Onset of Secondary Instabilities and Plasma Heating during Magnetic Reconnection in Strongly Magnetized Regions of the Low Solar Atmosphere

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Abstract

We numerically study magnetic reconnection on different spatial scales and at different heights in the weakly ionized plasma of the low solar atmosphere (around 300–800 km above the solar surface) within a reactive 2.5D multifluid plasma–neutral model. We consider a strongly magnetized plasma ($\beta \sim 6\%$) evolving from a force-free magnetic configuration and perturbed to initialize formation of a reconnection current sheet. On large scales, the resulting current sheets are observed to undergo a secondary “plasmoid” instability. A series of simulations at different scales demonstrates a cascading current sheet formation process that terminates for current sheets with width of $\sim \text{2 m}$ and length of $\sim \text{100 m}$, corresponding to the critical current sheet aspect ratio of $\sim \text{50}$. We also observe that the plasmoid instability is the primary physical mechanism accelerating the magnetic reconnection in this plasma parameter regime. After plasmoid instabilities appear, the reconnection rate sharply increases to a value of $\sim \text{0.035}$, observed to be independent of the Lundquist number. These characteristics are very similar to magnetic reconnection fully ionized plasmas. In this low-$\beta$ guide-field reconnection regime, both the recombination and collisionless effects are observed to have a small contribution to the reconnection rate. The simulations show that it is difficult to heat the dense weakly ionized photospheric plasmas to above $2 \times 10^4 \text{ K}$ during the magnetic reconnection process. However, the plasmas in the low solar chromosphere can be heated above $3 \times 10^4 \text{ K}$ with reconnection magnetic fields of 500 G or stronger.

Key words: magnetic reconnection – magnetohydrodynamics (MHD) – methods: numerical – Sun: chromosphere

1. Introduction

Magnetic reconnection is a fundamental magnetized plasma process in the heliosphere. In the partially ionized low solar atmosphere, numerous small-scale magnetic reconnection events have been observed by both the ground- and space-based high-resolution solar telescopes. Examples of such observations include microflares (e.g., Ding et al. 1999; Brosius & Holman 2009), chromospheric jets (e.g., Liu et al. 2009; Bharti et al. 2013; Tian et al. 2014; Huang et al. 2017), type II spicules (e.g., Sterling 2000; Sekse et al. 2012), Ellermann bombs (EBs; e.g., Fang et al. 2006; Hong et al. 2014; Nelson et al. 2015), and the high-temperature IRIS bursts (IBs; e.g., Peter et al. 2014; Vissers et al. 2015; Grubecka et al. 2016; Tian et al. 2016; Rouppe van der Voort et al. 2017; Tian et al. 2018b). At present, there are at least two challenges for a magnetic reconnection model in the low solar atmosphere: the plasma heating and the reconnection rate in a viable model should match the observational data.

One of the most puzzling problems of the low solar atmosphere reconnection events is whether or not the IBs are related to the traditional EBs. The traditional EBs are usually observed in the upper solar photosphere or the low chromosphere. A semi-empirical model of a variety of combinations of the spectral lines originally led to estimates of the temperature increases in EBs ranging from $\sim \text{400}$ to $\text{2000 K}$ (e.g., Fang et al. 2006; Berlicki & Heinzel 2014; Grubecka et al. 2016). The IBs observed by the IRIS Si IV “Transition Region” (TR) lines imply that the temperature in IBs can be above $8 \times 10^4 \text{ K}$ (De Pontieu et al. 2014; Peter et al. 2014). Peter et al. (2014) suggested that such IBs could be evidence of heating within photospheric EBs with temperatures an order of magnitude higher than those predicted by previous semi-empirical models. The recent joint observations from ground- and space-based telescopes also support the assertion that the plasmas within some of the photospheric active region EBs and quiet Sun Ellermann-like brightenings can be strongly heated to reach TR temperatures (e.g., Vissers et al. 2015; Tian et al. 2016, 2018a; Nelson et al. 2017). However, Rutten (2016) suggested that temperatures as low as $\sim 2 \times 10^4 \text{ K}$ could account for the observed increased emission in the Si IV line. In their work, they assumed local thermodynamic equilibrium for line extinctions during the onsets of EBs.

Several previous single-fluid MHD simulations showed that the maximum temperature increases in magnetic reconnection regions below the upper chromosphere are always only several thousand Kelvin (e.g., Isobe et al. 2007; Archontis & Hood 2009; Xu et al. 2011; Hansteen et al. 2017). However, the resolution in these simulations was not sufficient to resolve the current sheet substructure, and the artificial hyperdiffusivity operator that was included to prevent the collapse of the current sheets leaves open the possibility of formation of a hotter smaller scale. The high-resolution simulations with the more realistic magnetic diffusion in Ni et al. (2016) showed that the plasma can be heated from $\text{4200 K}$ to above $8 \times 10^4 \text{ K}$ inside the multiple magnetic islands of a reconnection process with strong magnetic fields (500 G) in the temperature minimum region (TMR) of the solar atmosphere. Ambipolar diffusion, temperature-dependent magnetic diffusion, heat conduction,
and optically thin radiative cooling were all included in Ni et al. (2016). However, some of the important interactions between ions and neutrals were ignored, and the non-equilibrium ionization–recombination effects were not considered in any of the single-fluid MHD simulations. The recent reactive multifluid plasma–neutral simulations (Ni et al. 2018a, 2018b) with the HiFi code showed that the plasma in TMR can be heated above $2 \times 10^4$ K when the reconnection magnetic field strength was greater than 500 G. Ni et al. (2018a) has also proved that the non-equilibrium ionization–recombination dynamics play a critical role in determining the structure of the reconnection region, and it leads to much lower temperature increases as compared to simulations that assume plasma to be in ionization–recombination equilibrium. The high temperature increases in the single-fluid simulations might be overestimated.

