Numerical Simulations of Solar Chromospheric Jets Associated with Emerging Flux

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Abstract

We studied the acceleration mechanisms of chromospheric jets associated with emerging flux using a two-dimensional magnetohydrodynamic (MHD) simulation. We found that slow-mode shock waves generated by magnetic reconnection in the chromosphere and the photosphere play key roles in the acceleration mechanisms of chromospheric jets. An important parameter is the height of magnetic reconnection. When magnetic reconnection takes place near the photosphere, the reconnection outflow collides with the region where the plasma beta is much larger than unity. Then, the plasma moves along a magnetic field. This motion generates a slow-mode wave. The slow-mode wave develops to a strong slow shock as it propagates upward. When the slow shock crosses the transition region, this region is lifted up. As a result, we obtain a chromospheric jet as the lifted transition region. When magnetic reconnection takes place in the upper chromosphere, the chromospheric plasma is accelerated due to the combination of the Lorentz force and the whip-like motion of the magnetic field. We found that the chromospheric plasma is further accelerated through the interaction between the transition region (steep density gradient) and a slow shock emanating from the reconnection point. In the process, the magnetic energy released by magnetic reconnection is efficiently converted into the kinetic energy of jets. This is an MHD effect that has not been discussed before.

Key words: Magnetic reconnection — Shock waves — Sun: chromosphere — Sun: photosphere

1. Introduction

Collimated ejections of chromospheric plasma into coronal heights have been studied for many years. Spicules are jet-like structures seen in the limb of the Sun (Sterling 2000; Anan et al. 2010). Surges are chromospheric jets seen dark in Hα images and their heights reach up to ~ 200 Mm (Roy 1973; Yoshimura et al. 2003). Recent observations have found tiny chromospheric jets with the loop structures at their foot-points (Shibata et al. 2007). They are called chromospheric anemone jets. Their typical height is 1–4 Mm (Nishizuka et al. 2011). Figure 1 shows an example of the jets observed by the Solar Optical Telescope (Tsuneta et al. 2008) on board the Hinode satellite (Kosugi et al. 2007).

Chromospheric jets are often observed in emerging flux regions and around sunspots. They are cool (~ 10^4 K) compared to coronal plasma (> 10^6 K). Many chromospheric jets disappear in the images of chromospheric lines before returning to the chromosphere. This may be due to some heating process, or rapid expansion of the chromospheric material of jets (the rapid expansion is shown in this paper). Therefore, for correct interpretations of observations we should understand how chromospheric plasma is ejected.

The triggering of chromospheric jets can lie at various heights in the solar atmosphere. X-ray observations found coronal hot jets (X-ray jets) associated with magnetic loops (Shibata et al. 1994), which implies that the energy release process takes place in the corona. Multi-wavelength coronal observations have shown that some cool jets (surges) coexist with hot jets (Schmieder et al. 1995; Canfield et al. 1996; Chae et al. 1999), which implies that the trigger is located in the lower corona, or in the upper chromosphere. It is known that some surges are associated with the Ellerman bombs or moustaches (Ellerman 1917; Rust 1968), where the Ellerman bombs are sudden energy releases in the low chromosphere or in the photosphere; their emissions are bright at both wings of the Hα line (Kitai 1983; Fang et al. 2006; Watanabe et al. 2008; Matsumoto et al. 2008). This implies that they are triggered in the low chromosphere or in the photosphere.

Using MHD simulations, Yokoyama and Shibata (1996) studied the acceleration mechanism of surges. They simulated the evolution of an emerging flux in a uniform coronal magnetic field. They obtained both hot and cool jets (which can correspond to the X-ray jets and surges, respectively) simultaneously as results of magnetic reconnection between the emerging flux and the pre-existing coronal field, where magnetic reconnection is a physical process in which a magnetic field in a highly conducting plasma changes its topology due to finite resistivity. Their results are consistent with multi-wavelength observations. They pointed out that the hot jets and the cool jets are accelerated mainly due to the gas-pressure gradient of the heated plasma and the Lorentz force, respectively. Note that chromospheric evaporation due to heat conduction is also important for the acceleration of X-ray jets (Shimojo & Shibata 2000). More elaborated modeling for coronal jets has been performed by other
Martínez-Sykora, Hansteen, and Moreno-Insertis (2011) proposed that a strong Lorentz force due to large magnetic field gradient squeezes chromospheric plasma, and produces spicule-like jets moving at the local sound speed. Moreno-Insertis, Galsgaard, and Ugarte-Urra (2008) compared a three-dimensional numerical experiment with an observed coronal jet, and showed that their numerical results are consistent with the observations.

When magnetic loops expand in the solar atmosphere, magnetic reconnection can take place at various heights. Now we should note that in the chromosphere the gravitational stratification is much stronger than in the corona. As a result, the plasma beta varies from much larger than unity to less than unity. Therefore, the physical processes associated with magnetic reconnection can be different at various heights in the solar atmosphere.

Can the lower atmospheric reconnection produce tall (∼1–10 Mm) jets? For example, observations have suggested that chromospheric anemone jets are considered to be results of magnetic reconnection in the chromosphere (Morita et al. 2010; Singh et al. 2011). Their speed is comparable to the expected local Alfvén speed (∼10 km s⁻¹) in the upper chromosphere. Suppose that magnetic reconnection takes place in the upper chromosphere. Then, it is highly possible that jets are reconnection outflows accelerated due to the whip-like motion of magnetic fields, considering that the outflow speed is almost the same as the local Alfvén speed. Next, suppose that magnetic reconnection occurred in the middle or low chromosphere where the plasma beta is expected to be larger than unity. Then, it is difficult to regard observed jets as being reconnection outflows, because the Alfvén velocity is lower than the observed velocity.

Let us estimate the maximum height where reconnection outflows can reach. Suppose that the kinetic energy is converted into potential energy. We then obtain \( pV_A^2/2 \sim \rho g H_{\text{jet}} \), where \( \rho \), \( V_A \), \( g \), and \( H_{\text{jet}} \) are the density, Alfvén velocity, gravitational acceleration, and height of the jet, respectively. We obtain \( H_{\text{jet}} \sim [p/(\rho g)][B^2/(8\pi p)] \sim H_{\beta}/\beta \), where \( p \), \( B \), \( H_{\beta} \), and \( \beta \) are the gas pressure, magnetic field strength, pressure scale height, and plasma beta in the chromosphere, respectively. Suppose that, for example, \( \beta \sim 1 \). Then, the height of the jet would be \( H \sim 200 \text{ km} \), which is much shorter than the typical height of observed jets (e.g., 1–4 \times 10^3 \text{ km} for the chromospheric anemone jets). Therefore, we need to consider another scenario in which only a fraction of plasma is accelerated.

