

Doctor Thesis

“A STUDY OF SOLAR FLARES BASED
ON COMPARISON BETWEEN THEORY
AND OBSERVATION”

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Extended Abstract

In this thesis, we study solar flares by comparing the theoretical results obtained by means of the MHD numerical simulation with the observational results to understand their driving mechanism and physical structure. Detailed accounts of the organization of this thesis are given below.

Chapter 1. General Introduction

In this chapter, we show a brief lead-in toward this thesis, that is, we mention both the background and the aim of our studies in this thesis.

Chapter 2. Physical Structure and Evolution of The Postflare Loops

In this chapter, on the basis of the Masuda's discovery of a hard X-ray loop-top source (see Masuda 1994; Masuda et al. 1994), we consider the physical structure of the postflare loops, which are those closed loops observed after the occurrence of flares (see Figure 1. 1). Masuda discovered that there existed a hard X-ray source above the soft X-ray postflare loops, which implies that a sufficient amount of energy is released outside these loops. In order to explain this observational fact, we use a flare model based on the magnetic reconnection and investigate the process of how such loop-top source is formed.

Chapter 3. Plasmoid Dynamics in Eruptive Flares

In this chapter, we cast our attention to ejected plasma blobs (plasmoids) observed frequently in solar flares (see Figure 1. 1). By referring to the observational results presented by Ohyama & Shibata (1997) who analyzed several *Yohkoh* soft X-ray images and showed the evolution of plasmoids, we study the dynamics of plasmoids in detail.

Chapter 4. 3-Dimensionality of Solar Phenomena

In this chapter, we consider some 3-dimensional solar phenomena. To do researches on that topic, we develop a numerical code for the 3-dimensional MHD numerical simulation. We explain this code briefly as well as show some interesting results obtained by means of the 3-dimensional MHD numerical simulation.

Appendix A. Summary of The Important Physical Processes in Solar Flares

In this appendix, we summarize several important physical processes relevant to solar flares, such as the magnetic reconnection, the tearing instability, and the anomalous resistivity.

Appendix B. Specification of A Numerical Code

In this appendix, we explain a numerical code used in this thesis. A family of the MHD partial-differential equations are converted into a set of difference equations, which are solved on the basis of the *modified Lax-Wendroff scheme*.

Most of the studies in this thesis (except for Chap. 4) are done by using a code for the 2.5-dimensional numerical simulation developed originally by Dr. Shibata and Dr. Yokoyama (see Shibata et al. 1989; Yokoyama 1995) with some modifications. The author is responsible for such modifications and the analyses of numerical data, while the interpretations of results and the discussions have been made cooperatively among the author, Dr. Shibata, and Dr. Yokoyama. Also the author has developed a code for the 3-dimensional MHD numerical simulation which is shown in Chap. 4. Some of the results in this thesis have been already reported in Magara et al. (1996a, b), Magara & Shibata (1997), and Magara, Shibata, & Yokoyama (1997).

Chapter 1. General Introduction

As is well known to us, the sun is the nearest and the most familiar star in the universe. Many interesting phenomena on the sun, therefore, have attracted people widely and deeply for a long time. Also, many efforts to try to understand the physical processes running behind these phenomena have been made continuously. Out of such interesting solar phenomena, I take up solar flares as a subject of my thesis and study their physical aspects by comparing the theoretical results with the observational one.

Solar flares are one of the most prominent phenomena observed in the outer atmosphere of the sun. Before the occurrence of flares, some area in an active region slowly enhances several kinds of radiation. This is called the preflare phase, which is followed by an violent phase when all sorts of electromagnetic wave are strongly emitted, plasma blobs are ejected, and highly energetic particles are generated. This violent phase is often called the impulsive phase, continuing in a fairly short time (typically $10\text{--}10^2$ s). After that violent phase, the radiation gradually decreases to its normal level. This phase is called the gradual phase (see Figure 1. 2).

Since they were first observed in the photosphere by Carrington in 1859, solar flares have been one of the most important topics in the solar physics. However, it had been quite difficult to understand their whole image as long as only the photospheric information was usable, because solar flares are the events observed mainly in the upper atmosphere of the sun, such as the chromosphere and the corona, and its behavior could not be understood well by means of the white light observation.

A turning point of the study of solar flares came when one spacecraft, called Skylab, succeeded to take many beautiful pictures of the corona through 1973 to 1974. This operation has been memorized as a milestone which taught solar physicists the important role that the magnetic field plays in various coronal activities, and solar flare is not exception. Before this operation, several primitive flare models including the effect of the magnetic field were proposed (Carmichael 1964, Sturrock 1966). After the operation, not only these magnetic-field-based flare models were widely accepted but also it began to be believed that various coronal activities had a direct relationship to the release of the magnetic energy stored by the coronal magnetic field. From this viewpoint, Hirayama (1974) and Kopp & Pneuman (1976) presented a flare model based on the magnetic reconnection. This model, called CSHKP model by including the results of Carmichael and Sturrock, suggested that the magnetic reconnection occurs in a current sheet formed in the corona so that the magnetic energy was converted into both the thermal and the kinetic energy, associated with the formation of closed loops and the eruption of plasma blobs (plasmoids) (see Figure 1. 1).

In addition to that qualitative study of solar flares, more fundamental researches started, that is, solar physicists began to investigate the possible equilibrium configuration of the coronal magnetic field and its stabilities. Zweibel & Hundhausen (1982) found several 2-dimensional magnetohydrostatic solutions analytically. This work was extended by Melville, Hood, & Priest (1984) and Su (1985, 1990). Cargill, Hood, & Migliuolo (1986), Velli & Hood (1986), and Hood & Anzer (1987) studied the MHD stability of the cylindrically symmetric magnetic arcade. The effect of the viscosity on that kind of arcade was discussed by van der Linden, Goossens, & Hood (1988) and Bogaert & Goossens (1991). Recently, Neukirch (1997) succeeded to find out a family of 3-dimensional analytic solutions of the magnetohydrostatic equations.

Researches on the evolution of the coronal magnetic field causing eruptive phenomena have also become an important subject in terms of studying the onset condition of solar flares. Historically, this kind of work started when we regarded the preflare evolution as the transition process in a series of force-free equilibria. Barnes & Sturrock (1972) studied the behavior of the cylindrically symmetric force-free field subject to a particular photospheric twist motion. Low (1977) introduced the generating-function method and obtained a family of force-free equilibria characterized by a certain parameter, which was later reconsidered by Jockers (1978) and Klimchuk & Sturrock (1989). Zwingmann (1987) showed a series of magnetohydrostatic solutions of the coronal magnetic field and explained the onset condition of eruptive phenomena, which was reconsidered by Platt & Neukirch (1994). Recently, Aly (1995) showed the possibility that an axisymmetric force-free magnetic field subject to a shearing motion reached an open-field state in a finite time.

In this respect, the recent development of computers enabled us to trace the temporal evolution of the coronal magnetic field by means of numerical simulations. These studies are found in Mikic, Barnes, & Schnack (1988), Biskamp & Welter (1989), Finn, Guzdar, & Chen (1992), Inhester, Birn, & Hesse (1992), Kusano, Suzuki, & Nishikawa (1995), Choe & Lee (1996), and so on. All of these successfully confirmed the formation of a current sheet within a magnetic arcade, though there has been one controversy about the topological configuration causing the formation of a current sheet, that is, in order to form a current sheet within a magnetic arcade, it was suggested that at least two neighboring arcades were needed (Biskamp & Welter 1989) or a converging motion toward the neutral line had to be added to a shearing motion (Inhester et al. 1992). Lately Amari et al. (1996) obtained an interesting result on this problem, in which they showed that a current sheet could be formed within a single magnetic arcade by imposing a long-lasting shearing motion alone.

Turning to the solar observation, the surprising development of observational tools has greatly contributed to our understanding of the magnetic and gaseous fine structures of solar flares. For example, Kurokawa (1987, 1989) made high-resolution H α observations and concluded that the twisted flux tube emerging from the subphotosphere played a crucial role

in the occurrence of flares. Tanaka (1991) studied the complex-subsurface magnetic-rope structure of a very flare-active isolated δ group by means of high-resolution evolutionary data from BBSO magnetic & velocity data. Wiik et al. (1996) used MSDP (*Multichannel Subtractive Double-Pass Spectrograph*) and derived the bulk-flow velocity field along a postflare loop. As for the X-ray observation, several space operations such as SMM (*Solar Maximum Mission*, starting in 1980) and Hinotori (starting in 1981) took over the role of Skylab and presented many interesting results. Sakurai (1985) used some hard X-ray images obtained by Hinotori and studied the magnetic structures causing flares.

After the launch of solar-A satellite, called *Yohkoh*, in 1991, the researches in solar flares entered on the next stage (Ogawara et al. 1991). Since the spacial and time resolutions of this spacecraft are both high, it becomes possible to make a detailed comparison between the *Yohkoh* X-ray images and other data. For examples, Schmieder et al. (1995, 1996) and van Driel-Gesztelyi (1997) studied the features of the postflare loops by comparing H α data with soft X-ray data. Yoshimura et al. (1996) studied several emerging flux regions by using soft X-ray data and H α data. Aschwanden & Benz (1997) used both radio data and soft X-ray data to estimate the gas density inside the postflare loops and at a region where the particle acceleration occurs.

Not only the comparison among several kinds of observational data but also the comparison of theoretically predicted results with fine observational results makes a large contribution to our understanding of the real image of solar flares. Since solar flares are fairly complicated and nonlinear phenomena, it is appropriate to use the method of numerical simulations for the direct comparison between theory and observation. This is why we perform the MHD numerical simulation and compare the results to some observational results in this thesis.

Chapter 2. Physical Structure and Evolution of The Postflare Loops

Abstract

In this chapter, we investigate the physical structure and evolution of the postflare loops by performing the 2-dimensional MHD numerical simulation with a particular attention paid to the observational results brought by *Yohkoh*. By using a flare model based on the magnetic reconnection, we confirm that a hot and dense region appears behind a fast MHD shock which is formed through the interaction between the reconnection jets and the postflare loops. That dense and hot region is maintained by the *magnetic pressure walls*, that is, the enhanced magnetic pressure confines a mass of energetic gases into a small region. Since such energetic region exists above the soft X-ray closed loops in the impulsive flares, we conclude that it corresponds to the so-called Masuda source (see Masuda et al. 1994).

As for the evolution of the postflare loops, we plot the time variation of the height of the reconnected magnetic-field line and confirm its decrease tendency with time, which is consistent with observational results. Moreover, the overall expansion behavior of the postflare loops is also reproduced in our simulation.

2. 1. Introduction

There are several different ways to classify solar flares (see Table 2. 1). Here, we distinguish between the LDE (long duration event) flares and the impulsive flares. Morphologically, the former has a cusp-shaped loop structure, while the latter does not, having only a simple loop structure. Since these two types of flares are different in appearance, it has been long thought that a different mechanism may operate for each of them. For example, the LDE flares (hereafter referred to as *cusp-type* flares) has been understood by the classical model for two-ribbon flares, that is, it was long suspected that the magnetic reconnection occurred above a cusp-shaped loop and the stored energy was almost released outside this loop (see Carmichael 1964, Sturrock 1966, Hirayama 1974, and Kopp & Pneuman 1976). *Yohkoh* actually confirmed this hypothesis by discovering a cusp-shaped loop in the LDE flares (see Figure 2. 1 and also Tsuneta et al. 1992). Impulsive flares (hereafter referred to as *non-cusp-type* flares), on the other hand, have been explained by a loop flare model in which the main energy release occurs inside the loop and the magnetic reconnection does not take place outside this loop (Alfvén & Carlqvist 1967, Spicer 1977, Uchida and Shibata 1988). To sum up, it has been thought that there was a clear distinction between *cusp-type* and *non-cusp-type* flares.

A doubt has been cast on this interpretation, however, when an intense hard X-ray emission was observed from above the soft X-ray closed loops in some impulsive limb flares (see Figure 2. 2 and also Masuda et al. 1994). Previously, this type of flares was classified as *non-cusp-type* flares, because only closed loops could be seen in soft X-ray. The presence of a hard X-ray source above the closed loops has now established that an impulsive energy release occurs not inside but outside the loops in at least some of *non-cusp-type* flares. Moreover, in the same impulsive limb flares as Masuda studied, Shibata et al. (1995) found the soft X-ray signature of plasmoids moving upward at about 50–400 km / s, which is about 10 % of the coronal Alfvén velocity, v_A , because $v_A \sim 3000 (B / 100 \text{ G}) (n / 10^{10} \text{ cm}^{-3})^{-1/2} \text{ km} / \text{s}$ in a typical active region in the corona. Since the plasmoid eruption is one of the key features associated with *cusp-type* flares, Shibata's discovery also supports the idea that at least part of *non-cusp-type* flares may have a close similarity with *cusp-type* flares (see Shibata 1995).

Now a question arises as to the relationship between *non-cusp-type* and *cusp-type* flares. To answer this question, we apply a flare model based on the fast magnetic reconnection originally developed by Ugai (1986) to solar flares and study the physical structure and evolution of the postflare loops. In this study, we perform the 2-dimensional MHD numerical simulation and investigate some key characteristics of *cusp-type* flares, keeping in mind the application of them to *non-cusp-type* flares. We give our physical assumptions, basic equations, and method of calculation in § 2. 2. The results are presented in § 2. 3. We then introduce some observational results relevant to our study in § 2. 4. In § 2. 5, we discuss the relationship between the impulsive flares and the LDE flares, trying to generalize the *cusp-type* flare model based on the fast magnetic reconnection. The final section is devoted to a summary.

