MAGNETOHYDRODYNAMIC MODELING FOR A FORMATION PROCESS OF CORONAL MASS EJECTIONS: INTERACTION BETWEEN AN EJECTING FLUX ROPE AND AN AMBIENT FIELD

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ABSTRACT

We performed a magnetohydrodynamic simulation of a formation process of coronal mass ejections (CMEs), focusing on the interaction (reconnection) between an ejecting flux rope and its ambient field. We examined three cases with different ambient fields: one had no ambient field, while the other two had dipole fields with opposite directions, parallel and anti-parallel to that of the flux rope surface. We found that while the flux rope disappears in the anti-parallel case, in the other cases the flux ropes can evolve to CMEs and show different amounts of flux rope rotation. The results imply that the interaction between an ejecting flux rope and its ambient field is an important process for determining CME formation and CME orientation, and also show that the amount and direction of the magnetic flux within the flux rope and the ambient field are key parameters for CME formation. The interaction (reconnection) plays a significant role in the rotation of the flux rope especially with a process similar to “tilting instability” in a spheromak-type experiment of laboratory plasma.

Key words: magnetic reconnection – magnetohydrodynamics (MHD) – Sun: corona – Sun: coronal mass ejections (CMEs)

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1. INTRODUCTION

Coronal mass ejections (CMEs) are one of the most spectacular explosive phenomena in which large amounts of mass and magnetic flux are ejected into interplanetary space. CMEs are believed to be the consequences of a sudden release of magnetic energy, i.e., the disruptions of a coronal magnetic field, the same as flares, which are strong brightenings observed in the soft X-ray. The energy released upward can be converted to kinetic energy as ejected flux ropes and/or CMEs, while the energy released downward can be converted to thermal energy or high energy particles. Therefore, larger flares show a higher association rate with CMEs (Kahler 1992; Kim et al. 2005; Yashiro et al. 2005), but they do not necessarily accompany each other. An association of flares can be explained by an energy release rate that depends on the condition of the coronal plasma. If the energy release is fast enough to heat the chromospheric plasma, the heated plasma fills the coronal loops above, which are accompanied with strong brightening in the soft X-ray observed as flares. If the energy release is not fast enough to overcome the radiative energy loss, strong brightening does not occur. At the same time, if the ejecting structures successfully escape into interplanetary space, the event can be understood as CMEs without flares. On the other hand, if an ejection associated with a strong soft X-ray brightening cannot escape into interplanetary space, the event should be considered a flare without CME (Yashiro et al. 2006).

The disruptions of a coronal magnetic field are observed to commonly occur in a wide spatial range. The Yohkoh satellite reveals that faint giant coronal arcades, which are similar to flares, are formed in association with CMEs (Hiei et al. 1993; McAllister et al. 1996). On the other hand, similar evolutions where flares occur with mass (plasmoid) ejections are observed accompanying much smaller scale emerging flux (Sakajiri et al. 2004). Based on such observational evidence, Shibata (1999) proposed that giant arcades, flares, mass ejections, and CMEs be understood with a unified view, i.e., the so-called “flares”. The model is supported by the statistical characteristics of flares. Solar flares show a “power-law” distribution between their energy scales and the occurrence rate (Drake 1971; Datlowe et al. 1974; Dennis 1985; Shimizu 1995).

In contrast to flares, CMEs show a log-normal distribution between the energy scales and the event number (Aoki et al. 2004; Yurcheyshyn et al. 2005; Lara et al. 2006). The statistical characteristics of CMEs and flares are different in spite of their common origin. This difference in statistical distribution implies the existence of filter effects that selectively prevent small-scale mass ejections from evolving into CMEs. Candidates for such filter effects are the interactions between eruptions and their ambient magnetic field whose spatial scale is larger than that of the eruption regions. Important points in the formation process of CMEs include not only the formation of ejecting structures (flux rope) but also whether they can overcome these obstacles.

The trigger processes of such eruptions have been numerically investigated for a few decades (Forbes & Priest 1995; Antiochos et al. 1999; Chen & Shibata 2000; Kusano et al. 2004; Kliem & Török 2006). In a “breakout” model (Antiochos et al. 1999), the interaction between a sheared arcade and its ambient field is considered to play a fundamental role in the triggering. In this model, the field configuration is assumed to be favorable to reconnection between the eroding sheared arcade and its ambient field. During the eruption, the reconnection can reduce the poloidal magnetic flux of the flux rope which is supplied by another reconnection occurring inside the sheared arcade. As a result of competition between the magnetic flux reduction and the supply, an eruption can fail to produce a CME. In some other
models (Forbes & Priest 1995; Kusano et al. 2004; Inoue & Kusano 2006; Kliem & Török 2006), ideal magnetohydrodynamic (MHD) processes (the loss of equilibrium or instability) of detached flux ropes are dominant in the trigger processes. In such models, the direction of the ambient field can be independent of the trigger processes. If the direction of the ambient field is favorable to reconnection, the reconnection can reduce the magnetic flux as mentioned above. In contrast, if the direction is not favorable to reconnection, the reconnection can reduce the magnetic flux.

