SIMULATIONS OF EMERGING MAGNETIC FLUX. I. THE FORMATION OF STABLE CORONAL FLUX ROPES

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ABSTRACT

We present results from three-dimensional visco-resistive magnetohydrodynamic simulations of the emergence of a convection zone magnetic flux tube into a solar atmosphere containing a pre-existing dipole coronal field, which is oriented to minimize reconnection with the emerging field. We observe that the emergence process is capable of producing a coronal flux rope by the transfer of twist from the convection zone, as found in previous simulations. We find that this flux rope is stable, with no evidence of a fast rise, and that its ultimate height in the corona is determined by the strength of the pre-existing dipole field. We also find that although the electric currents in the initial convection zone flux tube are almost perfectly neutralized, the resultant coronal flux rope carries a significant net current. These results suggest that flux tube emergence is capable of creating non-current-neutralized stable flux ropes in the corona, tethered by overlying potential fields, a magnetic configuration that is believed to be the source of coronal mass ejections.

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

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1. INTRODUCTION

Coronal mass ejections (CMEs) are a primary source of space weather and almost all theoretical models of CMEs require the presence or formation of a coronal magnetic flux rope (e.g., Forbes 2000). There exists observational evidence that many CMEs, particularly those originating from quiet Sun regions, are composed of a bright core associated with an erupting prominence and a relatively dark cavity that is associated with a magnetic flux rope (e.g., Gibson & Fan 2006). Moreover, a magnetic flux rope geometry has been fit to coronagraphic observations of propagating CMEs (e.g., Vourlidas et al. 2013). In addition, there is also growing evidence that these flux ropes are formed before the eruption. This evidence exists for both quiet Sun regions (Robbrecht et al. 2009; Vourlidas et al. 2013) and active regions (e.g., Green & Kliem 2009; Green et al. 2011; Pantoulakis et al. 2013). Although identification of flux rope magnetic geometry is more difficult for active regions than for the quiet Sun due to differences in size and complexity, the existence of active region flux tubes has been supported by recent non-linear force free extrapolations (e.g., Canou & Amari 2010; Guo et al. 2010; Yelles Chauuche et al. 2012).

It has long been postulated that the source of the magnetic field in the corona is a magnetic field in the deep convection zone, created by dynamo action (Parker 1979), and that the process by which this field arrives in the corona is the buoyant rise of twisted flux tubes to the surface and their subsequent emergence. The partial emergence of twisted flux tubes into the solar atmosphere has been extensively studied and reviewed by Archontis (2008) summarizes the various types of theoretical investigations. Early three-dimensional (3D) simulations found that an emerging sub-surface flux tube does not rise bodily into the corona, but that only the upper portion of the tube emerges, while the tube axis remains near the solar surface (Fan 2001; Magara 2001; Archontis et al. 2004; Murray et al. 2006). More recent simulations found that a new flux rope structure forms in the corona within the partially emerged flux tube and that this flux rope rises slowly in the corona (Manchester et al. 2004; Fan 2009). Archontis & Török (2008) and Archontis & Hood (2012) demonstrated that the rise of the flux rope came to a halt, due to the stabilizing magnetic tension of the surrounding (envelope) flux tube field. Fan (2009) associated the coronal flux rope formation mechanism with the transfer of twist from the convection zone, while Manchester et al. (2004) suggested that the mechanism is due to the reconnection of sheared magnetic fields. By imposing a pre-existing strong horizontal field in the corona, Archontis et al. (2006) found that reconnection between the emerging flux tube and the coronal field can create horizontal jets and plasmoids at relatively low heights in the corona. Later simulations also found that favorably orientated and sufficiently weak horizontal fields can remove part of the envelope field constraining the newly formed flux rope, allowing a strong upward acceleration of the rope resembling eruptive behavior (Archontis & Török 2008; Archontis & Hood 2012).

In this paper, we focus on the formation of stable coronal flux ropes as a result of flux emergence. Creating such configurations is an important step in improving initial equilibrium magnetic field configurations for models of CMEs. A primary example of a pre-eruption configuration is the flux rope model of Titov & Démoulin (1999, hereafter the TD model), in which a coronal flux rope is confined by an overlying potential field. This model has been successfully applied as the initial condition of a number of CME simulations (Roussev et al. 2003; Török & Kliem 2005; Manchester et al. 2008). A specific property of the TD model is that the coronal flux rope carries a net current, since there is no return current in the configuration (see Török et al. 2013 for a detailed discussion of return currents in active regions).