2. 2 Basic Formulations

2. 2. 1. Basic Equations

In this study, we use the standard set of the MHD equations except that the effects of gravity and viscosity are both neglected. These equations are

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \mathbf{v} = 0, \quad (2. 1)$$

$$\left[\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \mathbf{v} \right] = - \nabla P + \frac{1}{4\pi} (\nabla \times \mathbf{B}) \times \mathbf{B}, \quad (2. 2)$$

$$\frac{\partial P}{\partial t} + (\mathbf{v} \cdot \nabla) P = \frac{1}{4\pi} |\nabla \times \mathbf{B}|^2, \quad (2. 3)$$

$$\frac{\mathbf{B}}{t} = \nabla \times (\mathbf{v} \times \mathbf{B}) - \nabla \times \left(\frac{\mathbf{v}}{\mu} \times \mathbf{B} \right), \quad (2.4)$$

$$P = \frac{R T}{\mu}. \quad (2.5)$$

Additionally, we use $\mathbf{B} = 0$ as the initial condition for equation (2.4). Here all the symbols, such as P , T , \mathbf{v} , and \mathbf{B} have their usual meanings, γ is the adiabatic index, R is the gas constant, μ is the mean molecular weight, and η is the electrical resistivity. On the basis of these equations, we perform the 2-dimensional MHD numerical simulation in the Cartesian coordinate. All the physical variables are dependent on both the x -coordinate and the z -coordinate but independent of the y -coordinate. In practice, all the calculations are carried out in the nondimensional form derived by some particular units. These units are summarized in Table 2.2.

2.2.2. Initial Configuration & Boundary Conditions

Figure 2.3 shows the initial configuration of the present numerical simulation in the (x, z) -plane. Initially, the magnetic field is set to be antiparallel with respect to the z -axis. A hatched area represents the region where the artificial resistivity is turned on as the initial perturbation (during $0 < t < 0.6$), while a gray area represents the region where the initial density is enhanced for modeling on the lower dense atmosphere of the sun. In the actual calculations, we solve the basic equations only over a half domain of the calculation region ($x > 0$), assuming a symmetry with respect to the z -axis.

The physical variables are initially set to be in equilibrium and described as

$$\mathbf{v} = \mathbf{0}, \quad (2.6)$$

$$\mathbf{B}(x, z) = -B_0 \tanh\left(\frac{3}{2}x\right) \hat{z}, \quad (2.7)$$

$$P(x, z) = P_0 \left[\frac{1}{\cosh^2\left(\frac{3}{2}x\right)} + 1 \right], \quad (2.8)$$

$$T(x, z) = T_0 \left\{ 0.45 \tanh\left[100(z - 0.1)\right] + 0.55 \right\}, \quad (2.9)$$

$$\rho(x, z) = \frac{P(x, z)}{T(x, z)}, \quad (2.10)$$

where we set $\gamma = 0.2$, $\mu = 5/3$, $P_0 = 1/\mu = 0.6$, $B_0 = [8/\mu]^{1/2} = 8.68$, and $T_0 = 1$ in this calculations. Since the initial magnetic field is antiparallel, a current sheet is formed along the z -axis, where the gas pressure is enhanced to obtain the balance of forces. As we mentioned,

for modeling on the massive layers such as the chromosphere and the photosphere, we put a dense region between $z = 0$ and $z = 0.1$. This region plays an important role in line-tying the magnetic field at the base.

To carry out our calculations, we have to set four boundary conditions around the calculation region. At the top boundary (at $z = 24$), we set a free boundary condition:

$$\frac{B_z}{z} = \frac{v_x}{z} = \frac{v_z}{z} = \frac{P}{z} = \frac{P}{z} = 0, \text{ and } \mathbf{B} = 0. \quad (2.11)$$

Along the z -axis (at $x = 0$), at the side boundary (at $x = 8$), and at the bottom boundary (at $z = 0$), we set

$$v_x = B_z = \frac{v_z}{x} = \frac{B_x}{x} = \frac{P}{x} = \frac{P}{x} = 0, \quad (2.12)$$

$$v_x = B_x = \frac{v_z}{x} = \frac{B_z}{x} = \frac{P}{x} = \frac{P}{x} = 0, \quad (2.13)$$

and

$$v_z = B_x = \frac{v_x}{z} = \frac{B_z}{z} = \frac{P}{z} = \frac{P}{z} = 0, \quad (2.14)$$

respectively. We perform the present numerical simulation on the basis of the *modified Lax-Wendroff scheme*. The number of grid points is $(N_x \times N_z) = (100 \times 100)$ for the low-resolution case and (137×400) for the high-resolution case. In the former case, grid points are distributed uniformly along both the x -coordinate and the z -coordinate, and hence the mesh size is $(x, z) = (0.08, 0.24)$, while, in the latter case, grid points are distributed uniformly along the z -coordinate but nonuniformly along the x -coordinate with fine meshes assigned to the region of a current sheet ($0 < x < 0.2$). Accordingly, the mesh size is $(x, z) = (0.02, 0.06)$ inside the sheet and $(x, z) = (0.1, 0.06)$ in the distant region from the sheet.

2. 2. 3. Initial Perturbation

Initially, the electrical resistivity is set to be zero everywhere except in a small region of $\left[x^2 + (z - 4.0)^2 \right]^{1/2} < 0.8$ (gray area in Figure 2. 3) where $\eta = 1 / 30$ is assigned in a finite time, $0 < t < 0.6$. Owing to this resistivity, the magnetic field can begin to dissipate, which causes inflows toward this resistive region. These flows drag the magnetic field, making a X-type neutral point. Consequently, the current density increases and turns on the anomalous resistivity (see below).

2. 2. 4. Anomalous Resistivity

We may safely assume that the fast magnetic reconnection is responsible for the violent energy release observed in solar flares. Among several physical processes capable of causing the fast magnetic reconnection, the anomalous resistivity is probably one of the most important processes. We, therefore, give a brief explanation of this process in Appendix. A. 3. The anomalous resistivity is known to have a close relationship with plasma microturbulences (see Parker 1979; Treumann & Baumjohann 1997). Suppose that the ion-electron drift velocity exceeds its sound velocity, the mean flow of electrons excites the plasma microturbulences, which provides a seed for the anomalous resistivity. It is important to note that the anomalous resistivity can make the diffusion region localized into a fairly small area and thereby the Petschek-like configuration is formed around the diffusion region. That configuration is actually favored in terms of causing the fast magnetic reconnection (see Yokoyama and Shibata 1994). Ugai pointed out the essence of the fast magnetic reconnection by using a spontaneous fast-magnetic-reconnection model. He described a new-type nonlinear instability that grows by the self-consistent interaction (or feedback) between the microscopic anomalous resistivity and the macroscopic global reconnection flows (Ugai 1996). The anomalous resistivity is also used to understand other active phenomena in the universe. For example, Borovsky (1986) explained the extragalactic jet by a hybrid double-layer/anomalous resistivity model.

In the present study, we assume the following form for the anomalous resistivity except that η is over 1. In that case η is constantly set to be 1.

$$\eta = \begin{cases} \frac{1}{15} \left(\left| \frac{v_d}{v_c} \right| - 1 \right), & \text{for } |v_d| \geq \left| \frac{j_y}{\omega} \right| \geq |v_c|, \\ 0, & \text{for } |v_d| < |v_c|, \end{cases} \quad (2. 15)$$

where $\mathbf{j} = \nabla \times \mathbf{B}$, v_d represents the ion-electron drift velocity, and v_c is a threshold velocity (see Ugai 1986). In the present study, we take $v_c = 12$.

2. 3. Structure and Evolution of The Postflare Loops

2. 3. 1. Overview

Figure 2. 4 shows how the magnetic reconnection proceeds in this simulation. Here, contour lines and arrows represent the magnetic field lines and velocity field, respectively. Elapsed times are $t = 0$ (top-left), 6 (top-right), 12 (bottom-left), and 18 (bottom-right). Note that part of the calculation region, $(-4, 0) \leq (x, z) \leq (4, 20)$, is displayed here.

The fast magnetic reconnection sets in at $t = 6$ (see the top-right panel) and then a plasmoid begins to be formed. As the magnetic reconnection still proceeds, the plasmoid moves upward, getting larger in size. The reconnected field lines flow downward and form closed loops in the range of $0 < z < 2$ (see the bottom-left panel). At $t = 18$ (the bottom-right panel), a cusp-like structure is formed near $z = 3$. Basically, these evolutionary features are consistent with the results of previous works (e.g., Forbes & Priest 1983). In the following, we study the physical structure of such closed loops in detail, while the investigation of the dynamics of plasmoids is left to Chap. 3.

2.3.2. Physical Structure of The Postflare Loops

We use the results of the high-resolution case to clarify the physical structure around the loop-top region. Figure 2.5a is composed of four color maps, each of which shows the temperature (top-left), gas density (top-right), gas pressure (bottom-left), and current density (bottom-right), respectively. An elapsed time is $t = 10$. Solid lines and arrows represent the magnetic field lines and velocity field, respectively. Figures 2.5b–5c show the spacial distributions around the loop-top region of such physical variables as the fast-mode Mach number (M_a), temperature (T), y -component of the current density (j_y), gas pressure (P_g), magnetic pressure (P_m), gas density (ρ), and the absolute value of the x -component of the magnetic field ($|B_x|$). Figure 2.5b is a vertical distribution along the z -axis and Figure 2.5c is a horizontal distribution along a particular horizontal line ($z = 1.6$). An elapsed time is $t = 10$.

The region of the maximum temperature appears around a neutral point, where the temperature is about 8 times higher than in the surrounding region. There is little matter around this point, because a lot of plasma has been flowed away along the z -axis due to the strong reconnection outflows generated by the magnetic reconnection. Casting eyes to the loop-top region, we find that a fast MHD shock is formed at $z \sim 2$ in Figure 2.5b. Behind this shock, a region which is hotter and denser than the ambient region appears at $z \sim 1.6$. Moreover, the horizontal distribution given in Figure 2.5c gives us other information about this region. The profile of the current density implies the existence of a pair of slow MHD shocks by which the surrounded gases are heated. In addition, when we look at the profile of the magnetic pressure, it is found that such hot and dense gases are effectively confined by the enhanced magnetic pressure, which we call the *magnetic pressure wall*. Consequently, it is clearly understood that a hot and dense loop-top source is formed and survives by the help of the fast MHD shock and a pair of slow MHD shocks.

Figure 2.6 is a picture showing the structure of the loop-top source schematically. If the slow MHD shock is assumed to be the so-called switch-off shock with the upstream magnetic field nearly tangential to the shock surface, the enhancement of the gas density and temperature by this shock is expressed as

$$\frac{\rho_{II}}{\rho_I} = \frac{1 + \beta_I}{2 + \beta_I}, \quad (2.16)$$

and

$$\frac{T_{II}}{T_I} = \frac{2}{5 + \beta_I} + 1, \quad (2.17)$$

where subscripts I and II means the value measured in the region I and II , respectively. Here we take $\beta = 5/3$ (see Forbes, Malherbe, & Priest 1989). Similarly, the enhancement of the gas density and temperature by the fast MHD shock is expressed as

$$\frac{\rho_{III}}{\rho_{II}} = \frac{-U + [U^2 + V]^{1/2}}{4[2 - \beta]}, \quad (2.18)$$

and

$$\frac{T_{III}}{T_{II}} = \left[\frac{\frac{\rho_{III}}{\rho_{II}} - 1}{\beta} \frac{(\beta - 1) \left(\frac{\rho_{III}}{\rho_{II}} \right)^2 + 2(1 - \beta) \frac{\rho_{III}}{\rho_{II}} + (2\beta - \beta + 1) - 1}{\left(1 - \frac{\rho_{III}}{\rho_{II}} \right) + 1 + \frac{\rho_{III}}{\rho_{II}}} + 1 \right] \frac{\rho_{II}}{\rho_{III}}, \quad (2.19)$$

where

$$U = \left[2\beta - \beta + \frac{(\beta - 1)(\beta + 2)M_a^2}{\beta} + 2 \right], \quad (2.20)$$

and

$$V = 8[2 - \beta][\beta + 1][\beta - \beta + 2]M_a^2 \quad (2.21)$$

(see, for example, Priest 1982 or Shu 1992). Here subscript III means the value measured in the region III . In the present simulation, we have $\beta = 5/3$, $\beta_I = 0.2$, $\beta_{II} = 143$, and $M_a = 1.25$ so that we obtain $\rho_{III}/\rho_{II} = 1.37$ and $T_{III}/T_{II} = 1.24$, which is consistent with the results shown in Figure 2. 5b. Actually, an active region in the solar corona has such a plasma beta as is much lower than the present numerical simulation so that those enhancements described above become stronger in the real circumstances of the corona. For example, when we take $\beta_I = 0.01$, equations (2. 16) and (2. 17) give $\rho_{II}/\rho_I = 2.5$ and $T_{II}/T_I = 40$. Then if we assume $M_a = 1.25$ and $\beta_{II} = 100$, we finally obtain $\rho_{III}/\rho_{II} = 3.43$ and $T_{III}/T_{II} = 50$ by using equations (2. 18) and (2. 19). These results imply that the temperature of the loop-top source is about one hundred million degrees because T_I is about 2×10^6 K. Observationally, Masuda et al. (1994) and recently Sato (J. Sato 1997, private communication) found out that the temperature of

this source was about one hundred million degrees by assuming that the source has a thermal emission. Hence our estimation is in good accordance with their result.

Turning our eyes to the inside of the postflare loops in Figure 2. 5a, there remains hot regions because the present numerical simulation does not include the radiative cooling effect. However, in the real circumstances of the postflare loops, this effect does play an important role in reducing the temperature inside the postflare loops, which causes the so-called H_β postflare loops (see § 2. 5 and Figure 2. 9). Forbes & Malherbe (1985, 1991) considered that effect and studied the structure of the postflare loops.

2. 3. 3. Evolution of The Postflare Loops

Next, we turn our attention to the evolution of the postflare loops. In Figure 2. 7, each dotted line represents the time variation of the height of the reconnected magnetic-field line measured along the z -axis. Asterisks show the inflection points of dotted lines. These points are considered to indicate the loop-top region which can be made through the collision between the downflow of the reconnected magnetic-field line and the pre-existed closed loops. A thick curve line is drawn on the basis of the power-law fitting applied to such inflection points and this line is considered to represent an evolutionary line of the postflare loops. From this figure, we find that all the reconnected magnetic-field lines have a downward movement with time. On the other hand, as to the overall behavior of the postflare loops, they rise rapidly at first and then their rising speed gradually decreases.