From the point of view of space weather science, when CMEs containing southward magnetic flux impact the Earth’s magnetosphere, they have a strong influence on the space environment near the Earth. Hence, it is also important for space weather forecasting to understand the whole evolutionary process of CMEs, including how they are formed and how much southward magnetic flux they contain when they reach the orbit of the Earth. Estimation of the southward magnetic flux is a complicated task that requires understanding of the initial conditions and the whole evolutionary process of eruptions. For example, in the nonlinear evolution of kink instability of a strongly twisted flux rope, its untwisting (writhing) motion can rotate the flux rope (Kliem et al. 2004; Török & Kliem 2005). The ejecting twisted flux ropes could change their directions due to the writhing rotations after their onset or formation if the ropes become unstable in the mode.

Rotations of ejecting flux ropes are reported in CME observations. Yurchyshyn (2008) investigated the angles of magnetic clouds, halo CME axes derived from coronagraph images, and extreme-ultraviolet (EUV) post-eruptive arcade axes in CME source regions. He found that there is a good correlation of angles between the axis of EUV arcades and CMEs, and between the axes of magnetic clouds and coronal neutral lines (Heliospheric current sheets), while there is a low correlation of angles between CMEs and magnetic clouds. These facts mean that some fraction of ejected magnetic structures are significantly rotated during their travel from the solar corona to 1 AU. Along the way, the ambient field should have two types of structures: closed and open. In the outer corona or interplanetary space, all magnetic fields should be open due to the dragging of the solar wind, while most of the magnetic field in the lower corona should be closed. The rotation may result from interactions with both types of ambient fields. Yurchyshyn (2008) reported that a few events show deviation between the neutral line angles of post-eruptive arcades (in the lower corona) and the axes of CMEs (in a few solar radii). This means that the rotation during the evolution in the lower corona, i.e., the interaction with a closed field, is not negligible.

In this study, we performed a three-dimensional MHD simulation of a twisted flux rope ejected from an active region surrounded by a global ambient magnetic field. According to the numerical results, we discuss the conditions for the formation of CMEs and the evolutions of CME magnetic field structures, focusing on the interactions between ejecting flux ropes and their ambient closed fields.

In the following section, we describe the detailed methodology of the numerical simulation. In Section 3, the numerical results for different ambient fields are shown. We discuss the quantitative relations suggested by the numerical results in Section 4. Finally, we summarize this paper in Section 5.

2. NUMERICAL MODEL

2.1. Numerical Scheme

We performed a three-dimensional MHD simulation which solves the time variation of eight physical quantities: density, gas pressure, velocity, and magnetic field (ρ, p, v, and B, respectively). In order to simplify the MHD equations, the eight physical quantities are normalized with the following typical values in the corona: \( L_0 = R_\odot = 6.99 \times 10^8 \) km, \( B_0 = 30 \) G, \( \rho_0 = m_H \times 10^9 = 1.67 \times 10^{-16} \) g cm\(^{-3}\), and then, using these, we derive \( v_0 = v_{0A} = B_0/\sqrt{4\pi \rho_0} = 2071 \) km s\(^{-1}\), \( t_0 = L_0/v_{0A} = 350 \) s, and \( \rho_0 = \rho_0 v_{0A}^2 = 7.16 \times 10^1 \) erg cm\(^{-3}\). The normalized MHD equations (the equation of continuity, the equation of motion, the induction equation, and the equation of energy with gravity) are expressed as follows:

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \quad \frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot (\rho \mathbf{v} \mathbf{v} - \mathbf{B} \mathbf{B} + \rho T \mathbf{I}) = \rho g, \quad \frac{\partial \mathbf{B}}{\partial t} + \nabla \cdot (\mathbf{B} \mathbf{v} - \mathbf{v} \mathbf{B} + \psi \mathbf{I}) = -\nabla \times (\eta \mathbf{J}), \quad \frac{\partial e}{\partial t} + \nabla \cdot [(e + \rho T) \mathbf{v} - (\mathbf{v} \cdot \mathbf{B}) \mathbf{B}] = 0,
\]

where

\[
\mathbf{J} = \nabla \times \mathbf{B}, \quad e = \frac{\rho v^2}{2} + \frac{p}{\gamma - 1} + \frac{\mathbf{B}^2}{2} + \frac{G_0}{r \rho}, \quad p_T = p + \frac{\mathbf{B}^2}{2}, \quad g = -G_0 \frac{r}{r^3}, \quad G_0 = \frac{GM_\odot}{R_\odot^2 v_\odot^2}.
\]

In these equations, \( G \) is the gravitational constant, \( \mathbf{I} \) is the unit matrix in the Cartesian coordinate, \( \psi \) is an additional variable described below, and \( M_\odot \) is the mass of the Sun. The divergence operator is solved by a finite-volume method with the Harten–Lax–van Leer Discontinuities nonlinear Riemann solver (Miyoshi & Kusano 2005), which is combined with the third-order Monotone Upstream-centered Schemes for Conservation Laws (MUSCL) developed by van Leer (1979) and the second-order Runge–Kutta time integration. The specific heat ratio \( \gamma \) is set to be 1.05, which means an additional heat source. Resistivity \( \eta \) is set to be 0, which means the reconnection is caused by numerical diffusion.