The spatial scales, lifetimes, plasma flow velocities, and emission intensity of the low solar atmospheric reconnection events are usually much smaller than the eruptions in the solar corona. According to the length scale, lifetime, and the estimated Alfvén speed (or the reconnection outflow speed), the reconnection rate in low solar atmosphere events can be estimated as $M \approx L/(V_{A,\text{finite}}) \approx L/(V_{out,\text{finite}})$, which is about 0.01–0.5 (e.g., Nishizuka et al. 2011; Leake et al. 2012; Ni et al. 2015). The magnetic diffusion coefficient including both the electron–ion and electron–neutral collisions (Leake et al. 2012) is estimated order of $10^3–10^4$ m$^2$ s$^{-1}$ from the upper photosphere to the middle chromosphere. The corresponding Lundquist number (magnetic Reynolds number) for a length scale $L \sim 1$ Mm and Alfvén speed $V_A \sim (10–100)$ km s$^{-1}$ is of order $S = L V_A / \eta \sim 10^3–10^4$. Thus, the reconnection rate for the classical Sweet–Parker current sheet is calculated as $M_{SP} \sim S^{-1/2} \sim 10^{-4}–10^{-3}$ from the one-fluid model, which is far smaller than the reconnection rate required in low solar atmosphere reconnection events.

The plasmas around the TMR are weakly ionized, and the plasma density is very high. Thus, it is important to include the interactions between ions and neutral particles as well as the radiative cooling in any realistic model of magnetic reconnection around the TMR. Previous 1D analytical works (e.g., Vishniac & Lazarian 1999; Heitsch & Zweibel 2003; Lazarian et al. 2004) found that the high recombination can induce a loss of ions in the reconnection region that prevents ion pressure from building up, which leads to faster magnetic reconnection independent of magnetic diffusivity. Leake et al. (2012, 2013) used the reactive multifluid plasma–neutral module within the HiFi modeling framework to study null-point magnetic reconnection in the solar chromosphere. They showed that even for a Lundquist number as high as $10^2$, the reconnection rates can reach above 0.05, which is much faster than that predicted from the one-fluid Sweet–Parker model. They concluded that the strong ion recombination in the reconnection region, combined with Alfvénic outflows, can lead to fast null-point reconnection within a 2D numerical model. The plasmoid instability only increased the reconnection rate by approximately 15%, and the Hall effects were negligible in their simulations. The recent 2.5D simulations by Ni et al. (2018a, 2018b) used the same reactive multifluid plasma–neutral module within the HiFi modeling framework to model guide-field, low-$\beta$ reconnection within a weakly ionized plasma. Their results showed that the decoupling of the ionized and neutral fluids was not obvious in the strongly magnetized regions with much lower plasma $\beta$, and the ionization rate was much faster than the recombination rate, which was very different from the previous high-$\beta$ null-point simulations by Leake et al. (2012, 2013). Though the strong radiative cooling was shown to result in loss of most of the generated thermal energy, it also did not noticeably accelerate magnetic reconnection (Ni et al. 2018a, 2018b), which is different from the high-$\beta$ simulations with radiative cooling in Alvarez Laguna et al. (2017).

In this work, we present 2.5D simulations of magnetic reconnection on different spatial scales and at different heights above the solar surface by continuing to use the reactive multifluid plasma–neutral module within the HiFi modeling framework. Compared with our previous works (Ni et al. 2018a, 2018b), the simulations are performed on much larger scales. We find the critical Lundquist number and current sheet aspect ratio for the onset of the plasmoid instability in an initially strongly magnetized, weakly ionized plasma in the solar TMR.

2. Numerical Model and Initial Conditions

The same reactive multifluid plasma–neutral module of the HiFi modeling framework as used by Ni et al. (2018a) is applied in this work. We still normalize the equations by using the characteristic plasma density and magnetic field around the strongly magnetized solar TMR. The characteristic plasma number density is $n_s = 10^{21}$ m$^{-3}$, and the characteristic magnetic field is $B_s = 0.05$ T = 500 G. From these quantities, we derive additional normalizing values to be $v_A \equiv B_s / \sqrt{\mu_0 m_p n_s} = 34.613$ km s$^{-1}$ and $T_s \equiv B_s^2 / (k_B n_s) = 1.441 \times 10^4$ K. There are three characteristic length scales in our simulations: $L_1 = 200$ m, $L_2 = 1000$ m, and $L_3 = 10^3$ m. The three corresponding characteristic times are $t_1 \equiv L_1 / v_A = 0.0058$ s, $t_2 \equiv L_2 / v_A = 0.029$ s, and $t_3 \equiv L_3 / v_A = 0.29$ s, respectively.

The simulations are initialized with a force-free Harris sheet magnetic equilibria, where the reconnecting magnetic field is along the $y$-direction and the guide $B$-field is in the $z$-direction. Specifically, the initial dimensionless magnetic flux in the $z$-direction is given by

$$A_{\theta}(y) = -b_p \lambda_{\psi} \ln \left[ \cosh \left( \frac{y}{\lambda_{\psi}} \right) \right],$$

where $b_p$ is the strength of the the magnetic field, and $\lambda_{\psi}$ is the initial thickness of the current sheet. The initial magnetic field in the $z$-direction is

$$B_{\theta}(y) = b_p \left[ \frac{1}{\cosh \left( \frac{y}{\lambda_{\psi}} \right)} \right].$$

Since the ion inertial length is much smaller than the width of the current sheets in this work, the Hall effect is not important and ignored. The numerical results in our previous paper (Ni et al. 2018a) demonstrate that the collisions between electrons and neutrals are not important for magnetic reconnection in low-$\beta$ plasmas around the solar TMR. We
also ignore the collisions between electrons and neutral particles, and the dimensionless magnetic diffusivity in this work is

\[ \eta = \eta_{ei} = \eta_{el} T_e^{-1.5}, \]

where \( \eta_{ei} \) is a normalization constant derived from the characteristic values \( n_e, B_x, \) and \( L_x \). One can get \( \eta_{ei+1} = 3.729 \times 10^{-6}, \eta_{ei+2} = 7.457 \times 10^{-7}, \) and \( \eta_{ei+3} = 7.457 \times 10^{-8} \) by using the three different length scales. \( T_e \) is the dimensionless electron temperature. The electron and ion fluids are assumed to be well coupled and only the hydrogen gas is considered in our model. Thus, we assume \( T_e = T_i, n_e = n_i, \) and the pressure of the ionized component is twice the ion (or electron) pressure, \( P_p = P_e + P_i = 2P_e \).

