1}} = \frac{1}{X}, \quad (30)
\]
\[
\frac{B_{z2}}{B_{z1}} = X, \quad (31)
\]
\[
\frac{p_2}{\rho_1} = \gamma \dot{M}_t \left( 1 - \frac{1}{X} \right) + \frac{1 - X^2}{\beta_1} + 1, \quad (32)
\]
where the physical quantities in front of and behind the shock are denoted by 1 and 2, respectively, \( \dot{M}_t = v_{x1}/c_{x1} \) is the Mach number, and \( X = \rho_2/\rho_1 \) is the density ratio, which is the positive solution of
\[
2(2 - \gamma)X^2 + \left[ 2\beta_1 + (\gamma - 1)\beta_1\dot{M}_t^2 + 2 \right]\gamma X
- \gamma(\gamma + 1)\beta_1\dot{M}_t^2 = 0. \quad (33)
\]
From Equations (31) and (33), the compression ratio is limited in the range of
\[
1 < \frac{B_2}{B_1} < \frac{\gamma + 1}{\gamma - 1}. \quad (34)
\]
where the upper limit is 4 for \( \gamma = 5/3 \). Therefore, the maximum value of the strength of the magnetic field should weakly depend on \( \beta_0 \) (\( B_{\text{max}}/B_0 \sim \text{const} \) with respect to \( \beta_0 \), Equation (27)). The current sheet, which is pinched by the shocks, is in the MHD equilibrium (i.e., \( v_c = 0 \)). This means that the gas pressure at the center of the current sheet, \( p_{\text{max}} \), should be comparable to the magnetic pressure outside of the current sheet:
\[
p_{\text{max}} \sim \frac{B_{\text{max}}^2}{8\pi} \sim \frac{B_0^2}{8\pi}. \quad (35)
\]
which gives the relation (26)
\[
\frac{p_{\text{max}}}{p_0} \sim \beta_0^{-1}. \quad (36)
\]
where \( p_0 \) is the initial gas pressure. From Equation (36) and the adiabatic condition, the density at the center (which takes the
maximum value, Equation (25)) can be written by
\[ \frac{\rho_{\text{max}}}{\rho_0} \sim \left( \frac{\rho_{\text{max}}}{\rho_0} \right)^{1/\gamma} \sim \beta_0^{-3/5}, \]
where \( \rho_0 \) is the initial plasma density and \( \gamma = 3/5 \) is the adiabatic index.

To derive the \( \beta_0 \) dependence of \( L_{\text{min}} \), relation (29), we defined \( B_1 \) as the magnetic field strength of the maximum value in the pinched current sheet. As shown in Lagrangian particle trajectories, note that the pinched current sheet is formed from not the whole of the initial current but a part of it. We also defined the length of the part in the initial current sheet as \( L_0^* \) and the maximum value of the magnetic field strength in the part as \( B_0^* \). In our numerical simulations, we confirmed \( L_0^* \ll L_0 \) by the Lagrangian particle motions. Near the center of the initial current sheet, the magnetic field strength linearly depends on \( x \),
\[ \frac{L_0^*}{L_0} \sim \frac{B_0^*}{B_0}. \]

We considered the conservation law of a magnetic flux,
\[ B_0^* L_0^* = B_1 L_{\text{min}}, \]
\[ \Rightarrow B_0^* L_0^* = B_0 L_{\text{min}}. \]

From Equations (38) and (39), we obtained
\[ \frac{L_{\text{min}}}{L_0} \sim \left( \frac{L_0^*}{L_0} \right)^2. \]
Finally, from the mass conservation law,

\[ \frac{L_0}{L_{\text{min}}} = \frac{\rho_{\text{max}}}{\rho_0} \sim \beta_0^{-3/5}. \]  

(41)

From Equations (40) and (41), we derived

\[ \frac{L_{\text{min}}}{L_0} \sim \beta_0^{6/5}, \]  

(42)

which is consistent with Equation (29).

6. DISCUSSION

In this paper, we studied the implosion process of current sheets in a low-\(\beta\) plasma using one-dimensional ideal MHD simulations. We confirmed that the self-similar solution by Tajima et al. (1987) holds before the explosion time. The plasma adiabatically evolves until the explosion time. At \(t \sim t_0\), the behavior of the current sheet deviates from the analytical self-similar solution, and the MHD fast-mode shocks are formed because of the increase in the gas pressure at the center of the current sheet. After a pair of MHD fast-mode shocks propagate away from the current sheet, the pinched current sheet is formed between a pair of the shocks. We studied the \(\beta_0\) dependence of the maximum values of the physical quantities inside the current sheet and the minimum thickness of the current sheet (Equations (36), (37), and (42)).

In contrast to our study, the \(\eta\) dependence of the physical quantities and of the reconnection rate was investigated by McClymont & Craig (1996) in MHD simulations similar to ours, where \(\eta\) is the magnetic diffusivity. They found that the effect of the gas pressure does not significantly change the dependence of the thickness of the current sheet on \(\eta\), which means that our scaling laws can be easily extended to the cases with the resistivity.