2. 4. Observational Results of The Postflare Loops

2. 4. 1. Loop-top Source

In this section, we show two types of loop-top source observed by *Yohkoh*. One of them is observed in the impulsive flares. This is known as the Masuda's source and an example of this is shown in Figure 2. 2. The character of this source is that it exists above the soft X-ray closed loops and can be observed in such a high-energy electromagnetic wave as hard X-ray. Another type is observed in the LDE flares and an example of this can be seen in Figure 2. 1. Compared with the Masuda's source, this source is formed inside the soft X-ray closed loops and can be observed in soft X-ray.

2. 4. 2. Evolution of The Postflare Loops

Recently, Hiei & Hundhausen (1996) reported the shrinking behavior of the postflare loops observed by *Yohkoh*. They showed that one loop appeared at a certain height and shrank with time, then another loop appeared at the greater height than the previous one (see

Figure 2. 8). From Figure 2. 8, it is found that the downward velocity of each loop is about 3 km / s and the overall expansion speed is about 2 km / s. Such shrinking behaviors of the postflare loops are also discussed by Forbes & Acton (1996), who compared the potential magnetic field with the postflare loops observed by *Yohkoh*.

2. 5. Discussion

2. 5. 1. Physical Structure of The Postflare Loops

In this section, we discuss the physical structure of the postflare loops in both the impulsive flares and the LDE flares. According to Figure 2. 5a, the temperature around a neutral point and a pair of slow MHD shocks is so high that the thermal conduction can efficiently transport a sufficient amount of energy to the lower atmosphere along magnetic field lines. This is followed by heating of the chromospheric gases and then these heated gases begin to evaporate into the corona. In this way, such evaporated gases eventually fill the coronal loops in the postflare phase, which are observed as the soft X-ray closed loops. On the basis of this consideration, we then consider why there are two types of loop-top source. Here we pay attention to the efficiency of the magnetic reconnection in the impulsive flares and the LDE flares. As to the impulsive flares, the magnetic reconnection proceeds so rapidly that many closed loops can be formed. The evaporated gases try to fill these loops, though they do not have enough time to fill the outermost closed loop because of the quickness of the magnetic reconnection. As is mentioned in § 2. 3. 2, the loop-top source is formed just behind the fast MHD shock which is located at the top of the outermost closed loop, so that the loop-top source in the impulsive flares exists above the soft X-ray closed loops. On the other hand, the rate of the magnetic reconnection in the LDE flares is so slow that the evaporated gases can always fill even the outermost closed loop, which makes the loop-top source surrounded by the soft X-ray flaring gases.

Followed by the above discussion, we summarize our flare model based on the magnetic reconnection which can be applied to not only the LDE flares but also the impulsive flares. Figure 2. 9 is a schematic illustration of this model.

In the impulsive flares, there occur the downward flow of energetic electrons and the thermal conduction, each of which originates from a neutral point and a pair of slow MHD shocks. The electron flow and the thermal conduction direct toward the chromosphere along magnetic field lines, while the latter forms the conduction front, whereby the neighboring plasmas are heated (see Forbes & Malherbe 1985). The chromospheric gases are heated by such electron flow and the thermal conduction so that the hard X-ray footpoint sources and the H β bright ribbons are generated (see Sakao 1994). The heated gases evaporate into the corona, forming the dense soft X-ray closed loops. In this respect, it must be taken in mind

that a hot but low-density region cannot always be observed in soft X-ray because the emissivity is more sensitive to the gas density than the temperature. The reconnection jets collide with the outermost closed loop to form a fast MHD shock, which makes a hot and dense core in cooperation with a pair of slow MHD shocks (see § 2. 3. 2). Since the evaporated gases cannot reach this energetic core, it is observed that there is a hard X-ray source above the soft X-ray closed loops.

In the LDE flares, the closed loops under a neutral point completely fill with the evaporated gases. In this case, the evaporated gases contact with the conduction front so that they are sufficiently heated up, forming a soft X-ray cusp-like structure, which gives the reason to classify these flares into *cusp-type* flares (see Figure 2. 1). The loop-top source is observed to be embedded by the evaporated gases. Inside the closed loops, heating by the conduction front is no longer effective so that the plasma in the closed loops can cool via the radiative cooling process, whereby the postflare loops can be observed in H α . Van Driel-Gesztelyi et al. (1997) investigated the formation process of the postflare loops in detail and Heinzel (1994) studied the physical structure of the cool component of the postflare loops.

We have said that the difference between the impulsive flares and the LDE flares comes from the difference in the rate of the magnetic reconnection, which can also give appropriate explanations on other aspects of such two types of flares. First, when the rate of the magnetic reconnection is high, the released energy per unit time is so large as to make very energetic sources, such as hard X-ray loop-top and footpoint sources observed in the impulsive flares. Second, if we assume that the total released energy is equal, the smaller the rate of the magnetic reconnection is, the longer the duration of flares is, which supports the long lifetime of the LDE flares. Finally, the topological aspect of flares is also influenced by the rate of the magnetic reconnection. Since the impulsive flares are usually small in size, a small diffusion region is expected in this type of flares, which implies that the fast magnetic reconnection occurs in that configuration because the fast magnetic reconnection favors a small diffusion region, known as the Petschek-type reconnection (see Appendix. A. 1). On the other hand, when a diffusion region is long, the magnetic reconnection is not very effective, which can be applied to the case of large-scale flares such as the LDE flares.

Those physical aspects seen in the impulsive flares and the LDE flares also can be seen in the impulsive phase and the gradual phase of a typical flare. In the impulsive phase, the energy-release rate dramatically increases, which means that the magnetic reconnection proceeds rapidly. After that, the reconnected magnetic-field lines pile up and the size of the postflare loops becomes large. Moreover, an ejected plasmoid makes a diffusion region extended long, which decreases the rate of the magnetic reconnection. In this way, the evolution enters on a not so energetic but long-lived stage, which is called the gradual phase.

2. 5. 2. Evolution of The Postflare Loops

Comparing Figure 2. 7 with Figure 2. 8, it is found that the height of a particular postflare loop is always decreasing with time in both the theoretical results and the observational results. Here, our attention is paid to the evolution after the inflection points in Figure 2. 7, because a reconnected magnetic-field line can be recognized as a new postflare loop after it collides with the pre-existed closed loops. By looking at Figure 2. 8, we find that the observation suggests a very slow downward velocity of a postflare loop (about 3 km / s). This is consistent with the theoretical results in Figure 2. 7, because the evolution after the inflection points shows a very gradual descent.

The overall expansion of the postflare loops shown in Figure 2. 7 indicates that the rate of expansion is high at first and then gradually decreases, which is consistent with the time variation of the height of the postflare loops in *cusp-type* flares presented in Figure 7. 6 of Pneuman (1981).

2. 6. Summary of Chapter 2

Finally, we summarize main results of this chapter as follows:

1. By using a flare model based on the fast magnetic reconnection, we have studied the physical structure of *cusp-type* flares in detail. We find that a hot and dense region is formed and survives by the help of both a fast MHD shock and a pair of slow MHD shocks.
2. Taking into consideration the rate of the magnetic reconnection, we conclude that a hot and dense region can be formed above the soft X-ray closed loops in the impulsive flares and inside these loops in the LDE flares. This implies that at least some of *non-cusp-type* flares can also be explained by the model developed for *cusp-type* flares. The validity of this model is also considered in Aschwanden & Benz (1997) from observational viewpoints.
3. We find that each of the reconnected magnetic-field lines has a tendency of downward motions, which is consistent with observational results. The downward velocity is quite small.
4. As for the overall expansion of the postflare loops, we succeed to reproduce it by using a flare model based on the fast magnetic reconnection. This model suggests that the expansion is caused by the successive piles of the reconnected magnetic-field lines rather than the rising motion of every reconnected magnetic-field line.

Chapter 3. Plasmoid Dynamics in Eruptive Flares

Abstract

In this chapter, we investigate the plasmoid dynamics in eruptive flares by performing the 2.5-dimensional MHD numerical simulation. We start with a linear force-free magnetic arcade and impose the localized resistive perturbation on the axis of symmetry of the arcade. Owing to the resistivity, the magnetic field begins to dissipate, producing inflows toward the region where the resistivity is imposed. These inflows make the magnetic field convex to the axis of symmetry and hence a X-type neutral point is formed on that axis and a magnetic island appears. We find that the evolution of the magnetic island consists of three stages. In the first stage, it is slowly lifted by those upflows produced by the initial resistive perturbation. Then once the anomalous resistivity sets in, it begins to be accelerated. This acceleration stops after a fast MHD shock is formed at the bottom of the magnetic island, which implies that the upflows around the center part of the magnetic island are no longer strong. These three stages are also confirmed to exist in observational results (Ohyama & Shibata 1997). Moreover, we find that there is a time-lag between the start time of the acceleration of a magnetic island and the peak time of the electric field at a neutral point. This fact can be understood by considering that the component of the magnetic field which is perpendicular to a current sheet (perpendicular magnetic field) inhibits the efficient magnetic reconnection. Such time-lag is also found in the observational results which Ohyama & Shibata (1997) presented. In addition, we compare the initial rise velocity of a magnetic island in our simulation to the initial rise velocity of a plasmoid in their observational results, whereby we derive a conclusion that in the actual situations, the initial resistive perturbation proceeds very weakly and at many positions within a magnetic arcade.

3.1. Introduction

Solar flares are highly energetic and complicated phenomena in which mass eruptions occur, energetic particles are generated, and high energy radiations are emitted. In this chapter, we study the mass eruption in solar flares, which is sometimes called the plasmoid eruption.

Theoretically, these eruptive phenomena have been considered from a viewpoint of instabilities or a loss of equilibrium of the coronal magnetic field. Since the magnetic force is much stronger than other forces in the corona, any coronal structure is mainly controlled by the magnetic field. Therefore, unless magnetically driven events occur, coronal structures evolve in a series of quasi-static states. However, once such events occur, the coronal structures no longer stay in a static state but begin to make a dynamical evolution. This scenario has

made it important to study the equilibria and stabilities of the coronal magnetic field. Zweibel (1981, 1982) investigated the stabilities of 2-dimensional magnetohydrostatic equilibria. Priest (1988) and Steele et al. (1989) discussed the problem of a loss of equilibrium in terms of CMEs (coronal mass ejections). Forbes, Priest, & Isenberg (1994) estimated the amount of energy that is released when a loss of equilibrium occurs. Priest & Forbes (1990), Forbes (1990) used the method of complex analyses and studied the prominence eruption.

Recently, the rapid development of computers enabled us to trace the temporal evolution of the coronal magnetic field directly by means of numerical simulations. Such studies are found in Mikic, Barnes, & Schnack (1988), Biskamp & Welter (1989), Finn, Guzdar, & Chen (1992), Inhester, Birn, & Hesse (1992), Kusano, Suzuki, & Nishikawa (1995), Choe & Lee (1996), and so on. These authors investigate how the coronal magnetic field evolves when a particular photospheric motion is imposed.

Turning to the observation, the satellite *Yohkoh* has brought us many excellent data of the corona since its launch in 1991, which have contributed to our understanding of the coronal phenomena (see Masuda 1994; Sakao 1994; Hara 1996; Shimizu 1997). Owing to these data, some pre-proposed problems of solar phenomena were solved, such as the reliability of the magnetic reconnection in solar flares (Tsuneta et al. 1992; Magara et al. 1996b), while some new predictions have been proposed (Shibata et al. 1994, 1995). As for the plasmoid eruption, Ohyama & Shibata (1997) analyzed the evolution of plasmoids in detail and show their dynamical properties. The aim of this chapter is to clarify the basic physical processes involved in the plasmoid eruption on the basis of those new observational results presented by Ohyama & Shibata (1997). For this purpose, we perform the 2.5-dimensional MHD numerical simulation and compare the simulation results with the observational one.

An organization of this chapter is as follows. We show basic formulations in the next section. Main results are presented in § 3. 3, while § 3. 4 is used to introduce some observational results relevant to the plasmoid eruption. In § 3. 5, after discussing our results and making comments on several problems of the present study, we give our conclusion. The final section is devoted to a summary.

3. 2. Basic Formulations

3. 2. 1. Basic Equations

Basic equations are the same one as are used in Chap. 2; namely, equations (2. 1)–(2. 5). On the basis of these equations, we perform the 2.5-dimensional MHD numerical simulation in the Cartesian coordinate. All the physical variables are dependent on both the x -coordinate and the z -coordinate but independent of the y -coordinate, while the y -components of the

velocity and magnetic field are included in our calculations, which means the 2.5-dimensional numerical simulation.

3. 2. 2. Initial Configuration

Initially we assume a linear force-free field described as

$$B_x = -\frac{2L}{H} B_0 \cos\left(\frac{x}{2L}\right) e^{-\frac{z}{H}}, \quad (3.1)$$

$$B_y = -\sqrt{1 - \left(\frac{2L}{H}\right)^2} B_0 \cos\left(\frac{x}{2L}\right) e^{-\frac{z}{H}}, \quad (3.2)$$

$$B_z = B_0 \sin\left(\frac{x}{2L}\right) e^{-\frac{z}{H}}, \quad (3.3)$$

where L is a horizontal scale length and is used as a normalization length unit ($L = 1$). H means a vertical scale height of the magnetic field. For the present study, H ranges from $2L/$ to , where $H = 2L/$ corresponds to the case of the potential field and $H =$ corresponds to the case of the open field. Usually, a linear force-free field is characterized by a parameter , which relates the magnetic field to its rotation in the following way.

$$\nabla \times \mathbf{B} = \alpha \mathbf{B}. \quad (3.4)$$

In our present formulation, this value is described as $\alpha = \left[\left(\frac{1}{2L} \right)^2 - \left(\frac{1}{H} \right)^2 \right]^{1/2}$. Therefore, α ranges from zero (potential field) to $1/2L$ (open field). A linear force-free field is the lowest energy state for given boundary conditions with a prescribed helicity (see Heyvaerts & Priest 1984; Biskamp 1993), but the coronal magnetic field does not always lie in this state (Schmieder et al. 1996). This is because the relaxation time to this state is not so short as the dynamical evolution time (see Browning & Priest 1986). However, in the present study, we start with this state for simplicity.