In this work, we have separately tested our simulations by using two different radiative cooling functions in the plasma internal energy equation as presented in our previous papers (Ni et al. 2018a, 2018b). The first simple radiative cooling model (Meier 2011; Leake et al. 2012) represents the radiative losses that are due to atomic physics such as radiative recombination, and the excited states are not tracked. The formula of this simple model is given by

\[ L_{\text{rad1}} = \Gamma_{i}^{\text{ion}} \phi_{\text{eff}}, \tag{4} \]

where \( \phi_{\text{eff}} = 33 \text{ eV} \times 5.28 \times 10^{-18} \text{J} \). The ionization rate \( \Gamma_{i}^{\text{ion}} \) is defined as

\[ \Gamma_{i}^{\text{ion}} = \frac{n_n n_e A}{X + \phi_{\text{ion}}/T_e^2} \left( \frac{\phi_{\text{ion}}}{T_e^2} \right)^K \exp \left( -\frac{\phi_{\text{ion}}}{T_e^2} \right), \tag{5} \]

using the values \( A = 2.91 \times 10^{-14}, K = 0.39, X = 0.232, \) and the hydrogen ionization potential \( \phi_{\text{ion}} = 13.6 \text{ eV} \). The units for the neutral and electron number density are both m\(^{-3}\), and \( T_e^2 \) is the electron temperature specified in eV. Then, the units for \( \Gamma_{i}^{\text{ion}} \) is m\(^{-3}\) s\(^{-1}\) and the units for \( L_{\text{rad}} \) is J m\(^{-3}\) s\(^{-1}\). The radiative cooling goes to zero if the plasma is fully ionized. As shown in our previous work (Ni et al. 2018a), the plasmas will be fully ionized if the reconnection magnetic fields are strong enough. This simple radiative cooling model becomes invalid in this situation. The second radiative cooling function is computed by using the OPACITY project and the CHIANTI databases (Goodman & Judge 2012). This radiative cooling model is considered to be a more realistic cooling model for plasmas in the solar chromosphere (Alvarez Laguna et al. 2017). The expression of this radiative model in units of J m\(^{-3}\) s\(^{-1}\) is

\[ L_{\text{rad2}} = C_F n_e (n_n + n_i) 8.63 \times 10^{-6} T_e^{-1/2} \]

\[ \times \sum_{i=1}^{2} E_i T_i \exp (-e E_i / k_B T), \tag{6} \]

where \( C_F = 1.6022 \times 10^{-25} \text{ J eV}^{-1}, E_1 = 3.54 \text{ eV} \) and \( E_2 = 8.28 \text{ eV} \), and \( T_1 = 0.15 \times 10^{-3} \) and \( T_2 = 0.065 \). A three-level hydrogen atom with two excited levels is included in this function, and \( E_1 \) and \( E_2 \) are the excited level energies. The units for the temperature in this equation is Kelvin, and the units for the number density is m\(^{-3}\). The background constant heating is included to balance the initial radiative cooling when the second radiative cooling model is applied. The heating function \( H_0 \) is equal to the initial value of \( L_{\text{rad2}} \). In our simulations, we have normalized \( L_{\text{rad1}} \) in Equation (4) and \( L_{\text{rad2}} \) in Equation (6) by using the characteristic values presented above.

The two radiative cooling models with the remaining simulation setup being identical for each case have been tested, and we find that the main conclusions in this work are not dependent on the two different radiative cooling terms applied in the code. We will present the numerical results of five cases in this work, Cases I, II, III, IV, and V. In order to compare with Case A in our previous work (Ni et al. 2018a), Cases I, II, III, and IV show the results with the simple radiative cooling model \( L_{\text{rad1}} \). Since the plasmas in Case V will be fully ionized if the simple radiative cooling model is applied, Case V shows the results with the more realistic radiative cooling model \( L_{\text{rad2}} \). Both a higher resolution run and a lower resolution run have been carried out for each case.

Cases I, II, and III represent magnetic reconnection around the solar TMR (\( \sim 600 \text{ km above the solar surface} \)) on three different scales. All of the initial parameters in Cases I, II, and III are the same, except for the simulation length scales. The characteristic length scales in Cases I, II, and III are \( L_{s1}, L_{s2}, \) and \( L_{s3} \) as presented above, respectively. The initial neutral particle number density and ionization fraction in these three cases are \( n_{n0} = 0.5 n_n \) and \( f_{i0} = 0.01 \% \). Since the hydrogen gas only has the ground state and the ionized state in this model, the initial temperatures of the ionized and neutral fluids are set to be uniform at \( T_{i0} = T_{n0} = 0.05829 T_e \) in Cases I, II, and III to keep the initial ionization fraction \( f_{i0} = 0.01 \% \), which is the same as that around the solar TMR. The strength of the initial magnetic field in dimensionless form is \( b_p = 1 \) in Cases I, II, and III.

