Numerical Study of Magnetic Reconnection Processes in Solar Flares

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W. F. Ames (1965) said “fluid mechanics is especially rich in nonlinearities”. In this sense, I might as well say “Magnetohydrodynamics (MHD) is much richer”. Such rich nonlinearities bring about not only infinite research interest, but also enormous difficulty in mathematically treating them. Numerical simulation provides a powerful approach to overcome the difficulty and to investigate MHD processes of interest. This method differs from both the theoretical and experimental methods. Just like what M. A. Biot (1956) has said: In this field, there is at least as much artistry as science. However, it has become to be a more and more significant branch in MHD research. Its development will undoubtedly dispel some possible prejudices of some people against it step by step.

This thesis consists of the work I have done in simulating solar flares during recent years. To me, the solar flares are also erupting in the computers, as well as in the telescopes and on the Sun.

The text of this thesis is compiled by \LaTeX in English for convenience. I am sorry for not writing in my native language—Chinese, which I love so deeply.

I would feel the greatest happiness if this study I present here may contribute to the development of solar physics, even only like a drop to the sea.

——— To the coming Solar Maximum XXIII

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Extended Abstract

The solar flare represents a sudden release of energy ($10^{29} - 10^{33}$ erg within 100–1000 s), which is accompanied by plasma heating, radiations at multiwavelengths, and particle acceleration. It is of great interest not only to solar physics research, but also to space science, and solar-terrestrial study. Based on the early researches both in observations and in theory, magnetic reconnection model for solar flares was proposed, which was further supported by the Yohkoh observations. This thesis is aimed to study the magnetic reconnection processes in the solar flares by numerically solving magnetohydrodynamic (MHD) equations.

Since the MHD equations are highly nonlinear, it is very difficult to treat them analytically, and only a few exact solutions for reconnection are known. Part of the efforts are put into linear analysis, which is limited by the simplifying assumptions. To detailedly study the MHD problems, people have to resort to numerical simulation.

This thesis consists of the following 7 chapters:

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Chapter 1. Introductory review

The MHD numerical studies of magnetic reconnection in solar flares in the past thirty years are reviewed, with a small part describing the numerical methods used in the simulations, most part is devoted to introduce the progress in numerical results. In this chapter, the numerical simulations are classified into three categories, i.e., checking theoretical models, explaining observations, and studying the effect of some factors on magnetic reconnection. Then three crucial points are addressed, i.e., numerical instability, numerical resistivity, and driven/spontaneous magnetic reconnection. Finally, the aims of this thesis are put forward.

Chapter 2. Numerical procedure

This chapter depicts the equations for the 2.5-dimensional resistive MHD process, and how they are transformed into the dimensionless form. The multistep implicit scheme adopted for numerically solving the MHD equations is described briefly. Finally we define a physical quantity, the magnetic reconnection rate $R$.

Chapter 3. Pseudo-reconnection in MHD simulations

Symmetry boundary appears widely in numerical simulations. There are two methods to specify the symmetry conditions, which correspond to two different grid mesh systems, one with the boundary on the symmetry axis, the other with the symmetry axis centered between the first two columns of grid points. The simulations in this chapter illustrate that the first mesh system can produce physically acceptable numerical results, on the contrary, the second mesh system, i.e., the shifted mesh system, will introduce strong numerical resistivity, and lead to pseudo-reconnection. The first mesh system is strongly recommended.

Chapter 4. Two-ribbon flares and magnetic reconnection with heat conduction

The rise of the flaring loop and the separation of its two footpoints constitute one of the clearest signatures of magnetic reconnection in the two-ribbon flares. The results in this chapter do present such apparent motions. It is shown that a magnetic loop, with
heated plasma embedded in, is formed below the reconnection X-point. As the reconnection
proceeds, the field lines pile up and the bright loop rises, with its two footpoints separating.
As indicated by observations, the flaring loop motions are apparent rather than mass motions.
In fact, the reconnected field lines themselves shrink weakly. Our results are qualitatively
consistent with the CSHKP model for two-ribbon flares. However, we suggested that the
correlation between the rises of the loop and the reconnection X-point as mentioned in the
CSHKP model should be not valid. The bright loop can rise even if the X-point is fixed.

The role of the slow shocks in magnetic reconnection is discussed. The results are in favor
of the theory of Cargill & Priest (1982), i.e., slow shocks emanating from the reconnection
X-point contribute to the soft X-ray loop heating.

Finally, the effect of field-aligned heat conduction on magnetic reconnection is studied. It
shows that heat conduction may accelerate the magnetic reconnection. When the time-scale
of heat conduction is less than Alfvén time-scale, heat conduction dissociates the adiabatic
slow shock into an isothermal slow shock and a heat conduction front.

Chapter 5. Flaring loop motion and a unified model for solar flares

This chapter studies how the speed of the flaring loop motions, as seen in Chapter 4, is
influenced by the physical parameters. Many cases have been studied with different \( T_0, \rho_0, \)
magnetic field \( B_0, \) spatial length \( L_0, \) and resistivity \( \eta. \) It is found that the rise speed of the
loop and the separation speed of its footpoints are strongly dependent on \( B_0, \) and mediumly
on \( \rho_0, \) but weakly on \( T_0, L_0, \) and \( \eta. \) The dimensionless speeds (scaled by the Alfvén speed \( v_A) \) is around \( 4.4 \times 10^{-2}. \)

Extended simulations show that when the bright loop rises close to the reconnection
region, magnetic reconnection stops. The lifetime of magnetic reconnection is proportional
to the height of the reconnection X-point. When the X-point is high, the reconnection lasts
long, and reveals many features comparable to the two-ribbon flares as detailedly described
in Chapter 4. On the contrary, when the X-point is low, the reconnection is slowed down
after the impulsive phase, presenting a rather stable single loop configuration, which shows
a large similarity to the compact flares. Therefore, in this chapter, we put forward a unified
model for solar flares, where the different heights of the X-point lead to the bifurcation.

Chapter 6. Type II white-light flares and magnetic reconnection in low solar
atmosphere

This chapter attempts to explain Type II white-light flares by the low-layer magnetic
reconnection in four aspects: the lifetime \( \sim 10 \) min, the temperature rise near TMR \( \sim 150–
250 \) K, the radiation surface flux of the white-light kernel \( 1–2 \times 10^{10} \) erg cm\(^{-2} \) s\(^{-1} \), and the
required velocity profile to account for the red asymmetry of Ca II K line. Moreover, the effect
of ionization and radiation is investigated. It is shown that they both affect the magnetic
reconnection rate \( R \) weakly. However, they are effective mechanisms for cooling the heated
plasma.
Chapter 7. Summary
The numerical simulations performed in this thesis are summarized with a brief discussion.
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Chapter 1

INTRODUCTORY REVIEW

1.1 MHD Numerical Simulation of Magnetic Reconnection in Solar Atmosphere

Magnetic Reconnection is developed from the resistive tearing-mode instability in magnetized plasma, which can efficiently convert a large amount of magnetic energy into thermal and kinetic energies, and can also accelerate particles into high energy. Since it was put forward in the 1940s (Giovanelli 1946), the magnetic reconnection theory has been widely applied to account for a lot of eruptive phenomena in astrophysics (e.g., Parker 1979), interplanetary space (e.g., Vasylleunas 1975), and laboratory plasma (e.g., Hosea 1971), such as the accretion disk, magnetic tail, and tokamak device. Particularly for the Sun, from the cancelling magnetic features in the lower atmosphere (Wang & Shi 1992) to the coronal heating (Priest et al. 1998), from the small scale bright points to the large scale solar flares and CMEs, the magnetic reconnection plays an important role in these events. Lots of efforts are put to study the magnetic reconnection, such as theoretical analysis, laboratory experiments (Van Hoven 1976), and numerical simulation.

In the cases with dynamical length-scale greater than the ion Larmor radius and time-scale greater than the ion cyclotron period (sometimes they are not necessary), the dynamics of the magnetized plasma can be described by the magnetohydrodynamics (MHD) equations. Due to the high nonlinearity of the equations (especially with resistivity), it is very difficult to treat them analytically, and only a few exact solutions for reconnection are known. Otherwise, some models represent the diffusion region as a cut in the complex plane (e.g., Syrovatsky 1971), or they assume that the upstream magnetic field is potential and describe the field at the entrance to the diffusion region and shocks by a series of monopoles (e.g., Petschek 1964). Some other models use a linearization of the MHD equations in powers of the reconnection rate and so become less valid at high reconnection rates. Numerical simulation, however, can overcome these problems (see Jardine 1991 for a review). In fact, in many areas of theoretical physics, analytical theories and numerical experiments go hand in hand, complementing one another and together providing a much deeper understanding of a problem than either one alone (Priest & Forbes 1992).

Since 1970, people have begun to study magnetic reconnection by MHD numerical simulation (e.g., Burn 1970). Much progress has been made thereafter in both the numerical methods and the numerical results.
1.1.1 Progress in Numerical Methods

MHD numerical simulation can be regarded as solving an initial-and-boundary-value problem which is described by a set of partial differential equations. There are many numerical methods, such as the difference method, the finite element method, etc. Among them, the difference method is the most popular. Besides, some attempts are put at the spectral method (cf. Fu & Hu 1995).

Before performing simulation by the difference method, a numerical scheme should be appointed, such as the widely-used two-step Lax-Wendroff scheme (Hayashi & Sato 1978; Yokoyama & Shibata 1998), ADI scheme (Schanck & Killeen 1980), and so on. Based on the ADI scheme, Hu (1989) developed a “multi-step implicit scheme”, which is suitable for low $\beta$ (the ratio of gas to magnetic pressures) cases. Runge-Kutta method is also visible in some papers (Suzuki et al. 1997; Wang et al. 1997). Much attention has been paid to the physical models, and a lot of new schemes are seldom used in the numerical simulations of magnetic reconnection.

The development of the computational technology is reflected in two aspects: one is the high-resolution mesh, from $83 \times 72$ in 1970s (Hayashi & Sato 1978) to $440 \times 340$ at present (Yokoyama & Shibata 1998), the other is the three-dimensional simulations (e.g., Suzuki et al. 1997).

The finite-difference may introduce numerical diffusion, whose accuracy increases with smaller grid-spacing. Simulation with high-resolution mesh can avoid some numerical instabilities and unphysical results, therefore, the numerical results are closer to the real case. Otherwise, for instance, numerical instabilities will limit the value range of some quantities, e.g., the magnetic Reynolds number $R_m$ can not be too large.

In the two-dimensional (including 2D and 2.5D) simulations, usually one kind of symmetry is presumed, e.g., translational invariance or rotational invariance, which the actual events seldom possess. Due to the complexity of the magnetic configurations in the solar atmosphere and the development of the theory of 3D magnetic reconnection, 3D numerical simulations have been tried in recent years. On the one hand, some results in the 2D simulations can not necessarily be obtained in 3D simulations; on the other hand, some unsolved problems in 2D may be settled in 3D simulations.

1.1.2 Progress in Numerical Results

During the past thirty years, a large amount of numerical simulations have been performed to study the magnetic reconnection occurring in the solar atmosphere. Roughly they fall into three categories: checking theoretical models, explaining observational features, and studying the effect of some factors on magnetic reconnection.

• Checking theoretical models

There are a lot of theoretical models related to magnetic reconnection, whose validity remains to be checked by further observations and numerical simulations.

(a) **fast magnetic reconnection**: Fast magnetic reconnection is a fundamental problem in the magnetic reconnection theory (Biskamp 1993). It is now generally accepted that solar flares are triggered by magnetic reconnection, with their time-scale as short as 100-1000 s. Observations also indicate that fast reconnection processes are universal in the solar activities, with time-scale about 10-1000 $\tau_A$, where $\tau_A$ is the Alfvén transit time
However, according to the classical theory, magnetic Reynolds number ($R_m$) in the solar atmosphere is very high, $R_m = \tau_D/\tau_A \sim 10^6$-$10^{12}$, where $\tau_D$ is the magnetic diffusion time-scale. In order to merge the enormous gap between the observational and the classical theory, “fast magnetic reconnection” is proposed. In another word, the magnetic reconnection rate ($R$) is scaled by the resistivity ($\eta$) in such a relation: $R \sim \eta^\alpha$, where $\alpha \gg 1$ or $\alpha = 0$. Petschek (1964) firstly put forward a fast reconnection mechanism. He pointed out that slow MHD shocks will emanate from the small diffusion region, and stand at the surface between the inflow and the outflow. Most of the energy conversion, from magnetic energy to thermal and kinetic energies, takes place at the shocks. It is such a fascinating model that Jardine (1991) thought one of the main aims of numerical simulations is to reproduce fast Petschek-like reconnection.

By performing numerical simulations, Ugai & Tsuda (1977) and Scholer (1989) showed that locally enhanced resistivity will produce Petschek reconnection. Hayashi & Sato (1978) and Sato & Hayashi (1979) thought that the driving inflow is essential to the fast reconnection, the physical conditions in the diffusion region, on the contrary, have little effect on the reconnection rate. However, in the simulations by Biskamp (1986), the fast reconnection does not occur even when the driving inflow is given.

In view of the diversity in numerical results, Priest & Forbes (1986, 1992) proposed a unified theory for magnetic reconnection, which emphasizes the importance of the inflow boundary conditions in determining the reconnection regime, of which Sweet-Parker’s, Petschek’s, and Sonnerup’s solutions are special cases. They explained some conflicting numerical result. Moreover, they pointed out that Petschek-like reconnection is likely to occur when the inflow boundary conditions are free and the fast reconnection has developed nonlinearly from the tearing mode instability. The validity of such a unified theory is testified by the numerical simulations of Yan et al. (1992), and the Petschek-type reconnection was identified later by many simulations, e.g., Yokoyama & Shibata (1994) showed that Petschek-like reconnection and the slow mode shocks indeed occur in the simulation with free boundaries and locally enhanced anomalous resistivity.

Though, the mechanism of the fast magnetic reconnection is still controversial. In the simulations in favor of the Petschek mechanism, the anomalous resistivity ($\eta$) should be large enough to maintain the slow MHD shock structure and fast reconnection. For instance, Biskamp & Weltr (1980) and DeLuca & Craig (1992) found that the reconnection rate $R$ is almost independent of $\eta$, say, $R \sim \eta^{1/4}$ in the latter. However, such a scale relation breaks down when $\eta$ is too small. Therefore Biskamp (1986, 1993) thought that the Petschek mechanism is not suitable for high magnetic Reynolds number (i.e., small $\eta$), and the Petschek-type reconnection is only one special case of magnetic reconnection rather than universal rule. On the other hand, as pointed out by Richard & Craig (1993), such a saturation of $R$-$\eta$ relation may be a physical characteristic of reconnection, or it is merely an artifact of lack of resolution.

Besides, the numerical study by Ugai (1986) shed light on the possibility of fast reconnection with small $\eta$. He introduced a disturbance resistivity in the static magnetized plasma, then he found the current sheet is shortened, and the current density is enlarged, which may initiate the plasma micro-instability to give rise to anomalous resistivity $\eta$ near the X-point. With the assumption that $\eta$ is proportional to the electron drift speed ($v_d$), he found that the reconnection rate $R$ is related to $\partial \eta/\partial v_d$, rather than $\eta$. 

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If so, fast reconnection can proceed even if \( \eta \) is very small. Yokoyama & Shibata (1994) also indicated that nonuniform resistivity (compared with uniform resistivity) and the ejection of magnetic island are in favor of the occurrence of fast reconnection.