A promising model is the acceleration mechanism through the interaction between a slow shock and the transition region (Osterbrock 1961; Shibata et al. 1982, 2007; Suematsu et al. 1982; De Pontieu et al. 2004; Heggland et al. 2007). The transition region can be launched when a slow shock passes through it. The ejected chromospheric plasma behind the contact discontinuous layer is regarded as being chromospheric jets. Shibata et al. (2007) pointed out that slow-mode waves can be generated by magnetic reconnection between small emerging flux and a pre-existing ambient field. They subsequently proposed that the slow-mode waves become shocks for some reason, and produce chromospheric jets. Magnetic
reconnection can generate waves of various kinds (Isobe et al. 2008; Nishizuka et al. 2008; He et al. 2009; Kigure et al. 2010). The wave can become shocks through non-linear processes. Carlsson and Stein (1997) compared observations with simulations, and found that slow shocks are ubiquitously generated in the chromosphere. As supported in the above studies, the slow-shock acceleration mechanism has the potential to systematically account for the observational characteristics of chromospheric jets.

The slow-shock acceleration mechanism has been studied by many authors. Suematsu et al. (1982) and Shibata et al. (1982) studied the mechanism using one-dimensional hydrodynamic simulations. Their results are both qualitatively and quantitatively consistent with the observational properties of spicules and Hz surges. They also proposed the possibility that the Ellerman bombs generate slow shocks, which lead to surges. Saito, Kudoh, and Shibata (2001) showed that the Alfvén waves also produce slow shocks due to a non-linear coupling. Although both slow and fast-mode waves are generated through a non-linear process of the Alfvén waves, their results suggest that slow-mode waves play more important roles in the generation of jets than the fast mode waves. Some authors discussed the slow-shock mechanism by including other effects (Sterling et al. 1993; Heggland et al. 2007; Martínez-Sykora et al. 2009; Murawski et al. 2011). Hansteen et al. (2006) and De Pontieu et al. (2007) discussed the relation between chromospheric jets and the slow-mode shocks generated by convective flows and global p-mode oscillations.

In the previous studies, waves were found to be generated in the chromosphere or in the photosphere, artificially (e.g., Shibata et al. 1982), by using observational data of the photospheric motions (e.g., Matsumoto et al. 2008), and convective flows in a self-consistent manner (e.g., Hansteen et al. 2006). Although Martínez-Sykora et al. (2009) suggested the possibility that a sudden energy release due to reconnection in the low atmosphere can lead to chromospheric jets, they did not identify the connection. The relation between the acceleration of chromospheric jets and magnetic reconnection is still unclear.

Even if magnetic reconnection takes place in the upper chromosphere, the interaction between slow shocks and the transition region can play an important role in accelerating chromospheric jets. Standing slow shocks will emanate from a reconnection point if the resistivity is localized in some way [the Petschek-type reconnection: Petschek (1964)]. For a detailed discussion on the possibility for localization of the resistivity, see subsection 4.6). The units of length, velocity, time, and density in the simulation are $H_0$, $C_{\infty}$, $H_0/C_{\infty}$, and $\rho_0$, respectively, where $H_0 = k_B T_0/(m g_0)$ is the pressure scale height ($k_B$ is the Boltzmann constant and $m$ is the mean molecular mass), $C_{\infty}$ the sound speed, and $\rho_0$ the density at the base of the photosphere ($z = 0$). The gas pressure, temperature, magnetic field strength, and energy are normalized by those units, i.e., $p_0 = \rho_0 C_{\infty}^2$, $T_0 = m C_{\infty}^2/(y k_B)$, $B_0 = (\rho_0 C_{\infty}^2)^{1/2}$, and $E_0 = \rho_0 C_{\infty}^2 H_0^3$, respectively. The gravity is normalized by $g_0 = C_{\infty}^2/(y H_0)$. The values of the normalization units are given as the typical values in the photosphere: $H_0 = 170 \text{ km}$, $C_{\infty} = 6.8 \text{ km s}^{-1}$, $\tau = H_0/C_{\infty} = 25 \text{s}$, and $\rho_0 = 1.4 \times 10^{-7} \text{ g cm}^{-3}$. Thus, we obtain particular, we focused on the relation between magnetic reconnection and chromospheric jets. Magnetic reconnection occurred between emerging magnetic loops and the pre-existing field in the chromosphere, and at the foot-points of the emerging flux. Slow shocks are generated by magnetic reconnection in the lower atmosphere (near the photosphere) and upper chromosphere. The slow shocks accelerate the chromospheric plasma along a magnetic field, producing chromospheric jets. The key process considered in this paper is the interaction between slow shocks and the transition region. An important point we found is that the reconnection height determines the acceleration mechanism of chromospheric jets. To investigate the interaction process in detail, we also performed numerous 1.5D simulations, as presented in the Appendix.

We note that the evolution of emerging flux in a 3D space is similar to that in a 2D space. For the emerging process, compare Shibata et al. (1989) (2D) and Matsumoto et al. (1993) (3D); for Ellerman bombs, compare Isobe, Tripathi, and Archontis (2007) (2D) and Archontis and Hood (2009) (3D). In 3D we will find a wider variety of physics than in 2D, including the interchange instability (Isobe et al. 2006) and magnetic reconnection between not perfectly anti-parallel field lines (Jiang et al. 2011; Nakamura et al. 2012). Here, we focus on fundamental physics in a 2D space.

In section 2, we introduce the numerical model adopted here. In section 3, we present results of our 2D numerical simulation. Here, we separately describe the acceleration processes of two jets, Jet A and B, to clarify that the height of magnetic reconnection determines the acceleration mechanism. To investigate the acceleration process of Jet B in detail, we performed numerous 1.5D simulations, which are summarized in the Appendix. In section 4 we present a schematic diagram summarizing the acceleration mechanisms of chromospheric jets. Here, we also give implications for observations, and discuss the validity of the adopted assumptions.