Case IV represents magnetic reconnection in the photosphere (around \( 300 \text{ km above the solar surface} \)) with a length scale of \( L_{s1} = 200 \text{ m} \). The initial neutral particle number density and ionization fraction in Case IV are \( n_{n0} = 12.5 n_n \) and \( f_{i0} = 0.01 \% \). The initial temperatures of the ionized and neutral fluids are also set to be uniform at \( T_{i0} = T_{n0} = 0.05829 T_e \) in this case. The initial magnetic field in dimensionless form in Case IV is \( b_p = 5 \), which is five times higher than that in Case I.

Case V represents magnetic reconnection in the low solar chromosphere (around \( 800 \text{ km above the solar surface} \)) with a length scale of \( L_{s1} = 200 \text{ m} \). The initial neutral particle number density and ionization fraction in Case V are \( n_{n0} = 0.1 n_n \) and \( f_{i0} = 0.1 \% \). The initial temperatures of the ionized and neutral fluids are also set to be uniform at \( T_{i0} = T_{n0} = 0.06559 T_e \) in this case. The initial magnetic field in dimensionless form in Case V is \( b_p = 1 \), which is the same as in Case I.

The initial parameters in the five cases are presented in Table 1. We only simulate one-quarter of the domain \( (0 < x < 2L_x, 0 < y < L_y) \) in all five cases; \( L_x \) corresponding to \( L_{s1}, L_{s2}, \) and \( L_{s3} \) in Cases I, II, and III, respectively. We use the same outer boundary conditions at \( \vert y \vert = L_y \) and the same form of the initial electric field perturbations as Murphy & Lukin (2015) to initiate magnetic reconnection in all of the cases in this work. The perturbation electric field is applied for \( 0 \leq t \leq t_e \). The perturbation magnitude is proportional to the value of \( b_p \) in each of the cases with the amplitude of \( \delta E = 10^{-3} b_p \). Periodicity of the physical system is imposed in the \( x \)-direction at \( \vert x \vert = 2L_x \). The presented simulation results from all five cases have been tested by using a lower resolution and a higher resolution. The highest number of grid elements in Case III is \( m_x = 288 \) elements in the \( x \)-direction and \( m_y = 240 \).
The Astrophysical Journal, 868:144 (10pp), 2018 December 1

Ni & Lukin

The Different Characteristic Length Scales \( L_c \) and Times \( t_c \), and the Initial Parameters in the Five Cases

| Case | \( L_c (\text{m}) \) | \( t_c (s) \) | \( n_{\alpha} \) | \( j_\alpha \) | \( T_\alpha \) | \( N_{\beta} \) | \( T_{\beta} \) | \( f_\alpha \) | Radiative Cooling |
|------|------------------|------------|-----------|----------|----------|---------|---------|----------|----------------|
| A    | 100              | 0.0029     | 0.5 \( n_e \) | \( 10^{-4} \) | 0.05829 \( T_F \) | 1        | 0.058    | \( L_{\text{rad1}} \) |
| I    | 200              | 0.0058     | 0.5 \( n_e \) | \( 10^{-4} \) | 0.05829 \( T_F \) | 1        | 0.058    | \( L_{\text{rad1}} \) |
| II   | \( 10^3 \)       | 0.029      | 0.5 \( n_e \) | \( 10^{-4} \) | 0.05829 \( T_F \) | 1        | 0.058    | \( L_{\text{rad1}} \) |
| III  | \( 10^4 \)       | 0.29       | 0.5 \( n_e \) | \( 10^{-4} \) | 0.05829 \( T_F \) | 1        | 0.058    | \( L_{\text{rad1}} \) |
| IV   | 200              | 0.0058     | 12.5 \( n_e \) | \( 10^{-4} \) | 0.05829 \( T_F \) | 5        | 0.058    | \( L_{\text{rad1}} \) |
| V    | 200              | 0.0058     | 0.1 \( n_e \) | \( 10^{-3} \) | 0.06559 \( T_F \) | 1        | 0.01305  | \( L_{\text{rad2}} \) |

Note. The characteristic plasma number density \( n_e = 10^{21} \text{ m}^{-3} \), the characteristic magnetic field \( B_e = 0.05 \text{ T} \), the characteristic velocity \( V_e = 34.613 \text{ km s}^{-1} \), and the characteristic temperature \( T_e = 1.441 \times 10^6 \text{ K} \) are the same in all five cases.

The Smallest Normalized Grid Sizes Along the \( y \)-direction Near \( y = 0 \) in Both the Higher and Lower Resolution Runs of Cases I, II, III, IV, and V

| Case | Lower resolution | Higher resolution |
|------|------------------|-------------------|
| I    | 3.095 \( \times 10^{-4} \) | 1.855 \( \times 10^{-4} \) |
| II   | 1.284 \( \times 10^{-4} \) | 3.309 \( \times 10^{-4} \) |
| III  | 7.950 \( \times 10^{-5} \) | 3.095 \( \times 10^{-4} \) |
| IV   | 3.095 \( \times 10^{-4} \) | 1.855 \( \times 10^{-4} \) |
| V    | 3.095 \( \times 10^{-4} \) | 1.855 \( \times 10^{-4} \) |