Our aim in this paper is to use numerical magnetohydrodynamic (MHD) simulations to investigate how flux emergence from the convection zone into the solar atmosphere can create a stable coronal flux rope with a net current that is confined by an overlying magnetic field. To do this, we model the partial emergence of a buoyant convection zone twisted flux tube into a pre-existing dipole coronal magnetic field, the strength of which
we vary. The coronal dipole field is intended to represent the remnant field of an old, dispersed active region, into which a new magnetic field emerges. This simulation of flux emergence into a pre-existing dipole follows a paradigm similar to that of the simulation presented in MacTaggart (2011). That study focused on reconnection between the emerging field and the pre-existing field in the corona, whereas in this study we focus on the formation and stability of a coronal flux rope formed within the emerging field. In addition, for reference, we also model flux emergence into a field-free corona, to compare with the simulations of Archontis & Hood (2012) and Manchester et al. (2004).

In Section 2, the model is described. The results are presented in Section 3 and the consequences of these simulations for the theory of coronal flux rope formation and CME initiation are discussed in Section 4.

2. NUMERICAL METHOD

2.1. Equations

The evolution of a magnetic field in a plasma domain that includes the upper layers of the solar convection zone plus a photosphere/chromosphere, transition region, and corona is modeled using the visco-resistive MHD Lagrangian-remap code Lare3D (Arber et al. 2001). The equations solved by Lare3D are presented here in Lagrangian form:

\[ \frac{D\rho}{Dt} = -\rho \nabla \cdot \mathbf{v}, \]  
\[ \frac{D\mathbf{v}}{Dt} = -\frac{1}{\rho} [\nabla P + \mathbf{j} \times \mathbf{B} + \rho g + \nabla \cdot S], \]  
\[ \frac{D\mathbf{B}}{Dt} = (\mathbf{B} \cdot \nabla) \mathbf{v} - \mathbf{B}(\nabla \cdot \mathbf{v}) - \nabla \times (\eta \mathbf{j}), \]  
\[ \frac{De}{Dt} = \frac{1}{\rho} [ -P \nabla \cdot \mathbf{v} + \zeta_{ij} S_{ij} + \eta j^2]. \]

Here, \( \rho \) is the mass density, \( \mathbf{v} \) is the velocity, \( \mathbf{B} \) is the magnetic field, and \( \epsilon \) is the specific energy density. The current density is given by \( \mathbf{j} = \nabla \times \mathbf{B}/\mu_0 \), where \( \mu_0 \) is the permeability of free space and the resistivity \( \eta = 14.6 \Omega \). The gravitational acceleration is denoted by \( g \) and is set to the gravity at the mean solar surface (\( g_{\text{sun}} = -274 \text{ m s}^{-2} \)). \( S \) is the stress tensor that has components \( S_{ij} = \nu(\zeta_{ij} - (1/3) \delta_{ij} \nabla \cdot \mathbf{v}) \), with \( \zeta_{ij} = (1/2)((\partial v_i/\partial x_j) + (\partial v_j/\partial x_i)) \). The viscosity \( \nu \) is set to \( 3.35 \times 10^3 \text{ kg m}^{-1} \text{s}^{-1} \) and \( \delta_{ij} \) is the Kronecker delta function. Assuming an ideal gas law, the gas pressure, \( P \), and the specific internal energy density, \( \epsilon \), are

\[ P = \rho k_B T/\mu_m, \]  
\[ \epsilon = \frac{k_B T}{\mu_m(\gamma - 1)}, \]

respectively, where \( k_B \) is Boltzmann’s constant and \( \gamma = 5/3 \). The reduced mass, \( \mu_m \), is given by \( \mu_m = m_i m_p / m_p \), where \( m_p \) is the mass of a proton and \( m_i = 1.25 \) is a pre-factor designed to include the effect of elements heavier than hydrogen. In this study, as in many previous simulations of flux emergence, we assume that the plasma is fully ionized and so the reduced mass is spatially independent. In the partially ionized plasma of the Sun, the reduced mass changes as the ionization changes, as discussed in Leake & Linton (2013). In this study, however, we use the average value of \( \mu_m = m_i m_p \), which was shown in Leake & Linton (2013) to give the best constant-\( \mu_m \) match to one-dimensional (1D) models of the solar atmosphere that include partial ionization effects (e.g., Vernazza et al. 1981; Fontenla et al. 2006). The plasma variables \( \epsilon \) and \( \rho \) are defined at cell centers. The magnetic field is defined at cell faces and the velocity is defined at cell vertices. The staggered grid preserves \( \nabla \cdot \mathbf{B} \) during the simulation.