(b) CSHKP (or Kopp-Pneuman) model: To account for the flaring loop rise and the separation of two H\(\alpha\) ribbons in two-ribbon flares, Carmichael (1964), Sturrock (1968), Hirayama (1974), and Kopp & Pneuman (1976) developed the so-called CSHKP model (or Kopp-Pneuman model). In this model, magnetic reconnection is thought to occur in the current sheet which is formed by filament eruption. This model is improved later by Cargill & Priest (1982), and Forbes, Malherbe & Priest (1989). Forbes & Priest (1982, 1983a) were the first to perform the complete 2D dynamical simulation on the two-ribbon flare eruption. In their simulations, the initial state is presumed to be a current sheet configuration in magnetostatic equilibrium, the resistivity is uniform and constant. The results showed that as the reconnection proceeds, the magnetic neutral line rises, and the closed magnetic loop rises accordingly as in the CSHKP model. They also presented a slow mode MHD shock in their results, which however, heats the small area near the loop top. In their simulations, the reconnection rate is consistent with the prediction by the Petschek mechanism, and deviates from that by the Sweet-Parker mechanism. Their further study indicated that the inflow of the magnetic reconnection in the corona is not enough to provide the necessary mass for the postflare loop (only 1/1000-1/100). Therefore, they concluded that the mass in the postflare loop comes mainly from chromospheric evaporation (Forbes & Priest 1983b).

The numerical simulation by Magara et al. (1996) showed that although single magnetic line may move downward, the apparent rise motion is found by the continual pile-up of reconnected field lines. And the speed of the apparent motion decreases with time.

Yokoyama & Shibata (1997, 1998) coupled field-aligned heat conduction into the magnetic reconnection simulation and found that a cusp structure is seen below the reconnection region. The cusp structure is widely observed by Yohkoh satellite. For the first time in 2D simulations, they obtained the chromospheric evaporation incurred by the heat conduction front, which presents a strong support for the CSHKP model, although the speed of the evaporating plasma is about one order of magnitude smaller than the observational values at the impulsive phase of a solar flare.

Besides the widely studied CSHKP model, some other models related to the magnetic reconnection are also investigated. For instance, by 2.5D numerical simulation, Hu et al. (1997) illustrated that the approximate conservation of magnetic helicity is valid for the magnetic reconnection process. They pointed out as well that magnetic reconnection in the lower solar atmosphere will lead to a transfer of magnetic helicity to the upper corona.

- Explaining observational features

The solar atmosphere is mainly controlled by the magnetic field, and a lot of activities may related to magnetic reconnection. In fact, many observational evidences are found for the existence of the magnetic reconnection. With appropriate initial and boundary conditions prescribed, one can simulate the following evolution and compare it with observations. Based on the comparison, it can be investigated how the magnetic reconnection can account for the observations in the shape, the speed, the lifetime of the events, and some other dynamic features.
(a) Chromospheric Bursts: The chromospheric bursts include spicules, surges, UV microflares, and so on. By numerical simulation, Karpen et al. (1995) showed that magnetic reconnection in the chromosphere can explain many observational features, such as intermittency and large velocities. The simultaneous appearance of oppositely directed flows, both in and transverse to the computational plane, is suggestive of the velocity pattern observed for some explosive events. They also suggested the explanation for the fact that both red- and blue-shifts can characterize each event, but there is a preponderance of blue-shifts by a 3:2 ratio.

(b) Soft X-ray Jet: One of the interesting discoveries of Yohkoh is coronal soft X-ray (SXR) jets, which were observed as transitory X-ray enhancements with an apparent collimated motion. Some of the jets are accompanied by Hα surges. Yokoyama & Shibata (1995) simulated a magnetic sheet in the solar convection zone, which produces the emerging magnetic flux due to the Parker instability. When the emerging magnetic loop rises up to the corona and interacts with the pre-existing anti-parallel field, magnetic field lines are cut and reconnection occurs. As a result, the plasma is heated to form the X-ray brightening, and is accelerated to high speed (∼100 km s⁻¹) by the magnetic tension force. Two different types of magnetic configurations (horizontal and oblique) in the corona correspond to the observed two types of SXR jets, respectively. In the oblique case the SXR jet is accompanied by the motion of cold plasma, i.e., Hα surge.

(c) CMEs: In recent years, a large number of papers are devoted to the relation between the helmet streamer, coronal mass ejections (CMEs), and the eruptive filament. Mikić & Linker (1994) studied the dynamical evolution of an isolated arcade which is subjected to a quasi-static shear motion in axisymmetric spherical geometry. They found that past a critical level of shear, the configuration experiences ideal MHD magnetic nonequilibrium, leading to an opening of the field and the formation of a tangential discontinuity in the magnetic field. Finite resistivity resolves the tangential discontinuity into a current sheet at which there is a rapid reconnection of the magnetic field, leading to the release of the magnetic energy, fast flows, and the ejection of large length-scale plasmoid. They pointed out that for the partially open helmet streamer configuration, a small amount of shear is required, and it is not necessary for the corona arcade to have dimensions comparable to the solar radius for the disruption to occur, active-region-sized arcade may be enough. The multi-current-sheet inside the CME-related streamer and the role of emerging flux or of a current in CMEs eruptions were also investigated in detail (e.g., Guo, Wu, and Tandberg-hanssen 1996).

The numerical simulations of plasmoid ejection (Magara et al. 1997; Zhang, Wang & Zheng 1995) and coronal bright points (Waldron & Mullan 1987) revealed a large similarity to the corresponding observational features.

- Studying the effect of some factors

Researches in this aspect are meaningful only when coupled with the previous two aspects. On the one hand, by changing some terms in the MHD equations or the values of some parameters, the required conditions to satisfy the observations can be deduced; on the other hand, since numerical simulations are repeatable and ‘cheap’ experiments, one can do a large quantity of simulation to find some rules in the magnetic reconnection, which are helpful to improve theoretical models and to direct the further observations.
Forbes & Malherbe (1991) investigated the effect of radiation on magnetic reconnection. They found that the compressibility strongly affects the condensation in the radiative reconnection mode. Since the important parameter range, which is crucial to the trigger of radiative tearing mode, is difficult to simulate in 2-3 dimensions, they had to modify the optical-thin radiation formula to be suitable for their simulation.

Magara, Shibata & Yokoyama (1996) studied the parameter dependence of the eruption process in solar flares. Their simulations indicated that if the speed and the time are scaled by Alfvén speed and Alfvén transit time respectively, the dimensionless results are almost independent of $\beta$, while the density inside the current sheet strongly affects the ejection speed of the plasmoid.

Since heat conduction in the magnetized plasma is field-aligned and therefore highly nonlinear, the effect of heat conduction on 2D magnetic reconnection has been seldom studied. Recently, Yokoyama & Shibata (1997) studied the 2D magnetic reconnection coupled with field-aligned heat conduction. It was shown that when the time-scale of heat conduction is smaller than the Alfvén transit time, an adiabatic slow MHD shock is dissociated into an isothermal slow shock and a heat conduction front. Their results also showed in their case, heat conduction has very weak effect on the magnetic reconnection rate and the magnetic energy releasing rate.

1.1.3 Discussion

With the development of computer technology, scientific researchers become more and more interested in numerical simulations. Several kinds of simulations are performed in different time-scale and length-scale, such as MHD simulation, particle simulation, and hybrid simulation. Numerical simulation has become one separate research method from theoretical analysis and laboratory experiments. It has been interesting perhaps not so much for the ways in which it has agreed with the analytical results, but in several ways in which the results have been conflicting.

(1) Numerical Instability

Numerical simulation treats the interesting region by finite discretized numerical mesh. For one thing, it will introduce computational boundaries. Secondly, the finite difference may give rise to truncation error. Therefore, numerical instability is universal in any differencing form (Matthaeus & Montgomery 1981), particularly for long-term simulation. The numerical instability is due to that the Kolmogoroff dissipation scale-length $\lambda_\eta$ becomes less than the numerical grid spacing, where $\lambda_\eta = (d\epsilon/dtR_m^3)^{-1/4}$, $d\epsilon/dt$ is the rate of magnetic energy dissipation, $R_m$ is the magnetic Reynolds number. To avoid the instability, $R_m$ should be large enough (compared to the classical Spitzer type) in numerical simulation. In fact, some large resistivity models are suggested, e.g., hyperresistivity, turbulence resistivity (see Karpen et al. 1995).

(2) Numerical Resistivity

The differentiation of the MHD equations will introduce numerical diffusion. For magnetic reconnection, numerical resistivity is prominent, which may lead to numerical reconnection. So it is important to dynamically debug the results. For instance, in the simulations of driven reconnection by Hayashi & Sato (1978), they checked their numerical mode by setting the resistivity to be zero before formal simulation, and the results show that the driven field lines only accumulate near the magnetic inversion line without reconnection, which is physically consistent. Generally, it is acceptable when the effect of finite resistivity is much greater than
that of numerical resistivity. On the other hand, the numerical simulation means that such a reconnection would occur if the finite resistivity is introduced the same as the numerical resistivity. For example, Karpen et al. (1995) simulated magnetic reconnection incurred by numerical resistivity. However, the numerical resistivity is determined dynamically, and not controllable.

(3) Driven and Spontaneous Reconnections

The driven and spontaneous reconnection are distinguished in some literature. However, Biskamp (1986) pointed out, such a division is overemphasized, these two types are the same in nature. Besides, some driving inflows are too artificial. The meaningful driving process should be introduced self-consistently (Yokoyama & Shibata 1994), e.g., Pritchett & Wu (1979) and Richard & Craig (1993) investigated the driving process produced by the coalescence of magnetic islands; Yokoyama & Shibata (1994, 1995) studied the emerging magnetic flux produced by Parker instability; and Mikić & Linker (1994) simulated the driving force over coronal magnetic field by the photospheric shear or twist motions. Besides, shock waves initiated by nearby active region might become a driving force (e.g., Odstrcil & Larlicky 1997); the eruptive ejections of plasmoids (magnetic islands) and of filaments will decrease greatly the pressure below the eruption, therefore the outside plasma are driven by the strong pressure gradient.

1.2 Study in This Thesis

The solar flare represents a sudden release of energy, which is accompanied by plasma heating, radiations from wide wavelengths, particle acceleration, and so on. Such a magnificent while mysterious phenomenon is of great interest to both the observational and theoretical researchers due to its diversity and its complexity. This thesis is aimed to numerically study the MHD processes during magnetic reconnection and to show how solar flares can be explained by magnetic reconnection. The ensuing chapters are organized as follows:

Chapter 2: Numerical Methods. In this chapter, MHD equations in both the original and the dimensionless forms are presented. It is also described how to numerically solve the dimensionless MHD equations. The initial magnetic configuration used in this thesis is presented as well.

Chapter 3: Pseudo-reconnection in MHD Numerical Simulations. This chapter compares the numerical results obtained in two different numerical mesh systems, which both are widely used. We intend to explore the best numerical mesh system suitable for MHD numerical simulations.

Chapter 4: Two-ribbon Solar Flares and Magnetic Reconnection with Heat Conduction. This chapter studies the magnetic reconnection process similar to that in the Kopp-Pneuman model for two-ribbon flares. Both the adiabatic and the conductive cases are simulated to check how two-ribbon flares can be explained by magnetic reconnection, and to what extent the Kopp-Pneuman model and its later revised versions are correct.

Chapter 5: Flaring Loop Motion and a Unified Model for Solar Flares. Based on the observational fact that the usually classified two types of solar flares may share the same model, we performed 2 dimensional numerical simulations to reproduce their common feature—bright loop, and further explore the possibility that the two types of solar flares can be unified into the same model.

Chapter 6: Type II White-light Flares and Magnetic Reconnection in low solar
atmosphere. The simulations in this chapter are aimed to study the magnetic reconnection process occurring in the lower atmosphere, with emphasis on its application to Type II white-light flares.

Chapter 7: Summary. The numerical simulations in this thesis are summarized, some crucial points are discussed.
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Chapter 2

NUMERICAL PROCEDURE

2.1 Introduction to MHD Equations

In some circumstances, the magnetized plasma can be regarded as conducting fluid, and its dynamics can be described by the magnetohydrodynamic (MHD) equations, as mentioned in Chapter 1. Plasma consists of electrons, ions, and sometimes neutral atoms. Unless the thermodynamic equilibrium among these components breaks down, or the plasma changes in a discontinuous way, the multi-component plasma can be treated by the one-fluid model (Xu & Tang 1987). For cosmical plasma not in extreme cases, its evolution is governed by the following equations:

Continuity equation:
\[ \frac{d\rho^*}{dt^*} + \rho^* \nabla \cdot \mathbf{v}^* = 0, \quad \text{(2.1)} \]

Momentum equations:
\[ \rho^* \frac{d\mathbf{v}^*}{dt^*} = - \nabla P^* + \mathbf{j}^* \times \mathbf{B}^* + \mathbf{f}^*, \quad \text{(2.2)} \]

Energy equation:
\[ \rho^* \frac{d(\varepsilon^* + \mathbf{v}^* \cdot \mathbf{v}^*)}{dt^*} = \nabla \cdot (P^* \mathbf{v}^*) + \mathbf{E}^* \cdot \mathbf{j}^* + \mathbf{f}^* \cdot \mathbf{v}^* + Q^*, \quad \text{(2.3)} \]

Maxwell equations:
\[ \nabla \times \mathbf{E}^* = - \frac{\partial \mathbf{B}^*}{\partial t^*}, \quad \text{(2.4)} \]
\[ \nabla \times \mathbf{B}^* = \mu_0 \mathbf{j}^*, \quad \text{(2.5)} \]
\[ \eta^* \mathbf{j}^* = \mathbf{E}^* + \mathbf{v}^* \times \mathbf{B}^*, \quad \text{(2.6)} \]

where viscosity and radiative gain/loss are neglected and isotropic fluid is assumed; \( \frac{d}{dt} = \frac{\partial}{\partial t} + \mathbf{v}^* \cdot \nabla \); \( \rho^*, \mathbf{v}^*, \mathbf{B}^* \), and \( P^* \) are the density, the velocity, the magnetic field, and the gas pressure, respectively; Generally, the plasma is considered as perfect gas, i.e., \( P^* = \rho^* R_0 T^* \), where \( R_0 \) is the gas constant, \( T^* \) is the temperature; \( \mathbf{j}^* \) the current density, \( \mathbf{E}^* \) the electric field, \( \eta^* \) the resistivity, \( \mathbf{f}^* \) the external forces besides the electromagnetic force and internal friction; \( \varepsilon^* \) is the internal energy density; \( \gamma = 5/3 \) is the polytropic exponent; the quantity \( Q^* \) in Equation (2.3) is the heat conduction function.
2.2 Dimensionless Form of the MHD Equations

In numerical simulations, it makes convenience to solve the dimensionless equations. From Equations (2.1)–(2.6) it can be seen that the independent physical quantities are \( \rho^*, v^*, B^*, \) and \( T \), which are the functions of space (\( x^* \)) and time (\( t^* \)). We substitute them by \( \rho_0, v_0, B, T_0, L_0 \) and \( t_0 \), where \( \rho_0, v_0, B_0, T_0, L_0 \) and \( t_0 \) are characteristic values of the density, isothermal sound speed (\( v_0 = \sqrt{\gamma R_0 T_0} \)), magnetic field, temperature, length and time; \( \rho, v, B, T, x, \) and \( t \) are corresponding dimensionless quantities. Then the above MHD equations are reduced into

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \tag{2.7}
\]

\[
\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla)\mathbf{v} + \frac{1}{\rho} \nabla P - \frac{1}{\rho} \mathbf{j} \times \mathbf{B} - fL_0/(\rho \rho_0 v_0) = 0, \tag{2.8}
\]

\[
\frac{\partial \mathbf{B}}{\partial t} - \nabla \times (\mathbf{v} \times \mathbf{B}) + \nabla \times (\eta \nabla \times \mathbf{B}) = 0, \tag{2.9}
\]

\[
\frac{\partial T}{\partial t} + \mathbf{v} \cdot \nabla T + (\gamma - 1) T \nabla \cdot \mathbf{v} - 2(\gamma - 1) \eta \frac{\mathbf{j} \cdot \mathbf{j}}{\rho_0 \rho_0 \rho_0} - \frac{(\gamma - 1) L_0 Q}{\rho \rho_0 v_0^3} = 0, \tag{2.10}
\]

where dimensionless quantities \( \mathbf{j} = \nabla \times \mathbf{B} \).