Generally, a multi-dimensional MHD simulation may violate the divergence-free (solenoidal) condition of the magnetic field without any divergence cleaning. Equations (2), (3), and (4) are derived with the solenoidal condition, and therefore such finite numerical \( \nabla \cdot \mathbf{B} \) can break down the simulation. In order to keep the numerical \( \nabla \cdot \mathbf{B} \) at a sufficiently small value, we applied an additional variable \( \psi \) and an additional equation

\[
\frac{\partial \psi}{\partial t} + c_h \nabla \cdot \mathbf{B} + \frac{c_h^2}{c_p} \nabla^2 \psi = 0
\]
based on the hyperbolic divergence $\nabla \cdot \mathbf{B}$ cleaning method (Dedner et al. 2002). The coefficient $c_b$ in Equation (10) is the propagation speed of the numerical $\nabla \cdot \mathbf{B}$ and the coefficient $c_p$ is the diffusion coefficient of $\psi$. These numerical methods are described in detail in Shiota et al. (2008).

The numerical domain is a spherical shell of $(r_{\text{min}}, r_{\text{max}}) = (1.01, 5.0)$. The domain is discretized with a Yin-Yang grid (Kageyama & Sato 2004), a chimera grid composed of two congruent spherical partial shells ($\pi/4 \leq \theta \leq 3\pi/4$, $-3\pi/4 \leq \phi \leq 3\pi/4$), each of which is installed in a different direction. The physical quantities on the boundary of one partial shell are copied from the values at the same points on the other partial shell that are interpolated by a trigonometric function.

Thus, we have to specify the boundary condition in this grid system only at the inner and outer boundaries for the radial direction. The inner boundary is assumed to be the line-tying wall, where $B_{\text{norm}}$ is a constant and $v = 0$. The outer boundary is assumed to be the free boundary, where the radial gradients of all quantities are zero.

Note that in this simulation, the test particles, which are initially distributed in some manner, are advected with plasma bulk flows. The magnetic field lines in all figures in this paper are obtained with the integration of the magnetic field from the position of the test particles of each time step. Thus, we capture the evolution of individual magnetic field lines with this technique.

2.2. Initial Condition

The subject of the present work is to understand how a flux rope ejected from an active region evolves to a CME. The initial magnetic field consists of the ejecting flux rope ($\mathbf{B}_S$) and the ambient global field ($\mathbf{B}_D$). We performed three cases of simulation where only the ambient field $\mathbf{B}_D$ is different. The ambient global field ($\mathbf{B}_D$) is chosen to be a simple dipole field:

$$B_{r,D}(r, \theta, \phi) = \frac{2B_D \sin \theta}{r^3},$$

$$B_{\theta,D}(r, \theta, \phi) = \frac{B_D \cos \theta}{r^3}.$$  
(11)

In order to study the interaction between an ejected flux rope and the ambient field, we carried out three cases of simulations with different strengths and directions of the ambient fields, i.e., their amplitude $B_D$ of $0.0, -B_{D,0},$ and $B_{D,0}$, which are labeled as cases A, B, and C, respectively. $B_{D,0} = 0.124$ (3.72 G) is the base field strength of the corona where the Alfvén speed is equal to the sound speed of $T = 2 \times 10^6$ K.

We mimic the ejected flux rope formed as a result of eruption in an active region with a spheromak-type magnetic field which is a linear force-free field in a completely isolated sphere (see Figure 1(a)):

$$\mathbf{B}_S(r, \theta, \phi) = \mathbf{B}(\tilde{r}, \tilde{\theta}, \tilde{\phi}),$$

(13)

where $\tilde{\mathbf{r}} = \alpha(r - \mathbf{r}_S)$ is the translated and rescaled local spherical coordinate whose origin corresponds to the center of the spheromak. The spheromak-type field is expressed as follows:

$$\tilde{B}_r(\tilde{r}, \tilde{\theta}, \tilde{\phi}) = -2B_S0 \frac{j_1(\tilde{r})}{\tilde{r}} \cos \tilde{\phi},$$

$$\tilde{B}_\theta(\tilde{r}, \tilde{\theta}, \tilde{\phi}) = B_S0 \left( \frac{j_1(\tilde{r})}{\tilde{r}} + \frac{j'_1(\tilde{r})}{\tilde{r}} \right) \sin \tilde{\phi},$$

(14)

$$\tilde{B}_z(\tilde{r}, \tilde{\theta}, \tilde{\phi}) = B_S0j_1(\tilde{r}) \sin \tilde{\theta},$$

(15)

where $B_{S0}$ is the field strength of the spheromak and $j_1$ and $j'_1$ are the first-order spherical Bessel function and its derivative, respectively:

$$j_1(x) = \frac{\sin x - x \cos x}{x^2},$$

$$j'_1(x) = \frac{2x \cos x - (x^2 - 2) \sin x}{x^3}.$$  
(16)

We only adopt the field within the isolated spherical shell $\tilde{r} \leq \tilde{a}$ (transparent sphere in Figure 1(a)) while $\tilde{B}_S = 0$ outside of the shell. As the definition of the spherical Bessel function, $\tilde{a} = 4.493409458$ and $\alpha = \tilde{a}/\alpha_{S0}$. The toroidal flux ($\Phi_S$) and the poloidal flux ($\Phi_P$) of this field are obtained by numerical integrations:

$$\Phi_S = \int_{0}^{\pi} d\tilde{\theta} \int_{0}^{\tilde{a}} \tilde{r} \tilde{B}_s d\tilde{r} = 0.261\alpha_{S0}^3 B_{S0},$$