The gas pressure P is uniform ($P = P_0$) and the ratio of this to the magnetic pressure is defined as $\beta = 8 P_0 / B^2$ (plasma beta). The gas density ρ is uniform ($\rho = \rho_0$) except in the bottom region where it is 10 times higher than elsewhere. This dense region is modeled on the massive layers of the solar atmosphere, such as the chromosphere and the photosphere. Therefore, the gas pressure and density are expressed as

$$P = P_0, \quad (3.5)$$

and

$$\frac{\rho}{\rho_0} = 4.5 \left[\tanh \left[-50(z - 0.1) \right] + 1 \right] + 1, \quad (3.6)$$

respectively. The plasma beta is also expressed as

$$\begin{aligned}
&= \frac{P_0}{\frac{B^2}{8}} \\
&= \frac{P_0}{\frac{B_0^2}{8} e^{-\frac{2z}{H}}} \\
&= \frac{P_0}{B_0^2} e^{\frac{2z}{H}}.
\end{aligned} \tag{3.7}$$

In the present study, we set $\beta_0 = 0.2$, $H = 5 / 3$, $P_0 = 1 / 3 = 0.6$, $B_0 = 1$, and $B_0 = \left[8 P_0 / \beta_0 \right]^{1/2} = \left[8 / 0.2 \right]^{1/2} = 8.68$. From now on, all the physical variables presented in this chapter are normalized by the units in Table 2. 2.

Finally, the temperature is defined as

$$T = \frac{P}{\beta}. \tag{3.8}$$

3. 2. 3. Boundary Conditions

Figure 3. 1 illustrates the domain of the present numerical simulation. This figure also shows the initial configuration of the magnetic-field lines projected onto the (x, z) -plane. We set a free boundary condition at the top boundary (at $z = 40$):

$$\frac{B_x}{z} = \frac{B_y}{z} = \frac{v_x}{z} = \frac{v_y}{z} = \frac{v_z}{z} = \frac{P}{z} = \frac{\beta}{z} = 0, \quad \mathbf{B} = 0. \tag{3.9}$$

We set antisymmetric boundary conditions both along the z -axis (at $x = 0$) and along the side boundary (at $x = 8$):

$$v_x = v_y = B_z = \frac{v_z}{x} = \frac{B_x}{x} = \frac{B_y}{x} = \frac{P}{x} = \frac{\beta}{x} = 0. \tag{3.10}$$

We set a rigid boundary condition at the bottom boundary (at $z = 0$):

$$v_y = v_z = \frac{v_x}{z} = \frac{B_x}{z} = \frac{B_y}{z} = \frac{B_z}{z} = \frac{P}{z} = \frac{\beta}{z} = 0. \tag{3.11}$$

In the actual calculations, we use the *modified Lax-Wendroff scheme* and solve the equations only over a half domain ($0 \leq x \leq 8$), assuming a symmetry with respect to the z -axis. The number of grid points is $(N_x \times N_z) = (160 \times 200)$, where grid points are distributed uniformly along both the x -coordinate and the z -coordinate. Accordingly, the mesh size is $(\Delta x, \Delta z) = (0.05, 0.2)$.

Although many arcades exist in the calculation region (see Figure 3. 1), our attention is concentrated on the central arcade located within the range of $0 < x < 1$. The other arcades are set in order to make a smooth boundary condition at $x = 1$, that is, if we set a free boundary condition at $x = 1$, the effect of those numerical flows generated in this boundary could not be negligible. From this point of view, we impose the initial perturbation on the central arcade alone and investigate its evolution (see below).

3. 2. 4. Initial Perturbation

Initially the electrical resistivity is set to be zero everywhere except in a small region of $\left[x^2 + (z - h)^2\right]^{1/2} < r$, where η_{init} is assigned for a finite time ($0 < t < 2$). Here η_{init} is a normalized value and defined as the reciprocal of the magnetic Reynolds number, $R_m = C_s L / \eta_{\text{init}}$. Here h and r are those parameters which give the height and radius of the region where the initial perturbation is imposed. In the present study we take various models, each of which has different values of η_{init} , h , and r . All the models in this study are summarized in Table 3. 1. As far as r is concerned, this value is always less than 0.9 because we focus our concentration on the central arcade. Models r1–r4 and models 1– 3 are different in r and η_{init} , respectively. Those models are used in order to investigate the effects of the initial perturbation on the subsequent evolution (see Section 3. 3. 3). Models H1–H4 are different in a vertical scale height of the magnetic field, H . In model M, unlike those models described above, we impose the initial perturbation on four different positions, all of which are distributed along the z -axis.

Since a finite value of the resistivity is initially assigned, the magnetic field begins to dissipate, causing the inflows toward the region subject to the initial perturbation. Such inflows make the magnetic field convex to the axis of symmetry (z -axis) and hence a X-type neutral point (X-point) is formed on this axis. Then, not only the current density increases but also the gas density decreases around this point because a lot of gases are ejected away by the reconnection outflows. In this way, the anomalous resistivity eventually occurs (see § 2. 2. 4 and Appendix A. 3). In the present study, we assume the following form for the anomalous resistivity:

$$\eta = \begin{cases} \frac{1}{150} \left(\left| \frac{v_d}{v_c} \right| - 1 \right) & \text{for } |v_d| > \left| \frac{j_y}{c} \right| |v_c|, \\ 0 & \text{for } |v_d| < |v_c| \end{cases}, \quad (3. 12)$$

unless $\left| \frac{j_y}{c} \right| |v_c|$ exceeds 1. In that case η is fixed to 1. Here $\mathbf{j} = \nabla \times \mathbf{B}$, v_d and v_c are used as the ion-electron drift velocity and a threshold velocity (see Ugai 1986; Yokoyama 1995). The values of v_c used in the present numerical simulation can be seen in Table 3. 1.

3.3. Main Results

3.3.1. Overviews

First, we show the typical evolution of eruptive flares by using model r3. In Figures 3. 2a–2c, we display the time variations of the y -component of the magnetic field (hereafter, this is called the perpendicular magnetic field), temperature, and gas density, respectively. Contour lines and arrows represent the magnetic field lines and velocity field projected onto the (x, z) -plane. Only the area of $0 \leq z \leq 30$ is displayed in these figures, because we take into account the limitation of neglecting the gravity effect. According to Tsuneta (1996), the temperature of active regions in the corona is typically $2 \sim 4$ MK for the background component of the corona. If we assume 3 MK, the pressure scale height is given by about 1.5×10^8 m (see Priest 1982), which is 30 times longer than a half length between the footpoints of a coronal loop in the present study (5000 km, see Table 2. 2). Therefore, a reasonable vertical extent of a coronal loop is at most $0 \leq z \leq 30$ in the present study. In addition, we can avoid the problem of the numerical flows generated at the top free boundary (at $z = 40$) by setting a tentative boundary far from this free boundary.

From these figures, we can find that a magnetic island is formed and rising upward with time. By looking at the panel of $t = 12$ in Figure 3. 2b, it is found that there is a very hot region both at the bottom of the magnetic island and at the top of the closed loops formed in the lower part of the panel. The gas density map (Figure 3. 2c) shows that a high density region of V shape appears in the lower part of the magnetic island at $t = 12$.

Figure 3. 3 shows the time variations of the y -component of the electric field at a neutral point (hereafter, this is called the neutral-point electric field) and the height of a magnetic island of model r3, each of which is represented by a thin solid line and crosses (+), respectively. The height of the magnetic island is determined by measuring the position of the O-point within the magnetic island, where the x -component of the magnetic field changes its sign. Two thick solid lines represent the result of the line-fittings applied to the time variation of the height of the magnetic island. The time ranges of such line-fittings are $0 \leq t \leq 5$ for the first one and $15 \leq t \leq 17$ for the second one, both of which are determined on the basis of the temporal behavior of the neutral-point electric field, that is, the time-range in the first case corresponds to the phase when the electric field does not arise and in the second case corresponds to the phase when it rapidly decreases.

This figure indicates that after the initial perturbation phase ($0 \leq t \leq 2$) the anomalous resistivity is turned on at $t = 5$, which gives rise to the neutral-point electric field. This is because that field is proportional to the rate of the magnetic reconnection at a neutral point (see Forbes & Priest 1983). It reaches its maximum at $t = 14$ when the anomalous resistivity

is quite large, and then it rapidly decreases, because the amount of magnetic fields inflowing toward the neutral point becomes small.

Turning to the dynamics of a magnetic island, we find that the magnetic island is at first slowly going upward ($0 < t < 5$), then it is accelerated ($5 < t < 12$), and finally it rises at a constant rate ($12 < t$). In order to clarify what kinds of mechanism play an important role in such three stages, we pick up three different times ($t = 3, 6, \text{ and } 12$) and investigate the physical situations at such three times, each of which corresponds to the time when the magnetic island is slowly rising, when it is accelerated, and when it rises at a constant rate, respectively. Figures 3.4a–4c show the spacial distributions along the z -axis of the perpendicular magnetic field (B_y , dotted line), vertical velocity (v_z , solid line), and neutral-point electric field (E_y , broken line) at $t = 3, 6, \text{ and } 12$, respectively. The position of an O-point is also represented by a vertical thick solid line in these figures.

At $t = 3$ (see Figure 3.4a), the neutral-point electric field is zero because the anomalous resistivity does not arise yet. By looking at the position of the O-point, we find that the magnetic island rises upward slowly, moving with the upflows produced by the initial resistive perturbation. As this process proceeds, the matter around the neutral point is ejected away and hence the gas density around this point decreases, leading to the occurrence of the anomalous resistivity (at $t = 5$). Owing to the anomalous resistivity, the reconnection outflows become so strong that the magnetic island begins to be accelerated effectively. At $t = 6$ (see Figure 3.4b), we can find that the upward velocity around the O-point is about 1.5 times higher than that at $t = 3$ (see Figure 3.4a). Although the start time of the acceleration of the magnetic island corresponds to the occurrence time of the anomalous resistivity, the electric field at the neutral point does not take its maximum value at that time, as long as a sufficient amount of perpendicular magnetic fields exists around this point (see Figure 3.4b). According to some experimental researches in the magnetic reconnection, the rate of the magnetic reconnection coupled with the perpendicular magnetic field is lower than that without it (Ono, Morita, & Katsurai 1993). When this field is completely lost around the neutral point, then the rate of the magnetic reconnection becomes so high that the neutral-point electric field also takes a large value. This process can be confirmed in Figure 3.4c (at $t = 12.0$). In this stage, the upward velocity of the reconnection outflows becomes very large (reconnection jets) because of the efficient magnetic reconnection so that they form a fast MHD shock at the bottom of the magnetic island. (The position of this shock is shown in Figure 3.4c by a vertical thick broken line.) It is found that the O-point is now behind this shock and therefore the magnetic island is no longer accelerated effectively, rising at a constant rate. (The behavior of the perpendicular magnetic field described above is also recognized in Figure 3.2a. Looking at the panel of $t = 12$, the region around the neutral point (at $z \sim 5$) is colored in white, which means that little perpendicular magnetic fields exists around this point.)

3.3.2. Parameter Dependences

Following the results of § 3.3.1 where we show some basic features of the plasmoid eruption, we then ask whether such features are generally acceptable or not, that is, how the change of the initial resistive perturbation affects the dynamics of plasmoids. To answer this question, we use those models in which we vary the radius of the region where the initial perturbation is imposed (models r1–r4) or the strength of the initially assigned resistivity (models 1–3) as well as we assign the initial resistive perturbation at four different positions (model M) instead of only one position. Then, we make the figures similar to Figure 3.3 for these models and compare such figures with Figure 3.3. As a consequence, we find that there are some similarities among them, such as the existence of three stages in the evolution of a magnetic island and a time-lag between the start time of the acceleration of a magnetic island and the peak time of the neutral-point electric field. However, we find that two things are significantly changed, one of which is the upward velocity of a magnetic island and the other is the occurrence of the coalescence process between several magnetic islands.

Figures 3.5a–5b indicate the variation of the upward velocity of a magnetic island when r and η_{init} are changed. Here the upward velocity is derived from the inclination of lines obtained by means of the line-fittings similar to that in Figure 3.3. In both figures asterisks (*) and crosses (+) represent the upward velocity in the first stage (before the acceleration) and in the third stage (after the acceleration), respectively. Thick curves are the result of such curve-fittings as $v_{upward} = a r^b$ and $c \eta_{init}^d$, where a , b , c , and d are all constant values. In the case of $r = 0$ or $\eta_{init} = 0$, we obtain $v_{upward} = 0$, which means that no initial perturbation is imposed.

These figures tell us that the upward velocity in both the first stage and the third stage decreases as the initial resistive perturbation becomes weak (in a sense that r and η_{init} become 0). However, there is a difference between these two stages in a manner of such decreasing. By looking at the power term in the formula of curve-fittings, it is found that the variation in the third stage is gentler than that in the first stage. We consider that this difference reflects the sensitivity of both stages to the initial resistive perturbation, that is, the upward velocity in the first stage is directly connected with the initial perturbation, while that in the third stage is measured under the situation where the anomalous resistivity already arises so that the effect of the initial perturbation on the third stage is relatively weak.

Figure 3.6a shows the evolution of model M. Top and bottom panels are the temperature and gas density maps, respectively. Contour lines and arrows represent the magnetic field lines and the velocity field projected onto the (x, z) -plane. At $t = 5.0$ we can find four neutral points (at $z = 5, 10, 15,$ and 20), each of which corresponds to the position of the initial perturbation. Figure 3.6b is a similar figure to Figure 3.3 except for model M instead of model r3, showing the time variations of the heights of four magnetic islands and the neutral-point electric field at the lowest neutral point (at $z = 5$) which always has a larger value than at any other neutral points (at $z = 10, 15,$ or 20). Crosses, asterisks, dots, and diamonds represent the

heights of four magnetic islands, while a thin solid line shows the time variation of the neutral-point electric field at the lowest neutral point. A line-fitting is carried out only for the lowest magnetic island, represented by two thick solid lines.