3. Numerical Results

3.1. Magnetic Reconnection on Different Length Scales

In our previous papers (Ni et al. 2018a, 2018b), magnetic reconnection on a very small scale \( (L_c = 100 \text{ m}) \) with low Lundquist number \( (\sim 2000) \) has been studied. The plasmoid instability did not appear in those simulations. In this section, we will show studies of magnetic reconnection on three larger scales: Case I with \( L_{x1} = 200 \text{ m} \), Case II with \( L_{x2} = 1000 \text{ m} \), and Case III with \( L_{x3} = 10^4 \text{ m} \). The plasmoid instability appears in all three cases. Case A in Ni et al. (2018a) is used to compare with the three cases in this work. The only difference among the four cases (Cases A, I, II, and III) is the characteristic length scale. In this work, the half-length \( L_{\text{sim}} \) and half-width \( \delta_{\text{sim}} \) of the current sheet are defined by using the half-width at half-maximum and half-length at half-maximum in the current density \( J_{\alpha}/J_{\beta} \). The method for calculating the reconnection rate \( M_{\text{sim}} \) in this work is the same as that in the previous papers (Leake et al. 2013; Ni et al. 2018a, 2018b), \( M_{\text{sim}} = \eta j_{\text{max}} \sqrt{V_{\alpha}^2 B_{\text{up}}} \), where \( j_{\text{max}} \) is the maximum value of the out-of-plane current density \( J_{\alpha}/J_{\beta} \) at the reconnection X-point \( (x_{\beta}, y_{\beta}) \), \( B_{\text{up}} \) is \( B_e \) evaluated at \( (x_{\beta}, \delta_{\text{sim}}) \). \( V_{\alpha}^2 \) is the relevant Alfvén velocity defined using \( B_{\text{up}} \) and the total number density \( n^e \) at the location of \( j_{\text{max}} \). \( \eta \) is the magnetic diffusion coefficient defined by Equation (3) at the location of \( j_{\text{max}} \). The calculated Lundquist number \( S_{\text{sim}} \) is defined as \( S_{\text{sim}} = V_{\alpha}^2 L_{\text{sim}}/\eta^2 \).

Figure 1 presents the distributions of the current density and ionization fraction during the later stage of the reconnection process in Cases A, I, II, and III. This figure reveals the cascading process of the plasmoid instability. Several magnetic islands appear inside the current sheet in Case III. The outflow region is very turbulent in Case III—many small-scale current sheet fragments and small magnetic islands even appear in the outflow region. The small-scale current sheet fragments also appear in the outflow regions of Cases II and I, but this region is relatively smoother in the case with smaller length scale. We only identify one magnetic island in Case I. There is no magnetic island inside the current sheet of Case A. The outflow region of Case A is also very smooth; no current sheet fragments appear there. Therefore, the plasmoid instabilities terminate at the reconnection current sheet length of 100 m.

Figures 2(a) and (b) show the calculated aspect ratio of the current sheets (CSs) \( L_{\text{sim}}/\delta_{\text{sim}} \) and Lundquist number \( S_{\text{sim}} \) versus time before magnetic islands appear in Cases A, I, II, and III; both the higher and lower resolution results of Cases I, II, and III are presented. One should note that the characteristic length scale \( L_c \) and time \( t_c \) are different in the different cases. \( f_{\beta1} \) is the dimensionless time point when the first magnetic island appears in each run. For the same case, one can find that magnetic islands appear later in the run with the higher resolution. The reason is that the higher resolution results in the lower level of numerical noises, which allows the Lundquist number to reach higher values before the plasmoid instability enters the nonlinear regime. The sensitivity of plasmoid-unstable high Lundquist number current sheets to noise is well known in the single-fluid MHD regime, e.g., see Huang & Bhattacharjee (2010), and is apparent here for Cases II and III. In Case I, the aspect ratio reaches about 70, and the Lundquist number is about 3400 just before the first magnetic island appears in both the higher and the lower resolution runs. The aspect ratio reaches 45, and the Lundquist number is about 2000 at the end of the simulation in Case A, but no magnetic island appears in this case. Therefore, we can conclude that the critical aspect ratio is between 45–70, and the critical Lundquist number is between 2000 and 3400 for the plasmoid instability to appear in our simulations. These results are different from the previous work by Leake et al. (2012). In their higher \( \beta \) simulations with a 2D Harris sheet as initial magnetic field configurations, they find that the critical aspect ratio is about 200, and the critical Lundquist number is about \( 10^7 \) for the plasmoid instability. They point out that the magnetic islands that form due to the plasmoid instability in their simulations are losing ions due to recombination, and so their evolution is potentially very different from those in fully ionized conditions.
simulations. In their later work with the more realistic magnetic diffusion (Leake et al. 2013), they find that the plasmoid instability does not appear even when the aspect ratio is greater than 500, which they argue may be due to the lower magnetic Prandtl number used in Leake et al. (2013), allowing for stronger shearing of the reconnection outflow and leading to the stabilization of the plasmoid instability. However, we should point out that the critical Lundquist number and aspect ratio depend on many different factors, e.g., the amplitudes of viscosity and perturbation noises (Comisso & Grasso 2016; Comisso et al. 2016, 2017; Huang et al. 2017), and the plasma $\beta$ and the initial current sheet configuration (Ni et al. 2012, 2013) can impact the onset of plasmoid instability. In previous works on magnetic reconnection in fully ionized plasmas (Ni et al. 2012, 2013), it was found that the critical Lundquist number in the 2D case with an initial Harris current sheet and nonuniform plasma pressure is much higher than that in the 2.5D case with an initial force-free current sheet and uniform plasma pressure. The results presented here are consistent with these prior findings.

The prior work by Huang et al. (2017) indicates that the aspect ratio scales with the Lundquist number as $L_{\text{sim}}/\delta_{\text{sim}} \sim S^{1/2}$ for the lower Lundquist cases ($S < 2 \times 10^5$), and this scaling is closer to $L_{\text{sim}}/\delta_{\text{sim}} \sim S^{1/3}$ for the higher Lundquist number cases.