(19)

$$\Phi_P = \int_{-\pi}^{\pi} d\tilde{\phi} \int_{0}^{\tilde{a}} \tilde{r} \tilde{B}_p d\tilde{r} = 0.0736\alpha_{S0}^3 B_{S0}.$$  
(20)

Although a spheromak is a linear force-free field, we adopt it only within the boundary $\tilde{r} \leq \tilde{a}$. At the spherical boundary $\tilde{r} = \tilde{a}$, the local radial ($\tilde{B}_r$) and azimuthal ($\tilde{B}_\theta$) components vanish and only the local zenithal component $\tilde{B}_z$ exists, while $\tilde{B} = 0$ outside of the boundary. Therefore, the forces do not balance at the spherical boundary, i.e., a strong outward gradient of the magnetic pressure exists there. The Lorentz force due to the current, which corresponds to a steep magnetic pressure gradient, causes the strong expansion. The spheromak field naturally swells in the simulation.

Figure 1 shows schematic pictures of the magnetic field direction on the meridional plane through the center in cases B (Figure 1(b)) and C (Figure 1(c)). The spheromak local zenithal field on its surface $\tilde{r} = \tilde{a}$ satisfies $\tilde{B}_z(\tilde{r} = \tilde{a}) > 0$ (cf. Equation (15)), which is southward. On the other hand, the ambipolar field is northward ($\tilde{B}_\theta < 0$) in case B, while it is southward in case C. On the equator, $\tilde{B}_z(\tilde{r} = \tilde{a})$ and $\tilde{B}_D > 0$ becomes anti-parallel in case B and parallel in case C. Hence, cases A, B, and C are referred to as “no ambient,” “anti-parallel,” and “parallel” cases, respectively, in this paper.

Density and pressure are determined by hydrostatic equilibrium with uniform temperature $T = 2 \times 10^6$ K,

$$\rho = \rho_b \exp \left[ \frac{G_0 T}{p_b} \left( \frac{1}{r} - 1 \right) \right],$$

(21)

$$\rho = \gamma \rho_b,$$

(22)

$$v = 0,$$

(23)

where $p_b = 2\rho_b k_b T/(m_H p_b)$ is the normalized pressure on the solar surface ($r = 1$). In the present study, we did not apply the
solar wind, which is also an important factor for the formation of a CME. This is because we intend to investigate only the interaction between the ejecting flux rope and the ambient fields in the early phase of CME formation.

3. NUMERICAL RESULTS

3.1. Common Evolution

In all cases, the spheromak (flux rope) initially swells outward due to the non-equilibrium initial condition and the line-tying boundary as described in the previous section. At the same time, the expanding flux rope also shows similar rotation around the line of ejection in all cases. This rotation is caused by an untwisting (writhing) motion of the highly twisted initial spheromak structure (see Figure 1(a)).

In order to show the height–time evolution, the right panels of Figure 2 illustrate the time variation of the plasma $\beta \equiv 0.5 p / B^2$ along the $x$-axis. The low plasma $\beta$ regions along the $x$-axis express the positions of the flux rope. Therefore, the low plasma $\beta$ in the height–time diagrams illustrated in Figure 2 corresponds to the trajectories of the flux ropes. In all the cases, the flux rope rapidly expands just after the start, followed by a rising and expanding motion with some deceleration.

In all cases, many small features appear around the altitude $r \sim 1.1$ (below the flux rope) and rise as shown in the right panels of Figure 2. These structures are small plasmoids formed by the tearing of the current sheet between the anchored axial flux. Once such plasmoids are formed in the current sheet, due to the gradient of magnetic pressure, they are ejected into the upper direction, and finally collide and coalesce into the flux rope above. However, the momentum and energy additions by the plasmoids are too small to accelerate the flux rope. The evolution is similar to the results of two-dimensional simulations by Shiota et al. (2008).

3.2. “No Ambient” Case

In the “no ambient” case (case A), the flux rope can expand into an unmagnetized plasma. We simulated this case as a reference case with no magnetic obstacle. The expanding flux rope drags outward its anchored axial flux. The dragged axial flux forms an $\Omega$ shape and a current sheet is formed at the center. Although the reconnection occurs between the dragged fluxes, a part of the flux remains anchored.

Figures 2(a) and (b) show the magnetic field structure at $t = 50$ (well-developed phase) and the height–time diagrams in this case. Figure 2(b) expresses the time evolution of the flux rope in this case. The flux rope rapidly expands just after the start ($t \lesssim 10$: early phase), and then shows a rising and expanding motion during the middle phase ($10 \lesssim t \lesssim 30$). The flux rope is slightly decelerated. In the late phase ($t \gtrsim 30$), the flux rope expands self-similarly at an almost constant speed, and finally reaches the outer boundary of the numerical domain.