2.2. Normalization

The equations are non-dimensionalized by dividing each variable (\( \mathcal{C} \)) by its normalizing value (\( \mathcal{C}_0 \)). The set of equations requires a choice of three normalizing values. We choose normalizing values for the length (\( L_0 = 1.7 \times 10^3 \text{ m} \)), magnetic field (\( B_0 = 0.12 \text{ T} \)), and gravitational acceleration (\( g_0 = g_{\text{sun}} = 274 \text{ m s}^{-2} \)). From these values, the normalizing values for the gas pressure (\( P_0 = B_0^2/\mu_0 L_0 g_0 = 2.46 \times 10^{-4} \text{ kg m}^{-2} \)), velocity (\( v_0 = L_0 g_0 = 6.82 \times 10^3 \text{ m s}^{-1} \)), temperature (\( T_0 = m_p g_0/\kappa_b = 5.64 \times 10^3 \text{ K} \)), current density (\( j_0 = B_0/(\mu_0 L_0) = 0.56 \text{ Am}^2 \)), viscosity (\( \nu_0 = B_0^2/\eta L_0 g_0 = 2.85 \times 10^3 \text{ kg m}^{-2} \text{s}^{-1} \)), and resistivity (\( \eta_0 = \mu_0 L_0^3/8 \eta_0 = 1.46 \times 10^3 \Omega \)) can be derived. With these values of normalization and the values of \( \nu \) and \( \eta \) given above, the Reynolds number \( \mathcal{R}_e = (P_0 L_0 v_0)/\nu \) and magnetic Reynolds number \( \mathcal{R}_m = (\mu_0 L_0 v_0)/\eta \) in this simulation are both 100. The value of resistivity used in this simulation (0.01\( \eta_0 = 14.6 \Omega \text{m} \)), although comparable to upper estimates of the resistivity in the lower chromosphere, is much larger than typical values in the corona. We use this large value to ensure that the explicit resistivity is larger than the numerical value for the scheme used in regions where electric current densities build up. The normalized numerical resistivity is \( \hat{\nu}_A \hat{\Delta}_x/L \), where \( \Delta_x \) is the normalized grid size, \( \hat{\nu}_A \) is the normalized local Alfven speed, and \( \hat{\Delta}_x \) is a typical normalized length scale over which the magnetic field varies (Arber et al. 2007). In regions of increased current density, we find \( \hat{\Delta}_x \approx 0.66, \hat{\nu}_A \approx 0.05, \) and \( \hat{\Delta}_x \approx 5 \), which give a value for the normalized numerical resistivity of \( \eta_0/\hat{\eta}_0 = 0.0044 \).

2.3. Numerical Domain

The simulations use an irregular Cartesian grid with 304 cells in each direction. In the vertical direction, \( z \), the grid extends from \(-30 L_0 \) to \( 210.45 L_0 \) with a resolution of 0.428 \( L_0 \) at the bottom boundary and 1.99 \( L_0 \) at the top boundary. In the horizontal directions, \( x \) and \( y \), the grid is centered on 0 and has side boundaries at \( \pm 126.85 L_0 \). The resolution at \( x = y = 0 \) is 0.658 \( L_0 \) and the resolution at the side boundaries is 2.61 \( L_0 \). This irregular grid has the following form:

\[ x, y = \pm \left[ (1 + f_h) \chi_h + f_h w \ln \left( \frac{\cosh \left( \frac{w - L_0}{w} \right)}{\cosh \left( \frac{w - L_0}{w} \right)} \right) \right], \]
\[ z = -30 L_0 + \left[ (1 + f_v) \chi_v + f_v w \ln \left( \frac{\cosh \left( \frac{w - L_0}{w} \right)}{\cosh \left( \frac{w - L_0}{w} \right)} \right) \right], \]

where \( \chi_h = [0, 1, 2, ..., 152] L_0/152, \chi_v = [0, 1, 2, ..., 304] L_0/304, f_h = 2.1, f_v = 1.83, L_h = 95 L_0, L_v = 100 L_0, \) and \( w = 100 \).
and \( w = 10 L_0 \). We also perform one additional simulation (called ND1), which has a higher top boundary at 270 \( L_0 \), with \( f_v = 2.83 \).