2. Numerical Setup

2.1. Basic Equations and Assumptions

We here describe how to solve the two-dimensional, resistive, and compressible MHD equations in Cartesian coordinates $(x, z)$, where the $z$-direction is antiparallel to the gravitational acceleration. We assume that the medium is an inviscid, perfect gas with a ratio of specific heats of $\gamma = 5/3$. We include the radiative cooling effect, as described below, but neglect the heat conduction (for a discussion on this validity, see subsection 4.6). The units of length, velocity, time, and density in the simulation are $H_0$, $C_{\infty}$, $H_0/C_{\infty}$, and $\rho_0$, respectively, where $H_0 = k_B T_0/(m g_0)$ is the pressure scale height ($k_B$ is the Boltzmann constant and $m$ is the mean molecular mass), $C_{\infty}$ the sound speed, and $\rho_0$ the density at the base of the photosphere ($z = 0$). The gas pressure, temperature, magnetic field strength, and energy are normalized by those units, i.e., $p_0 = \rho_0 C_{\infty}^2$, $T_0 = m C_{\infty}^2/(y k_B)$, $B_0 = (\rho_0 C_{\infty}^2)^{1/2}$, and $E_0 = \rho_0 C_{\infty}^2 H_0^3$, respectively. The gravity is normalized by $g_0 = C_{\infty}^2/(y H_0)$. The values of the normalization units are given as the typical values in the photosphere: $H_0 = 170 \text{ km}$, $C_{\infty} = 6.8 \text{ km s}^{-1}$, $\tau = H_0/C_{\infty} = 25 \text{s}$, and $\rho_0 = 1.4 \times 10^{-7} \text{ g cm}^{-3}$. Thus, we obtain
\( p_0 = 6.3 \times 10^4 \text{ dyn cm}^{-2}, \ T_0 = 5600 \text{ K}, \ B_0 = 250 \text{ G}, \) and \( E_0 = 3.1 \times 10^{26} \text{ erg}. \)

The basic equations are as follows:

\[ \frac{\partial \rho}{\partial t} + \rho (v \cdot \nabla) v = -\rho \nabla \cdot v, \tag{1} \]

\[ \frac{\partial v}{\partial t} + (v \cdot \nabla) v = -\frac{1}{\rho} \nabla p + \frac{1}{4\pi \rho} (\nabla \times B) \times B + g. \tag{2} \]

\[ \frac{\partial T}{\partial t} + (v \cdot \nabla) T = -(\gamma - 1) T \nabla \cdot v - \frac{T}{\tau_{\text{cooling}}}. \tag{3} \]

\[ p = \frac{k_B}{m} T, \tag{4} \]

\[ \frac{\partial B}{\partial t} - c \nabla \times E = 0, \tag{5} \]

\[ E = \eta J - \frac{v}{c} \times B, \tag{6} \]

\[ J = \frac{c}{4\pi} \nabla \times B. \tag{7} \]

Here, \( g = (0, -g_0) \) is the gravitational acceleration, \( \eta \) the resistivity, and \( J \) the current density. \( \tau_{\text{cooling}} \) represents the radiative cooling time, given as a function of \( z \). We include the radiative cooling effect in a simple manner: we assume that the radiative cooling time varies linearly in the photosphere and is only applied to those regions where the pressure fluctuation is positive compared to these initial pressure (therefore we do not apply the radiative cooling to the expanded regions). We do not consider radiative cooling in the convectively unstable zone \( (z < 0) \).

\[ \tau_{\text{cooling}}(z) = \begin{cases} \infty & (z < 0) \\
(\tau_{c1} - \tau_{c0}) z / z_{tr} + \tau_{c0} & (0 \leq z \leq z_{tr}) \\
\tau_{c1} & (z > z_{tr}) \end{cases} \tag{8} \]

where \( \tau_{c0} (= 0.04 t_0 = 1 \text{ s}) \) and \( \tau_{c1} (= 40 t_0 = 1000 \text{ s}) \) are the radiative cooling times at the bottom and top of the chromosphere, respectively. Here, we roughly mimic the average cooling time given in the VAL-C model (Vernazza et al. 1981); \( z_{tr} \) is the height of the transition region. We also assume the anomalous resistivity model in the simulation box:

\[ \eta = \begin{cases} 0 & \text{for } v_{id} < v_{c} \\
\alpha (v_{id} / v_{c} - 1)^2 & \text{for } v_{id} \geq v_{c} \end{cases} \tag{9} \]

\[ v_{id} = |J| / \rho, \tag{10} \]

\[ v_{c} = v_{c1} + (v_{c2} - v_{c1}) \left\{ \frac{1}{2} \left[ \tanh \left( \frac{z - z_{tr}}{w_{tr}} \right) + 1 \right] \right\}, \tag{11} \]

where \( \alpha (= 0.01 H_0^2 / \tau) \), \( v_{id} \), and \( v_{c} \) are a constant, the (relative ion–electron) drift velocity, and the threshold above which the anomalous resistivity sets in, respectively; \( v_{c1} (= 50 C_{ad}) \) and \( v_{c2} (= 3000 C_{ad}) \) are the threshold in the chromosphere and the corona, respectively, and \( w_{tr} = 0.5 H_0 \).

### 2.2. Initial and Boundary Conditions

We simply model the solar atmosphere: a superadiabatically stratified layer representing the upper convection zone \((Z_{min} \leq z < 0)\), an isothermal cool \((T = T_0)\) layer representing the photosphere/chromosphere \((0 \leq z \leq z_{tr})\), and an isothermal hot \((T = 150 T_0)\) layer representing the corona \((z > z_{tr})\). We take \( z = 0 \) to be the base height of the photosphere.

The initial distribution of temperature is given as

\[ T(z) = T_{\text{pho}} - \left( a \left| \frac{dT}{dz} \right|_{ad} \right) z \quad (Z_{min} \leq z \leq 0) \quad (12) \]

for the convectively unstable zone, and

\[ T(z) = T_{\text{pho}} + (T_{\text{cor}} - T_{\text{pho}}) \left\{ \frac{1}{2} \left[ \tanh \left( \frac{z - z_{tr}}{w_{ad}} \right) + 1 \right] \right\} \quad (z \geq 0) \quad (13) \]

for the upper atmosphere, where \( T_{\text{pho}} \) and \( T_{\text{cor}} \) are the temperatures in the photosphere/chromosphere and the corona, respectively. In calculations we took \( z_{tr} = 10 H_0 \) (1700 km), \( w_{tr} = 0.2 H_0 \) (34 km). \( |dT/dz|_{ad} \equiv (\gamma - 1) / \gamma \) is the adiabatic temperature gradient and \( a \) is a dimensionless constant. The layer is convectively unstable in the case that \( a > 1 \). We used \( a = 1.5 \) to model the convection zone. The two isothermal layers are joined through the transition region with steep temperature gradient, whose width is \( \sim 0.4 H_0 \). As a result, the transition region physically corresponds to a density discontinuous layer.

For the initial magnetic field, we put a magnetic flux sheet in the convectively unstable zone and uniform background field. The former represents the flux sheet that produces emerging flux loops, and the latter represents a pre-existing magnetic field in the chromosphere and the corona. A schematic illustration of the model is shown in figure 2.