3.3. “Anti-parallel” Case

We simulated the “anti-parallel” case (case B) as a condition where the directions of poloidal fluxes of the ejection and the ambient field are opposite. Such a condition is probable if the flux rope is formed and ejected in a newly emerged flux system whose direction is opposite to the pre-existing ambient field. This situation is modeled as a candidate system for the rapid eruption for a CME (Antiochos et al. 1999). The time evolution of the flux rope is significantly different from that in case A. Because the direction of the ambient field is opposite to the poloidal field of the ejecting flux rope, there is a strong current sheet where magnetic reconnection occurs. Similar to the scenario of Antiochos et al. (1999), the rising motion of the flux rope presses the current sheet, and then enhances its current density, i.e., the reconnection rate there. However, the poloidal flux of the flux rope is peeling off due to the reconnection with the ambient field as it rises. Finally, as shown in Figure 2(c), after most of the poloidal flux is reconnected, the flux rope no longer exists. The magnetic flux becomes just coronal loops within which the shear Alfvén waves propagate. Such a result may be a candidate condition for plasmoid ejections (flares) without CMEs.

We can see the time evolution of the flux rope in the height–time diagram (Figure 2(d)). In the early phase ($t \lesssim 10$), the trajectory and the vertical size are similar to those in case A. After this phase, the top edge of the flux rope becomes lower than that in case A because of the flux peeling ($10 \lesssim t \lesssim 35$: middle phase). Finally ($t \gtrsim 35$: late phase), the plasma $\beta$ inside the flux rope becomes comparable to unity, and stays almost at the same height. In this phase, the structure is recognized as twisted loops rather than a flux rope.

3.4. “Parallel” Case

We simulated the “parallel” case (case C) as a condition where the directions of poloidal fluxes of the ejection and the ambient field are parallel. The CME formation process in a similar but axisymmetric situation is studied by Shiota et al. (2008). The present study is an extension of Shiota et al. (2008) to study the three dimensionality of the CME formation process.
The time evolution of the flux rope is a little different from that in case A as shown in Figures 2(e) and (f). In the early phase ($t \lesssim 10$) the flux rope rapidly expands and shows a rising and expanding motion in the middle phase ($10 \lesssim t \lesssim 45$), and then a self-similar evolution in the late phase ($t \gtrsim 35$). The expanding flux rope and its central current sheet where magnetic reconnection occurs form an $\Omega$-shaped loop, and the flux rope finally reaches the outer boundary. This evolution is similar to that in case A. In contrast to the results in case A, the ejecting flux rope is significantly deformed (Figures 2(a) and (e)), although the flux rope is successfully ejected to the outer boundary in both cases. The difference could be caused by reconnection between the flux rope and the ambient magnetic field. The height evolution of the flux rope is similar to that in case A but the flux rope continues to be decelerated (Figures 2(b) and (f)).

We investigated the time evolution of the magnetic field configuration. Figures 3 and 4 show the time evolution of the magnetic field configuration seen from the side and the face of the ejecting flux rope. Some characteristic field lines are highlighted with colors (green, light blue, blue, and red) in the figures. These lines are drawn by integrating the magnetic field from the same points on the inner boundary in each time step.

The green field line represents the initial axial field line. As the flux rope expands, the axial field lines are dragged and stretched as shown in Figures 3(a)–(e). Finally, the field line reconnects at the $\Omega$-shaped center as described in Section 3.2.

The light blue field line indicates the interchanged field line. The line is initially an ambient field line (Figures 3(a) and 4(a)). As the flux rope expands, the line is pushed aside by the flux rope (Figures 3(b) and 4(b)). Because the direction of the magnetic field is parallel, the ambient field line does not reconnect with
Figure 3. Time evolution of the magnetic field and velocity in case C. Panels (a)–(f) show $t = 0, 2, 5, 10, 20,$ and $40$, respectively. The inner boundary is shown with the sphere in the left part of each panel whose colors indicate the $B_z$ map. The same field lines are displayed with tubes in each panel. Characteristic field lines are colored with green, light blue, blue, and red; see the text for details. The velocity map at $y = 0$ is displayed with colors on the vertical transparent background.

the flux rope. However, as the flux rope evolves, it rotates due to the writhing motion. The field lines are connected to the inside of the flux rope (Figures 3(c) and 4(c)). The result implies that the field lines reconnect with the flux rope at the current sheets formed on the flux rope surface. The location of the reconnection is examined later (Section 4.2).

The blue field line also indicates the interchanged field line. Initially, the line is also an ambient field line further out than the light blue line. The blue field line reconnects with the ejecting flux rope at the current sheet formed on its surface in the same way as the light blue field line.

Reconnection also occurs in the front of the flux rope. In the initial condition, the poloidal field of the surface of the flux rope is parallel to the ambient field. As mentioned in Section 3.1, however, the flux rope rotates due to its writhing motion. The magnetic field on the flux rope front rotates counterclockwise as a result of the rotation of the flux rope, as shown in Figure 4. The rotation causes the formation of a current sheet on the front and then magnetic reconnection occurs there.

The red field line in Figures 3 and 4 was initially an ambient dipole field but finally connected to the inside of the flux rope, i.e., became the new axial field of the ejecting flux rope (Figures 3(f) and 4(f)).