![Figure 1](image_url) Figure 1. Magnetic reconnection during the later stage at different length scales. The left side shows the distributions of the dimensionless current density $J_z/J_\star$, and the right side shows the ionization fraction $f_i$. The values of the characteristic length scale $L_\star$ correspond to $L_{\text{sim}} = 200$ m in Case I, $L_{\text{sim}} = 1000$ m in Case II, and $L_{\text{sim}} = 10^4$ m in Case III, respectively. One should note that the values of the characteristic current density $J_\star$ are also different in each case because of the different length scales, but the same color bar is used for all cases. Since the length scale in Case A is the smallest, the value of the characteristic current density $J_\star$ in Case A is greatest, which makes the distribution of the normalized current density $J_z/J_\star$ in Case A hardly visible. The green boxes correspond to the same size in each case with different length scales.
From Figures 2a and (b) in this work, we can find that the aspect ratios are 70, 116, and 120, and the corresponding Lundquist numbers are 3400, 11,000, and 17,000 at \( t_{p1} \) just before the first magnetic island appears in Cases I, II, and III, respectively. The scaling in this work is closer to \( \frac{L_{sim}}{\delta_{sim}} \sim S_{1}^{1/2} \). Therefore, our numerical results are consistent with Huang et al. (2017).

Figures 3(a) and (b) show the half-width of the current sheet \( \delta_{sim} \) and the calculated reconnection rate \( M_{sim} \) vs. time in Case A in our previous work and in Cases I, II, and III in this work before and after the plasmoid instability appear. \( t_{p1} \) represents the time point when the first magnetic island appeared inside the current sheet in each run; \( t_{p2} \) represents the time point when the second magnetic island appeared inside the current sheet in the lower resolution run of Case III. The values of \( t_{\alpha} \) are \( t_{\alpha 1} = 0.0058 \) s in Case I, \( t_{\alpha 2} = 0.029 \) s in Case II, and \( t_{\alpha 3} = 0.29 \) s in Case III, respectively.

From Figures 2(a) and (b) in this work, we can find that the aspect ratios are 70, 116, and 120, and the corresponding Lundquist numbers are 3400, 11,000, and 17,000 at \( t_{p1} \) just before the first magnetic island appears in Cases I, II, and III, respectively. The scaling in this work is closer to \( \frac{L_{sim}}{\delta_{sim}} \sim S_{1}^{1/2} \). Therefore, our numerical results are consistent with Huang et al. (2017).

Figures 3(a) and (b) show the half-width of the current sheet and reconnection rate versus time before and after plasmoid instabilities appear; both the higher and lower resolution results of Cases I, II, and III are presented. In both the higher and lower resolution runs, one can find that the half-widths of current sheets (CSs) drop to about 2 m, 5 m, and 20 m at \( t_{p1} \) just before the first magnetic island appears inside the current sheet in Cases I, II, and III, respectively. \( t_{p2} \) is the time point when the second magnetic island begins to appear inside the current sheet in the lower resolution run of Case III. After plasmoid instabilities appear, the current sheet width around the main reconnection X-point continues to drop sharply in all of the three cases. In both the higher and lower resolution runs, the thinnest half-width eventually reaches about 1 m and 2 m in Cases I and II, respectively. It reaches about 7 m in the lower resolution run of Case III. One should note that the higher resolution simulation of Case III is terminated when the plasmoid instability just develops into the nonlinear stage, and the sharp decrease of the current sheet width does not happen yet. The resolution in the lower resolution run of Case III is obviously not high enough for the half-width of the current sheet to reach the 1 m scale. In Case I, the result in the higher resolution run confirms that the half-width of the current sheet thins to 1 m without undergoing the higher order plasmoid instability. When the half-width of the current sheet thins to 1 m in Case A, the plasmoid instability does not appear. Therefore, we can conclude that the plasmoid instability terminates at the reconnection current sheet width of 2 m.

From Figure 3(b), one can find that the reconnection rate in Case III is the lowest in all three cases before the plasmoid...
instabilities appear, with the reconnection rate approximately scaling with the Lundquist number as \( M_{\text{sim}} \sim S_{\text{sim}}^{-1/2} \) at \( t_{p1} \). After plasmoid instabilities appear, the reconnection rates increase sharply in all three cases. The first peak of the reconnection rate in the lower resolution run of Case III is about 0.015, which is about two times higher than that before the first magnetic island appears. The reconnection rate reaches above 0.035 after the second magnetic island appears, resulting in a five-fold increase in the reconnection rate in the lower resolution run of Case III during a very short period. Though the highest reconnection rate in this work and that in Leake et al. (2013) are close, the physical mechanisms for increasing the reconnection rates are very different in the two studies. In Leake et al. (2013), the main mechanism for increasing the reconnection rate is the strong recombination effect. The plasmoid instabilities in their simulations only increase the reconnection rate by approximately 15%. In this work, the strong ionization rate in our low-\( \beta \) simulations makes the neutral–ion collisional mean free path within the current sheet much smaller than the width of the current sheet, and the ionized and neutral fluid flows are well coupled throughout the reconnection region. Consequently, the ionization–recombination non-equilibrium effects do not lead to fast magnetic reconnection in this regime. The plasmoid instability is still the main mechanism to lead the fast magnetic reconnection in Case III in this work. As mentioned in the previous paragraph, the higher resolution simulation of Case III is terminated when the plasmoid instability just develops into the nonlinear stage; the sharp increase of the reconnection rate also does not appear yet. However, one can find that the reconnection rate sharply increases to a similarly high value in both the higher and the lower resolution runs of Case II. Therefore, we can expect that the reconnection rate in the higher resolution run of Case III will also behave similarly to that in the lower resolution run. Still higher resolution simulations of Case III will be pursued in future work to track multiscale plasmoid evolution further into the nonlinear regime.