In McLaughlin et al. (2009), the collapse of a null point to a current sheet triggered by MHD fast-mode waves was investigated. In their simulations, successive current sheet pinching was observed. The pinching is done by the gas pressure of the heated plasma in the outflow region, while the pinching in our study is done by a Lorentz force. To extend this study, we are currently performing two-dimensional MHD simulations of the coalescence process of plasmoids and studying the pinching process of the current sheet formed between plasmoids in detail.

We analyzed the \(\beta_0\) dependence of the fast-mode Mach number and the compression ratio, \(p_2/p_1\) (see Figure 7).

Figure 7. \(\beta_0\) dependence of the fast-mode Mach number and \(p_2/p_1\), where the gas pressures behind and in front of the shock are denoted by the subscripts 1 and 2, respectively.

Figure 8. (a) Schematic picture of the time evolution of the pinching process and the formation of MHD fast-mode shocks. (b) Schematic picture of our conjecture about an actual high Lundquist number reconnection process, where particles could be accelerated at fast-mode shocks generated by the coalescence of plasmoids.
This $\beta_0$ dependence of the compression ratio, $p_2/p_1$, is similar to Equation (43) with $\beta_0 < 10^{-2}$. In recent studies, it is indicated that particles can be accelerated in contracting plasmoids (Drake et al. 2006) and in plasmoids crossing fast-mode shocks (Nishizuka & Shibata 2013). Figure 8(a) shows an overview of the implosion process of the current sheet between plasmoids and the formation of the MHD fast-mode shocks. In our study, we found that fast-mode shocks can be formed during the implosion process in low-$\beta$ plasmas. Applying the shock formation in the implosion process to a coalescence process, we conjecture that particles in the plasmoids are efficiently accelerated at shocks generated by the coalescence of plasmoids through a Fermi acceleration process (see Figure 8(b)).

We investigated the parameter region where our scaling laws are valid. Since we assume the framework of non-relativistic ideal MHD, (1) the drift velocity of the electrons needs to be sufficiently smaller than the light speed and (2) the thickness of the current sheet needs to be larger than the ion skin depth. The electron drift velocity can be written as follows

$$J = nev_{\text{drift, e}} = \frac{c}{4\pi} \nabla \times B \sim \frac{c}{4\pi L} B,$$

where $L$ is the thickness of the current sheet. From Equation (45), $v_{\text{drift, e}}/c$, the drift velocity of electrons, is given as

$$v_{\text{drift, e}} \sim \frac{1}{4\pi e n L} B.$$

Applying the scaling laws of relations (36), (37), and (42) to Equation (45), we obtain

$$\frac{v_{\text{drift, e}}}{c} \sim 10^{-4} \left( \frac{\beta_0}{\beta_2} \right)^{-3/5} \times \left( \frac{B_0}{B_1} \right)^{-1} \left( \frac{n_0}{n_9} \right)^{-1} \left( \frac{L_0}{L_6} \right)^{-1},$$

where $\beta_2 = 10^{-2}$, $B_1 = 10$ G, $n_9 = 10^9$ cm$^{-3}$, and $L_6 = 10^6$ cm. Since we study non-relativistic MHD processes, $v_{\text{drift, e}}/c < 1$ is required. This leads to

$$\log_{10} \left( \frac{L_0}{L_6} \right) > - \log_{10} \left( \frac{n_0}{n_9} \right) - 3 \log_{10} \left( \frac{\beta_0}{\beta_2} \right) + \log_{10} \left( \frac{B_0}{B_1} \right) - 4.$$  (47)

Since the thickness of the current sheet must be larger than the ion skin depth, we similarly obtain

$$\frac{\lambda_i}{d} = \frac{1}{c} \sqrt{\frac{4\pi e^2}{m_i} \frac{\sqrt{\pi}}{d}} \sim 10^{-4} \left( \frac{\beta_0}{\beta_2} \right)^{-3/2} \left( \frac{n_0}{n_9} \right)^{1/2} \left( \frac{L_0}{L_6} \right)^{-1} < 1.$$  (48)
where $\lambda_i$ is the ion skin depth. We can transform Equation (48) as

$$
\log_{10} \left( \frac{L_0}{L_6} \right) > \frac{1}{2} \log_{10} \left( \frac{n_0}{n_9} \right) - \frac{3}{2} \log_{10} \left( \frac{\beta_0}{\beta_-} \right) - 4.
$$

In Figure 9, we show the parameter region using inequalities (47) and (49) with the values in the solar corona, $\beta_0 = 10^{-2}$ and $B_0 = 10$ G. It is found that our scaling laws are applicable in the solar corona. These scaling laws will help us estimate physical quantities in a current sheet formed in a low-$\beta$ plasma. Describing the evolution of the pinching process in an extremely low-$\beta$ plasma like the atmosphere of the neutron stars will be our future work.

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