The magnetic flux sheet in the convectively unstable zone is given as

\[ B_{\text{fs}}(z) = \left[ \frac{8\pi p(z)}{\beta(z)} \right]^{1/2} \hat{x}, \tag{14} \]

where

\[ \beta(z) = \frac{B_{\text{fs}}}{f(z)} \tag{15} \]

and
The size of the simulation box was \(-70H_0 \leq x \leq 70H_0\) and \(-10H_0 \leq z \leq 140H_0\), namely 23800 km \(\times\) 25500 km, which is large enough to include the observed chromospheric jets (\(~1000\) km). The grid spacing in the horizontal direction was taken to be uniform with \(\Delta x = 0.1H_0\) in \(-40H_0 \leq x \leq 40H_0\), and increases up to \(1.5H_0\) in \(x < -40H_0\) and \(40H_0 < x\). In the vertical direction \(\Delta z = 0.05H_0\) in \(z \leq 30H_0\), and increases up to \(1.0H_0\) in \(z > 30H_0\). The total grid number was \(N_x \times N_y = 828 \times 906\).

The numerical scheme that we adopted is based on the CIMP-MOCCT method (Kudoh et al. 1999). The scheme could capture the profiles of physical parameters with contact discontinuities. Therefore, using this method we could investigate the dynamics of the transition region in detail. We used some of the routines contained in CANS (Coordinated Astronomical Numerical Software) maintained by Yokoyama et al.

### 3. Numerical Results

#### 3.1. Evolution of Emerging Flux

Magnetic loops expand as a result of the convective Parker instability (Parker 1966; Nozawa et al. 1992). The time evolution of the magnetic loops is shown in figure 4. In figure 4, two jets are indicated by white arrows (Jet A and B). We focus on these jets, and consider how magnetic reconnection creates these jets.

#### 3.2. Jet Resulting from Magnetic Reconnection near Photosphere: Jet A

The foot-points of the emerging loops sink into the convective zone because the plasma in the lifted flux sheets falls down along the magnetic field and the foot-points become heavy. The sinking motion naturally creates an antiparallel field (U-shape structure) and a current sheet. This behavior is also found in the models presented by Isobe, Tripathi, and Archontis (2007) and Archontis and Hood (2009). As the result of the current sheet formation, magnetic reconnection takes place just below the photosphere \((z \sim -4H_0)\) (see figure 5). In our simulation, the numerical resolution is insufficient to resolve the reconnection region, or the current sheet. Although we cannot capture detailed structures of the reconnection region, we infer that the simulation describes the global dynamics, like the acceleration of reconnection outflow. The reconnection rate of the numerical reconnection is \(\sim 0.1\). If fast reconnection takes place near the actual photosphere, the reconnection outflow will generate slow-mode waves, as found in this simulation.

The reconnection outflow is drastically slowed down by the dense and high-beta photospheric plasma within a short travel distance. The outflow moves upward along a magnetic field due to both the inertia effect and squeezing. The upward flow collides with the downward flow along the field. As a result, a compressed region is formed, leading to the generation of a slow-mode wave. This behavior is also found in the model presented by Takeuchi and Shibata (2001).

The slow-mode wave shows wave amplification when it propagates upward through the gravitationally stratified atmosphere as a result of the law of conservation of energy. Eventually, the slow-mode wave becomes a slow shock. This is analyzed in detail in section 4. The slow shock produces...
a chromospheric jet through the interaction with the transition region. Note that to guide the slow-mode wave generated by magnetic reconnection in the lower atmosphere, an ambient field plays an important role. The chromospheric material of the jet is expanded (see div \( V \) distribution of figure 5) and the density is decreased. This may be important for the disappearance of jets in images of chromospheric lines.

We obtained the physical parameters along the field line indicated in figure 6. Figure 7 shows the time evolution of the velocity parallel to the magnetic field line, the density, and the pressure measured along the field line.

The growth of the shock, i.e., the growth of the amplitude of the velocity parallel to the field line, can be seen in figure 7b. The velocity of the jet is \( \sim 6C_{s0} \sim 40\text{ km s}^{-1} \). This is much larger than the local sound speed. Note that the upward velocity of the reconnection outflow in the photosphere is \( 0.6C_{s0} \sim 4\text{ km s}^{-1} \). The maximum height of the jet is \( \sim 25H_0 \sim 4300\text{ km} \) from the photosphere. The difference between the maximum height and the initial height is \( \sim 17H_0 \sim 2900\text{ km} \).

This example shows that even if magnetic reconnection takes place below the photosphere, tall (\( \sim 4\text{ Mm} \)) jets can be produced through the interaction between a slow shock

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**Fig. 4.** Time evolution of the density (left) and the temperature (right). The black solid lines show the magnetic field lines. Black arrows show the velocity vectors. \( \rho_0 = 1.4 \times 10^{-7}\text{ g cm}^{-3}, T_0 = 5600\text{ K}, C_{s0} = 6.8\text{ km s}^{-1}, \) and \( \tau = 25\text{ s} \). (Color online)

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**Fig. 5.** Jet A: Time evolution of the density (left), the temperature (middle), and the divergence of the velocity field (div \( V \); right). The black solid lines show the magnetic field lines. The arrows show the velocity vectors. The green solid lines show the transition region. Note that in the images of div \( V \) the blue and red regions correspond to the compressed and expanded regions, respectively. \( \rho_0 = 1.4 \times 10^{-7}\text{ g cm}^{-3}, T_0 = 5600\text{ K}, C_{s0} = 6.8\text{ km s}^{-1}, \) and \( \tau = 25\text{ s} \). An animation is available at [LINK1](http://www.kwasan.kyoto-u.ac.jp/~takasao/movie/chrom_jet/recon_foot.gif). (Color online)
and the transition region. It should be noted that the jet is not the reconnection outflow, but the transition region lifted by the slow shock.

3.3. Jet Resulting from Magnetic Reconnection in Upper Chromosphere: Jet B

Magnetic loops push an ambient field, and start to reconnect with it as magnetic loops expand. Chromospheric material is accelerated after magnetic reconnection. The accelerated chromospheric plasma eventually start to rise upward along the magnetic field, becoming a chromospheric jet. No hot jet driven by the pressure gradient, as in Yokoyama and Shibata (1996), coexists with the chromospheric jet.

Figure 8 shows the time evolution of the density, temperature, and divergence of the velocity field of the emerging flux region. The density distributions show a shape similar to that of the observed jets, i.e., a jet with a loop structure (see figure 1).

We assumed the anomalous resistivity model in the simulation (the justification of the anomalous resistivity model will be discussed in section 4). As a result, the Petschek-type reconnection, in which the standing slow shocks emanate from the magnetic reconnection point, occurs between the emerging flux and the pre-existing ambient field. In the following, we call the standing slow shocks emanating from the reconnection point the Petschek slow shocks.