In this thesis, we study only 2.5-dimensional (two-dimensional and three-component) problems. The Cartesian coordinates are taken such that, the \( x \) and \( z \) axes are parallel with the surface of the photosphere, while the \( y \)-axis is upward and perpendicular to the photosphere. All physical quantities depend only on \( x, y, \) and \( t \), i.e., \( \partial / \partial z = 0 \). Considering the gravity as the only external force, the MHD Equations (2.7)–(2.10) are cast into the following form

\[
\frac{\partial \rho}{\partial t} + v_x \frac{\partial \rho}{\partial x} + v_y \frac{\partial \rho}{\partial y} + \rho \frac{\partial v_x}{\partial x} + \rho \frac{\partial v_y}{\partial y} = 0, \tag{2.11}
\]

\[
\frac{\partial v_x}{\partial t} + v_x \frac{\partial v_x}{\partial x} + v_y \frac{\partial v_x}{\partial y} + \frac{\partial T}{\partial x} + \frac{2}{\rho \beta_0} \frac{\partial \psi}{\partial x} \Delta \psi + \frac{2 B_z \partial B_z}{\rho \beta_0} = 0, \tag{2.12}
\]

\[
\frac{\partial v_y}{\partial t} + v_x \frac{\partial v_y}{\partial x} + v_y \frac{\partial v_y}{\partial y} + \frac{\partial T}{\partial y} + \frac{2}{\rho \beta_0} \frac{\partial \psi}{\partial y} \Delta \psi + \frac{2 B_z \partial B_z}{\rho \beta_0} \partial y = 0, \tag{2.13}
\]

\[
\frac{\partial v_z}{\partial t} + v_x \frac{\partial v_z}{\partial x} + v_y \frac{\partial v_z}{\partial y} + \frac{2}{\rho \beta_0} \frac{\partial \psi}{\partial x} \partial B_z - \frac{2}{\rho \beta_0} \frac{\partial \psi}{\partial y} \partial B_z = 0, \tag{2.14}
\]

\[
\frac{\partial \psi}{\partial t} + v_x \frac{\partial \psi}{\partial x} + v_y \frac{\partial \psi}{\partial y} - \frac{1}{R_m} \Delta \psi = 0, \tag{2.15}
\]

\[
\frac{\partial B_z}{\partial t} + v_x \frac{\partial B_z}{\partial x} + v_y \frac{\partial B_z}{\partial y} + B_z \frac{\partial v_x}{\partial x} + B_z \frac{\partial v_y}{\partial y} - \frac{\partial v_x}{\partial x} \partial \psi + \frac{\partial v_y}{\partial y} \partial \psi = \frac{\partial}{\partial x} \left( \frac{1}{R_m} \frac{\partial B_z}{\partial x} \right) - \frac{\partial}{\partial y} \left( \frac{1}{R_m} \frac{\partial B_z}{\partial y} \right) = 0, \tag{2.16}
\]
\[
\frac{\partial T}{\partial t} + v_x \frac{\partial T}{\partial x} + v_y \frac{\partial T}{\partial y} + (\gamma - 1)T \frac{\partial v_x}{\partial x} + (\gamma - 1)T \frac{\partial v_y}{\partial y} - \\
\frac{2(\gamma - 1)}{\rho \beta_0 R_m} \left[ (\nabla \psi)^2 + \left( \frac{\partial B_z}{\partial x} \right)^2 + \left( \frac{\partial B_z}{\partial y} \right)^2 \right] - \frac{CQ}{\rho} = 0,
\]

where variables \(v_x, v_y, v_z\) are the three components of the velocity; \(B_z\) is the \(z\)-component of the magnetic field; \(\psi\) is the flux function, which is related to magnetic field by

\[
B = \nabla \times (\psi \hat{e}_z) + B_z \hat{e}_z.
\]

The quantity \(Q\) in Equation (2.17) is the dimensionless heat conduction function, with the form as

\[
Q = \begin{cases} 
0, & \text{in the case without conduction,} \\
\nabla \cdot \left( \frac{\nabla \psi}{B} \right), & \text{in the case with conduction.}
\end{cases}
\]

The dimensionless coefficients \(\beta_0\) (the ratio of gas to magnetic pressures), \(R_m\) (magnetic Reynolds number), \(g\) (the acceleration of gravity) and \(C\) (the heat conduction coefficient) are expressed as

\[
\beta_0 = \frac{2 \mu_0 \rho_0 v_0^2 L_0^2}{\psi_0^2},
\]

\[
R_m = \frac{1}{\eta} = \frac{\mu_0 v_0 L_0}{\eta^*},
\]

\[
g = \frac{L_0 g_s}{v_0^2},
\]

\[
C = \frac{(\gamma - 1) \kappa_0 T_0^2}{\rho_0 L_0 v_0^3},
\]

where \(R_m\) is scaled by the sound speed \((v_0)\) rather than the characteristic Alfvén wave speed \((v_A)\) (see Hu et al. 1997). \(g_s\) is the acceleration of gravity at the surface of the Sun. \(\kappa_0\), the heat conduction coefficient, is taken to be the classical Spitzer type, i.e., \(\kappa_0 = 10^{-11}\) W m\(^{-1}\) K\(^{-\frac{2}{3}}\). \(v_0 = \sqrt{RT}\), where \(R\) is the gas constant. \(\mu_0\) is the magnetic permeability. \(v_A = \psi_0 / L_0 (\mu_0 \rho_0)^{1/2}\). Parameters \(\rho_0, T_0, v_0, \psi_0, B_0, L_0\) and \(t_0\) are characteristic values of the density, temperature, velocity \((v_x, v_y, v_z)\), magnetic flux function, magnetic field, length and time, where \(\psi_0 = B_0 L_0\).

It is noted that the dimensionless Equations (2.11)–(2.17) are determined only by four quantities: \(\beta, R_m, g\) and \(C\). In the absence of the knowledge of resistivity \((\eta^*)\), the independent parameters are \(\rho_0, T_0, \psi_0, L_0\) and \(\eta\). So any given characteristic values are only one choice. In another word, the dimensionless results in the simulation can be applied to other physical conditions with different \(\rho_0, T_0, \psi_0, L_0\) and \(\eta\), only if the values given by Equations (2.20)–(2.23) remain unchanged.
2.3 Initial Conditions

Since $\beta$ (the ratio of gas to magnetic pressures) is very low in the solar corona, it is thought that the widely existing magnetic fields therein are force-free. In this thesis, similar to that given in Hu et al. (1997), the initial magnetic configuration is chosen as a force-free current sheet surrounded by a potential field, as shown in Figure 2.1,

$$\psi = \begin{cases} 
\frac{2w}{\pi} - \frac{2w}{\pi} \cos\left(\frac{\pi}{2w}\right), & (|x| < w), \\
\frac{2w}{\pi} - w, & (|x| \geq w), 
\end{cases} \quad (2.24)$$

$$B_z = \begin{cases} 
\cos\left(\frac{\pi}{2w}\right), & (|x| < w), \\
0, & (|x| \geq w), 
\end{cases} \quad (2.25)$$

where $w$, the half width of the current sheet, is set to be 0.1. The initial atmosphere is isothermal ($T=1$ everywhere) and in hydrostatic equilibrium, i.e., $v = 0$, $\rho = e^{(4-y)/g}$. Here $\rho$ at $y = 4$ is set to be unit.

2.4 Space Discretization

The computational region in this paper is a Cartesian plane, with the $y$-axis always being symmetry axis. The region is discretized into 41-61 uniformly spaced grid points along the $x$-axis (so as to guarantee enough grid points in the reconnection region) and uniformly spaced grid points along the $y$-axis. The grid spacing increases according to a geometric series along the $x$-axis.
2.5 Numerical Scheme

The numerical method used in this thesis is a multistep implicit scheme, which is suitable for strong magnetic field cases.

In numerical simulations, explicit schemes are preferred for transient phenomena. However, for long-term events (compared to the Alfvén time-scale), implicit schemes are helpful to procure a high stability. Among the implicit schemes, the alternating direction implicit method, or ADI method, is introduced by Peaceman & Rachford (1955), and Douglas (1955). This method uses a splitting of the time step to obtain a multi-dimensional implicit method which requires only the inversion of a tridiagonal matrix.

For the MHD simulations, the conventional ADI techniques seem passable, however, some ripples in the numerical results may still appear for a long-term computation. Therefore, Hu (1989) proposed a multistep implicit algorithm which is applicable to time-dependent MHD flows in a strong magnetic field.

This scheme is implemented in three steps, which are outlined as follows. Note that the first two steps correspond to the $x$- and $y$-direction split-implicit treatment of the MHD equations, similar to those steps in Lindemuth & Killeen (1973):

The set of MHD Equations (2.11)–(2.17) may be rewritten in the compact form

\[
\frac{\partial U}{\partial t} + W \left( U, \frac{\partial U}{\partial x}, \frac{\partial U}{\partial y}, \frac{\partial^2 U}{\partial x^2}, \frac{\partial^2 U}{\partial y^2} \right),
\]

where $U = (\rho, v_x, v_y, v_z, \psi, B_z, T)$ and $W$ are vectors, the seven elements of $W$ are omitted for conciseness.

(1) The first step: At the $n$-th time step and for any point $(x_i, y_j)$, Equation (2.26) is discretized into

\[
\frac{U_{ij}^{n+1} - U_{ij}^n}{\Delta t} + W_{ij}^{n+1} = 0,
\]

where

\[
W_{ij}^{n+1} = W^n + \left( \frac{\partial W}{\partial U} \right)^n [U_{ij}^{n+1} - U^n] + \left( \frac{\partial W}{\partial (\partial U/\partial x)} \right)^n \left[ \left( \frac{\partial U}{\partial x} \right)^{n+1} - \left( \frac{\partial U}{\partial x} \right)^n \right] +
\]

\[
+ \left( \frac{\partial W}{\partial (\partial^2 U/\partial x^2)} \right)^n \left[ \left( \frac{\partial^2 U}{\partial x^2} \right)^{n+1} - \left( \frac{\partial^2 U}{\partial x^2} \right)^n \right].
\]

Then, $U_{ij}^{n+1}$ is obtained.

(2) The second step: Equation (2.26) is discretized again,

\[
\frac{U_{ij}^{n+2} - U_{ij}^{n+1}}{\Delta t} + W_{ij}^{n+2} = 0,
\]

where
where
\[
W_{ij}^{n+2} = W_{ij}^{n+1} + \left( \frac{\partial W}{\partial U} \right)^{(n+1)} \left[ \mathbf{U}_{ij}^{n+2} - \mathbf{U}_{ij}^{n+1} \right] + \\
+ \left( \frac{\partial W}{\partial \left( \partial U / \partial y \right)} \right)^{n+1} \left[ \left( \frac{\partial U}{\partial y} \right)^{n+2} - \left( \frac{\partial U}{\partial y} \right)^{n+1} \right] + \\
+ \left( \frac{\partial W}{\partial \left( \partial^2 U / \partial y^2 \right)} \right)^{n+1} \left[ \left( \frac{\partial^2 U}{\partial y^2} \right)^{n+2} - \left( \frac{\partial^2 U}{\partial y^2} \right)^{n+1} \right].
\] (2.30)

Then, \( \mathbf{U}_{ij}^{n+2} \) is obtained.

(3) The third step: The numerical results in the new time step \( U_{ij}^{n+1} \) is calculated by taking an arithmetic average between \( U_{ij}^n \) and \( U_{ij}^{n+2} \), i.e.,
\[
U_{ij}^{n+1} = \left( U_{ij}^n + U_{ij}^{n+2} \right) / 2. 
\] (2.31)

Besides, artificial diffusion is introduced into all the MHD equations except that of \( \psi \), as indicated by Hu (1989). Note that the multistep implicit scheme is not an energy conserved scheme. There may be an error up to 10%~20% in energy conservation.

2.6 Magnetic Reconnection Rate

The magnetic reconnection rate (\( R \)) is a fundamental quantity in magnetic reconnection. When reconnection occurs, there will be a neutral reconnection point (X-point). The reconnection rate \( R \) is defined by the closing rate of the field lines at the X-point, i.e., \( R = d\psi_n / dt \), where the subscript “n” means it is evaluated at the neutral point. From Faraday’s equation, the electric field at the X-point, \( E_n \), is related to \( R \) by
\[
E_n - E_0 = -\frac{d}{dt} \int_0^X B_x(0, y) dy = -\frac{d\psi_n}{dt} = -R, 
\] (2.32)

where \( E_0 \) is the electric field at the origin \((x = 0, y = 0)\). In the most parts of this thesis, the bottom of the computational region as in Figure 2.1 is a line-tying boundary. Therefore, \( E_0 \) equals to 0. Considering the general Ohm’s law, \( \eta j = E + v \times B \), and omitting the sign of \( E \), the relation between the reconnection rate \( R \), the electric field at the X-point \( E \), and the resistivity \( \eta \) is reduced:
\[
R = E_n = \eta j_n. 
\] (2.33)

Bibliography
Chapter 3

PSEUDO-RECONNECTION IN MHD SIMULATIONS

3.1 Introduction

In MHD numerical simulations, MHD equations are differentially solved in a finite region which is discretized into a numerical mesh. Due to the truncation error in the differencing form, some unphysical results may be introduced, e.g., numerical instability (Matthaeus & Montgomery 1981), etc. For the first time, we point out in this chapter that the widely used “shifted grid mesh” may introduce strong resistivity.

Grid meshes are classified into two types, i.e., Lagrangian and Eulerian meshes. Among the Eulerian meshes, two kinds of meshes can be distinguished in dealing with the symmetry boundary, one with the first line of grid points on the symmetry axis, the other with the symmetry axis centred between the first two lines of grid points. This chapter shows that physically acceptable results are obtained in the first kind of mesh (“normal mesh” afterward), while the second mesh (i.e., shifted mesh) will introduce strong numerical resistivity and lead to pseudo-reconnection. Sect. 3.2 describes the numerical procedure. Numerical results are presented in Sect. 3.3. We conclude in Sect. 3.4 with brief discussions.

3.2 Numerical Procedure

The dimensionless MHD equations adopted in this chapter are given in Chapter 2. For convenience and generality, the gravity and the heat conduction in the MHD Equations (2.11–2.17) are neglected, and then the only independent parameter in the equations is $\beta$ (the ratio of gas to magnetic pressures), which is chosen to be 0.1, a typical value for solar corona.