From Figure 3. 6a, it can be seen that the lowest magnetic island continues to be merged with other magnetic islands with time, except for the highest one. At $t = 20$, there appears a well-developed magnetic island between $z = 10$ and $z = 20$, at the bottom of which a hot and dense region is formed. This configuration is quite similar to that at $t = 12$ in Figures 3. 2b–2c. Referring to Figure 3. 6b, we can again find that the coalescence among magnetic islands proceeds and there are some features similar to the case of model r3: the existence of three stages in the evolution of the magnetic island and a time-lag between the start time of the acceleration of the magnetic island and the peak time of the neutral-point electric field.

On the basis of these results, we consider that those two features—the existence of three stages in the evolution of the magnetic island and a time-lag between the start time of the acceleration of the magnetic island and the peak time of the neutral-point electric field—are important factors in understanding the dynamics of plasmoids. In § 3. 5 we discuss such features in detail, but before doing that, several important observational results related to the present study are introduced briefly in the next section.

3. 4. Observational Results

In this section, we show some observational data brought by *Yohkoh*, which help us to understand the mass eruption in solar flares. In addition, we exhibit some interesting results of the plasmoid eruption which is originally analyzed by Ohyama & Shibata (1997), and then make some comments on them.

Figure 3. 7a shows the eruptive process running in a typical LDE (long duration event) flare observed by *Yohkoh* (Tsuneta et al. 1992; Hudson 1994). White arrows indicate the position of an ejected mass (plasmoid). Figure 3. 7b shows a GOES X-ray plot of this event. These figures clearly tell us how the mass eruption proceeds from the preflare phase to the rise phase of solar flares. By looking at the bottom-right panel in Figure 3. 7a, we can find that a mass of V shape is ejected upward. This shape is quite similar to a hot and dense region formed at the bottom of a magnetic island shown in Figures 3. 2b–2c, and Figure 3. 6a.

Figure 3. 8 (from Ohyama & Shibata 1997) indicates the time variations of the height of a plasmoid and the hard X-ray intensity. “Core” and “top” represent the positions of the highest and its 1/e-fold values of the soft X-ray intensity within the plasmoid. This figure

tells us two important facts. One of them is that the plasmoid evolves through several different stages, that is, it rises slowly from 11:04 to 11:15, then it is accelerated in a relatively short time (from 11:15 to 11:18), and finally it rises almost at a constant rate (after 11:18). Moreover, the start time of the acceleration of the plasmoid (11:15) is certainly before the peak time of the hard X-ray intensity (11:18). These features are consistent with the results derived from our numerical simulation. (In this comparison, we assume that the time variation of the neutral-point electric field in our simulation has a direct connection to that of the hard X-ray intensity. That assumption stands on the ground that the observed hard X-ray radiation is produced by the high energy electrons accelerated by the neutral-point electric field, although the precise acceleration mechanism is still not understood completely.)

3. 5. Discussion & Conclusion

In this section we compare the simulation results with the observational one, considering several problems related to the present study and giving our conclusion.

3. 5. 1. Comparison between the Simulation Results and The Observational One

By comparing the simulation results with the observational one, it is found that the evolution of a magnetic island in the former is quite similar to the evolution of a plasmoid in the latter. Comparing Figure 3. 3 to 3. 8, we find that both the magnetic island and the plasmoid have three different stages in their evolutions, that is, the first stage of rising slowly, the second stage of acceleration, and the third stage of rising at a constant rate. Having a constant upward velocity in the third stage is also consistent with the result of Steele & Priest (1989) who studied the prominence eruption. Precisely speaking, Figure 3. 8 indicates that the plasmoid rises very slowly in the first stage, the velocity of which is 20 times smaller than that in the third stage. On the other hand, Figure 3. 3 shows that the initial upward velocity of the magnetic island is only a half smaller than the final one. We explain this discrepancy by considering the effect of the initial resistive perturbation on the flare evolution. From Figures 3. 5a and 3. 5b, as long as the anomalous resistivity occurs, weak initial resistive perturbations reduce the initial upward velocity of the magnetic island, while it does not change the final upward velocity so much. Therefore, one possibility is that the initial resistive perturbation proceeds very weakly in the real circumstances. Another possibility is derived from Figure 3. 6b. In this figure there are four magnetic islands, but we think that the lowest one corresponds to the observed plasmoid because a hot and dense region is formed only inside the lowest magnetic island (see Figure 3. 6a). Model r3 in Figure 3. 3 and model M in Figure 3. 6 have the common parameters characterizing the initial resistive perturbation ($r = 0.7$, $v_{init} = 1 / 15$) except for the number of the regions where the initial perturbation is imposed. However the initial upward velocity in Figure 3. 3 is 60 % of that in Figure 3. 6. This is explained by the fact that the lowest magnetic island of model M is prevented from rising freely by other

upper magnetic islands. This implies that the initial resistive perturbation probably occurs at many positions within a magnetic arcade.

Next, when it comes to the final upward velocity of the magnetic island, it is about $0.4 v_{A0}$, where $v_{A0} = B_0 / [4 \mu_0]^{1/2}$ and in the present study this value is given by $v_{A0} / c_{S0} = [2 / \mu_0]^{1/2} \sim 2.45$. In Magara et al. (1996b), we discussed that the observed ejection velocity of a plasmoid was always quite smaller than v_{A0} , and hence the present result is along this tendency.

3. 5. 2. Does the Anomalous Resistivity Really Arise?

In § 3. 5. 1 we consider the evolution of eruptive flares, but there is one crucial point in that consideration. The question is whether the anomalous resistivity always sets in or not. In order to answer this question, we perform those numerical simulations in which we do not assign the anomalous resistivity after the initial resistive perturbation finishes, and investigate the time variation of v_d (the ion-electron drift velocity). These simulations are represented by models H1–H4, each of which has a different vertical scale height (see Table 3. 1). Figure 3. 9 shows the time variation of v_d at the neutral point ($v_{d, neutral}$) of model H1. Here $v_{d, neutral}$ is always the maximum value of v_d within the range of $(0, 0) < (x, z) < (1, 30)$.

Figure 3. 9 indicates that $v_{d, neutral}$ is initially increasing with time, then rapidly decreasing. We think that this abrupt decreasing is due to the effect of finite mesh size in our numerical simulation. In the actual situations in the corona, we expect that $v_{d, neutral}$ continues to increase with time, which probably leads to the occurrence of the anomalous resistivity. Moreover, according to Schumacher & Kliem (1996), the current density at a neutral point takes a larger maximum value as the Lundquist number is getting larger, which is suitable to the real circumstances of the corona. Hence their result also supports the idea that the rapid increase of v_d occurs in the corona and the anomalous resistivity eventually sets in. From this point of view, we do such a curve-fitting as $v_{d0} e^{kt}$ within the range of $0 < t < 8$, where v_{d0} and k are constant values. The result of this fitting is represented by a thick solid curve in Figure 3. 9.

Next, we do the same curve-fittings as in Figure 3. 9 for other models (models H2–H4). In this way the (k, H) -relation is obtained, displayed in Figure 3. 10a. This relation is characterized by $H = 5.53 e^{6.05k}$, which means that $(k, H) = (0, 5.5)$ is the limiting case where it takes an infinite time for $v_{d, neutral}$ to reach its rapid increase phase. In Figure 3. 10b, we make a similar plot to that in Figure 3. 10a for other models, all of which have a different height ($h = 3$) of the region where the initial perturbation is imposed, compared to models H1–H4. As for these models, the (k, H) -relation is given by $H = 4.55 e^{6.79k}$, which shows that $(k, H) = (0, 4.55)$ is the limiting case. Therefore, in order to cause the anomalous resistivity, we need those magnetic arcades whose vertical scale height is several times larger

than its horizontal extent, which is consistent with the conclusion of Kusano et al. (1995).

3. 5. 3. Boundary Effects

As we mention in § 3. 1, a lot of investigations of the evolution of eruptive flares by means of numerical simulations has been made recently. This paper belongs to this kind of work, where there has been a controversy about the effect of boundary conditions. Mikic et al. (1988) set a symmetric boundary condition at the side boundary, which was reconsidered to have an artificial effect on the evolution of a magnetic arcade by Biskamp & Welter (1989), who suggested that the effect of the neighboring arcades could cause eruptive phenomena. Instead of the effect of the neighboring arcades, van Ballegooijen & Martens (1989) and Inhester et al. (1992) considered the effect of a photospheric converging flow toward the neutral line of a magnetic arcade. Finn et al. (1992) discussed the effect of the time-varying flow pattern imposed on a double arcade system. As for our present study, we set an antisymmetric boundary at $x = 8$, which is far from a central arcade ($0 < x < 1$), so that we can safely neglect the effect of this boundary. However, taking a periodic arcade system as our model is not suitable for studying the evolution of a single arcade system in the sun. This weak point becomes crucial in the third stage of the evolution, when a lot of magnetic fields in the neighboring arcades begins to flow into the region of the central arcade.

3. 5. 4. Limitation of The Present Study

As mentioned above, we start with a periodic linear force-free magnetic field in order to study the evolution of a single arcade system. This means that we assume the distribution of the photospheric magnetic field to be too simple, compared to its more complicated distribution in the real circumstances. That simplification brings several limitations to our present study. First, since we take only one mode of the solution of equation (3. 4), we must use a very tall arcade which has a tremendous amount of magnetic energy in order to construct a current sheet within a magnetic arcade. Thanks to this, we cannot make a quantitative estimate about the released energy during the flare evolution. Second, as we mention in § 3. 5. 2, since the effect of the neighboring arcades cannot be neglected in the third stage of the flare evolution, we do not discuss the feature of this stage in detail, that is, although it is found that a magnetic island rises at a constant rate in this stage, it is still an open question whether such motion continues all through this stage. As a matter of fact, taking a careful look at Figure 3. 3 tells us that the upward velocity is gradually decreasing after $t = 20$, unlike the results of our previous study (Magara et al. 1996b). In Magara et al. (1996b), we studied the magnetic reconnection by assuming the initial magnetic field to be antiparallel, where a magnetic island continues to rise at a constant rate. In our present study, on the other hand, the closed structure of the overlying magnetic field exerts the downward force to the magnetic island, which contributes to the reduction of its upward velocity. In addition, we cannot neglect other important effects, such as the viscosity, the radiative cooling, and crossing other

magnetic fields (S. Koutchmy 1996, private communication; C. Delannée 1996, private communication). The researches including these effects should be done in the future.

Another uncertainty is the origin of the localization of the initial resistive perturbation. Biskamp & Welter (1989) mentioned the possibility that some resistive structures develop inside a magnetic arcade. Priest (1988) suggested that the rise of the overlying magnetic arcade induced the resistive processes beneath it. Wiechen, Büchner, & Otto (1996) studied the features of the resistive instability developing within a magnetic arcade. They emphasized the importance of the localization of the resistivity. Schumacher & Kliem (1996) showed the dynamical features of a current sheet subject to the anomalous kinetic instability.

3. 5. 5. Conclusion

Considering the results of our present simulation, now we describe how those flares associated with mass eruptions evolve from the preflare phase to the gradual phase. We find three different stages in the mass eruption, corresponding to the preflare phase, the impulsive or rise phase, and the gradual phase of solar flares, and then we investigate the dynamical features of each stage in detail. Figure 3. 11 shows a schematic illustration of the evolution of eruptive flares. Crosses (+) and a thick curve represent the time variations of the height of a plasmoid and the hard X-ray intensity, respectively. In this figure the initial resistive instability starts inside a magnetic arcade at $t = t_1$, forming several small plasmoids. The lowest plasmoid is lifted slowly by the outflows from a neutral point, which is produced by the initial resistive perturbation. As the plasmoid rises, the gas density decreases and the spacial gradient of the magnetic field increases around the neutral point, which means that the ion-electron drift velocity increases, leading to the occurrence of the anomalous resistivity (at $t = t_2$). Then those reconnection outflows from the neutral point are rapidly enhanced, beginning to accelerate the plasmoid (from t_2 to t_3). At the same time when the anomalous resistivity sets in, the neutral-point electric field arises, but it does not reach its maximum value as long as a sufficient amount of perpendicular magnetic fields exists around the neutral point, because that field inhibits the efficient magnetic reconnection. At $t = t_3$, the neutral-point electric field reaches its maximum, which implies that the high energy electrons can be generated and hence the high-energy electromagnetic wave such as hard X-ray can be emitted. At this time, the magnetic reconnection becomes so lively that it produces the reconnection jets, which form a fast MHD shock at the bottom of the plasmoid. After that, since an O-point (center part of the plasmoid) is located behind this shock, the strong acceleration ends and the plasmoid rises almost at a constant rate.

Since the role that the perpendicular magnetic field plays in this scenario of the flare evolution is very important, we make a detailed study about this topic in the future. This kind of study is also seen in Birn & Hesse (1991), where they discussed the magnetic reconnection in the Earth's magnetotail and found the concentration of the electric field around a X-type

neutral point.

3. 6. Summary of Chapter 3

Finally, we summarize main results of this chapter as follows:

1. We investigate the dynamics of plasmoids in eruptive flares. There are three different stages in the mass eruption. In the first stage, a magnetic island is slowly lifted by the upflows produced by the initial resistive perturbation, and then begins to be accelerated when the anomalous resistivity sets in. This acceleration stops after a fast MHD shock is formed at the bottom of the magnetic island, which implies that the upflows around the center part of it are no longer strong. However, it is still an open question whether it continues to rise at a constant rate after that.
2. We find that a time-lag between the start time of the acceleration of a magnetic island and the peak time of the neutral-point electric field exists. This is explained by considering the inhibition of the efficient magnetic reconnection by the perpendicular magnetic field. We think that this time-lag reflects the observed time-lag between the start time of the acceleration of a plasmoid and the peak time of the hard X-ray intensity, reported by Ohyama & Shibata (1997).
3. Comparing the simulation results with the observational one, we find that the observed initial upward velocity of a plasmoid is quite slower than that in our simulation. By studying the effects of the initial perturbation, we conjecture that the initial resistive perturbation occurs very weakly and at many positions within a coronal magnetic arcade.
4. The case where the vertical extent of a magnetic arcade is several times larger than its horizontal extent is expected to lead to the occurrence of the anomalous resistivity. This is consistent with the conclusion of Kusano et al. (1995).
5. We do not investigate how the initial localization of the resistivity proceeds, which is discussed in our future study.