Figure 9 shows the time evolution of the resistive electric current $|\eta J_z|$ at the neutral point as a measure of the reconnection rate, where the reconnection rate is defined as the reconnected magnetic flux per unit time. This is measured in the region $-10H_0 < x < 5H_0$ and $H_0 < z < 10H_0$. As shown, magnetic reconnection starts at $t \sim 46\tau$.

Many shocks are formed in the chromosphere. The blue and red regions in the div $V$ distribution shown in figure 8 correspond to the compressed and expanded regions, respectively. The shock structures can be discerned as sharp blue regions in figure 8 (blue regions emanating from the reconnection point are the Petschek slow shocks). The enlarged images before and after one of the Petschek slow shocks crosses the transition region are shown in figure 10.

We obtained the physical parameters along the field line shown in figure 11. The black thick field lines are the same field line. Figure 12 shows the evolution of the velocity parallel to the magnetic field line, the density, and the pressure measured along the field line. The enhanced density, and the pressure in
Fig. 7. Jet A: (a) Time evolution of the velocity parallel to the field line, the density, and the pressure. These data were obtained along the field line shown in figure 6. (b) Close-up images of the period between $t = 47t$ and $51t$. SS and TR stand for the slow shock and the transition region, respectively.
the outflow region appear as the peaks at a distance of about $10H_0$ in figure 12 (positions where the distance is $10H_0$ are marked with the black points on the field line in figure 11). We can see that the height of the transition region (i.e., the steep gradient in the density profile at $15H_0$ distance from the start point of the line at $t = 59\tau$) decreases with time during the period between $t = 59\tau$ and $60\tau$. This is because the reconnection inflow reduces the height of the transition region (see figures 8 and 10). Eventually, at $t \sim 60\tau$ the transition region crosses one of the Petschek slow shocks attached to the reconnection point.

After reconnection starts, the inflow and outflow regions are separated by stationary Petschek slow shocks. We express such a situation in words “the slow shock is attached to the outflow region.” When one of the Petschek slow shocks crosses the transition region, it starts to propagate upward. Then, the shock is no longer attached to the chromospheric material accelerated due to magnetic tension (see the bottom panels of figure 10) (the propagating speed of a slow shock in the low-beta region is almost the same as the local sound speed, which is much larger in the corona than that in the chromosphere). In div $V$ distributions at $t = 59.8\tau$ and later times shown in figure 8, we can see the slow shock (blue region) propagating upward. After the interaction, we can find that the steep gradient in the magnetic field strength distribution is smoothed out.

After the interaction between the Petschek slow shock and the transition region, the reconnection outflow, which is almost in the horizontal direction, is accelerated, although the transition region is not lifted upward yet (see figure 12b). Considering that the outflow speed is comparable to the local

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**Fig. 8.** Jet B: Time evolution of the density (left), the temperature (middle), and the divergence of the velocity field (div $V$; right). The black solid lines show the magnetic field lines. The arrows show the velocity vectors. The green solid lines show the transition region. Note that in the images of div $V$ the blue and the red regions correspond to the compressed and the expanded regions, respectively. $\rho_0 = 1.4 \times 10^{-7} \text{ g cm}^{-3}$, $T_0 = 5600 \text{ K}$, $C_s = 6.8 \text{ km s}^{-1}$, and $\tau = 25 \text{ s}$. An animation is available at LINK2.² (Color online)
Alfvén speed, it is worth showing the Alfvén speed distribution of the reconnection region. Figure 13 shows that the Alfvén speed in the outflow region becomes large after the interaction. In the region where the Alfvén speed increases, the outflow velocity has also increased.

To investigate the interaction process in detail, we performed 1.5D MHD simulations, and analyzed the interaction between one of Petschek slow shocks and the transition region. We summarize the simulation results in the Appendix. We found that the outflow is further accelerated due to magnetic tension after the interaction between the Petschek slow shock and the transition region. When the slow shock crosses the transition region, a slow-mode rarefaction wave is generated, and then propagates into the outflow region. The slow-mode rarefaction wave causes the density and the magnetic field strength to become low and high, respectively. Thus, the Alfvén speed becomes large (see figure 13).

When magnetic reconnection takes place, a magnetic field tends to be relaxed, and re-configured to a relatively straight structure. Reconnection outflows are accelerated as a consequence of the relaxation process. The Alfvén velocity near the outflow region determines the outflow speeds. Therefore,
the relaxation process will proceed quickly when the Alfvén velocity increases. Therefore, the high Alfvén speed leads to quick relaxation of the configuration of the magnetic field. The re-configuration process results in further acceleration of the reconnection outflow.

How is the chromospheric plasma accelerated upward due to magnetic reconnection? We noted that the upward motion of the transition region starts at around \( t = 62 \). To clarify how the jet is accelerated, we show the forces acting on the chromospheric plasma just below the transition region in figure 14. The plasma in the outflow region feels no strong upward acceleration during \( t = 60 \) and \( 62 \). The \( z \)-component of the velocity shows the gradual acceleration during this period. Note that the chromospheric plasma in the outflow feels no strong acceleration due to the pressure gradient force when the outflow collides with the ambient field. The chromospheric plasma (reconnection outflow), which is almost horizontally accelerated, eventually starts to rise upward along the magnetic field after colliding with the ambient field.

A schematic diagram of the acceleration process is shown in figure 15. After the outflow collides with the ambient field, the outflow plasma moves in the direction along the ambient field. Consequently, we obtain the chromospheric jet. This process, which we call the “whip-like acceleration,” is essentially the centrifugal force acceleration.

The upward velocity \( (V_z) \) of the chromospheric jet is
Fig. 12. Jet A: (a) Time evolution of the velocity parallel to the field line, the density, and the pressure. These data were obtained along the field line shown in figure 11. (b) Close-up images of the period between $t = 59\tau$ and $61\tau$. SS and TR stand for the slow shock and the transition region, respectively.
\[ t = 58.8 \tau \]
\[ t = 60.2 \tau \]
\[ t = 62.0 \tau \]

\[ C_{s0} = 6.8 \text{ km s}^{-1} \] and \[ \tau = 25 \text{ s} \].

We can discern the transition region corrugated at the top of the jet (see figures 8 and 13). Stone and Edelman (1995) found that parallel slow shocks are unconditionally unstable to the corrugation instability. In our simulation, the slow shock crosses the transition region and propagates almost parallel to the magnetic field. Therefore, the corrugated structure can be a result of the corrugation instability.

4. Discussion

We studied the acceleration mechanisms of chromospheric jets associated with emerging flux using a 2D MHD simulation. It is found that the height of magnetic reconnection determines the acceleration mechanism of chromospheric jets. Here, we present a schematic diagram summarizing the acceleration mechanisms of chromospheric jets. We also give implications for observations, and discuss the validity of the adopted assumptions.