In the initial state, the plasma is isothermal and uniform ($T = 1$, and $\rho = 1$ everywhere). The simulated region is $-1 \leq x \leq 1$, $-4 \leq y \leq 4$. Due to the symmetry, simulation is made only in the first quadrant of the $x$-$y$ plane, i.e., $0 \leq x \leq 1$, $0 \leq y \leq 4$, where the magnetic inversion line coincides with one symmetry axis $x = 0$. The top ($y = 4$) and the right-hand side ($x = 1$) are treated as open boundaries, while symmetry conditions are applied to the left-hand side ($x = 0$) and the bottom ($y = 0$). The simulated region is discretized into 41 nonuniformly spaced grid points in the $x$-direction (over ten grid points are within the half current sheet) and 91 uniformly spaced grid points in the $y$-direction.
Figure 3.1: Two kinds of mesh systems used to treat the symmetry boundary at \( x = 0 \). a Model A, where symmetry axis lies at the first column of the grid points. The symmetry conditions are set by \( f_0 = f_2 \), where \( f = \rho, v_y, \psi, B_z, \) and \( T \); \( f_0 = -f_2 \), where \( f = v_x \) and \( v_z \); \( f_0 \) represents the physical quantity at the image grid points (\( \times \)). b Model B, where grid system is shifted to the left by half mesh. The symmetry conditions are set by \( f_1 = f_2 \), where \( f = \rho, v_y, \psi, B_z, \) and \( T \); \( f_1 = -f_2 \), where \( f = v_x \) and \( v_z \).

Generally, there are two methods for setting the symmetry conditions (Roache 1972), which correspond to two different kinds of mesh systems, i.e., the normal mesh system and the shifted mesh system. Taking the left-hand symmetry axis as an example, in the normal mesh, the first column of grid points are just on the symmetry axis, as shown in Figure 3.1(a). The boundary values of physical quantities at the axis are determined in the following way: after introducing the image grid points of the second column of grid points (which are marked by \( \times \)), interior-point differencing method is applied to the boundary points. In the shifted mesh, the first grid-node is shifted to the left by half mesh-spacing across the symmetry axis, which is centred between the first two columns of grid points as shown in Figure 3.1(b). Symmetry conditions are applied to the first two columns of grid points. In this chapter, numerical simulations are performed in two models, i.e., Models A and B. In Model A, the symmetry conditions at \( x = 0 \) are specified by the first method (normal mesh), while Model B corresponds to the second model (shifted mesh). In both models, the symmetry conditions at the bottom (\( y = 0 \)) are specified by the first method (normal mesh). The different numerical results between the two models will be compared in three cases.

### 3.3 Numerical Results

Firstly, we keep the localized anomalous resistivity (\( \eta \)) constant all the time, where \( \eta = 0.02 \cos(5\pi x) \cos(2.5\pi y) \) for \(|x| \leq 0.1, |y| \leq 0.2\), and \( \eta = 0 \) for elsewhere. The numerical results show that an X-type configuration appears at the origin \((x = 0 \) and \( y = 0)) in both models. Two symmetrical inflows move into the current sheet, the plasma therein is ejected after the field lines reconnect rapidly. Two jets are formed along the vertical direction. As shown in Figure 3.2, the time profiles of the magnetic reconnection rate \( (R) \) in Models A and B are almost coincident, with very weak relative deviation of \( \sim 5\% \) during the simulated \( 40 \tau_A \). Apart from the accordance in the reconnection rates, large difference is found in the
Figure 3.2: Time profile of reconnection rate $R$ in the two models when anomalous resistivity is kept all the time. The solid line corresponds to Model A, and the dotted line to Model B.

Figure 3.3: Distributions of the $\rho$, $v_y$, and $T$ in the two models along the $y$-axis at $t = 20 \tau_A$. The solid line corresponds to Model A, and the dotted line to Model B.

distributions of some physical quantities for the two models. Figure 3.3 gives the $y$-plots of the density ($\rho$), the $y$-component of velocity ($v_y$), and temperature ($T$). From the figure it can be seen that the discrepancy is overwhelming for $\rho$ and $T$.

Secondly, the above anomalous resistivity is introduced only before $t = 10\tau_A$, thereafter $\eta$ equals to 0. According to the MHD theory about frozen-in effect, the field lines will not reconnect after $t = 10 \tau_A$, but accumulate near the symmetry axis $x = 0$ due to the existing inflows. The results in Model A are well in agreement with the theoretical prediction, with its reconnection rate ($R$) falling to 0 rapidly after $t = 10 \tau_A$. Note that $R$ in Model A after $t = 10 \tau_A$ is not equal to 0 exactly due to the unavoidable numerical errors. However, the reconnection rate in Model B keeps limited value rather than 0 after $t = 10 \tau_A$, as shown in Figure 3.4.

Finally, we simulate an ideal MHD process ($\eta = 0$) with driving inflows. Symmetric inflows lie at $|x| = 1$, with a steady velocity distribution $v = -0.5v_0 \text{sgn}(x) \cos(2.5(y - 2)\pi)$, where $v_0$ is the sound speed, $1.8 < y < 2.2$, and $\text{sgn}(x)$ equals 1 at $x = 1$, and $-1$ at $x = -1$. The numerical results indicate that as predicted by the frozen-in theory, the field lines in Model A do not reconnect, merely accumulate near the $y$-axis. On the contrary, magnetic
Figure 3.4: Time profile of reconnection rate $R$ in the two models when anomalous resistivity is removed after $t = 10 \tau_A$. The solid line corresponds to Model A, and the dotted line to Model B.

Figure 3.5: Time profile of reconnection rate $R$ in the two models when the ideal magnetized plasma is driven by symmetric inflows. The dashed line corresponds to Model B. $R$ keeps zero in Model A all the time.

Figure 3.6: Magnetic configurations of the two models at $t = 20 \tau_A$. The left panel corresponds to Model A, and the right panel to Model B.
reconnection occurs rapidly in Model B. Figure 3.5 plots the time profiles of $R$ for the two models, where $R$ in Model A is equal to 0 exactly, while $R$ for Model B is very large. The magnetic configurations of Models A and B are illustrated in Figure 3.6. In this figure, it is clearly seen that the field lines in Model A just accumulate near the $y$-axis. In Model B, however, they are cut and reconnect. In particular, small magnetic islands appear in Model B.

### 3.4 Discussions and conclusions

In this chapter, we perform 2-dimensional MHD simulations, with emphasis on the effect of two different symmetry-boundary-treating methods on the numerical results, where the symmetry boundary is located at the magnetic inversion line. The two methods correspond to two kinds of mesh systems, denoted by Models A and B, respectively. Model A is to the normal mesh system, with its first column of grid points on the boundary, while Model B corresponds to the shifted mesh system, with its symmetry axis centred between the first two columns of grid points, as illustrated in Figure 3.1. Roache (1972) pointed out that for the numerical simulation of fluid-dynamics, these two methods are both acceptable mathematically, only that the second method gives a residual truncation error in the $v$-flux of $v$-momentum across the symmetry axis, which is only $O(\Delta x^3)$. However, numerical results in this chapter show that in the MHD simulations with high magnetic Reynolds number, pseudo-reconnection may be introduced when the shifted mesh system is used to specify the symmetry conditions at the magnetic inversion line: 1) the magnetic reconnection rate ($R$) keeps large in Model B even after all the resistivity terms in the MHD equations are removed from $t = 10 \tau_A$; 2) field lines in Model B reconnect in the ideal MHD process with driving inflows. According to the MHD frozen-in theory, magnetic reconnection can not occur when the magnetic Reynolds number ($R_m$) is infinite (i.e., $\eta = 0$). This is just the situation in Model A. However, pseudo-reconnection is present in Model B. Since the two models have the same numerical resolution with the only difference in treating the symmetry boundary conditions, the above nonphysical results in Model B arise mainly from the strong numerical resistivity at the magnetic inversion line, which is intrinsically processed by the shifted mesh system. In the shifted mesh system, the maximum of the magnetic flux function ($\psi$), which corresponds to the magnetic inversion line, lies between the first two columns of grid points, rather than on a grid point. Therefore the extreme point of $\psi$-$x$ profile can not be resolved by the grid, as shown in Figure 3.7, which leads to the “wave-cut-off” effect, i.e., the numerical results near $x = 0$, as represented by the dashed line, are always smaller than the actual distribution represented by the solid line. Such “wave-cut-off” effect gives rise to strong numerical magnetic diffusion, and leads to unphysical numerical results, e.g., the pseudo-reconnection (i.e., numerical reconnection). In this sense, the different distributions of $\rho$, $v_y$, and $T$ between Models A and B are also due to the unphysical results arising from the numerical resistivity in the shifted mesh system.

Similarly, in any nonsymmetry case, such a “wave-cut-off” effect also exists when the magnetic inversion line does not lie just on one line of grid points. In this circumstance, unphysical results will be produced as well.

In summary, the following conclusions can be deduced by the numerical study in this chapter:

1. The shifted mesh system as in Model B, when used to specify the symmetry conditions
at the magnetic inversion line, will introduce strong numerical resistivity, and give rise to some unphysical results, such as pseudo-reconnection. One had better be careful while using the shifted mesh system.

(2) The first mesh system (i.e., the normal mesh system as in Model A) can effectively avoid the pseudo-reconnection, and therefore, is strongly recommended.

3) In any case of MHD numerical simulations, either symmetric or nonsymmetric, strong numerical resistivity will be introduced when the magnetic inversion line is not located on one column of grid points.

It should be noted that, if the magnetic flux function \( \psi \) is replaced by \( B_x \) and \( B_y \) in the magnetic diffusion Equation (2.9), the “wave-cut-off” effect discussed above will vanish, the unphysical results in Model B will reduce greatly. Besides, the discrepancy between the numerical results of Model B and Model A will become smaller when the numerical resolution is higher (e.g., more grid points near the symmetry line).

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Chapter 4

TWO-RIBBON FLARES AND MAGNETIC RECONNECTION WITH HEAT CONDUCTION

4.1 Introduction

The role of magnetic reconnection in solar flares is of great importance. In order to account for the phenomena observed in the two-ribbon flares, Carmichael (1964), Sturrock (1968), Hirayama (1974), and Kopp & Pneuman (1976) developed reconnection models in which a neutral sheet is considered as the preflare configuration. Based on the Yohkoh observations, Tsuneta (1993) also suggested that in active regions the global instability of the magnetic field prior to the flares creates the neutral sheet, and the subsequent magnetic reconnection plays a key physical role for the flare energy release. This model can explain some observational features, such as the rise of postflare loops and the separation of Hα two ribbons. However, the details need to be studied further.

Based on the models mentioned above, a lot of numerical simulations are performed to study the magnetic reconnection process, and in return, the models are improved by the numerical simulations (see Forbes & Acton 1996). Numerical simulations have shown many features comparable to those of observations, such as the postflare loops (Forbes & Priest 1982), cusp structures (Yokoyama & Shibata 1997), and plasmoid motions (Magara, Shibata, & Yokoyama 1997), etc. Some authors studied the parameter dependence of magnetic reconnection (e.g. Magara et al. 1996). However, there is still a long way to go before direct comparison between the numerical simulations and the observations.

Recently, Yokoyama & Shibata (1997) studied the effects of field-aligned heat conduction on magnetic reconnection. They showed that an adiabatic slow MHD shock is dissociated into a heat conduction front and an isothermal slow MHD shock. Another conclusion drawn by them is that heat conduction has weak effect on the reconnection rate. This may be due to the too far bottom boundary, which makes its wall effect on the outflow very weak. Besides, since in their case, the temperatures of the two outflows are symmetrical, there is only weak conduction across the reconnection region. In this chapter, we study further the magnetic reconnection coupled with anisotropic heat conduction, and some results are compared with corresponding adiabatic case. Section 4.2 describes the MHD equations, initial and boundary conditions. The numerical results are given in Section 4.3, some of them are
discussed in Section 4.4. Finally, conclusions are drawn in Section 4.5.

4.2 Numerical Method

In this chapter, two sets of parameters, which are denoted by Case A and Case B, are adopted to study the effect of heat conduction on magnetic reconnection respectively (as shown in Table 4.1). In Case A the initial timescale of heat conduction $t_{\text{cond}}$ is two orders of magnitude larger than the Alfvén timescale $\tau_A$, while in Case B, $t_{\text{cond}}$ is about five times larger than $\tau_A$. Hereafter, Model A0 means Case A without heat conduction, Model A1 means Case A with heat conduction. The same is to Case B.

Table 4.1. Characteristic values of physical quantities

<table>
<thead>
<tr>
<th>Case</th>
<th>$\rho_0$ (kg m$^{-3}$)</th>
<th>$T_0$ (10$^6$ K)</th>
<th>$v_0$ (km s$^{-1}$)</th>
<th>$\psi_0$ (Wb m$^{-1}$)</th>
<th>$B_0$ (10$^{-3}$ T)</th>
<th>$L_0$ (10$^4$ km)</th>
<th>$t_0$ (s)</th>
</tr>
</thead>
<tbody>
<tr>
<td>A</td>
<td>3.34 x 10$^{-11}$</td>
<td>1</td>
<td>128.6</td>
<td>7.45 x 10$^4$</td>
<td>3.72</td>
<td>2</td>
<td>155.5</td>
</tr>
<tr>
<td>B</td>
<td>0.334 x 10$^{-11}$</td>
<td>3</td>
<td>222.7</td>
<td>4.9 x 10$^4$</td>
<td>2.43</td>
<td>2</td>
<td>89.8</td>
</tr>
</tbody>
</table>

Then, in Case A, $\beta=0.1$, $v_A=575$ km s$^{-1}$, $\tau_A = L_0/v_A = 34.78$ s, the initial timescale of heat conduction $t_{\text{cond}} = L_0^2 n_0 k T_0^{5/2}$, where $n_0$ is the number density of gas, $k$ is the Boltzmann constant; in Case B, $\beta=0.07$, $v_A=1190$ km s$^{-1}$, $\tau_A = 16.8$ s, $t_{\text{cond}} = 77$ s.

As mentioned in Chapter 2, the characteristic values listed above are only one choice. In another word, the dimensionless results in the simulation can be applied to other physical conditions with different $\rho_0$, $T_0$, $\psi_0$, $L_0$ and $\eta$, only if the values given by Equations (2.20)–(2.23) remain unchanged.

The domain of our simulation is taken to be $-3 \leq x \leq 3$, $0 \leq y \leq 8$. The initial atmosphere is isothermal ($T=1$ everywhere) and in hydrostatic equilibrium, i.e., $v = 0$, $\rho = e^{(4-y)g}$. Here $\rho$ at $y=4$ is set to be unit. Note that the magnetic flux function $\psi$ in this chapter is multiplied a minus for some convenience.

In this chapter the resistivity is localized in a rectangular region $|x| \leq 0.1$, $3.75 \leq y \leq 4.25$, where $\eta/(\mu_0 v_0 L_0)$ equals $0.02 \cos(5\pi x) \cos(2(y-4)\pi)$. The purpose of adopting the two trigonometrical functions is to avoid discontinuity in the $\eta$ distribution so as to increase the numerical stability.