As shown in Figures 1 and 3, the plasmoid instabilities terminate at the reconnection current sheet length of 100 m and width of 2 m. The plasmas inside the current sheet become strongly ionized during the reconnection process, and the ion density can reach about \( 0.5 \times 10^{21} \text{ m}^{-3} \). Then, inside the current sheet, the neutral–ion collisional mean free path and the ion inertial length are \( \lambda_n \approx 2.37 \times 10^{-3} \text{ m} \) and \( d_i \approx 0.01 \text{ m} \), which are both much smaller than the width of the thinnest current sheet. Therefore, collisionless effects can also be ignored within reconnection current sheets in strongly magnetized regions around the solar TMR. Though the recombination effect does not significantly increase the reconnection rate in such strongly magnetized regions, the non-equilibrium ionization–recombination effect is still important. As discussed in Ni et al. (2018a), the non-equilibrium ionization–recombination factor leads to much lower temperature increases.

By comparing the numerical results in Cases A, I, II, and III, we find that the ionized and neutral fluid flows are well coupled throughout the reconnection region even after plasmoid instabilities appear in all cases (not shown). The ionization rate is also higher than the recombination rate in the reconnection process, and most of the generated thermal energy is also radiated away through radiative cooling in all cases (not shown).

### 3.2. Magnetic Reconnection at Different Heights in the Low Solar Atmosphere

The plasma density in Case IV is 25 times higher than that in Case I, and the plasma density in Case V is five times lower than that in Case I. According to the VAL-C solar atmosphere model (Vernazza et al. 1981), Cases IV, I, and V represent magnetic reconnection 300 km above the solar surface in the photosphere, 600 km above the solar surface around the TMR, and 800 km above the solar surface in the low solar chromosphere, respectively. The characteristic length scale in Cases I, IV, and V is the same with \( L_{\text{si}} = 200 \text{ m} \). The strength of reconnection magnetic field in Case IV is 2500 G, which is five times higher than that in Case I and Case V. The initial plasma \( \beta \) in both Case I and Case IV is 0.058.

Figure 4(a) shows the evolution of the calculated Lundquist number \( S_{\text{sim}} \) with time before plasmoid instabilities appear, and Figure 4(b) shows the reconnection rate \( M_{\text{sim}} \) versus time before and after plasmoid instabilities appear. Magnetic reconnection starts at about \( t = 14t_{\text{si}} \) in Case I and Case IV, and it starts at about \( t = 8.5t_{\text{si}} \) in Case V. \( t_{\text{si}} \) corresponds to the time point when the first magnetic island starts to appear in

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**Figure 4.** Panels (a) and (b) show the calculated Lundquist number \( S_{\text{sim}} \) and reconnection rate \( M_{\text{sim}} \) vs. time in Cases I, IV, and V. \( t_{p1} \) represents the time point when the first magnetic island appears inside the current sheet in Cases I and V.
Cases I and V. The Lundquist number is about 3400 just before the first magnetic island appears in Case I. At the end of the simulation in Case IV, the Lundquist number is about 2000, which is too low to result in plasmoid instabilities, and no magnetic islands appear in this case. After $t = 174t_{i,i}$, the quick increase of the plasma temperature at the reconnection X-point in Case V leads to the decrease of magnetic diffusion and the increase of Lundquist number $S_{\text{sim}}$, the value of $S_{\text{sim}}$ eventually reaches about $1.4 \times 10^4$ in Case V before the first magnetic island appears. As shown in Figure 4(b), magnetic reconnection basically follows the Sweet–Parker model with $M_{\text{sim}} \sim S_{\text{sim}}^{-1/2}$ before plasmoid instabilities appear. After the first magnetic islands appear, the reconnection rate starts to increase with time in Cases I and V. From these results, we can also conclude that the plasmoid cascading process will terminate at a larger scale in an environment with higher plasma density or weaker reconnection magnetic fields.

Figure 5 shows the distributions of current density $J_z/J_{\text{a,i}}$ and ionization fraction $f_i$ at three different time points in Cases I, IV, and V. The maximum current density during the later stage of the reconnection process reaches the greatest value in Case IV with the strongest reconnection magnetic fields. Though the strength of the reconnection magnetic fields is the same in Case I and V, the maximum current density in Case V is about two times greater than that in Case I. Therefore, the current density generated in the reconnection process relates to both the strength of the reconnection magnetic fields and plasma density. The stronger reconnection magnetic fields and lower plasma density both can result in the greater maximum current density in the reconnection process. In the case with higher plasma density, the temperature increase is lower, and the magnetic diffusion is higher. The higher plasma density also results in the stronger viscous diffusion. Therefore, the current is more strongly dissipated by higher magnetic and viscous diffusion in the case with higher plasma density. However, the results in Figure 5 show that the current density depends on the strength of the magnetic field more strongly than the plasma density. The initial values of the plasma $\beta$ and ionization fraction in Case I and Case IV are the same, but the maximum ionization fraction in Case IV is much smaller than that in Case I during the later stages of the magnetic reconnection process. The maximum ionization fraction reaches the greatest value in Case V with the minimum initial plasma density. The stronger reconnection magnetic fields also result in the higher maximum ionization fraction (Ni et al. 2018a). How high the maximum ionization fraction can reach during a reconnection process also depends on both the initial plasma density and the strength of the reconnection magnetic fields. Figure 5 also shows that the width of the current sheet during the later reconnection stage is a minimum in Case V, which is measured to be about 1.5 m by using the half-width at the half-maximum in the current density $J_z/J_{\text{a,i}}$. Though the plasma density in Case V is five times lower than that in Case I, we find that the width of the current sheet in Case V is still more than one order of magnitude larger than both the ion–neutral collisional mean free path $\lambda_{in}$ and the ion inertial length $d_i$ during the magnetic reconnection process.