4.1. Comparison between One-Dimensional Case and Two-Dimensional Case: Jet A

We found that a slow-mode wave grows to be a shock in the case that magnetic reconnection occurred just below the photosphere. The growth of shocks in one-dimensional cases has been studied by many authors (Osterbrock 1961; Shibata et al. 1982; Heggland et al. 2007). We consider how our 2D simulation results differ from the prediction of the idealized 1D theory in which a flux tube is rigid enough not to be disturbed by the plasma dynamics in it and the total energy is exactly conserved in a flux tube (so-called zero-beta approximation).

Here, we review the 1D theory. We consider the situation in which a slow-mode wave propagates upward in a rigid flux tube in a gravitationally stratified atmosphere. In the linear regime, we obtain the following relation from the energy conservation of the slow mode wave:

\[ A \nu_k^2 C_s = \text{constant}, \]  
\[ (20) \]

where \( A, \nu_k, \) and \( C_s \) are the cross-section of the flux tube, the velocity parallel to the magnetic field, and the sound speed, respectively. Considering the magnetic flux conservation,
BA = constant, and that \( C_s \) is almost constant in the chromosphere, we obtain
\[
v || B^{-0.5} \propto \rho^{-0.5},
\]  
where \( B \) is the magnetic field strength. In the non-linear regime, this relation is modified due to dissipation processes. Ono, Sakashita, and Yamazaki (1960) derived the relation in a strong-shock case,
\[
v || B^{-0.5} \propto \rho^{-0.236}.
\]  

Note that these relations will hold if the energy is exactly conserved in a flux tube.

We investigated the dependence of the growth rate \( v_0 B^{-0.5} \) on the density, \( \rho \), using our simulation results. We show the result in figure 16a [this figure corresponds to figure 6 in Shibata and Suematsu (1982)]. In figure 16a, the triangles were obtained from the numerical simulation. The dependence of the growth rate on the density changes approximately at \( -\log \rho = 2 \). We separately applied the least-squares method to the numerical results given below and above this density value to obtain the fitting lines. These lines correspond to the two solid lines. These lines have slopes of 0.36 and 0.15, respectively. The dotted and dashed lines correspond to the analytic relations (21) and (22), respectively.
Below $-\log \rho \sim 2$, the dependence of the growth rate on the density slightly deviates from the prediction of the idealized 1D theory (21). This is probably because the plasma beta is greater than unity in the low chromosphere, and the low-beta approximation is broken. This means that the energy of the shock wave in a flux tube leaks to other flux tubes. Above $-\log \rho \sim 2$, the discrepancy between the numerically obtained relation and the 1D analytic relation (22) becomes large. The energy loss of the wave along a flux tube is larger than that predicted with the aid of the one-dimensional analysis. This implies that the one-dimensional approximation, i.e., the low-beta approximation is broken in the non-linear regime.

To clarify the cause of the discrepancy, we investigated whether the assumption of the idealized 1D theory is valid. We plot the lines on which the plasma beta is unity in figure 16b. It should be noted that the plasma beta behind the shock is close to unity, even in the upper chromosphere. Therefore, the 1D approximation (zero-beta approximation) is severely broken in the regime, and the energy of the slow shock is carried to other flux tubes by fast-mode waves. If we consider the energy conversion rate from the magnetic energy liberated by magnetic reconnection into the kinetic energy of jets, we need to pay an attention to this effect.

4.2. Acceleration of Reconnection Outflow: Jet B

The reconnection outflow discussed in subsection 3.3 is basically accelerated due to the whip-like motion of the reconnected magnetic field. The reconnection outflow is accelerated after the interaction between one of the Petschek slow shocks and the transition region. Here, we focus on this point, which has not been addressed in previous studies.

We obtain the slow mode rarefaction wave propagating into the outflow region after the Petschek slow shock crosses the transition region. Behind the slow mode rarefaction wave, the magnetic field strength increases and the density decreases (see figure 12). Thus, the Alfvén speed increases, as shown in figure 13. The large Alfvén speed leads to a rapid re-configuration of the magnetic field. The re-configuration results in a further acceleration of the outflow plasma. For a detailed discussion, see the Appendix.

In the process, the magnetic energy released by magnetic reconnection is efficiently converted into the kinetic energy of jets. After magnetic reconnection takes place, the magnetic energy is converted into the kinetic (acceleration by the Lorentz force) and thermal (heating by the slow shocks) energies of the outflow plasma. When one of the slow shocks crosses the transition region, the slow shock accelerates the plasma along a magnetic field, i.e., the thermal energy is converted into the kinetic energy of the slow shock. Thus, a large fraction of the magnetic energy is converted into the kinetic energy of jets.

4.3. A Unified Scenario of Solar Chromospheric Jets

In the case of the jet associated with magnetic reconnection near the photosphere (Jet A), the slow-mode wave generated as a result of the collision of the reconnection outflow is amplified when it propagates upward. The slow-mode wave becomes a shock as a result of the law of conservation of energy. Through the interaction between the slow shock and
the transition region, a chromospheric jet is produced. This is essentially the same process as that discussed by Shibata et al. (1982), that is, a pure hydrodynamic process.

In the case of the jet associated with magnetic reconnection in the upper chromosphere (Jet B), the chromospheric plasma (reconnection outflow) is accelerated due to the whip-like motion of the reconnected magnetic field. The outflow is further accelerated through the interaction between the Petschek slow shock and the transition region. This is an MHD effect that cannot be found in pure hydrodynamic cases (see subsection 3.3 and the Appendix).

We self-consistently related the acceleration mechanisms of chromospheric jets with magnetic reconnection under the assumptions adopted in this study. The height of a reconnection point is found to determine the acceleration scenario. The acceleration processes are summarized in figure 17. We include the model found by Yokoyama and Shibata (1996) to unify the acceleration processes and the speculated scenario for the low and middle chromospheric cases.

4.4. Implication for Observations

4.4.1. Implication for recurrent acceleration of jets

Observations found the recurrent acceleration of chromospheric anemone jets (Nishizuka et al. 2011). Here we discuss a possible recurrent acceleration scenario in terms of the process found in this study.

We showed that slow-mode waves can be generated when reconnection outflows collide with an ambient field. Let us assume that the velocity of the reconnection outflows that hit the ambient field is time-dependent. This is possible, e.g., if the plasmoids are formed in the current sheet, where plasmoids are coherent structures of plasma and magnetic fields. The emergence of the plasmoids has been theoretically and observationally implied (Singh et al. 2011; Leake et al. 2012;
When the outflow hits the ambient magnetic field, slow-mode waves will be generated by squeezing. This process is essentially the same as a mode conversion process from fast-mode waves into slow-mode waves (e.g., Bogdan et al. 2003). The slow-mode waves (or slow shocks) will propagate upward and accelerate jets when passing through the transition region.