Since the problem is symmetrical about the $y$-axis, calculation was performed only in the right half region ($0 \leq x \leq 3$ and $0 \leq y \leq 8$). Therefore, the left-hand side ($x=0$) was treated as a symmetry boundary. The right-hand side ($x=3$) and the top ($y=8$) were treated to be open boundaries. An equivalent extrapolation was applied for all quantities at these two boundaries except that at the top, $\rho$ satisfies $\frac{\partial \rho}{\partial y} + g = 0$ (namely, the pressure gradient force compensates the gravity). The base ($y=0$) simulates the photosphere, where all quantities are fixed except $B_z$ and $T$, which are determined by equivalent extrapolation.

The computed region is discretized into 61 non-uniform grid points along the $x$-axis and 181 uniform grid points along the $y$-axis. The grid spacing increases according to a geometric series of common ratio 1.06 along the $x$-axis, and then 14 grid points are within the force-free
field region ($x < 0.1$). We coupled the heat conduction into the multistep implicit scheme (Hu, 1989), by which the MHD equations were numerically solved.

4.3 Numerical Results

4.3.1 General Description

The distributions of the temperature, the magnetic field and the velocity field for Model A1 in the $x$–$y$ plane at several times are shown in Figure 4.1 (only the part $|x| \leq 1$ is presented). After the localized resistivity is introduced, an X-type configuration is formed in the resistivity region near $y = 4$. The velocities of the upward and the downward jets increase rapidly to about Alfvén speeds. The jets are included by a pair of switch-off shocks. The downward jet collides with the line-tied magnetic loops, and a termination shock (Forbes & Priest 1983) appears at the loop top. The plasma near the outer flux loop is also heated. It means that a bright loop is formed. The height of the bright loop decreases rapidly until $t=10\tau_A$, when the loop top reaches its minimum height $y=0.9$. Thereafter, the bright loop rises almost uniformly with time, with its two footpoints separating. At the same time, the closed magnetic field lines below the loop top shrink weakly.

Other models show very similar process except a little difference. In Model A0, the loop top does not rise all the time; in Model B1, it is shown that the adiabatic slow MHD shock is dissociated into an isothermal slow shock and a conduction front. The heat conduction front is so fast that the numerical time step $\Delta t$ becomes too small after $t = 7.5\tau_A$. 

Figure 4.1: Evolutions of temperature (gray-scale), projected magnetic field (solid lines) and the velocity field (vector arrows) in Model A1 in the $x$–$y$ plane.
4.3.2 The Motion of the Flaring Loop

It is believed that the motions of the flare ribbons and loops constitute one of the clearest signatures of magnetic reconnection in the solar atmosphere, and the shrinkage effect of the reconnected field lines is a necessary property of any reconnection models (Forbes & Acton 1996). Indeed, such motions and the shrinkage appear in our simulation of magnetic reconnection.

As mentioned in subsection 4.3.1, a termination shock is formed at the loop top by the interaction between the downward jet and the closed magnetic loop. The loop top structure is characterized by the jumps in the distributions of the gas pressure, temperature ($T$), $v_y$ and $\partial^2 \psi / \partial y^2$, etc. Note that $\partial^2 \psi / \partial y^2$ contributes to $j_z$, the $z$-component of the current density. Figure 4.2 shows the profiles of $T$, $v_y$, and $\psi$ of Model A1 along the $y$-axis at two separate times. From it we can see that the loop top rises. During the simulated period, the loop top in Model A1 rises almost linearly with time, in Model A0, however, it sometimes drops down as shown in Figure 4.3. The fitted rise speeds of the loop tops in Models A0, A1 are 27 km s$^{-1}$ and 30 km s$^{-1}$, respectively.

Numerical results show that the plasma near the magnetic lines which link the loop top with the base is also heated to form a bright loop. The bright loop can be considered as the manifestation of SXR (soft X-ray) loop (H$\alpha$ ribbons are thought to be produced by the chromospheric heating at the both footpoints of the loop). Note that the two legs of the bright loop are on the inner side of the magnetic lines that link the loop top with the base due to the limited propagation speed of the heating mechanism as shown in Figure 4.1. As the loop top rises, the SXR loop rises simultaneously, and in the horizontal direction, the two bright footpoints of the loop separate. Figure 4.4 shows the horizontal distributions of the temperatures of Models A0 and A1 at several times (only the right half region is plotted),
where we can see the front-like structure propagates to the right. Since there is no effective cooling mechanism, the temperature of the plasma at the left side of the ‘front’ keeps high. Figure 4.5 shows the evolution of the footpoint position of the bright loop. For convenience, the position is determined by the coordinate of the plasma with \( T = 1.6 \). The fitted moving speeds of the leading loop footpoints in Models A0 and A1 are 31 km s\(^{-1}\), and 33 km s\(^{-1}\), respectively.

Shrinkage of the closed field lines is a phenomenon discovered in both solar flares (Švestka et al. 1987; Forbes & Acton 1996) and large-scale non-flare loop in the corona (Wang et al. 1997). Such effect is also clearly seen in our simulation. Figure 4.6 shows \( \psi \) (the flux function) distributions of Model A1 along the y-axis at two separate times. From the figure we see not only the rise of the loop top as indicated by the arrows (cf. Figure 4.2), but also the shrinkage of the closed field lines. It can be seen that the position with any given value of \( \psi \) below the loop top at \( t = 10\tau_A \) decreases at \( t = 20\tau_A \). Since one value of \( \psi \) corresponds to one field line, such height decrease means the shrinkage of the closed field lines below the loop top.

It should be noted that in our numerical results, the shrinkage is an integral effect. Figure 4.7 shows the time evolution of the vertical velocity at point \((0, 1)\) in Model A1. Because of the frozen-in effect, the velocity implies the motion of the closed field line at the point. The figure shows that the field line may also rise in one \( \tau_A \) timescale. This point needs to be checked by new observations.

The rise of the loop and the separation of the loop footpoints are not due to simple mass motion of plasma. In fact, they are formed by progressive appearance of newly-formed magnetic loop, which is consistent with observations (Bruzek 1964; Martin 1979; Engvold, Jensen, & Anderson 1979). The rise speed of the SXR loop and the separation speed of the ribbons are within the range of observational results (e.g. Nolte et al. 1979). Besides, the amount of shrinkage may depend on the geometry of the reconnecting field (Forbes & Acton 1996). How the motion of the SXR loops are affected by physical conditions will be discussed in the next chapter.
Figure 4.4: Distributions of temperature at the base for (a) Model A0, (b) Model A1.

Figure 4.5: Temporal variation of the footpoint position for Case A.
Figure 4.6: Local distributions of magnetic flux function ($\psi$) of Model A1 along the $y$-axis at $t=10\tau_A$ and $t=20\tau_A$.

Figure 4.7: Temporal variation of vertical velocity at the point (0, 1) in Model A1.
Figure 4.8: Temporal variations of the magnetic reconnection rate (\(R\)) for Case A and Case B. \(R = d\psi/dt\), where \(\psi\) is dimensionless, and the unit of \(t\) is \(\tau_A\).

4.3.3 Magnetic Reconnection Rate

Figure 4.8 compares the reconnection rates of the four models. The reconnection rate \(R\) of Model A1 is larger than that of Model A0 by 20%, and \(R\) of Model B1 is larger than that of Model B0 by 50%. It means that the heat conduction accelerates the reconnection, and such an effect is more significant in Case B than in Case A. It may be qualitatively explained as follows: When the reconnection goes on, part of the magnetic energy is converted into thermal energy, and the plasma temperature and the gas pressure near the reconnection region increase. The increased gas pressure will hinder the reconnection inflows so as to slow down the reconnection. If the heat conduction acts, the increased thermal energy can be taken away, so that the gas pressure decreases, and the reconnection evolves faster. Since the heat conduction timescale in Case B is much shorter than that in Case A, heat conduction is more effective in accelerating the reconnection.

4.4 Discussion

The numerical results in this chapter show some features similar to those obtained from observations and bear both similarities and differences compared with those of theoretical models (cf. Kopp & Pneuman 1976) and of published numerical simulations (e.g. Forbes & Priest 1983).

4.4.1 The Slow Shock and Its Role in Flaring Loop Heating

In the model of Kopp & Pneuman (1976), the loop is considered to be heated by gas-dynamic shocks. However, as pointed out by Heyvaerts, Priest, & Rust (1977) and Cargill & Priest (1982), the gas-dynamic shocks cannot heat the loop up to observed temperatures of \(5 \times 10^6\)–\(10^7\) K, and in a dynamical model, it should be replaced by slow MHD shocks. We have known that a pair of slow MHD shocks along the separatrix can effectively convert magnetic
energy into kinetic and thermal energy of the plasma in magnetic reconnection (Petschek 1964). Such shocks can be seen by the kinks of the magnetic lines along the jets shown in Figure 4.1.

Figure 4.9 shows the distributions of the temperature and the gas pressure in Case B at $t=7\tau_A$. In particular, its horizontal distributions of several quantities along the cut $y=1.6$ are presented in Figure 4.10(b). For comparison, $x$-plots of the counterparts in Case A along the cut $y=2$ at $t=10\tau_A$ are given in Figure 4.10(a). The slow MHD shocks, near which the velocity changes its direction and amplitude, are indicated by a significant increase of pressure as well as a decrease of magnetic field. From these two figures it can be seen that in Model B0, the slow shocks are adiabatic, while in Model B1, they are isothermal. The calculated propagation speed of heat conduction in Model B1 is several times larger than the local Alfvén speed, so adiabatic slow shock is dissociated into an isothermal slow shock and a heat conduction front, which is indicated by the broad high-temperature region in Figure 4.10(b). This does not occur in Model A1, whose slow shock is still adiabatic despite of heat conduction. Such effect was theoretically predicted by Forbes, Malherbe & Priest (1989) and numerically discussed by Yokoyama & Shibata (1997). The conduction front heats the plasma along the field lines, and may produce chromospheric evaporation if the chromosphere is included in the numerical region which is indicated by the numerical results of Yokoyama & Shibata (1998).

It is noticed that even in the models without heat conduction, such as Models A0 and B0, a bright loop is also formed. In such models, the heating mechanism is the slow MHD shocks, which trace the newly reconnected flux. As shown in Figure 4.1, a pair of slow MHD shocks move down along with magnetic lines until they reach the top of the loop system. Afterward, the field lines keep almost stationary except the shrinkage discussed in subsection 4.3.2, then the slow shocks propagate down along the field lines. Figure 4.11 shows the horizontal distributions of gas pressure ($P$), $y$-components of magnetic field and velocity field ($B_y$ and $v_y$) of Model A0 along $y=1.5$ cut at $t = 20\tau_A$. Slow MHD shock structures are roughly seen near $x = 0.6$ (the non-zero plateau of $B_y$ in the downstream of the slow shock.
Figure 4.10: Horizontal distributions of several physical quantities ($y$-component of velocity $v_y$, density $\rho$, temperature $T$, $y$-component of magnetic field $B_y$ and pressure $P$) (a) along the cut $y=2$ at $t=10\tau_A$ in Case A. (b) along the cut $y=1.6$ at $t=7\tau_A$ in Case B. All quantities are dimensionless except that $v_y$ is in km s$^{-1}$. 
Figure 4.11: Horizontal distributions of gas pressure ($P$), $y$-components of magnetic field and velocity field ($B_y$ and $v_y$) along the cut $y = 1.5$ across one leg of the bright loop at $t = 20\tau_A$ in Model A0.

means that the shock front is oblique). Since the propagation speed of the slow shocks is limited, the heated loop is not coincident with the flux loop that links the loop top with the base, but is on the inner side of it.

Based on the analyses above, we conclude that the pair of slow MHD shocks contribute to the heating of the SXR loop. This result is in favor of the theories of Heyvaerts et al. (1977) and Cargill & Priest (1982). However, it differs from the numerical result of Forbes & Priest (1983), who showed that such heating extends only a short distance near the loop top. This difference may result from that the reconnection point in their simulation was situated very low, and the slow shocks were weak, while in our simulation, the slow shocks can develop to be strong enough as the field lines, which convey the shocks, move from the reconnection site to the outer side of the loop system under the action of magnetic tension.

### 4.4.2 Is There Relation Between the Rises of the SXR Loop and of the Neutral Point?

During the rise of the flaring loops, the neutral point was believed to rise as the magnetic reconnection goes on (Kopp & Pneuman 1976; Tsuneta 1997). However, the numerical results in this chapter show that there is no direct relationship between them in our cases. For example, the neutral point in Model A1 keeps stationary at $y=4$ during all the simulated period, however, the SXR loop rises all the time. So our simulation shows that there is no direct relationship between the rise of the flaring loop and the rise of the reconnection point, although the neutral point may rise. The rise of the flaring loop is related to the rise of the loop top as described in subsection 4.3.2. The loop top corresponds to a position with $\partial^2\psi/\partial y^2$ maximum, which contributes to the $z$-component of the current density.

### 4.4.3 Other Points

We also note the following points:

1. In subsection 4.3.3, it is shown that the reconnection rate $R$ oscillates with an approximate period of $2\sim4$ times $\tau_A$. Such oscillation has also been found by Forbes & Priest (1982). They attributed it to wave reflection at the numerical boundaries. Here, we propose
a dynamical explanation, i.e., after resistivity sets in, $R$ increases so rapidly that the mass flux of the inflows is much more than that of the outflows, the plasma accumulates in the reconnection region, the increased gas pressure produced by which, in turn, inhibits the inflow, then reconnection is slowed down. When the mass flux of the inflows becomes less than the outflow, the inverse process occurs. Hereafter, the whole process repeats, and the oscillation period may be determined dynamically.

2. Observations show that initially the rise speed of the SXR loop and the separation speed of the footpoints are more than 50 km s$^{-1}$, but after several hours they decrease to less than 1 km s$^{-1}$ (Nolte et al. 1979). The corresponding speeds in this chapter are comparable with the initial ones of the observations. In actual solar atmosphere, the magnetic field is not vertically uniform, and the resistivity may be not constant, both of them will affect the rise speed of the SXR loop.

3. From Figure 4.10, it can be seen that in the models with heat conduction, the plasma accumulates in the downstream of the slow shock, however, in the adiabatic models, the plasma is concentrated around the shock front.

4.5 Conclusions

2.5–D magnetic reconnections are numerically studied in two cases. Case A, in which the timescale of heat conduction $t_{\text{cond}} \gg \tau_A$, may correspond to less intensive eruption process, while Case B, in which $t_{\text{cond}} \sim \tau_A$, may describe more violent eruption process. The reconnection starts after the anomalous resistivity is introduced into a hydrostatic solar atmosphere with a force-free magnetic configuration. The configuration with magnetic inversion line is similar to that of the classical model for long duration event (LDE). Recent research shows that the magnetic reconnection process of compact loop flares (or impulsive flares) may share the same model with the two-ribbon flares (or LDE flares) (Shibata et al. 1995), so some results in this simulation may be applied to the magnetic reconnection process occurring in both types of flares.

The main conclusions of this chapter are as follows:

1. After the anomalous resistivity is introduced, the magnetic reconnection starts, and two jets (upward and downward) appear. The downward jet collides with the closed field lines which are line-tied with the base, a termination shock (loop top) and a bright loop (SXR loop) are formed. As the reconnection goes on, the loop top, along with the SXR loop, rises almost uniformly, and the two footpoints of the loop separate. Besides the apparent loop motion, the closed field lines below the loop top shrink weakly. In this chapter, the rise speed of the SXR loop is $\sim 27$-30 km s$^{-1}$, the moving speed of the footpoints of the loop, which may correspond to the position of the H$\alpha$ 'ribbon', is $\sim 31$–33 km s$^{-1}$. In Case A, heat conduction increases these speeds weakly, but in Case B, the problem remains open since in Model B1, only the results before $t = 7.5 \tau_A$ are recorded.