4. DISCUSSION

4.1. Magnetic Flux Relation for CME Formation

The amounts of magnetic flux and magnetic helicity are key parameters for understanding the numerical results. Here we estimate the amounts of magnetic flux and helicity and discuss
Figure 4. Face-on view of the time evolution of the magnetic field structure in case C. Panels (a)–(f) show \( t = 0, 2, 5, 10, 20, \) and 40, respectively. The meaning of the colors on the sphere and the tubes are explained in the text.

the relationship between them and conditions for CME formation. In the initial conditions of cases B and C, though the spheromak field structure is superposed onto the ambient dipole field, the ambient dipole field component is so weak that it does not change the spheromak field. Hence, we can regard the time evolutions in cases B and C as the results from the interaction (reconnection) between the ejecting flux rope and the ambient field.

Whether an ejecting flux rope can evolve to a CME depends on the magnetic flux relation of the overlying ambient field and the ejecting flux rope. In the present study, the dipole field is assumed as the ambient field and therefore we can integrate the amount of flux with the following equation:

\[
\Phi_D = \int_{\phi_{\min}}^{\phi_{\max}} d\phi \int_{r_{\min}}^{r_{\max}} B_{\theta,D} \left( \theta = \frac{\pi}{2} \right) r dr.
\]  

(24)

Because only the magnetic flux above the spheromak field is relevant to the CME condition, the azimuthal width \((\phi_{\min}, \phi_{\max})\) of the overlying magnetic flux is determined by that of the spheromak field in the initial condition: \((\phi_{\min}, \phi_{\max}) = \pm \tan^{-1}(a_S/|r_S|)\). We then obtain the overlying magnetic flux \(\Phi_D = 1.19 \times 10^{-2}\). On the other hand, we can also integrate the poloidal flux of the spheromak field with Equation (20). The azimuthal width in the spheromak coordinate is set to \(\pm \pi/2\) when the upper half of the spheromak field is assumed to commit to their ejection. Then, we get \(\Phi_P = 2.45 \times 10^{-3}\). The relationship between the amount of magnetic flux is \(\Phi_P < \Phi_D\) in cases B and C. This condition naturally explains the disappearance of the flux rope due to flux peeling-off reconnection in case B. In case C, the reconnection is not completely anti-parallel and is so inefficient that the amount of the stripped magnetic flux is less than that of the flux rope. In case A, there is no obstacle ambient field. Hence, the flux rope in cases A and C is not dissipated.

Although only three typical cases are modeled in the present study, there can be many intermediate situations. For example, if the ambient field is anti-parallel but the total amount of magnetic flux is smaller than that of an ejecting flux rope, the flux rope will also retain its shape and escape to outer space. Such an anti-parallel weak ambient field condition is a candidate where explosive energy release has been studied as the breakout model (Antiochos et al. 1999). In the case of local simulation of emerging flux, Galsgaard et al. (2007) and Archontis & Török (2008) studied flux emerging into pre-existing and weak coronal fields. Galsgaard et al. (2007) showed that the rising motion of the emerging flux is not affected by the orientation between the emerging flux and the pre-existing field. Archontis & Török (2008) showed that in the anti-parallel
case the emerging flux rope successfully erupts while it remains confined in the parallel case. In such cases of emerging flux, the amount of the magnetic flux that emerged becomes more than that of the coronal field within the same spatial scale because the strength of the magnetic field is very strong in the lower atmosphere. However, since the magnetic field in the real Sun is very complicated, the small scale magnetic field does not necessarily have the same direction as the global scale field. Furthermore, the spatial scale of the emerging flux should be (sometimes very) small compared to the solar radius, so the flux of a global scale ambient field is possibly comparable to or (sometimes very) small compared to the solar radius, so the strength of the magnetic field is very strong in the lower atmosphere. The amount of magnetic flux open and decreases magnetic obstacles. Therefore, the condition of Equation (25) is possibly an important factor for CME formation, and therefore its detailed estimation with a realistic corona model is required for CME prediction for space weather forecasting. Note that we also did not consider the effect of solar wind, which makes a substantial amount of magnetic flux open and decreases magnetic obstacles. The detailed estimation of the amount of magnetic flux using observational data and a realistic corona model is an important task for a CME prediction.

Shiota et al. (2008) discussed a condition for CME formation with the two-dimensional simulation of a continuously sheared arcade which results in a CME or a confined flux rope formation. They suggested that conditions for CME formation depend on the balance between amounts of the magnetic helicity \(H\) within the flux rope and sheared arcade, and the overlying magnetic flux \(\Phi_{\text{ov}}\) which confines the sheared structure, as follows:

\[
\frac{\Phi_{\text{ov}}^2}{|H|} < 1.7. \tag{25}
\]

In this symmetric condition, the maximum energy for the flux rope to escape requires very much work to drag the overlying flux. Therefore, the condition of Equation (25) is possibly an upper limit for a three-dimensional condition.

The relative magnetic helicity of each case of the present study is estimated from the initial condition

\[
H = \int (A + A_r) \cdot (B - B_r) dV, \tag{26}
\]

where \(B_r\) and \(A_r\) are the potential magnetic field and the corresponding vector potential, respectively, which are determined by the magnetic field radial component on the inner boundary of the numerical domain. The estimated relative helicity is \(H = -3.65 \times 10^{-4}, -2.96 \times 10^{-4}, \text{ and } -4.34 \times 10^{-4}\) in cases A, B, and C, respectively. The relation of Equation (25) in this case becomes

\[
\frac{\Phi_{\text{ov}}^2}{|H|} = \frac{\Phi_0^2}{|H|} = \frac{1.19 \times 10^{-2}}{4.34 \times 10^{-4}} \sim 0.33 < 1.7. \tag{27}
\]

This relation suggests that the confinement due to the overlying field is inefficient in preventing flux rope escape in case C, and the results are consistent with the results of Shiota et al. (2008). As discussed above, the solar wind decreases the amount of the overlying magnetic flux and helps the flux rope to escape from the Sun.