Yokoyama and Shibata (1996) discussed direct acceleration due to the pressure enhancement due to the squeezing of the chromospheric plasma by reconnection outflows. Here, we note that slow-mode waves can be generated as a result of the collision of reconnection outflows, as in Takeuchi and Shibata (2001). This process will be preferable for the recurrent acceleration of chromospheric jets. A schematic diagram is shown in figure 18.

4.4.2. Relation between Ellerman bombs and surges
Observations have suggested that the Ellerman bombs result from the emergence of undulatory flux tubes (Pariat et al. 2004). Watanabe et al. (2011) observationally found morphological evidence for the magnetic reconnection in the deep photosphere. They also found that some Ellerman bombs were accompanied by dark Hα surges. Following their study, the apparent velocity of the ejections in the photosphere is \( \sim 10 \text{km s}^{-1} \), and the Doppler velocity of the surge is \( \sim 40 \text{km s}^{-1} \).

We obtained the tall (\( \sim 4 \text{Mm} \)) jet resulting from the reconnection just below the photosphere. The reconnection outflow velocity in the photosphere is \( 0.6C_{\|} \sim 4 \text{km s}^{-1} \). The velocity of the jet is \( 6C_{\|} \sim 40 \text{km s}^{-1} \), which is much larger than the reconnection outflow velocity. These values are consistent with the observations. The shock acceleration scenario can account for the difference between the velocity in the photosphere and the velocity of chromospheric jets.

4.4.3. Ejections from quiescent prominences
We found that the reconnection outflow speed increases through the interaction between one of the Petschek slow shocks and the steep density jump layer. Recent observations have implied that magnetic reconnection takes place in quiescent prominences, where the quiescent prominences are large structure composed of relatively cool (\( \sim 10^4 \text{K} \)) plasma in the corona of the quiet-sun regions (Hillier et al. 2011). There are density jump surfaces between the quiescent prominence and the corona. Therefore, if the Petschek-type reconnection occurs in quiescent prominences, the reconnection outflow will be accelerated through the interaction.

4.5. Justification of Anomalous Resistivity Model in Chromosphere
We adopt an anomalous resistivity model not only in the corona, but also in the chromosphere. The resistivity model is widely used to understand the rapid reconnection processes, namely that the reconnection rate is close to unity, in a fully ionized and collision-less plasma like the corona (e.g., Yokoyama & Shibata 1996). Using the resistivity model, many observational characteristics of explosive phenomena, like solar flares and coronal jets, have been successfully accounted for (for review, see Shibata & Magara 2011). If the anomalous resistivity model is adopted, the resistivity is spatially localized, which leads to rapid Petschek-type reconnection.

Although there is no clear evidence that the Petschek slow shocks are indeed formed during reconnection processes, we utilize the resistivity model to allow fast reconnection to occur. We should note that the chromosphere is fully collisional and partially ionized. For this reason, the characteristics of the resistivity in the chromosphere can be different from that in the corona, which makes it difficult for us to predict the reconnection physics there. Therefore, before Hinode era we had not known whether rapid reconnection processes would ubiquitously occur in the chromosphere. Recently, Hinode observed many rapid events, including the chromospheric anemone jets, which implies that rapid reconnection processes are taking place also in the chromosphere.

It is known that the collision between electron, ion, and neutral can play an important role in magnetic reconnection (Zweibel & Brandenburg 1997). Neutrals mainly contribute to the gas pressure in a current sheet. The collision between ions and neutrals leads to the diffusion of the neutrals from the current sheet. This diffusion, i.e., the ambipolar diffusion, can induce the current sheet to become thin. When the thickness of the current sheet becomes comparable to a microscopic scale, magnetic reconnection will set in. Recently, H. Isobe (in preparation) showed that even with uniform resistivity, a localized (or non-uniform) neutral distribution in the current sheet induces a local thinning of the current sheet, leading to the Petschek-like reconnection. This scenario may account for fast reconnection processes in the chromosphere. Because there is a possibility that the Petschek-type reconnection takes place in the chromosphere, we adopted the anomalous resistivity model to obtain it. We believe that we can mimic reconnection processes in the partially ionized and fully collisional plasma using this resistivity model.

4.6. Effect of Heat Conduction
Generally speaking, the heat conduction weakens shocks because the heat energy liberated at the shock fronts can escape in front of the shock surfaces. We do not include the effect of heat conduction in this paper. We consider whether or not the heat conduction significantly affects the simulation results.

The timescale of the heat conduction can be estimated as

\[
\tau_{\text{cond}} \sim 3 \times 10^{-10} n_e l^2 T^{-5/2} \text{s},
\]

where \( n_e, l, T \) are the electron density, typical length scale, and the temperature, respectively. The heat conduction is important at a shock front because the temperature gradient is steep. In our simulation, the minimum length scale was the smallest grid size. By taking \( l = 8.5 \text{km} \) (grid size), \( n_e = 10^{12} \text{cm}^{-3} \), and \( T = 10^4 \text{K} \), we estimated the heat conduction time, \( \tau_{\text{cond}} \), at \( \sim 2 \times 10^4 \text{s} \). Considering that the propagation time of shocks in the chromosphere was less than 100 s in our simulation, we can neglect the effect of the heat conduction on the shock propagation in the chromosphere.

Heat due to the thermal conduction and cooling due to the radiative loss control the location of the transition region. In the slow-shock acceleration scenario, the maximum height of a chromospheric jet is a function of the height of the transition region before a slow shock comes (Shibata et al. 1982). Therefore, the maximum height of jets will be affected if we include the thermal conduction. However, the acceleration...
scenario in which slow shocks lift the transition region will not be changed. The heating and cooling processes will affect the coronal EUV emission. Therefore, one should include the heat conduction as well as cooling terms into the energy equation if one tries to directly compare the results of numerical simulations with observations like those of De Pontieu et al. (2011) (see e.g., Heggland et al. 2009).

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Appendix. 1.5 Dimensional Study on Interaction between Slow Shock and Transition Region

The process described in subsection 3.3 can be considered to be a magnetic reconnection problem with a density contact discontinuous layer (transition region) in the inflow region (see figure 19). We investigated the interaction process between the Petschek slow shock and the transition region.