2. It is shown that the rise of the SXR loop does not necessarily mean the rise of the neutral point (neutral line in 3-dimensional space). In our simulation, the loop rises even when the neutral point keeps almost stable.

3. Heat conduction increases the reconnection rate $R$. Such effect is more significant in Case B than in Case A, since the timescale of heat conduction in Case B is much shorter. Besides, $R$ oscillates with a period of about 2-4 times $\tau_A$.

4. Slow MHD shock contributes to the heating of the SXR loop. When the propagation
speed of heat conduction is greater than that of the slow shock, the adiabatic slow shock is dissociated into an isothermal slow shock and a conduction front.
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Chapter 5

FLARING LOOP MOTION AND A UNIFIED MODEL FOR SOLAR FLARES

5.1 Introduction

Solar flares are usually classified into two types, i.e., compact and two-ribbon flares, which were thought to require quite different physical mechanisms (Priest 1981). Generally, a two-ribbon flare is related to the eruption of a dark filament overlying a magnetic polarity inversion line. During the event, dense plasma loops rise in the corona for over 10 hours, with their footpoints separating. The chromospheric plasma is heated up to form two Hα ribbons. Such apparent motions are due to progressive appearance of new loops and continual cooling of old ones. The high density of the loops is caused mainly by the chromospheric evaporation. To account for these observational features, Carmichael (1964), Sturrock (1968), Hirayama (1974), and Kopp & Pneuman (1976) developed the so-called CSHKP model (or Kopp-Pneuman model) (cf. Forbes, Malherbe & Priest 1989 for a review). Such a model was supported by the Yohkoh observations (e.g., Tsuneta et al. 1992), and the chromospheric evaporation was also illustrated by the numerical simulation based on this model (Yokoyama & Shibata 1998). We performed 2.5-dimensional numerical simulation of magnetic reconnection (Chen et al. 1999), which did show the apparent motions of the soft X-ray (SXR) loop, as well as the shrinkage effect of the magnetic loops below the SXR loop. The shrinkage effect was indicated by observations (Švestka et al. 1987; Forbes & Acton 1996). In contrast with the two-ribbon flares, the compact flares do not show expansion, obvious motion, or shape variation (e.g., Švestka 1981). They are usually attributed to the emerging flux reconnection model, (Heyvaerts, Priest & Rust 1977). As a hard X-ray source above the SXR loop was discovered by Yohkoh, it is suggested that the reconnection process in compact flares might be the same as that in two-ribbon flares (Masuda et al. 1994). Such an idea is supported by more observational evidences (e.g., Tsuneta et al. 1997), and it is proposed that there is a need to unify the two types of flares into the same reconnection model (cf. Shibata 1998, and references therein).

Following Chen et al. (1999), this chapter further investigates the dependence of the speed of the flaring loop motion on several physical quantities. Most of the quantities can be measured or derived from observations. Besides, a possibility is tentatively suggested to
unify the two types of solar flares into the same reconnection model.

5.2 Numerical Method

The 2.5D compressible resistive MHD equations are shown in Chapter 2. We note here again the dimensionless coefficients $\beta_0$ (the ratio of gas to magnetic pressures), $R_m$ (magnetic Reynolds number), $g$ (the acceleration of gravity) and $C$ (the heat conduction coefficient) are expressed as

$$\beta_0 = \frac{2\mu_0 \rho_0 v_0^2}{B_0^2},$$  \hfill (5.1)\n
$$R_m = \frac{\mu_0 v_0 L_0}{\eta},$$  \hfill (5.2)\n
$$g = \frac{L_0 g_s}{v_0^2},$$  \hfill (5.3)\n
$$C = \frac{(\gamma - 1)\kappa_0 T_0^\frac{7}{2}}{\rho_0 L_0 v_0^3},$$  \hfill (5.4)\n
where $g_s$ is the acceleration of gravity at the surface of the Sun. $\eta$ is the resistivity. It is noticed that all the coefficients of the MHD equations depend only on $\beta_0$, $T_0$, $L_0$, $\rho_0$, and $\eta_0$. Table 5.1 presents our different models, in which these five quantities, along with the height of reconnection point, $h_0$, serve as free parameters. There are two different values for each parameter in these models, so that the effect of each parameter on the results can be estimated.

The domain of simulation is $-3 \leq x \leq 3$, $0 \leq y \leq 2h_0$, $h_0$ is the altitude at which the localized resistivity will be introduced. Anomalous resistivity is introduced into a local region: $|x| \leq 0.1$, $|y - h_0| \leq 0.2$, where $\eta = R_m^{-1} = \eta_0 \cos(\frac{2\pi x}{w}) \cos(2.5(y - h_0)\pi)$. Since the neutral point is always within the localized resistivity region, $h_0$ then roughly represents the height of the reconnection point (the X-type point).

Due to the symmetry, calculation is made only in the right half region. The bottom ($y=0$) is a line-tying boundary. The top ($y = 2h_0$) and the right-hand side ($x = 3$) are treated as open boundaries. Symmetry conditions are used for the left-hand side ($x = 0$). The numerical mesh consists of 61 nonuniformly spaced grid points along the $x$-direction (13 points are within the half current sheet) and 181 uniformly spaced grid points along the $y$-direction (so the numerical resolution in Model B is higher than that in Models A1–A6).

Table 5.1. Characteristic values of some physical quantities
In order to check the effect of the boundary and the numerical resolution on the results, Model A1 is also calculated in two other cases, one with right-hand side at \( x = 1 \) (the domain size in the \( x \)-direction is smaller, denoted as Model A1S), the other with higher numerical resolution (75 × 281 grid points, denoted as Model A1H), where the horizontal grid spacing near the \( y \)-axis is half that in Model A1.

### 5.3 Numerical Results

As described in detail by Chen et al. (1999), the oppositely-directed field lines beside the neutral line begin to reconnect after a localized resistivity is introduced in Model A1. The magnetic tension of the reconnected field lines drives rapidly two outflows (upward and downward jets) to high speed, accompanied by two symmetrical inflows. At the separatrix between the inflow and the outflow, a pair of MHD slow-mode shocks (switch-off shocks) develop as predicted by Petschek (1964). The included angle of the shock pair is about 3°. The upward jet, along with the frozen-in field lines, is ejected out of the top approximately with the Alfvén speed \( v_A \); the downward jet collides with the line-tied magnetic loop to form a termination shock (Forbes & Priest, 1983). Along the magnetic loop, both the temperature and the density increase to form an SXR loop. As the reconnection goes on, the SXR loop along with the loop-top (the termination shock) rises, with its two footpoints separating as shown in Figure 5.1. As discussed in Chen et al. (1999), such loop motions are due to progressive appearance of newly-heated loop, and the footpoints may correspond to the positions of Hα ribbons. During the time period of simulation (40 \( \tau_A \)), the loop-top and the footpoints show uniform motions very well, with the rise and separation speeds being about 30 km s\(^{-1}\). The dynamic patterns of the upper half part and the lower half part show strong nonsymmetry, which results from the different boundary conditions, i.e., the top is free, but the base is a line-tying boundary.

Besides the apparent rise and separation motions, the magnetic loops below the outer SXR loop shrink weakly. Such an effect was discovered in observations (Švestka et al. 1987;
Forbes & Acton 1996). Figure 5.2 shows the spatial change of the field line with $\psi = 1.2$ from $t = 20 \tau_A$ to $t = 30 \tau_A$. The amount of the shrinkage is about 5%.

Figure 5.3 presents for Model A1 the $z$-component of the current density $j_z$ and the $y$-component of the magnetic force $(j \times B)_y$ along the $y$-axis ($x=0$) at $t = 30 \tau_A$. The loop-top corresponds to the sharp peaks of both $j_z$ and $(j \times B)_y$, below which the magnetic field is approximately potential except near the base. Besides, it can be seen that the outflows are significantly accelerated by the Lorentz force just near the X-type point ($B=0$), where two extreme values of $(j \times B)_y$ are coincident with those of $j_z$.

Models A2–A6 show a very similar process. Table 5.2 presents the corresponding rise speeds of the loop-top ($v_r$) and separation speeds of the footpoints ($v_s$), which all are fitted by the method of least squares since the flaring loop shows approximately uniform motions before $t = 40 \tau_A$ (Chen et al. 1999). In Model B, whose only difference from Model A1 is that the reconnection point is low (also is the top boundary), the SXR loop also appears. However, in three aspects it differs from Model A1. Firstly, the slow shock, which is characterized by $B_y \sim 0$ in the downstream, is not so obvious as in Model A1. Secondly, the loop-top (termination shock) is evident only before $t = 8 \tau_A$. Thirdly, the rise speed of the SXR loop drops down from 16 km s$^{-1}$ at $t \sim 10 \tau_A$, to 5 km s$^{-1}$ at $t \sim 30 \tau_A$, and even to 1 km s$^{-1}$ at $t \sim 40 \tau_A$. The SXR loop keeps rather stable after a short time of drastic reconnection process as seen in Figure 5.4, which shows the two-dimensional evolution of temperature and magnetic field for Model B. Correspondingly, the magnetic reconnection rate ($R$) in Model B becomes smaller and smaller after an impulsive increase. The time profile of $R$ for Model B is compared with that for Model A1 in Figure 5.5. It is noted that $R$ for Model B shown in this figure is calculated from a new simulation with the same numerical resolution as in Model A1, i.e., with 61 $\times$ 46 grid points due to the one-fourth height. It should be pointed that the $R$ profiles for Model B with different resolutions show very small deviation. The significant
Figure 5.2: Magnetic line with $\psi = 1.2$ of Model A1 at two separate times.

Figure 5.3: $y$-plots of $j_z$ ($z$-component of the current density) and $(j \times B)_y$ ($y$-component of the magnetic force) along the $y$-axis ($x=0$) in Model A1 at $t = 30 \tau_A$. The solid line corresponds to $j_z$, and the dotted line to $(j \times B)_y$. 

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decrease of the $R$ in Model B can be explained as follows. As illustrated in Figure 3, the magnetic field just below the loop-top is approximately potential. When the rising loop-top is close to the resistivity region, a balance between the forces around the Y-point (as shown in Figure 5.4) is slowly reached, then the magnetic reconnection is decelerated naturally unless the anomalous resistivity is re-introduced at a larger height. In another word, the height of the closed field can not exceed that of the resistivity region. Similarly, the SXR loop motions in Model A1 will also slow down when the loop is close to its reconnection point. It is then meaningful to show how the lifetime of the flare scales with the height of the X-point, $h_0$, in our model. Three models, Models B, C, and D with $h_0=1$, $2$, $3$, respectively, are simulated with the same numerical resolution as Model A1. The lifetime for each case is given in Figure 5.6. Here the lifetime is defined as the duration in which the reconnection rate $R$ is greater than 0.01. It can be seen that the lifetime is almost proportional to $h_0$.

<table>
<thead>
<tr>
<th>Model</th>
<th>$v_r$</th>
<th>$v_r/v_A$</th>
<th>$v_s$</th>
<th>$v_s/v_A$</th>
</tr>
</thead>
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<tr>
<td>A1</td>
<td>31.0</td>
<td>5.38</td>
<td>31.2</td>
<td>5.42</td>
</tr>
<tr>
<td>A2</td>
<td>18.0</td>
<td>3.14</td>
<td>22.4</td>
<td>3.89</td>
</tr>
<tr>
<td>A3</td>
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<td>3.74</td>
<td>33.9</td>
<td>4.17</td>
</tr>
<tr>
<td>A4</td>
<td>31.6</td>
<td>5.50</td>
<td>32.1</td>
<td>5.59</td>
</tr>
<tr>
<td>A5</td>
<td>32.1</td>
<td>5.59</td>
<td>31.2</td>
<td>5.43</td>
</tr>
<tr>
<td>A6</td>
<td>46.7</td>
<td>5.74</td>
<td>45.3</td>
<td>5.57</td>
</tr>
<tr>
<td>A1S</td>
<td>31.1</td>
<td>5.40</td>
<td>31.5</td>
<td>5.48</td>
</tr>
<tr>
<td>A1H</td>
<td>30.2</td>
<td>5.26</td>
<td>31.6</td>
<td>5.49</td>
</tr>
</tbody>
</table>

Models A1S and A1H are computed to check the results of Model A1. In Model A1S, which has a smaller domain size in the $x$-direction, the reconnection rate $R$ becomes a little larger and oscillates with a shorter quasi-period. In this sense, the oscillation of the saturated reconnection rate in some numerical simulations (e.g., Chen et al. 1999) should have a boundary effect. However, the fitted rise speed of the flaring loop before $t = 40$ $\tau_A$, $v_r=30.6$ km s$^{-1}$, shows little difference from that of Model A1. In the higher numerical resolution case, Model A1H, both $R$ and $v_r$ are approximately equal to those of Model A1 with deviations less than 6% and 1%, respectively. It implies a convergence in our computations, and there is a convergence of the numbers in Table 5.2 with increased numerical resolution.

5.4 Discussion

5.4.1 Flaring Loop Motion

The CSHKP model was developed phenomenologically to account for the loop motions in two-ribbon flares. It should still be checked to what extent the model is reliable. Self-consistent numerical simulations, as well as further detailed observations may be helpful. For instance, we have shown numerically that (1) the SXR loop rises with its footpoints separating, which
Figure 5.4: Local distributions of temperature (grey-scale) and magnetic field (lines) in Model B at three separate times.

Figure 5.5: Time profiles of magnetic reconnection rate ($R$) in Model A1 (solid line) and Model B (dotted line). $R = d\psi_n/dt$, where $\psi_n$ is dimensionless, and the unit of $t$ is $\tau_A$. 

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Figure 5.6: Lifetimes of the flare for different X-point heights, $h_0$, where the unit of lifetime is $\tau_A = 35s$.

is thought to constitute one of the clearest signatures of magnetic reconnection; (2) the slow-mode MHD shock contributes to the SXR loop heating; (3) there is no direct relationship between the loop rise and the motion of the neutral point if the neutral point is far above the loop. The SXR loop may rise even if the neutral point keeps stable (Chen et al. 1999).

This chapter numerically studies the dependence of the speed of the loop motion on the characteristic values of several physical quantities. The characteristic values adopted here are typical ones in the corona except that the anomalous resistivity is somewhat arbitrary. Our results may be checked by observations.

From Table 5.2, it is deduced that, among the physical quantities $\beta_0$, $T_0$, $\rho_0$, $L_0$, and $\eta_0$, the speeds of the loop motion ($v_r$ and $v_s$) are strongly dependent on $\beta_0$ (Models A2 vs. A3) and $T_0$ (A1 vs. A6), weakly dependent on $\eta_0$ (A1 vs. A2). $L_0$ and $\rho_0$ have little effect on the speeds (A1 vs. A4, and A1 vs. A5). For clarity and simplicity, the results in Table 5.2 can be written as the scaling law $v \sim f^{\alpha_i}$, where $v$ represents $v_r$ and $v_s$, $f_i$ represents $\beta_0$, $T_0$, $\eta_0$, $L_0$ and $\rho_0$, $\alpha_i$ is the corresponding power index, which is determined by two sets of data ($v$, $f_i$). Note that since only two sets of data are presented in this study, such a scale relation is rather qualitative. We choose such a fitting method so that the parameter dependence of the velocities can be compared directly by the values of the power index.