### 4.2. Rotation of an Ejecting Flux Rope

As described in Section 3, the rotations of ejecting flux ropes in the early phase in each case are caused by the writhing motion of the spheromak structure in the initial conditions. However, the final angles of the escaping flux ropes are quite different \(\sim 90°\) for cases A and C (Figures 2(a) and 2(e)). The difference is caused by the existence of the ambient field, i.e., the interaction between the ejecting flux rope and the ambient field as mentioned in Section 3.4. The interaction is the reconnection of the twisted field within the flux rope to the untwisted ambient field. As a result of the reconnection of the not completely anti-parallel fluxes, a strongly bent magnetic flux can be formed. As a result of the relaxation process of the flux, the flux rope can be rotated.

Angular momentum \((L_x)\) and torque \((\tau_x)\) around the \(x\)-axis are obtained from their integral in the space \(|r| \leq r_{\text{min}} \leq r \leq r_{\text{max}}\), \(\pi/4 \leq \theta \leq 3\pi/4\), \(-\pi/4 \leq \phi \leq \pi/4\)

\[
L_x = \iint_V dV (r \times \rho v) \cdot \mathbf{\hat{x}}, \tag{28}
\]

\[
\tau_x = \iint_V dV (r \times F) \cdot \mathbf{\hat{x}}. \tag{29}
\]

Here we calculated three kinds of torques where \(F\) is set to be either the Lorentz force, gradients of gas pressure, or dynamic pressure. The time variations of the torques around the \(x\)-axis are shown in Figure 5(a) and that of the angular momentum \(x\) component is shown in Figure 5(b). Figures 5(a) and (b) show that the structure continues to be rotated dominantly by the torque due to the Lorentz force. This result shows that the rotation is caused by the writhing motion of the initial flux rope. However, Figure 5(c) shows the three-dimensional distribution of the torque. The net torque mainly works at the surface of the flux rope (indicated by pink surfaces emphasized with a white circle in Figure 5(c)). It appears that the rotation is caused by interaction between the magnetic field inside the flux rope and the ambient magnetic field. The magnetic field line which penetrates the strong torque region is drawn with a purple line in Figure 5(c). The purple line appears strongly wound, i.e., just after reconnection. This result means that the strong torque regions correspond to the current sheets where the reconnection between the flux rope and the ambient field occurs.

Figure 5(d) shows a schematic picture of the evolution of the magnetic field lines in case C mentioned above and in Section 3.4. The expanding flux rope initially rotates counterclockwise due to its writhing motion. As a result, current sheets are formed at two points on the surface, where the anti-parallel fields are in contact (indicated with red lines in Figure 5(d)). Reconnection occurs in the current sheets, and then strongly wound field lines (indicated with green and light blue lines in Figure 5(d)) are formed. The untwisting motion of the winding field causes a strong torque (purple arrows in Figure 5(d)) that rotates the flux rope.

A few decades ago, a magnetic field configuration of a spheromak in an equilibrium state under a uniform field was applied to a laboratory plasma experiment, where plasma must be confined within a device. However, such a spheromak-type field configuration is unstable to tilting mode perturbation (Sato...
Figure 5. Time variation of torques (a) and angular momentum (b) around the $x$-axis in case C (parallel). Green, blue, red, and black lines in panel (a) show time variations of torques caused by dynamic pressure, gas pressure gradient, Lorentz force, and total force, respectively. (c) Three-dimensional distribution of the magnetic torque around the $x$-axis. Pink surfaces show positions where counterclockwise (negative $x$) torques work. The white circle indicates the position of the net counterclockwise torque which substantially rotates the flux rope. (d) Schematic pictures of magnetic field lines on a horizontal cross section in the ejecting structure. Due to its writhing motion, the twisted flux rope rotates counterclockwise (evolution from the left panel to the center panel). As a result, current sheets (indicated with red lines in the center panel) are formed on the partial surface of the flux rope. Reconnection in the sheets produce a steep bend magnetic field (green and light blue lines in the right panel of (d) and purple tube in panel (c)), which causes the torque due to the magnetic tension force.

The field configuration settles down where the magnetic moment of the spheromak is parallel to the background field. In the nonlinear stage of the instability, the magnetic field causes reconnection, which is the same as that explained in Figure 5(d). The numerical results in this study imply that such an effect in laboratory plasma also works in the coronal plasma. Such rotations of ejecting flux ropes are reported in some CME observations. Yurchyshyn (2008) investigated the axis angles of magnetic clouds, halo CME axes derived from coronagraph images, and EUV post-eruptive arcade axes in CME source regions. He found that there are good correlations between the axes of EUV arcades and CME angles within a deviation of $\pm 45^\circ$. This fact means that some fraction of ejected magnetic structures are significantly rotated during travel in the lower solar corona but the amount of the rotation is not much.