Chapter 4. 3-Dimensionality of Solar Phenomena

Abstract

In this chapter, we discuss those 3-dimensional solar phenomena which are enumerated as follows:

1. Emergence of twisted magnetic-flux tubes and the magnetic reconnection between the emerging magnetic field and the overlying magnetic field,
2. Interchange instability and the generation of turbulences in the loop-top region of the postflare loops,

and

3. Eruption of twisted magnetic ropes (plasmoids in a 3-dimensional viewpoint).

In addition, we explain a numerical code developed for the 3-dimensional MHD numerical simulation and show some interesting results obtained by using this code.

4. 1. Emergence of Twisted Magnetic-Flux Tubes

So far, many workers have made a theoretical consideration to the problem of the emerging subphotospheric magnetic field mainly from a 2-dimensional viewpoint. The concept of this is shown in Figure 4. 1. The subphotospheric magnetic field begins to rise because of the buoyancy effect (or Parker instability) and eventually emerges on the surface. Then, owing to the rapid decrease of the surrounding gas pressure, it begins to expand outward and finally interacts with the overlying coronal magnetic field. This interaction has been believed to cause various activities observed in the corona. For example, Heyvaerts, Priest, & Rust (1977) presented a loop flare model based on that interaction. Shibata et al. (1989, 1992) performed the 2-dimensional MHD numerical simulation and investigated the emergence and expansion of the emerging magnetic field. Yokoyama & Shibata (1996) also studied the interaction between the emerging magnetic field and the overlying coronal field by means of the 2-dimensional MHD numerical simulation. These studies fairly contributed to our understanding of the basic processes running in the emerging flux region (EFR) of the sun and clarified many aspects of those active phenomena observed in these regions.

However, the real image of the emergence of magnetic fields cannot be understood

completely as long as we stand on a 2-dimensional viewpoint. This will be easily acceptable if we look at Figure 4. 2, which shows how the emergence really proceeds in the sun (see Kurokawa 1989). The left panels are observed in white light and the right one are in H . From this figure, it is found that as the emergence proceeds, the direction of the emerging magnetic field does not remain constant but rotates, that is, the direction of the successive emerging magnetic field ($f_2 - p_2, f_3 - p_3$) is diverted from the initial direction ($f_1 - p_1$). Precisely speaking, as we goes to the later phase, we find that the magnetic field tends to emerge along the neutral line of the EFR. This observational fact strongly suggests that the emerging magnetic field has not only the poloidal component but also the toroidal one, the latter of which causes the rotation of the direction of the emerging magnetic field. The existence of the toroidal magnetic field causes the twisted structure of the emerging magnetic field and the formation of a sheared magnetic arcade on the solar surface (see Figure 4. 3).

Figure 4. 3 represents the emergence of the cylinder-like flux tube having the toroidal magnetic field as its axial component. This is an example of the emergence of twisted magnetic-flux tubes. Supposing that such flux tube emerges from the subphotosphere, we can expect that there appears the sheared EFR seen in Figure 4. 2. The relation between a 2-dimensional viewpoint and a 3-dimensional viewpoint is also indicated in this figure.

It is worth saying that when we try to do detailed researches on this kind of 3-dimensional phenomena, we need to know the information about the photospheric motion. Recently, some workers reported those interesting results where a small but significant flow pattern on the photosphere are clearly recognized (see November & Simon 1988; Kitai et al. 1997).

4. 2. Interchange Instability in The Loop-Top Region

In Chap. 2, we discuss a hot and dense region formed at the top of the postflare loops. This region is maintained by the *magnetic pressure walls*, which prevents a plasma from draining sidewise (see Figure 2. 5c and 2. 9). However, real flares have a 3-dimensional structure so that the consideration of the third direction is needed for the full comprehension of the loop-top region. We show the 3-dimensional structure around this region in Figure 4. 4. From a 3-dimensional viewpoint, the loop-top source is no longer regarded as a circle-like source but a cylindrical one. Since such configuration is similar to that of a magnetically confined plasma in laboratories, we expect that the interchange instability can occur in the loop-top region in the same way as in the laboratory case. This instability causes the bulk motion in the loop-top region, which is maybe detected as the turbulent component of line broadening when we observe a particular spectral line of this region.

4.3. Eruption of Twisted Magnetic Ropes

In Chap. 3, we make a detailed investigation on the plasmoid eruption and derive several dynamical properties of this phenomenon. Through that study, a plasmoid is identified with a magnetic island, or the region of a bundle of closed magnetic field lines. However, this image fails when we stand on a 3-dimensional viewpoint, in which a plasmoid is not regarded as such a disk-like thing but as a helical rope whose footpoints are possibly fixed on the surface of the sun. This situation is shown in Figure 4. 5, which also shows other typical aspects of eruptive phenomena in the sun. At first, a sheared magnetic arcade having the spacial variation along its neutral line appears on the surface, which is probably formed by the emergence of a twisted magnetic-flux tube described in § 4. 1. Then, a current sheet is formed within the arcade probably through some ideal MHD processes. Later the resistive process begins to run in the sheet, which eventually leads to the fast-magnetic-reconnection stage, in which the violent energy release occurs. This corresponds to the impulsive phase of solar flares (or rise phase of the LDE flares). Finally, there are formed the postflare loops on the surface of the sun, while a plasmoid of helical structure is ejected upward with its footpoints fixed on the surface.

4.4. Numerical Code for The 3-Dimensional MHD Numerical Simulation

In this section, we explain a numerical code for the 3-dimensional MHD numerical simulation based on the *modified Lax-Wendroff scheme*. First, we write the basic MHD equations in the conservational form:

$$-\frac{\partial}{\partial t} \mathbf{w} + \frac{\partial}{\partial x} \mathbf{E}(\mathbf{w}) + \frac{\partial}{\partial y} \mathbf{F}(\mathbf{w}) + \frac{\partial}{\partial z} \mathbf{G}(\mathbf{w}) + \mathbf{S}(\mathbf{w}) = \mathbf{0}. \quad (4.1)$$

Here

$$\mathbf{w} = \begin{pmatrix} v_x \\ v_z \\ B_x \\ B_z \\ \frac{P}{-1} + \frac{1}{2} [v_x^2 + v_z^2 + v_y^2] + \frac{1}{8} [B_x^2 + B_z^2 + B_y^2] \\ v_y \\ B_y \end{pmatrix}, \quad (4.2)$$

$$\mathbf{E}(\mathbf{w}) = \begin{pmatrix} v_x \\ v_x^2 + p_1 + \frac{1}{8} [B_{1y}^2 + B_{1z}^2 - B_{1x}^2] + \\ + \frac{1}{4} [B_{0y} B_{1y} + B_{0z} B_{1z} - B_{0x} B_{1x}] \\ v_z v_x - \frac{1}{4} [B_{0z} B_{1x} + B_{1z} B_{0x} + B_{1z} B_{1x}] \\ 0 \\ B_z v_x - B_x v_z + \overline{R}_m j_y \\ \left\{ \frac{1}{2} [v_x^2 + v_z^2 + v_y^2] + \frac{P}{-1} \right\} v_x + \\ + \frac{B_z}{4} [B_z v_x - B_x v_z + \overline{R}_m j_y] + \frac{B_y}{4} [B_y v_x - B_x v_y - \overline{R}_m j_z] \\ v_y v_x - \frac{1}{4} [B_{0y} B_{1x} + B_{1y} B_{0x} + B_{1y} B_{1x}] \\ B_y v_x - B_x v_y - \overline{R}_m j_z \end{pmatrix}, \quad (4.3)$$

$$\mathbf{F}(\mathbf{w}) = \begin{pmatrix} v_y \\ v_x v_y - \frac{1}{4} [B_{0x} B_{1y} + B_{1x} B_{0y} + B_{1x} B_{1y}] \\ v_z v_y - \frac{1}{4} [B_{0z} B_{1y} + B_{1z} B_{0y} + B_{1z} B_{1y}] \\ B_x v_y - B_y v_x + \overline{R}_m j_z \\ B_z v_y - B_y v_z - \overline{R}_m j_x \\ \left\{ \frac{1}{2} [v_x^2 + v_z^2 + v_y^2] + \frac{P}{-1} \right\} v_y + \\ + \frac{B_x}{4} [B_x v_y - B_y v_x + \overline{R}_m j_z] + \frac{B_z}{4} [B_z v_y - B_y v_z - \overline{R}_m j_x] \\ v_y^2 + p_1 + \frac{1}{8} [B_{1x}^2 + B_{1z}^2 - B_{1y}^2] + \\ + \frac{1}{4} [B_{0x} B_{1x} + B_{0z} B_{1z} - B_{0y} B_{1y}] \\ 0 \end{pmatrix}, \quad (4.4)$$

$$\mathbf{G}(\mathbf{w}) = \begin{pmatrix} v_z \\ v_x v_z - \frac{1}{4} [B_{0x} B_{1z} + B_{1x} B_{0z} + B_{1x} B_{1z}] \\ v_z^2 + p_1 + \frac{1}{8} [B_{1y}^2 + B_{1x}^2 - B_{1z}^2] + \\ + \frac{1}{4} [B_{0y} B_{1y} + B_{0x} B_{1x} - B_{0z} B_{1z}] \\ B_x v_z - B_z v_x - \overline{R}_m j_y \\ 0 \\ \left\{ \frac{1}{2} [v_x^2 + v_z^2 + v_y^2] + \frac{P}{-1} \right\} v_z + \\ + \frac{B_x}{4} [B_x v_z - B_z v_x - \overline{R}_m j_y] + \frac{B_y}{4} [B_y v_z - B_z v_y + \overline{R}_m j_x] \\ v_y v_z - \frac{1}{4} [B_{0y} B_{1z} + B_{1y} B_{0z} + B_{1y} B_{1z}] \\ B_y v_z - B_z v_y + \overline{R}_m j_x \end{pmatrix}, \quad (4.5)$$

and

$$\mathbf{S}(\mathbf{w}) = \begin{pmatrix} 0 \\ {}_1 g_x \\ {}_1 g_z \\ 0 \\ 0 \\ Q \\ 0 \\ {}_1 g_y \\ 0 \end{pmatrix}, \quad (4.6)$$

where g_i ($i = x, y, z$) and Q mean the rate of gravitational acceleration and the energy source such as the radiative cooling. In this formula, $X_1 = X - X_0$, where X_0 is the initial value of physical variables.

Next, we transform equation (4. 1) into a set of difference equations which can be solved by means of the 2-step *modified Lax-Wendroff scheme*. From now on we pick up one component from a vector equation (4. 1) as a representative and write this with plain characters.

<1st step>

$$\begin{aligned} \frac{1}{t} \left(w_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - w_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^n \right) &= -\frac{1}{x} \left(E_{i+1, j+\frac{1}{2}, k+\frac{1}{2}}^n - E_{i, j+\frac{1}{2}, k+\frac{1}{2}}^n \right) \\ &- \frac{1}{y} \left(F_{i+\frac{1}{2}, j+1, k+\frac{1}{2}}^n - F_{i+\frac{1}{2}, j, k+\frac{1}{2}}^n \right) \\ &- \frac{1}{z} \left(G_{i+\frac{1}{2}, j+\frac{1}{2}, k+1}^n - G_{i+\frac{1}{2}, j+\frac{1}{2}, k}^n \right) \\ &- S_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^n. \end{aligned} \quad (4.7)$$

Here t is the time spacing and x , y , and z are the mesh sizes along the x -coordinate,

y-coordinate, and z-coordinate, respectively, while n is the time-step number and i, j , and k mean the position number of a grid point measured along the x-coordinate, y-coordinate, and z-coordinate, respectively. Other characters are defined as follows.

$$\begin{aligned}
 w_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^n \dots & \quad \text{the value of } w^n \text{ at a virtual mid-mesh point, } \left(x_{i+\frac{1}{2}}, y_{j+\frac{1}{2}}, z_{k+\frac{1}{2}}\right), \\
 E_{i+1, j+\frac{1}{2}, k+\frac{1}{2}}^n \dots & \quad \text{the value of } E^n \text{ at a virtual mid-mesh point, } \left(x_{i+1}, y_{j+\frac{1}{2}}, z_{k+\frac{1}{2}}\right), \\
 F_{i+\frac{1}{2}, j+1, k+\frac{1}{2}}^n \dots & \quad \text{the value of } F^n \text{ at a virtual mid-mesh point, } \left(x_{i+\frac{1}{2}}, y_{j+1}, z_{k+\frac{1}{2}}\right), \\
 G_{i+\frac{1}{2}, j+\frac{1}{2}, k+1}^n \dots & \quad \text{the value of } G^n \text{ at a virtual mid-mesh point, } \left(x_{i+\frac{1}{2}}, y_{j+\frac{1}{2}}, z_{k+1}\right), \\
 S_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^n \dots & \quad \text{the value of } S^n \text{ at a virtual mid-mesh point, } \left(x_{i+\frac{1}{2}}, y_{j+\frac{1}{2}}, z_{k+\frac{1}{2}}\right),
 \end{aligned}$$

and so on.