(1) $v_r \sim \beta_0^{-0.76}$, $v_s \sim \beta_0^{-0.58}$; (2) $v_r \sim T_0^{0.16}$, $v_s \sim T_0^{0.54}$; (3) $v_r \sim L_0^{-0.03}$, $v_s \sim L_0^{-0.04}$; (4) $v_r \sim \rho_0^{0.02}$, $v_s \sim \rho_0^{0.02}$; and (5) $v_r \sim \eta_0^{0.23}$, $v_s \sim \eta_0^{0.15}$. The results are similar to those obtained in adiabatic MHD evolution (Chen et al. 1998a), except that in this chapter the effect of $T_0$ and $\eta_0$ is a little more important. The increased $T_0$ effect might be due to the heat conduction, while the increased $\eta_0$ effect might be due to the much smaller numerical resistivity in this chapter compared with that in Chen et al. (1998a). It is interesting to find that the dimensionless speeds ($v_r/v_A$ and $v_s/v_A$) are around $4.4 \pm 1.3 \times 10^{-2}$. During the quasi-steady state, $\Delta \psi \sim B_{\text{Loop}} \Delta y$, where $B_{\text{Loop}}$ is the dimensionless magnetic field.
near the SXR loop top, which is between 1 and 2 in Models A1–A6 before \( t = 40 \tau_A \). \( \Delta \psi \) and \( \Delta y \) are the reconnected flux and the spatial increase of the loop after time interval \( \Delta t \). Therefore, the reconnection rate \( R = \Delta \psi / \Delta t \sim B_{\text{Loop}} \Delta y / \Delta t \sim B_{\text{Loop}} v_r / v_A \), where the unit of \( R \) is \( \tau_A^{-1} \). It means that the dimensionless rise speed is related to the reconnection rate by \( v_r / v_A \sim R / B_{\text{Loop}} \). Our results then imply that \( R / B_{\text{Loop}} \) is around 4.4±1.3 × 10⁻², which is roughly in accord with the observed fast reconnection rate \( R \sim 0.001 - 0.1 \tau_A^{-1} \) (Dere 1996). Such an effect in the dimensionless velocity remains to be checked by observations in the early gradual phase of solar flares.

For convenience to see the parameter dependence, the above scaling laws may be written as

\[
v_r \sim \beta_0^{-0.76} T_0^{0.6} L_0^{-0.03} \rho_0^{-0.02} \eta_0^{0.23}, \quad \text{and} \quad v_s \sim \beta_0^{-0.58} T_0^{0.54} L_0^{-0.04} \rho_0^{0.15} \eta_0. \quad (5.5)
\]

Since \( \beta_0 \sim \rho_0 T_0 B_0^{-2} \) and \( \eta_0 \sim \eta L_0^{-1} T_0^{-1/2} \) are not original quantities, it is better to rewrite the scaling laws by

\[
v_r \sim B_0^{1.52} T_0^{-0.03} L_0^{-0.26} \rho_0^{-0.78} \eta^{0.23}, \quad \text{and} \quad v_s \sim B_0^{1.16} T_0^{-0.15} L_0^{-0.19} \rho_0^{-0.58} \eta^{0.15}. \quad (5.6)
\]

It indicates that, among the original physical quantities \( B_0, \rho_0, T_0, \eta, \) and \( L_0 \), the speeds are strongly dependent on the magnetic field \( (B_0) \), mediumly on \( \rho_0 \), and weakly on \( \eta, L_0, \) and \( T_0 \). The strong \( B_0 \) dependence can be explained by the change in the magnitude of the Lorentz force: stronger magnetic field leads to faster acceleration and therefore higher velocities; the moderate \( \rho_0 \) dependence may be due to the inertia of the plasma which is accelerated by the Lorentz force; the weak \( \eta \) effect may mean a fast reconnection; and the weak \( L_0 \) effect implies that the flaring loop motions have geometrical self-similarity, since the domain size, the width of the current sheet, and the height of the reconnection X-point are all scaled with \( L_0 \). Note that the power indices of \( T \) and \( \rho \) in Equation (5.6) are not so valid as those in Equation (5.5). However, our further simulations, where \( B_0, \rho_0, T_0, \eta, \) and \( L_0 \) serve as independent parameters, show the above conclusions are qualitatively right.

### 5.4.2 Possibility to Unify the Two Types of Solar Flares

Based on the fact that the usually classified two types of solar flares share many common features, e.g., the ejection of hot plasma, loop rise, X-type or Y-type morphology that suggests the presence of current sheets or neutral point, change of field configuration, and so on, Shibata (1997) proposed to unify the models for the two types of flares. The numerical results of this chapter evoke us to consider such a proposal.

In Model A1, where the reconnection point is high, a bright loop with high temperature and density is formed to appear as an SXR loop. The SXR loop keeps rising for a long time with its two footpoints separating, the magnetic loops below the loop-top shrink weakly, which are the typical features of two-ribbon flares. Note that the old SXR loops do not cool down rapidly to form Hα loops since radiation loss is not considered at present. On the contrary, in Model B, where the reconnection point is low, both the loop motion and the loop-top last only for a short time, the reconnection turns to be saturated naturally, and the rise speed of the SXR loop drops down rapidly so that the SXR loop seems rather stable, which shows a large similarity to compact flares.

In the CSHKP model for two-ribbon flares, the current sheet formed by the filament eruption is vertically long, and the reconnection point might be high; on the contrary, when
the newly-emerging magnetic flux from the photosphere interacts with the pre-existing oblique or nearly-vertical magnetic field in the atmosphere, the corresponding reconnection point may be low.

Here, we might as well suppose that the anomalous resistivity is caused by ion-acoustic current instability with the critical condition: \( v_d \sim B/(n_e l_0 e) \geq (kT_e/m_e)^{1/2} \), where \( v_d \) is the electron drift velocity, \( B \) the magnetic field, \( n_e \) the electron density, \( l_0 \) the characteristic width of the current sheet, \( e \) the electron charge, \( k \) the Boltzmann constant, \( T_e \) the electron temperature, and \( m_e \) the electron mass. From the formula, it can be deduced that the occurrence of anomalous resistivity in the lower atmosphere (with larger \( n_e \)) requires stronger magnetic field, and vice versa. Our model then implies that the reconnecting magnetic field is stronger for the compact flares, and weaker for the two-ribbon flares. Such a result is in agreement with that given by Shibata (1997, Table I therein). Because larger magnetic field may induce larger electric field, which accelerates the electron to high energy so as to form the hard X-ray source, it is then inferred that the compact flares should generally be stronger in hard X-ray emission than the two-ribbon flares.

It is noted that the shrinkage effect as shown in Figure 5.2 also appears in Model B. So our unified model further predicts that the shrinkage of the reconnected magnetic loop may also occur in the compact flares.

### 5.5 Conclusions

2.5-dimensional magnetic reconnection is numerically solved to study the flaring loop dynamics, with a force-free current sheet surrounded by a potential field as the initial configuration. The conclusions are as follows:

1. The loop motions (including loop rise, footpoint separation and shrinkage of the closed magnetic loop) are numerically shown self-consistently. Among the original physical quantities \( B_0, \rho_0, T_0, L_0 \), and \( \eta \), the rise speed of the loop and the separation speed of its footpoints are strongly dependent on the magnetic field \( B_0 \), mediumly on the density \( \rho_0 \), but weakly on the resistivity \( \eta \), the length scale \( L_0 \), and the temperature \( T_0 \). The strong \( B_0 \) dependence means that the Lorentz force is the dominant factor; the inertia of the plasma can account for the moderate \( \rho_0 \) dependence; the weak \( \eta \) dependence means the “fast-reconnection” occurs; while the weak \( L_0 \) effect may imply that the loop motions have geometrical self-similarity, since the domain size, the width of the current sheet, and the height of the reconnection X-point are all scaled with \( L_0 \).

2. It is tentatively suggested that the usually classified two types of solar flares can be unified into the same reconnection model. The essence is that the height of the reconnection point leads to the bifurcation, i.e., the reconnection point is high for the two-ribbon flares, and low for the compact flares. For our simplified model (vertically-uniform magnetic field, fixed local resistivity region, etc.), the lifetime of the flare is proportional to the height of the reconnection point.
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Chapter 6

TYPE II WHITE-LIGHT FLARES AND MAGNETIC RECONNECTION IN LOW SOLAR ATMOSPHERE

6.1 Introduction

Solar white-light flares (WLFs) are one of the strongest, also rare flaring events, with visible optical continuum. They are of great importance in flare research because they are similar in many aspects to stellar flares, and because they present a major challenge to the atmospheric models and energy transport mechanisms (Neidig 1989). It is suggested that there are two types of WLFs which show distinction in their emissions, i.e., Type I WLFs reveal the Balmer or Paschen jump, while Type II do not (Machado et al. 1986). Such distinction results from different continuum radiation mechanisms: hydrogen free-bound transitions for Type I WLFs, while negative hydrogen ($H^-$) radiation for Type II. Mauas, Machado, and Avrett (1990) firstly investigated the semiempirical atmospheric model for WLFs, which suggests that some white-light emission can take place in areas of the active region without chromospheric emissions and may correspond to the heating of the deep layer atmosphere. Further systematic studies on both the spectral characteristics and the atmospheric models for WLFs by Fang & Ding (1995) indicated that (1) for Type I WLFs, there is good time correlation between the maximum of continuum emission and the peaks of hard X-ray and microwave radiations. However, the correlation does not exist for Type II WLFs; (2) Type I are usually accompanied by strong and broadened hydrogen Balmer lines, which are absent in Type II; (3) for Type I, chromosphere is heated, but for Type II, the temperature of the photosphere increases greatly (e.g., hundreds of degrees) without obvious chromospheric heating. Those features for Type I are well explained by the conventional flare picture: energy is initially released in the corona, then is transported to the lower atmosphere and heats the plasma consecutively. However, for Type II, since the known mechanisms of energy transport are no longer effective (see Neidig 1989; Ding, Fang, & Yun 1999 for more references), an in situ heating mechanism deep in the photosphere is required.

Emslie & Machado (1979) and Mauas, Machado, and Avrett (1990) suggested that the required in situ heating near the temperature minimum region (TMR) may be due to the
local Joule dissipation of current. Recently, Li et al. (1997) proposed that Type II WLFs are produced by magnetic reconnection in a weakly ionized plasma. They try to account for some typical features of Type II WLFs: the length, the lifetime. However, their work is based on a linear analysis.

Magnetic reconnection is widely applied to explain many eruptive and/or heating events, particularly for the large flares occurring in the corona. Some chromospheric bursts are also thought to be incurred by the magnetic reconnection in the chromosphere (e.g., Karpen et al. 1995). From observations, more and more evidences are found for the magnetic reconnection occurring in the lower atmosphere (Wang & Shi 1993). In this chapter, we perform 2D numerical simulations to study the magnetic reconnection which occurs in the lower atmosphere, with emphasis on its application to Type II WLFs.

6.2 Numerical Method

For the magnetic reconnection in the deep atmosphere, ionization and radiation become important, while heat conduction is negligible, which is much different from the MHD process in the previous chapters. The weakly ionized plasma can as well be described by the one-fluid model, only that its ionicity changes with time. Another difference is the stratification, say, the maximum density is about 7 orders of magnitude larger than the minimum in this thin layer, which is hardly realizable in the present 2D simulations. In order to study the magnetic reconnection in this thin but complicated layer, we neglect the gravity and assume the plasma as uniform pure hydrogen atmosphere. The MHD equations are given as follows:

\[
\frac{\partial \rho^*}{\partial t^*} + \nabla \cdot (\rho \mathbf{v}^*) = 0, \tag{6.1}
\]

\[
\rho^* \frac{\partial \mathbf{v}^*}{\partial t^*} + \rho^* (\mathbf{v}^* \cdot \nabla) \mathbf{v}^* + \nabla P^* - \mathbf{j}^* \times \mathbf{B}^* = 0, \tag{6.2}
\]

\[
\frac{\partial \mathbf{B}^*}{\partial t^*} - \nabla \times (\mathbf{v}^* \times \mathbf{B}^*) + \nabla \times (\eta^* \nabla \times \mathbf{B}^*) = 0, \tag{6.3}
\]

\[
\frac{\partial}{\partial t^*} \left( \frac{P^*}{\gamma - 1} + n_e \chi_H + \rho \mathbf{v}^* \cdot \mathbf{v}^* \right) + \nabla \cdot \left[ \left( \frac{P^*}{\gamma - 1} + n_e \chi_H + \rho \mathbf{v}^* \cdot \mathbf{v}^* \right) \mathbf{v}^* \right] - \nabla \cdot (P^* \mathbf{v}^*) - \mathbf{E}^* \cdot \mathbf{j}^* + R - H = 0, \tag{6.4}
\]

where \(\rho^*, \mathbf{v}^*, \mathbf{B}^*,\) and \(P^*\) have the same meanings as in Chapter 2, however, \(P^* = (n_H + n_e)kT^*\), \(n_H\) and \(n_e\) are the number density of hydrogen atoms and electrons, respectively; \(\chi_H\) is the ionizing potential. \(n_e\) is deduced by modified Saha and Boltzman formula:

\[
n_e = \begin{cases} 
\frac{\left( \sqrt{\phi + 4n_H \phi} - \phi \right)^2}{n_H}, & T \leq 10^5 \text{ K}, \\
\frac{\left( \sqrt{\phi + 4n_H \phi} - \phi \right)^2}{n_H}, & T > 10^5 \text{ K},
\end{cases} \tag{6.5}
\]

where \(\phi = \frac{1}{b_1} \left( \frac{2\pi m_e kT}{\hbar^2} \right)^{3/2} e^{-\chi_H/kT}\), \(b_1\) is taken from Brown (1973).

Radiation is important for the lower atmosphere. Strictly speaking, it should be solved by the non-LTE theory, which is impossible at present for 2D simulations. Instead, it is substituted by the empirical formula given by Gan & Fang (1990),

\[
R = n_H n_e \alpha f'(T), \tag{6.6}
\]
where $f'(T) = 1.547 \times 10^{-23} \left( \frac{T}{10^6} \right)^{3/2}$, $\alpha(Z) = \alpha_1(Z) + \alpha_2(Z)$ is a function of $Z$, $\lg \alpha_1(Z) = 2.75 \times 10^{-3} Z - 5.445$, $\alpha_2(Z) = 2.3738 \times 10^{-4} e^{-Z/163}$, $Z$ is the height from the bottom of photosphere, whose unit is km. Since gravity is neglected, $\alpha$ in this chapter is set to be uniform correspondingly, which is done by fixing the value of $Z$, say, $Z = 0$. The radiative heating $H$ in Equation (6.4) is given by $H = n_e H$, where the constant $H = (n_e f')_{t=0}$.

Similar to Chapter 2, the MHD Equations (6.1)–(6.4) are rewritten in the corresponding dimensionless form with respect to $\rho$, $v$, $\psi$, $B_z$, and $T$. Since the density changes sharply in the deep atmosphere, three cases (Cases A, B, and C) are studied with different characteristic values as shown in Table 6.1. The density in Cases A and B is in the photosphere range, and in Case C, $\rho_0$ and $T_0$ are chosen from Table 3 of Fang & Ding (1995), which are the values for preflare chromosphere.