We simulated the time evolution of the reconnected magnetic field using 1.5D MHD simulations. Gravity was neglected, and therefore the results given here are scale-free. The basic equations are as follows:

\[ \frac{\partial \rho}{\partial t} + \rho \frac{\partial v_x}{\partial x} = -\frac{\partial p}{\partial x}, \]  
(A1)

\[ \frac{\partial v_x}{\partial t} + v_x \frac{\partial v_x}{\partial x} = -\frac{1}{\rho} \frac{\partial p}{\partial x} - \frac{1}{4\pi \rho} B_y \frac{\partial B_y}{\partial x}, \]  
(A2)

\[ \frac{\partial v_y}{\partial t} + v_x \frac{\partial v_y}{\partial x} = \frac{1}{4\pi \rho} B_x \frac{\partial B_x}{\partial x}, \]  
(A3)

\[ \frac{\partial T}{\partial t} + v_x \frac{\partial T}{\partial x} = -(\gamma - 1) T \frac{\partial v_x}{\partial x}. \]  
(A4)

\[ p = \frac{k_B}{m} \rho T. \]  
(A5)

\[ \frac{\partial B_y}{\partial t} - c \frac{\partial E_z}{\partial x} = 0, \]  
(A6)

\[ E_z = -\frac{v_x}{c} B_y + \frac{v_y}{c} B_x. \]  
(A7)

A.2. Initial and Boundary Conditions

Figure 20 shows the initial condition of the simulation in the case that \( \rho_{\text{CH}} / \rho_{\text{CR}} = 100 \), where \( \rho_{\text{CH}} \) and \( \rho_{\text{CR}} \) are the density in the chromosphere and the density in the corona, respectively.
Fig. 21. Time evolution of the velocity in the $x$-direction, the velocity in the $y$-direction, the gas pressure, the density, the magnetic field strength, and the Alfvén speed. TR, SS, and SR denote the transition region, the slow shock, and the slow-mode rarefaction wave, respectively.

Table 1. Parameters.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$p_{CH}/p_{CR}$</td>
<td>10, 100, 1000</td>
</tr>
<tr>
<td>$\beta$</td>
<td>0.01, 0.1, 1, 10</td>
</tr>
<tr>
<td>$\theta$ (°)</td>
<td>15, 20, 25, 30</td>
</tr>
</tbody>
</table>

The initial angle of the reconnected field line $\theta$, the plasma beta $\beta$, and the ratio of the density in the chromosphere and the density in the corona $p_{CH}/p_{CR}$ are parameters. The parameters used are summarized in table 1. We investigated 48 cases.

The initial distributions of the physical parameters are as follows:

$$
\rho = \rho_{CH} - (\rho_{CH} - \rho_{CR}) \left[ \frac{1}{2} \tanh \left( \frac{x - x_d}{w} \right) + 1 \right], \quad (A8)
$$

$$
p = p_0 \left( 1 + \frac{1}{\beta} \right) - \frac{B_x^2 + B_y^2}{8\pi}, \quad (A9)
$$

$$
B_x = B_0 \sin \theta, \quad (A10)
$$

$$
B_y = B_0 \cos \theta \tanh \left( \frac{x - x_c}{w} \right), \quad (A11)
$$

$$
v_x = 0, \quad (A12)
$$

$$
v_y = 0. \quad (A13)
$$
The size of the simulation box was $0 \leq x \leq 240$. The grid spacing was uniform with $\Delta x = 0.0027$ in $0 \leq x \leq 10$, and gradually increased in $10 < x$. The total grid number was 8200. We used free boundary conditions. The scheme used was a CIP-MOCCT scheme.

### A.3. Numerical Results

In figure 21, we show the time evolution of the physical parameters for the case that $\rho_{CH}/\rho_{CR} = 100$, $\beta = 0.1$, and $\theta = 20^\circ$. We can discern that two slow shocks are attached to the outflow region until $t \sim 0.6$. They correspond to the Petchek slow shocks. The transition region (TR) is advected by the inflow, which corresponds to a fast-mode rarefaction wave, into the outflow region. After one of the slow shocks crosses the transition region, the slow shock (SS) propagates in the positive $x$-direction. Simultaneously, a slow-mode rarefaction wave (SR), which corresponds to the reflection wave, is generated and propagates in the negative $x$-direction.

From figure 21, we find that the outflow velocity increases after the interaction between the slow shock and the transition
region. We obtained a relation between the outflow velocity before the interaction $v_{\text{outflow1}}$ and that after the interaction $v_{\text{outflow2}}$: $v_{\text{outflow2}} \sim 2v_{\text{outflow1}}$ (see figure 22).

A.4. Interpretation of Numerical Results

Why does the outflow velocity increase through the interaction between the Petschek slow shock and the transition region? We note that the slow-mode rarefaction wave propagates into the outflow region after the interaction. Behind the slow-mode rarefaction wave, the magnetic field strength increases and the density decreases. Thus, the Alfvén speed increases. The increase of the Alfvén speed can be found in our 2D simulation (see figure 13). This leads to a rapid relaxation of the re-configuration of the magnetic field. The re-configuration of the magnetic field causes the outflow plasma to be further accelerated. A schematic diagram is shown in figure 23.

Considering that the outflow plasma is further accelerated due to magnetic tension after the interaction, we estimate the outflow velocity after the interaction as follows:

$$\frac{dv}{dt} \sim \frac{1}{4\pi} (B \cdot \nabla) B,$$

(A14)

$$\Delta v \sim \frac{B_x}{4\pi \rho} \Delta t,$$

(A15)

$$\sim \frac{B_x}{4\pi \rho} W v_{\text{SMx}},$$

(A16)

where $\Delta v$ is the increase of the outflow velocity through the interaction, $\rho$ is the density in the outflow region, $W$ is the length scale (see figure 23), and $v_{\text{SMx}}$ is the slow-mode speed in the $x$-direction in the outflow region before the interaction. In the outflow region where the magnetic field is in the $x$-direction,

$$v_{\text{SMx}} = \begin{cases} C_s & (\beta < 2/\gamma) \\ V_{Ax} & (\beta > 2/\gamma). \end{cases}$$

(A17)

where $V_{Ax} = B_x/\sqrt{4\pi \rho}$ is the Alfvén speed in the $x$-direction. For all of the cases, the plasma beta in the outflow region is greater than $2/\gamma$. Thus, $v_{\text{SMx}} = V_{Ax}$. Therefore,

$$\Delta v \sim \frac{B_x}{4\pi \rho} W v_{\text{SMx}} = \frac{B_x}{4\pi \rho} W V_{Ax}$$

$$= \frac{B_y}{\sqrt{4\pi \rho}} = V_{Ay},$$

(A18)

where $V_{Ay} = B_y/\sqrt{4\pi \rho}$ is the Alfvén speed in the $y$-direction. Before the interaction, the reconnection outflow velocity $v_{\text{outflow1}}$ is $V_{Ay}$. After the interaction, $v_{\text{outflow2}} = V_{Ay} + \Delta v \sim 2V_{Ay} = 2v_{\text{outflow1}}$. This relation shows no strong dependence on the plasma beta and the ratio of the density jump, which is consistent with the numerical results.

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