The numerical results of the present work are consistent with observations in that flux ropes can be rotated significantly by the interaction with a closed field. However, the rotation angle in case C is more than that for most CMEs in Yurchyshyn (2008). The reason is that the effect of the ambient field in case C is much stronger than the real solar corona because of the following reasons: the dipole field assumed in this study decreases as $r^{-3}$, with $r$ being the distance from the Sun, and is the potential field whose influence can reach the furthest. Most of the dipole field flux is closed because the solar wind is not taken into account. The Alfvén speed in the simulation is very small because of the assumption of high plasma density. The slow evolutions allow the ejecting structure to show enough rotation. The condition in case C is an ideal case to examine the interaction between the ejecting flux rope and the ambient field. Furthermore, although the time evolution in case C is very slow (the terminal CME speed $\sim$40 km s$^{-1}$), the terminal speed is comparable to the Alfvén speed inside the flux rope.

The rotation process of ejecting flux ropes is also numerically studied by Gibson & Fan (2008) and Lynch et al. (2009). Gibson & Fan (2008) discussed the rotation process with the numerical result of eruption of a twisted flux rope within a potential arcade. They showed that the rotation of the flux rope axis ($\sim 115^\circ$) is caused by writhing motion. Lynch et al. (2009) discussed the rotation process with the results of two eruptions of arcades sheared opposite to each other. They showed that the opposite chirality of the sheared arcades leads to the same amount of opposite rotations of the resulting flux ropes. In contrast to these studies, the results of the present study show that the rotation of the flux rope is affected not only by the writhing motion but also by the interaction with the ambient field as described in Section 3. In contrast, too, to these previous works, the results of the present study suggest that an ejecting flux rope can be significantly rotated as a result of reconnection with its ambient field.

4.3. Evolution of the Anchored Axial Flux

In case C, because of the interchange reconnection mentioned in Section 3.4, the footpoints of the ejecting flux rope appear to move outside (colored lines in Figures 3 and 4). Initially, only the green line was included in the anchored axial flux of the ejecting flux rope, but all the light blue, blue, and red lines are included by the end. These changes of anchored footpoints are caused by a topological change of magnetic flux due to
reconnection in two regions: at the surface and behind the ejecting flux rope. As described in Section 3, the ejecting and expanding flux rope forms Ω-type loops of the initial axial flux. The first reconnection occurs in the current sheet at the center of the waist of the loop (see the green line in Figures 3 and 4). This reconnection severs the anchored axial flux and makes the ejecting flux rope structure detached.

At the same time, the ejecting flux rope reconnects with the ambient field on its side (light blue and blue lines) and front (the red line). In the initial condition of this case, the poloidal field of the flux rope is parallel and therefore unfavorable to reconnection. However, the flux rope rotates due to its untwisting (writhing) motion as described in Section 4.2. The reconnected ambient field is connected to the inside of the flux rope, i.e., it becomes the new axial flux of the ejecting flux rope. The interchange of the anchored axial field makes the distance between the legs larger and the footpoint field strength weaker. As a result, the magnetic tension force can decrease and it becomes easier for the flux rope to escape from the Sun.

The interchange of anchored flux explains the evolution of the dimming region observed in many flares and/or CMEs with soft X-ray or EUV. In many flares, strong dimmings are often observed on both sides of the post-eruptive arcades, followed by weak dimmings far from the flare site (Sterling & Hudson 1997). Such a strong dimming can be formed by a sudden expansion of the erupting core and the overlying flux dragged upward, coupled with coincident rarefaction due to reconnection (Shiota et al. 2005). The following weak dimmings are formed by some interaction (stretched or reconnected) with the erupting structure. This evolution of the magnetic field structure is similar to that mentioned by Gibson & Fan (2008). The results of case C in the present study also show that the weak dimmings are seen in new footpoints of axial flux formed due to the reconnection described in this section.

5. SUMMARY

In this study, we performed three global MHD simulations of interaction (magnetic reconnection) between an ejecting flux rope and its ambient field as a launch process of CMEs. The initial conditions are assumed to be the situations where an ejecting flux rope is embedded under three different ambient fields: a case without ambient field and cases with dipole fields of two opposite directions, which are parallel and anti-parallel to that of the flux rope surface. The flux rope can escape to the outer boundary of the numerical domain (evolve to CMEs) in the “no ambient case” and in the dipole case where the direction of the ambient field is parallel, while it cannot escape in the other dipole case.

The flux ropes show rotations perpendicular to the ejecting direction that are caused not only by the untwisting (writhing) motion of the flux rope but also by the interaction between the ejecting flux rope and the ambient field. The results especially show the possibility of rotation due to the relaxation of a complicated field structure which results from the reconnection between the flux rope and the ambient field. The reconnection and resulting rotation are similar to that seen in a spheromak experiment of laboratory plasma.

The results of this study show that the interaction between an ejecting flux rope and its ambient field is significant in determining CME formation and CME orientation. For space weather forecasting, it is necessary to predict the occurrence of a CME when a flare occurs. A powerful tool for such a prediction is the MHD simulation of the flare and the accompanying ejecting flux rope under a realistic coronal magnetic field.

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