<2nd step>

Owing to the 1st-step calculation described above, we obtain the next time-step values at every virtual mid-mesh point. Then we try to obtain the values at every real grid point.

$$\begin{aligned}
 w_{i,j,k}^{n+1} &= w_{i,j,k}^n \\
 & - \frac{t}{x} \left(\frac{1}{4} [E_{i+1,j,k}^n - E_{i-1,j,k}^n] + \frac{1}{8} \hat{E}^{n+1} \right) \\
 & - \frac{t}{y} \left(\frac{1}{4} [F_{i,j+1,k}^n - E_{i,j-1,k}^n] + \frac{1}{8} \hat{F}^{n+1} \right) \\
 & - \frac{t}{z} \left(\frac{1}{4} [G_{i,j,k+1}^n - G_{i,j,k-1}^n] + \frac{1}{8} \hat{G}^{n+1} \right) \\
 & - \left(\frac{1}{2} S_{i,j,k}^n + \frac{1}{16} \hat{S}^{n+1} \right).
 \end{aligned} \tag{4.8}$$

Here

$$\begin{aligned}
 \hat{E}^{n+1} &= E_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - E_{i-\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} \\
 & + E_{i+\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} - E_{i-\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} \\
 & + E_{i+\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} - E_{i-\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
 & + E_{i+\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1} - E_{i-\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1},
 \end{aligned}$$

$$\begin{aligned}
\hat{F}^{n+1} &= F_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - F_{i+\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} \\
&+ F_{i-\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - F_{i-\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} \\
&+ F_{i+\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} - F_{i+\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
&+ F_{i-\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} - F_{i-\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1},
\end{aligned}$$

$$\begin{aligned}
\hat{G}^{n+1} &= G_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - G_{i+\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
&+ G_{i-\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} - G_{i-\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
&+ G_{i+\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} - G_{i+\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
&+ G_{i-\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} - G_{i-\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1},
\end{aligned}$$

and

$$\begin{aligned}
\hat{S}^{n+1} &= S_{i+\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} + S_{i+\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} + S_{i+\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} + S_{i+\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1} \\
&+ S_{i-\frac{1}{2}, j+\frac{1}{2}, k+\frac{1}{2}}^{n+1} + S_{i-\frac{1}{2}, j-\frac{1}{2}, k+\frac{1}{2}}^{n+1} + S_{i-\frac{1}{2}, j+\frac{1}{2}, k-\frac{1}{2}}^{n+1} + S_{i-\frac{1}{2}, j-\frac{1}{2}, k-\frac{1}{2}}^{n+1}.
\end{aligned}$$

The code shown above is based on the case of uniformly distributed grid points. The case of non-uniformly distributed grid points is explained in Appendix B in this thesis.

By using this code, we investigate the maximum linear growth rate of the tearing instability and compare it to the analytically predicted value. The result is that the discrepancy between both values is quite small. For example, we assume such an antiparallel force-free configuration as is described in the following way:

$$B_x = 0, \quad (4.9)$$

$$B_y = B_0 \cosh^{-1}(4 - x), \quad (4.10)$$

$$B_z = -B_0 \tanh(4 - x), \quad (4.11)$$

$$P = P_0, \quad (4.12)$$

$$= 0, \quad (4.13)$$

and

$$T = T_0, \quad (4.14)$$

where $P_0 = 1 / \dots$, $B_0 = [8 P_0 /]^{1/2}$, $T_0 = P_0 / 0$. When $= 5 / 3$, $= 0.15$, and $R_m = 1000$,

we obtain the growth rate of 1.56 from the numerical simulation, while the predicted one is 1.49; consequently, the discrepancy is $(1.56 - 1.49) / 1.49 \sim 0.047$.

4.5. Results of The 3-Dimensional MHD Numerical Simulation

We then show some interesting results of the 3-dimensional MHD numerical simulation and make some comments on them. First of all, we explain the calculation configuration. Figure 4. 6 shows the calculation region and the initial configuration of the present numerical simulation. Initially, physical variables are set to be in the same force-free state as is given in the previous section. Grid points are distributed uniformly along the y -coordinate ($y = 1.0$) but non-uniformly along the x -coordinate (from 0.02 to 0.1, toward the side boundary) and z -coordinate (from 0.4 to 2.0, toward the top boundary). Although the actual calculation region is within the range of $(0, 0, 0) \leq (x, y, z) \leq (5, 25, 50)$, we show the region of $(-5, 0, 0) \leq (x, y, z) \leq (5, 25, 50)$ in this figure, by assuming a symmetry with respect to the (y, z) -plane. As the initial perturbation, we impose the locally enhanced resistivity constantly, that is, we set $R_m = 1000$ within the range of $x^2 + (z - 5)^2 \leq 0.4$ and $11 \leq y \leq 14$, while setting $R_m = 0$ elsewhere. The region of the enhanced resistivity lies in the center part of the calculation region, displayed as a shaded area in Figure 4. 6.

Next, we refer to some results of the present numerical simulation. Figure 4. 7 is a 3-dimensional viewgraph displaying how the magnetic field evolves under the locally enhanced resistivity described above. Red, blue, light blue, green, and yellow lines represent the magnetic field lines at different positions; red line initially lies far from the neutral line (y -axis) so that its tilt against the vertical direction is weak, while light blue, green, and yellow lines initially lie near the neutral line so that their tilts are strong. Blue line is initially set to be just above the neutral line so that its direction is completely horizontal. It should be mentioned that we draw these field lines on the basis of the Euler description. Hence, although the colors of field lines are common at different times, such field lines have no relation with each other. A color map displayed at the base represents the temperature. Elapsed times are $t = 0$ (top-left), 4 (top-right), 8 (bottom-left), and 12 (bottom-right). Figure 4. 8 is a picture similar to Figure 4. 7, apart from a different viewing angle.

Both of Figure 4. 7 and 4. 8 clearly indicate some prominent features of this simulation. First, as a result of the magnetic reconnection, a sheared arcade is formed near the base. Moreover, by looking at the temperature map at the base, we find that there are formed two hot sources near the footpoints of the arcade, which are produced by the hot plasma flows originating from the region where the resistivity is imposed (see the panels at $t = 8$, and 12 in both figures). Such two sources develop along the neutral line with time, which may correspond to H α ribbons observed in solar flares. (Actually, the thermal conduction probably plays a more dominant role in the formation of H α ribbons rather than the hot plasma flows.) Casting

our eyes to the upper part of the panel, we find that a twisted magnetic-field line appears (blue line), which corresponds to a helical rope, or an extended image of the magnetic island seen in many 2-dimensional MHD numerical simulations. (In the present study, an open-field configuration is assumed as the initial state so that the complete helical structure is not reproduced.)

Figure 4. 9 is a figure similar to Figure 4. 7 except that another plane parallel to the (x, z) -plane is added to Figure 4. 7. This plane lies at $y = 12$ and has a color map which represents the temperature. The color map on that plane at $t = 12$ gives similar features that are already seen in the 2-dimensional MHD numerical simulation in this thesis, that is, there is formed a superhot region around both the reconnection region (neutral-point region) and the loop-top region (see Figure 2. 5a).

Appendix A. Summary of The Important Physical Processes in Solar Flares

In this appendix, we make a brief explanation on some important physical processes in solar active phenomena. These are the magnetic reconnection, the resistive tearing process, and the anomalous resistivity, all of which have a close relationship with our studies in this thesis.

A. 1. Magnetic Reconnection

Solar flares are those phenomena that produce a huge amount of thermal, non-thermal, and kinetic energies, the resource of which is now believed to be the magnetic energy stored in the quasi-static coronal magnetic field. Since the corona is a good conductive medium but still has a finite resistivity, the magnetic energy can be released in the diffusion process. The region where the energy release works efficiently is known as a current sheet, in which a certain component of the magnetic field is abruptly changed across the sheet. Hence the current density in that region is so high that the Ohmic dissipation becomes very effective. The formation of current sheets in the coronal environment is studied in many works (e.g. Parker 1994; Karpen, Antiochos, & DeVore 1996).

The most fundamental time scale of such diffusion process in a current sheet is the so-called diffusion time τ_d defined as

$$\tau_d = \frac{l^2}{v_d}, \quad (\text{A. 1. 1})$$

where l is a typical thickness of the sheet. A typical velocity of the diffusion process is expressed as

$$v_d = \frac{l}{\tau_d} = \frac{l}{\tau}. \quad (\text{A. 1. 2})$$

The quickness of this process is defined as the ratio of v_d to another dynamical velocity, such as the Alfvén velocity v_A ,

$$\frac{v_d}{v_A} = \frac{l}{v_A \tau} = R_m^{-1} \frac{L}{l}, \quad (\text{A. 1. 3})$$

where L is a typical length of the sheet and $R_m = v_A L / \eta$ is the magnetic Reynolds number. Since the coronal value of the magnetic Reynolds number is so large that the diffusion process runs very slowly, which cannot explain the violent energy release observed in real flares typically proceeding at the rate of the Alfvén velocity.

The above simple estimation is followed by a more detailed consideration in which the coupling between the velocity field and the magnetic field is taken into account, that is, the behavior of the magnetic field under a particularly prescribed velocity field is studied. This approach is classified into the kinematic treatment and Parker (1979) did that study where he investigated the behavior of the magnetic field around a stagnation point (X-point), finding out the necessity of the localization of the resistivity around this point.

The dynamical treatment of a current sheet in which the equation of motion is taken into account has been also done in many works. The first model based on this treatment is the so-called *Sweet-Parker model* (Sweet 1958; Parker 1963), which has a relatively long diffusion region. In this model the energy release process is regarded as the magnetic annihilation, and the rate of it is given by

$$\frac{v_{inflow}}{v_A} \propto R_m^{-1/2}, \quad (\text{A. 1. 4})$$

where v_{inflow} represents the inflow velocity to the sheet. This rate is less dependent on the magnetic Reynolds number than the simple diffusion rate (see equation (A. 1. 3)), though it is still too small to satisfy the time scale of real flares.

After the *Sweet-Parker model* was found to be insufficient to explain the quickness of the flare evolution, Petschek (1964) proposed a more efficient model where the rate of the magnetic reconnection is almost independent of the magnetic Reynolds number. The feature of this model lies in the existence of a small diffusion region and a pair of slow MHD shocks, which enables the fast magnetic reconnection. The rate of this process is given by

$$\frac{v_{inflow}}{v_A} \propto \frac{1}{\ln R_m}. \quad (\text{A. 1. 5})$$

Moreover, Sonnerup added another discontinuity to the *Petschek model* and found out that the magnetic reconnection could proceed at the rate of the Alfvén velocity (Sonnerup 1970). Later, Priest & Forbes (1986) tried to unite several different models for the magnetic reconnection and came to the conclusion that the reconnection model was dependent on the flow pattern appearing in the system. According to their results, the *Petschek model* and the *Sonnerup-like model* have a connection with the converging flow pattern and the diverging flow pattern, respectively.

Those models described above are based on the assumption of steady state. The next problem is to investigate how the system evolves to that steady state naturally, that is, what is a crucial factor to form that particular steady state. In order to answer this question, we have to trace the temporal evolution of the system, which needs the help of numerical simulations.

According to Biskamp (1986, 1993), it is found that the Petschek's configuration does not occur but the Sweet-Parker's does under a small uniformly distributed resistivity. On the contrary, Ugai & Tsuda (1977) and Sato & Hayashi (1979) used a locally enhanced resistivity model and confirmed the occurrence of the Petschek's configuration. Their results suggest that the fast magnetic reconnection with the Petschek's configuration should have a close relationship to the local enhancement of the resistivity. The problem of how the localization of the resistivity arises in the flare evolution is fairly important so that this will be studied in our future study.

A 3-dimensional magnetic reconnection has been an important topic recently, especially because we must study the magnetic reconnection in such a complicated region as an active region of the sun, where the structure of the magnetic field is purely 3-dimensional. The fundamental researches in this field is briefly introduced in Priest (1991), Biskamp (1993), Ugai & Shimizu (1996), and so on. Démoulin et al. (1996) applied the concept of the 3-dimensional magnetic reconnection to the coronal magnetic field and discussed the concentration of currents on quasi-separatrix layers created by smooth photospheric motions.

A. 2. Resistive Tearing Process in A Current Sheet

Resistive tearing process is one of the most fundamental processes occurring spontaneously in a current sheet. If the length of the sheet is about ten times larger than its thickness, such sheet becomes unstable to the tearing instability and a lot of modes begin to develop at their own linear growth rates. Theoretical considerations of this subject began with the work of Furth, Killen, & Rosenbluth (1963) (FKR theory), who divided the region subject to that instability into two areas, that is, the external ideal MHD area and the internal resistive MHD area, with matching conditions imposed on the interface of these areas. Then they derived the dispersion relation in the linear stage. The maximum growth rate has a $R_m^{-1/2}$ -dependence and this is the same dependence on the magnetic Reynolds number as that of the *Sweet-Parker model* for the magnetic reconnection (apart from the length scale), so that the resistive tearing process develops slowly at the linear stage when the magnetic Reynolds number is quite large.

Although the tearing instability shows an exponential-like growth in the linear stage, such growth gradually slows down as the evolution enters on the nonlinear stage. This is one of the important features in the nonlinear tearing process, that is, an exponential growth in the linear stage is reduced and even replaced by an algebraic one in some cases (Rutherford 1973; Steinolfson & van Hoven 1984). As a theory of this stage, we know the *Rutherford regime*, which indicates that the width of a magnetic island increases with time at a constant rate (Rutherford 1973, Biskamp 1993). The saturation level of this width is also considered by White (1977). For example, the saturation width becomes about twice larger than the

initial thickness of the current sheet when the wavelength of the tearing mode is about 10 times larger than the initial thickness of the sheet (see Fig. 5. 12 in Biskamp 1993).

A coalescence process also plays a dominant role in the nonlinear stage. Through the coalescence process between magnetic islands, a sufficient amount of thermal energy is produced by the Joule heating, leading to the formation of a large magnetic island. Those studies of this topic are seen in Finn & Kaw (1977), Hayashi (1981), Biskamp (1982), Bhattacharjee, Brunel, & Tajima (1983), Steinolfson & van Hoven (1983, 1984), Schumacher & Kliem (1996), and so on. Turbulent effects in that process are also considered in Matthaeus and Montgomery (1981) and Matthaeus & Lamkin (1985).