<table>
<thead>
<tr>
<th>Case</th>
<th>$\rho_0$ $(10^{-5} \text{ kg m}^{-3})$</th>
<th>$T_0$ (K)</th>
<th>$v_0$ (km s$^{-1}$)</th>
<th>$L_0$ $(10^3 \text{ km})$</th>
<th>$t_0$ (s)</th>
<th>$\beta_0$</th>
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<tr>
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<tr>
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<td>7860</td>
<td>11</td>
<td>2.5</td>
<td>219</td>
<td>0.3</td>
</tr>
</tbody>
</table>

In this chapter, Model A means Case A without ionization and radiation, Model AI means Case A with ionization only, and Model AIR means Case A with both ionization and radiation considered. The similar notations are used for Cases B and C.

### 6.3 Numerical Results

The numerical results show that the evolution of the low-layer magnetic reconnection is similar to the low-X-point case in Chapter 5. As the anomalous resistivity sets in, two symmetrical convergent inflows move into the diffusion region. Meanwhile, two narrow jets are ejected rapidly. Note that along the surface between the inflow and the outflow (jet), the field lines show no obvious kink, which is one characteristic of slow-mode MHD shock. As indicated in our previous results (e.g., Chapter 5 and Chen et al. 1999), the reconnected field lines above the reconnection point (X-point) are ejected along with the upward jet, its counterpart below the X-point, however, pile up due to the line-tying effect at the bottom, so that the closed magnetic loop system rises as shown in Figure 6.1, which depicts the evolution of the temperature, the velocity and the magnetic configuration in Model B. When the loop system becomes close to the resistive region, the magnetic reconnection is slowed down. The saturation of $R$ was discussed in Chapter 5 and Chen et al. (1999), which show that the saturation time-scale is proportional to the height of the X-point. Although the reconnection stops after $t \sim 10 \tau_A$, a global upward flow is seen. Such motion can be explained by the melon-seed effect of the magnetic field. Since the field lines are anchored to the bottom,
the magnetic field $B$ becomes larger after the reconnection, on the contrary, the magnetic field above the small reconnection region is $\sim 0$ near the $y$-axis, as seen in Figure 6.1. The pressure gradient of the magnetic field accelerates the frozen-in plasma to move upward with maximum speed $\sim 0.5$ times $v_A$.

Figure 6.2 depicts the time profiles of the magnetic reconnection rate ($R$). It can be seen that the reconnection goes off self-consistently. Taking $R = 0.01$ as the critical level as in Chapter 5, the lifetime of magnetic reconnection are $7 \tau_A$, i.e., $\sim 600$ s, which is shown to be independent of the ionization and the radiation. It is also found that both ionization and radiation have very weak, but positive effect on the reconnection rate in these models.

Compared to the energy releasing rate, the hydrostatic quantities (e.g., the density, the temperature, etc.) are much more sensitive to the ionization and the radiation. The distributions of $\rho$ and $T$ in Case C along the two axes are plotted in Figure 6.3. It can be seen that the ionization and the radiation cool down the plasma, and lead to the plasma condensation. For instance, the maximum temperature is 3.2 for Model C, 1.2 for Model CI, and 1.02 for Model CIR.

### 6.4 Discussion and Conclusions

More and more evidences show that magnetic reconnection may occur in the lower-layer solar atmosphere. However, it is not yet studied detailedly in theory. One of the difficulties is the sharp transition of some physical quantities. With the uniform isothermal atmosphere assumed, this chapter numerically studied the magnetic reconnection occurring in the weakly ionized plasma.
Figure 6.2: Time profiles of magnetic reconnection in the three cases. The solid line corresponds to the model without ionization and radiation, the dashed line to the model with ionization, and the dotted line to the model with both ionization and radiation.

Figure 6.3: Distributions of density and temperature along the horizontal line $x = 0.33$ and along the $y$-axes. The solid line corresponds to the model without ionization and radiation, the dashed line to the model with ionization, and the dotted line to the model with both ionization and radiation.
6.4.1 Effect of Ionization and Radiation

(1) The results in Sect. 6.3 indicate that ionization decreases the plasma temperature. This is the natural result of the energy conversion from the thermal one to the ionizing potential. In fact, the ionizing potential contributes to a great part of the total energy, since in the deep atmosphere, the ionizing potential of the hydrogen atoms $\chi_H$ is greater than the thermal energy, i.e., $\chi_H > (\gamma - 1)kT$. From a naive point of view, such an energy conversion will lead to the decrease of the total gas pressure, as implied by the first term in Equation (6.4). In this sense, ionization will accelerate the magnetic reconnection (i.e., a positive effect), as clearly shown in Case C. In Cases A and B such an effect is much weaker. This is due to that the parameters in these two cases are not in a ionization-sensitive regime. In Cases A and B, the ionicity will rise to only 1.6% when the plasma is heated up to $T=2$, however, the ionicity will be near 100% in Case C. Therefore, much of the thermal energy will be converted into the ionizing energy in Case C when the plasma is heated by magnetic reconnection, which leads to a stronger positive effect on the reconnection rate.

(2) To avoid the complexity of non-LTE calculation, the radiation term in Equation (6.4) is replaced by the empirical formula for radiative loss derived by Gan & Fang (1990).

As shown in Figure 6.3, the radiation acts as an effective mechanism for cooling the plasma and strengthen the condensation. The temperature in Model CIR is close to the unity. However, the radiation weakly increases the reconnection rate ($R$) in these models. We further increase $Z$ in Equation (6.5) to exemplify the coefficient of the radiation term by 3 orders, the reconnection rate, however, increases only by 10% for Case A, 3% for Case B, and $\sim 0$ for Case C.

6.4.2 Applicability to Explaining Type II WLFs

Both the spectral observation and semi-empirical atmospheric model for Type II WLFs imply that an in situ heating mechanism is required in the chromosphere or photosphere (Fang & Ding 1995). Our numerical results of the deep-layer magnetic reconnection show that Type II WLFs can be well explained by such magnetic reconnection both dynamically and energetically.

(1) Saturation of the magnetic reconnection is self-consistently obtained by the simulations in this chapter. As shown in Figure 6.2, the reconnection rate ($R$) falls down rapidly just after it rises to its maximum. It means that the low-layer magnetic reconnection has a short “lifetime”. Taken $R = 0.01$ as the critical level for the saturation, the “lifetime” of such reconnection is about 600 s, which is the typical lifetime for a WLF.

(2) As shown in Figure 6.1, a large bright region is seen near the bottom, which may correspond to the bright kernel of WLFs. For Model CIR, the temperature rise in the bright region is $\sim 150$ K, which is enough to account for the strong negative hydrogen ion (H$^-$) radiation (Fang et al 1993). The energy releasing rate of the magnetic reconnection, with unit length in the $z$-direction, is approximately equal to $\frac{B_0^2 L_0^2 R}{\mu_0}$, where $B_0$ is the strength of the magnetic field. For $B_0 = 700$ G, $L_0 = 2500$ km, and $R = 0.1 \tau_A^{-1}$ as in Cases A and B. Assuming that half the energy is converted to radiate in the continuum of WLFs, the surface flux of the radiation in the kernel (with a size $\sim 0.4 L_0$ in the $x$-direction as shown in Figure 6.1) is reduced to be $\sim 1.2 \times 10^{16}$ erg cm$^{-2}$ s$^{-1}$. Such a value is consistent with that obtained by observations (Neidig 1989).

(3) As the Hα line was shown to have a strong red asymmetry during the impulsive phase
of flares, $K_1$ (intensity minimum) of the Ca II K line, which is formed near the temperature minimum region (TMR), also presents a red asymmetry, i.e., the intensity of the red wing at $K_{1r}$, $I_r$, is stronger than that of the blue wing at $K_{1b}$, $I_b$ (Fang et al. 1985, 1986; Gan & Fang 1987). For the first time, Fang et al (1992) systematically studied the red asymmetry of the Ca II K line for 12 flares (including one Type II flare). They showed that a downward-moving plasma in the region above TMR can produce the red asymmetry. However, as they pointed out, if there exists the downward motion only, it would be difficult to reproduce the observed $(I_r - I_b)/I_c$ in quantity, especially for the large Type II WLF (the 1979 September 19 WLF), where $I_c$ is the intensity of the related continuum. They suggested that an additional upward-moving plasma in the region below TMR, i.e., a contract velocity profile near the TMR, would increase $(I_r - I_b)/I_c$ to be compatible with the observations.

Our numerical results show that the plasma in the lower half region is seen to move upward except in the small outflow part (see Figure 6.1). Figure 6.4 plots the velocity profile along one line $y = x + 0.42$, which was marked by the dots in Figure 6.1. It is interesting to find that the velocity profile is consistent with the contract profile in the Figure 5 of Fang et al. (1992), i.e., the plasma in the upper part moves downward, while that in the lower part moves upward. In this sense, the reconnection X-point should be above the TMR to ensure that the zero velocity lies near the TMR. Note that the downward velocity in Figure 6.4 is contributed by the downward jet, and the upward part comes from the convergent inflow. The symmetry inflows themselves can also produce the contract velocity profile. However, it is very difficult for the velocity to reach the required values ($10 \sim 30$ km s$^{-1}$) as proposed by Fang et al. (1992), since the speed of the reconnection inflow is $\sim 0.1v_A$, where Alfvén speed $v_A$ is only $\sim 20$ km s$^{-1}$ near TMR.

In summary, the numerical simulation of the low-layer magnetic reconnection shows the following two points:
Both ionization and radiation have weak effect on magnetic reconnection rate. However, they greatly change the distributions of the temperature. In particular, most part of the thermal energy of the lower chromosphere, which is released by the magnetic reconnection, will be converted into the ionizing potential, leading to weak heating in this region;

(2) Saturation of magnetic reconnection is resulted self-consistently, which implies that the low-layer line-tying magnetic reconnection has a short lifetime. Our simulation shows that the lifetime is independent of both the ionization and the radiation;

(3) Magnetic reconnection in the deep atmosphere can account for Type II WLFs in many observational aspects, such as the lifetime (∼10 min), the temperature rise in the deep atmosphere, and the radiation surface flux in bright white-light kernels. Moreover, it can provide qualitatively the required contract velocity profile to produce the observed red asymmetry at $K_1$ of the Ca II K line.

Bibliography

Chapter 7

SUMMARY

2.5-dimensional magnetic reconnection was numerically simulated in this thesis to investigate how it can be applied to account for solar flares. The numerical method used in this study is the multistep implicit scheme developed by Hu (1989). The magnetic reconnection, which results from the saturation of nonlinear tearing instability (Furth, Killen, & Rosenbluth 1963) as shown in Figure 4.8, occurs in the current-sheet configuration after a localized resistivity is introduced.

The main numerical results in this thesis can be summarized into the following points:

1. Pseudo-reconnection produced by the shifted grid mesh system

Symmetry boundary appears widely in numerical simulations. There are two methods to specify the symmetry conditions, which correspond to two types of Eulerian grid mesh systems. Our results indicate that the normal mesh system as shown in Figure 3.1(a) is acceptable both mathematically and physically, and therefore is strongly recommended. On the contrary, the shifted grid mesh system produces strong numerical resistivity and leads to pseudo-reconnection and other unphysical results. It is further suggested that in any numerical simulations, symmetric or nonsymmetric, the absence of grid points along the magnetic inversion line will introduce a strong numerical resistivity and leads to pseudo-reconnection.

2. Checking the CSHKP model for two-ribbon flares

To account for the rise motion of the flaring loop and the separation of the loop footpoints, Carmichael (1964), Sturrock (1968), Hirayama (1974), and Kopp & Pneuman (1976) developed reconnection models, where a current-sheet is considered as the preflare configuration. With the similar configuration, this thesis studied the ensuing magnetic reconnection. It was found that the simulation can reproduce the above flaring loop motions. The results also indicate the shrinkage effect of the reconnected field lines as found in observations (Švestka et al. 1987; Forbes & Acton 1996). Our simulations qualitatively support the CSHKP model. However, our study implies the following two crucial points: (1) As theoretically pointed out by Cargill & Priest (1982), our numerical simulations show that the slow-mode MHD shocks emanating from the reconnection X-point contribute to the soft X-ray (SXR) loop heating; (2) As described by Kopp & Pneuman (1976), and other literature, the SXR loop rise was thought to imply the rise of the reconnection X-point. However, our results illustrate that there is no direct relation between the rises of the SXR loop and of the X-point.

3. A unified model for solar flares

Observations have suggested that the magnetic reconnection processes may be the same in the two types of solar flares, i.e., two-ribbon and compact flares (Masuda et al. 1994;
Shibata 1997). This thesis attempts in theory to explore the possibility to unify the two types of solar flares.

In Chapter 5, we simulated two cases with the small resistive region located at different heights in the same magnetic configuration. The height of the resistive region characterizes the height of reconnection X-point since the resistive region is small.

It is shown that in the high X-point case, a bright loop (SXR loop in observations) is formed below the X-point, and is seen to rise with a speed of tens of km s$^{-1}$. Meanwhile, the two footpoints of the loop separate. These are the typical features of two-ribbon flares. A bright loop is also seen in the low X-point case. However, the loop presents obvious motion only in the impulsive phase. Thereafter, it is rather stable. Such a pattern shows a large similarity to compact flares.

Based on this study, we put forward a unified theory for the usually classified two type of solar flares, where the height of the X-point is essential to the flare morphology, i.e., two-ribbon flares correspond to the magnetic reconnection with the X-point located high, while compact flares correspond to the low X-point case. This unified theory predicts that the shrinkage effect of the reconnected field lines may exist in both types of solar flares.

4. Low-layer magnetic reconnection — a possible mechanism for Type II WLFs

White-light flares (WLFs) fall into two types due to their different observational features, where Type II WLFs require an in situ (in the deep atmosphere) heating mechanism. The research in Chapter 6 indicates that low-layer magnetic reconnection can account for Type II WLFs in four important aspects:

- (1) the typical lifetime of WLFs: $\sim$10 min;
- (2) the temperature rise in the white-light kernel: $\sim$150–250 K;
- (3) the rate of energy releasing to maintain the radiation surface flux: $1-2 \times 10^{10}$ erg s$^{-1}$ cm$^{-2}$;
- (4) the contract velocity profile required to explain the red asymmetry of the Ca $II$ K line profile.

From the above discussion, it is then clearly shown that magnetic reconnection is an effective mechanism for solar flares. Most of the characteristics of solar flares can be reproduced in magnetic reconnection processes, such as the morphology, the motion, the heating and the lifetime, and even the spectral profiles. Now, the remaining problem is the anomalous resistivity, which our recent research shed a light.

Finally, we itemize the main conclusions of this thesis as follows:

- (1) Absence of grid points at magnetic inversion line will introduce strong numerical resistivity and leads to pseudo-reconnection;
- (2) Numerical simulations of magnetic reconnection can reproduce the flaring loop motions, including the shrinkage of the reconnected field lines;
- (3) Slow-mode MHD shocks contribute to flaring loop heating;
- (4) There is no direct relation between the rises of the flaring loop and of the reconnection X-point;
- (5) The usually-classified two types of solar flares can be unified into the same magnetic reconnection model, where the different heights of the reconnection X-point lead to the bifurcation.
- (6) Magnetic reconnection in the low-layer atmosphere can explain four observational features for Type II WLFs: the lifetime, the temperature rise near TMR, the radiation surface flux, and the red asymmetry of Ca $II$ K line.
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