Figure 6 shows the profiles of ion and neutral temperatures across the reconnection X-point along the y-direction in Cases I, IV, and V at the same three time instances for each of the Cases as in Figure 5. The plasma temperatures inside the current sheet increase with time during the reconnection process in each of the three cases. The maximum temperatures in Cases I, IV, and V reach $1.8 \times 10^4$ K, $1.3 \times 10^4$ K, and $3.1 \times 10^4$ K before the simulations are terminated due to lack of spatial resolution. The simulation results show that it is very difficult to heat the much denser photosphere plasmas to high temperatures. The plasma temperature over $2 \times 10^4$ K is not observed in Case IV even when the reconnection magnetic field is as strong as 2500 G. However, we cannot discard the possibility that the plasma temperature could exceed $2 \times 10^4$ K in Case IV if the simulation could run for a significantly longer period of time. In Case V with much lower number density, the plasma can be easily heated above $3 \times 10^4$ K as shown in our simulations.

4. Summary and Discussions

In this work, we have used the reactive multifluid plasma–neutral module of the HiFi modeling framework to study magnetic reconnection in initially weakly ionized plasmas
around the solar TMR on different spatial scales and also at three different heights from 300 to 800 km above the solar surface. The main conclusions are as follows:

1. In strongly magnetized solar TMR, the critical Lundquist number for the onset of the plasmoid instability is between 2000 and \( \sim 3400 \), and the corresponding critical aspect ratio is between 45 and \( \sim 70 \) in current sheets with the guide field evolving from a force-free magnetic configuration in 2.5D simulations. These critical values are about five times smaller than those in the previous works by Leake et al. (2012, 2013), which considered magnetic reconnection in a weakly magnetized middle chromosphere in a Harris Sheet resulting in null-point reconnection. The different plasma \( \beta \) and configurations of the reconnecting magnetic field are the main reasons for this substantial difference. The critical Lundquist number and aspect ratio in this work are very close to the results in previous work by Ni et al. (2013), in which magnetic reconnection in fully ionized plasmas was studied by using the same initial force-free magnetic field configuration.

2. The numerical results of magnetic reconnection at different scales show that the plasmoid cascading process terminates at a length scale of 100 m and width of 2 m in the strongly magnetized solar TMR with reconnection magnetic fields of 500 G and plasma \( \beta \) of 0.058. These length scales are much larger than the ion inertial length and the ion–neutral collisional mean free path within the resulting current sheets; the collisionless effects and the decoupling effects of ions and neutrals on magnetic reconnection rate are small. The plasmoid cascading process terminates at a larger scale in an environment with higher plasma number and weaker reconnection magnetic fields.

3. In strongly magnetized solar TMR region, the plasmoid instability plays an important role in accelerating magnetic reconnection. The reconnection rate scales with the Lundquist number as \( M_{\text{sim}} \sim S_{\text{sim}}^{-1/2} \) before plasmoid instabilities appear. After plasmoid instabilities appear, the reconnection rate sharply increases to a high value (about 0.035), which is independent of the Lundquist number. These characteristics are very similar to magnetic reconnection in fully ionized plasmas. The recombination effect does not significantly affect the reconnection rate in this region, which is very different from the previous weakly ionized higher plasma \( \beta \) simulations (e.g., Leake et al. 2012, 2013; Alvarez Laguna et al. 2017).

4. A weakly ionized low chromosphere plasma with the neutral density of \( 10^{20} \text{ m}^{-3} \) can be heated from several thousand Kelvin to above \( 3 \times 10^{4} K \) during a magnetic reconnection process with magnetic fields of 500 G. We find that the much denser photosphere plasmas with the neutral density of \( 1.25 \times 10^{22} \text{ m}^{-3} \) are heated to only about \( 1.3 \times 10^{4} K \) with strong reconnection magnetic fields of 2500 G.

5. A stronger reconnection magnetic field and lower plasma density can both result in a higher maximum current density and ionization fraction in a reconnection process with initially weakly ionized plasmas. However, the maximum current density depends on the strength of the magnetic field more strongly than the plasma number density.

The numerical results in this work indicate that the characteristics of guide-field magnetic reconnection in the strongly magnetized low solar atmosphere (300–800 km above the solar surface) are close to that of fully ionized plasmas. The initially weakly ionized dense plasmas in the current sheet eventually become strongly ionized during the reconnection process. Then, the high ion density makes the ion–neutral collisional mean free path much smaller than the width of the current sheet, the decoupling of ionized and neutral flows is not obvious, and the significant increase of the magnetic reconnection rate by the recombination effect as shown in previous weakly magnetized high plasma \( \beta \) simulations does not appear. The length scales of the observed EBs and IBs are usually about \( 10^{3} \) m, which is two orders of magnitude larger than the characteristic length scale of \( L_{\text{ts}} = 10^{4} \text{ m} \) in Case III. Therefore, the Lundquist numbers in these reconnection events are above \( 10^{6} \), which is much higher than the critical Lundquist number predicted from this work. We conclude that the plasmoid instability instead of the recombination effect is the dominant mechanism that results in fast magnetic reconnection in strongly magnetized regions around the solar TMR. Though the ionized and neutral fluids are coupled well in the reconnection site, it does not mean that the one-fluid model is good enough to study magnetic reconnection in this region. Ni et al. (2018a) have shown that ionization–recombination non-equilibrium dynamics are very important for studying the evolutions of the ionization fraction.
and plasma temperatures in the reconnection process. Therefore, the reactive multifluid plasma–neutral model applied in this work is likely necessary for studying magnetic reconnection in this region. Since the width of the current sheet is also much larger than the ion inertial length inside the current sheet during the reconnection process, the collisionless effects on magnetic reconnection are negligible in this strongly magnetized region.

The simulations in this work were terminated not long after the simulations enter the plasmoid regime. Since the plasma temperatures inside the current sheet increase with time during the reconnection process, one may expect that the plasmas can be heated to higher temperatures if the simulations can last for a longer period. In future work, we will aim to extend the large-scale simulations further into the cascading plasmoid regime to study plasma heating during the saturation phase of the multiscale current sheet evolution.

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