A. 3. Anomalous Resistivity

Normally, the diffusive effect is based on the interaction among particles and such effect plays an important role in a collision-dominated plasma. However, in a very hot and dilute plasma, collisions rarely occur so that the resistivity based on the particle interaction becomes quite small. In that case, the behavior of a plasma is usually described in the ideal MHD regime and the solar corona belongs to this case, where any structure in the corona makes an ideal MHD evolution. Accordingly, those transient activities observed in the corona are considered to be caused by another type of the resistivity, which is based on the interaction between particles and various kinds of wave. This is the origin of the anomalous resistivity.

Since the anomalous resistivity is caused by the interaction between particles and waves, the occurrence of this needs the existence of fluctuations in a plasma. Such fluctuations are mainly generated by various plasma instabilities. Usually, the anomalous resistivity is characterized by the effective collision frequency, defined as the ratio of fluctuation energy density to thermal energy density. For example, the effective collision frequency caused by the ion-acoustic instability is given by

$$\nu_{i,an} = \nu_{pe} \frac{W_i}{n_0 k_B T_e}, \quad (\text{A. 3. 1})$$

where ν_{pe} , n_0 , k_B , T_e , and W_i mean the plasma frequency, plasma density, Boltzmann constant, electron temperature, and ion-acoustic fluctuation energy density, respectively. This is called the *Sagdeev formula* for the ion-acoustic anomalous collision frequency (Treumann & Baumjohann 1997). By estimating W_i under the assumption of very large electron temperature ($T_e \gg T_i$), we get the following expression for that anomalous collision frequency:

$$\nu_{i,an} = 0.01 \nu_{pi} \frac{v_d}{v_s} \frac{T_e}{T_i}^{-2}. \quad (\text{A. 3. 2})$$

Here ν_{pi} , v_d , v_s , T_i , and ν_{pe} mean the ion plasma frequency, ion-electron drift velocity, plasma

sound velocity (ion-acoustic velocity), ion temperature, and wave scattering angle, respectively. Since the ion-acoustic instability arises in the case of $v_d \gg v_s$, the ion-acoustic anomalous collision frequency can become of the order of the ion plasma frequency assuming a very large drift velocity ($v_d \gg v_s$) and a narrow wave scattering angle ($\theta \ll 1$). On the other hand, the ion-electron collisional frequency is expressed as

$$\nu_{ei} = \frac{4\pi e^2 n_e}{3 n_0} \frac{1}{D} \sim 4 \times 10^{-6} \frac{n_0}{T_e^{3/2}} \ln \Lambda, \quad (\text{A. 3. 3})$$

where D and $\ln \Lambda$ mean the *Debye length* and the *Coulomb logarithm* (*Spitzer formula*). In the coronal environment, this collisional frequency is about 10^5 s^{-1} ($T_e = 10^6 \text{ K}$ and $n_0 = 10^9 \text{ cm}^{-3}$), while the ion-acoustic anomalous collision frequency is given by the ion plasma frequency, which is about 10^6 s^{-1} . Consequently, the effectiveness of the resistivity is enhanced about 10^5 times.

Another important instability relevant to the anomalous resistivity is the lower-hybrid drift instability, which plays an important role in a current sheet. According to Treumann & Baumjohann (1997), the anomalous collision frequency caused by this instability also becomes of the order of the ion plasma frequency, supposing a strong magnetic field with a steep density gradient of the order of the ion gyroradius.

Appendix B. Specification of A Numerical Code

In this appendix, we explain a numerical code for the 2.5-dimensional MHD numerical simulation used in this thesis. We start with such the basic MHD equations as are expressed in the conservational form:

$$-\frac{\partial}{\partial t} \mathbf{w} + \frac{\partial}{\partial x} \mathbf{E}(\mathbf{w}) + \frac{\partial}{\partial z} \mathbf{G}(\mathbf{w}) + \mathbf{S}(\mathbf{w}) = \mathbf{0}, \quad (\text{B. 1})$$

where \mathbf{w} , $\mathbf{E}(\mathbf{w})$, $\mathbf{G}(\mathbf{w})$, and $\mathbf{S}(\mathbf{w})$ are the 8-dimensional vectors given by (4. 2), (4. 3), (4. 5), and (4. 6), respectively. Then all the physical variables are discretised in terms of both space and time coordinates. In the following part, We pick up one component from vector equation (B. 1) and express it in discretised forms with plain characters.

$$w_{i,j}^n = w(t_n, x_i, z_j) \quad (\text{B. 2})$$

and

$$E_{i,j}^n = E(w_{i,j}^n), G_{i,j}^n = G(w_{i,j}^n), S_{i,j}^n = S(w_{i,j}^n). \quad (\text{B. 3})$$

Here t_n means the n -th time step and x_i and z_j mean the i -th and j -th grid points along the x -coordinate and the z -coordinate, respectively. These values are defined as

$$t_{n+1} = t_n + \Delta t, \quad (\text{B. 4})$$

$$x_{i+1} = x_i + \Delta x_{i+\frac{1}{2}}, \quad (\text{B. 5})$$

and

$$z_{j+1} = z_j + \Delta z_{j+\frac{1}{2}}, \quad (\text{B. 6})$$

where Δt is determined by the CFL (Courant-Friedrichs-Lewy) condition:

$$\Delta t \leq \Delta t_{CFL} \min \left[\frac{\min(x_i, z_j)}{v_{si,j}} \right] \text{ for all } (i, j), \quad (\text{B. 7})$$

where $v_s = (v^2 + c_s^2 + v_A^2)^{1/2}$. Here v , c_s , and v_A are the fluid velocity, adiabatic sound velocity, and Alfvén velocity, and we set $\Delta t = 0.4 \Delta t_{CFL}$ all through this thesis (see also Fletcher 1991). As for the spacial interval, we set

$$x_{i+\frac{1}{2}} = \text{const.}, \quad z_{j+\frac{1}{2}} = \text{const.}, \quad (\text{B. 8})$$

in the uniformly distributed grid-point case, while in the non-uniformly distributed grid-point case, we set

$$x_{i+\frac{1}{2}} = x \quad x_{i-\frac{1}{2}}, \quad z_{j+\frac{1}{2}} = z \quad z_{j-\frac{1}{2}}, \quad (\text{B. 9})$$

where x and z are constant values. All through this thesis, we set $x = z = 1.05$.

Next, we solve the basic equations on the basis of the 2-step *modified Lax-Wendroff scheme*.

<1st step>

$$\begin{aligned} w_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} &= \frac{1}{4} \left(w_{i+1, j+1}^n + w_{i+1, j}^n + w_{i, j+1}^n + w_{i, j}^n \right) \\ &\quad - \frac{t}{2} \frac{1}{x_{i+\frac{1}{2}}} \left(E_{i+1, j+1}^n - E_{i, j+1}^n + E_{i+1, j}^n - E_{i, j}^n \right) \\ &\quad - \frac{t}{2} \frac{1}{z_{j+\frac{1}{2}}} \left(G_{i+1, j+1}^n - G_{i+1, j}^n + G_{i, j+1}^n - G_{i, j}^n \right) \\ &\quad - \frac{t}{4} \left(S_{i+1, j+1}^n + S_{i+1, j}^n + S_{i, j+1}^n + S_{i, j}^n \right), \end{aligned} \quad (\text{B. 10})$$

where $w_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1}$ is the physical variable at a virtual mid-mesh point $(x_{i+\frac{1}{2}}, z_{j+\frac{1}{2}})$. Then we evaluate the values of flux and source at this point:

$$E_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} = E \left(w_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} \right), \quad (\text{B. 11})$$

$$G_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} = G \left(w_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} \right), \quad (\text{B. 12})$$

and

$$S_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} = S \left(w_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} \right). \quad (\text{B. 13})$$

Using these values, we determine the values at the interface of each grid point as follows.

$$\tilde{E}_{i+\frac{1}{2}, j}^{n+1} = \frac{z_{j-\frac{1}{2}}}{2} \frac{1}{z_j} E_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} + \frac{z_{j+\frac{1}{2}}}{2} \frac{1}{z_j} E_{i+\frac{1}{2}, j-\frac{1}{2}}^{n+1}, \quad (\text{B. 14})$$

$$\tilde{G}_{i, j+\frac{1}{2}}^{n+1} = \frac{x_{i-\frac{1}{2}}}{2} \frac{1}{x_i} G_{i+\frac{1}{2}, j+\frac{1}{2}}^{n+1} + \frac{x_{i+\frac{1}{2}}}{2} \frac{1}{x_i} G_{i-\frac{1}{2}, j+\frac{1}{2}}^{n+1}, \quad (\text{B. 15})$$

and

$$\begin{aligned}\tilde{S}_{i,j}^{n+1} &= \frac{z_{j-\frac{1}{2}}}{2} \frac{1}{z_j} \left(\frac{x_{i-\frac{1}{2}}}{2} \frac{1}{x_i} S_{i+\frac{1}{2},j+\frac{1}{2}}^{n+1} + \frac{x_{i+\frac{1}{2}}}{2} \frac{1}{x_i} S_{i-\frac{1}{2},j+\frac{1}{2}}^{n+1} \right) \\ &+ \frac{z_{j+\frac{1}{2}}}{2} \frac{1}{z_j} \left(\frac{x_{i-\frac{1}{2}}}{2} \frac{1}{x_i} S_{i+\frac{1}{2},j-\frac{1}{2}}^{n+1} + \frac{x_{i+\frac{1}{2}}}{2} \frac{1}{x_i} S_{i-\frac{1}{2},j-\frac{1}{2}}^{n+1} \right),\end{aligned}\quad (\text{B. 16})$$

where

$$x_i = \frac{1}{2} \left(x_{i+\frac{1}{2}} + x_{i-\frac{1}{2}} \right), \quad z_j = \frac{1}{2} \left(z_{j+\frac{1}{2}} + z_{j-\frac{1}{2}} \right). \quad (\text{B. 17})$$

<2nd step>

On the basis of the 1st-step calculation, we then try to obtain the next time-step physical variables at real grid points. This is given by

$$\begin{aligned}w_{i,j}^{n+1} &= w_{i,j}^n \\ &- \frac{t}{4} \frac{1}{x_i} \left(E_{i+1,j}^n - E_{i-1,j}^n + 2 \left[\tilde{E}_{i+\frac{1}{2},j}^{n+1} - \tilde{E}_{i-\frac{1}{2},j}^{n+1} \right] \right) \\ &- \frac{t}{4} \frac{1}{z_j} \left(G_{i,j+1}^n - G_{i,j-1}^n + 2 \left[\tilde{G}_{i,j+\frac{1}{2}}^{n+1} - \tilde{G}_{i,j-\frac{1}{2}}^{n+1} \right] \right) \\ &- \frac{t}{2} \left(S_{i,j}^n + \tilde{S}_{i,j}^{n+1} \right).\end{aligned}\quad (\text{B. 18})$$

Now we obtain the next time-step physical variables at real grid points. However in order to suppress the numerical overshoot behind those regions where the spacial gradient of physical variables are quite large (e.g. shock waves), we introduce the effect of the artificial viscosity. This effect is imposed on the results obtained after the 2nd-step calculation:

$${}^*w_{i,j}^{n+1} = w_{i,j}^{n+1} + q_{i,j}^{n+1} \quad (\text{B. 19})$$

Here ${}^*w_{i,j}^{n+1}$ is the physical value subject to the effect of the artificial viscosity and $q_{i,j}^{n+1}$ is the artificial viscous term. This viscous term is calculated by

$$\begin{aligned}q_{i,j}^{n+1} &= \frac{t}{x_i} \left(\tilde{q}_{xi+\frac{1}{2},j}^n \frac{w_{i+1,j}^n - w_{i,j}^n}{x_{i+\frac{1}{2}}} - \tilde{q}_{xi-\frac{1}{2},j}^n \frac{w_{i,j}^n - w_{i-1,j}^n}{x_{i-\frac{1}{2}}} \right) \\ &+ \frac{t}{z_j} \left(\tilde{q}_{zi,j+\frac{1}{2}}^n \frac{w_{i,j+1}^n - w_{i,j}^n}{z_{j+\frac{1}{2}}} - \tilde{q}_{zi,j-\frac{1}{2}}^n \frac{w_{i,j}^n - w_{i,j-1}^n}{z_{j-\frac{1}{2}}} \right),\end{aligned}\quad (\text{B. 20})$$

where

$$\tilde{q}_{x_{i+\frac{1}{2},j}}^n = \frac{1}{2} (q_{x_{i+1,j}}^n + q_{x_{i,j}}^n), \tilde{q}_{z_{i,j+\frac{1}{2}}}^n = \frac{1}{2} (q_{z_{i,j+1}}^n + q_{z_{i,j}}^n). \quad (\text{B. 21})$$

Here $q_{x_{i,j}}^n$ and $q_{z_{i,j}}^n$ are defined as

$$q_{x_{i,j}}^n = x_i \left[\max \left(\frac{1}{2} |v_{x_{i+1,j}}^n - v_{x_{i-1,j}}^n|, v_{\min} \right) - v_{\min} \right], \quad (\text{B. 22})$$

and

$$q_{z_{i,j}}^n = z_j \left[\max \left(\frac{1}{2} |v_{z_{i,j+1}}^n - v_{z_{i,j-1}}^n|, v_{\min} \right) - v_{\min} \right]. \quad (\text{B. 23})$$

In the above formulation for the artificial viscosity, the term x_i is a scale factor and the term v_{\min} is a critical velocity difference between the neighboring grid points, over which the artificial viscosity is set up. All through this thesis, we set $\alpha = 3$ and $v_{\min} = 0.01$.

By using this code, we investigate the maximum linear growth rate of the tearing instability. The comparison between the simulation result and analytically predicted one always gives a small discrepancy. For example, using the same force-free model as in § 4. 4, we obtain the simulation result of $\gamma_m = 1.43$ in the case of $\beta = 0.15$, $R_m = 1000$, while the analytical treatment gives $\gamma_m = 1.49$.

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