At the boundaries, all components of the velocity and the gradients of magnetic field, gas density, and specific energy density are set to zero. The resistivity is smoothly decreased to zero close to the side boundary to reduce the diffusion of the magnetic field at the boundary to its numerical value and ensure line-tied boundary conditions as much as possible:

\[
\eta = 0.01 \left[ \tanh\left( \frac{(r_\eta - L_\eta)}{w_\eta} \right) + 1 \right] \eta_0 \frac{z}{2},
\]

where \( r_\eta = \sqrt{x^2 + y^2} \), \( L_\eta = 100 L_0 \), and \( w_\eta = 5 L_0 \). In addition, a damping region is applied to the velocity at all four side boundaries and the top boundary. For a given coordinate \( (x = x, y, z) \), the velocity Equation (2) has an additional term when \( |\kappa| > \kappa_d \):

\[
\frac{Dv}{Dt} = \frac{1}{\rho} \left[ \nabla P + j \left( \mathbf{B} + \rho \mathbf{g} \right) + \nabla \cdot \mathbf{S} \right] - \mathbf{N} - \mathbf{v},
\]

with \( x_d = y_d = 96 L_0 \) and \( z_d = 170 L_0 \) (simulation ND1 has \( z_d = 254 L_0 \)). The parameter \( N \) is designed to increase linearly from 0 at \( \kappa_d \) to 1 at the boundary: \( N = (|\kappa| - \kappa_d)/(\max(\kappa) - \kappa_d) \). This approach is used to prevent any reflected waves from interfering with the solution in the interior.

### 2.4. Initial Conditions

The initial conditions consist of a hydrostatic background atmosphere that represents the upper solar convection zone \((-30 L_0 \leq z < 0)\), the photosphere/chromosphere \((0 \leq z < 10 L_0)\), the transition region \((10 L_0 \leq z < 20 L_0)\), and the corona \((20 L_0 \leq z)\). The transition region in this model is thinner than a typical width derived from semi-implicit models of the Sun, which compare the observed spectrum of the Sun with radiative transfer calculations (e.g., Fontenla et al. 2006). In those studies, the typical width is about 0.1 Mm (≈\( L_0 \)). As in previous simulations of flux emergence (see the review by Archontis 2008), we use a thicker transition region of 1.7 Mm. This artificial increase of the transition region is required to resolve the large changes in density and temperature that occur across this region for the given spatial resolution that is limited by the large total simulation domain and computational costs.

A magnetic field is imposed on this background atmosphere. This field consists of a background dipole field that permeates the entire domain and a localized twisted flux tube in the model convection zone. The flux tube’s pressure and density are perturbed to initiate its buoyant rise into the model solar atmosphere. The initial conditions are shown in Figures 1 and 2.

The initial hydrostatic atmosphere is created by first defining the temperature:

\[
\frac{dT}{dz} = a \left( \frac{dT}{dz} \right)_{\text{ad}} = -\frac{\gamma - 1}{\gamma} \frac{T}{L_0}, \quad z \leq 0, \tag{11}
\]

\[
T(z) = T_{ph}, \quad 0 < z < 10 L_0, \tag{12}
\]

\[
T(z) = T_{\text{cor}}^{(z - 10 L_0)/10 L_0}, \quad 10 L_0 \leq z < 20 L_0, \tag{13}
\]

\[
T(z) = T_{\text{cor}}, \quad z \geq 20 L_0, \tag{14}
\]

where \( T_{ph} = T_0 \) and \( T_{cor} = 150 T_0 \). The pre-factor \( a = 1 \) in Equation (11) ensures that the model convection zone is marginally stable to convective instability by setting the temperature gradient to its adiabatic value \( \frac{dT}{dz} = (\frac{dT}{dz})_{\text{ad}} \) (Stix 2004). The gas density profile is then obtained by solving the hydrostatic equilibrium equation \( \frac{dP}{dz} = -\rho g \) and using the ideal gas law and the condition that \( \rho(z = 0) = \rho_0 \).

The dipole field is translationally invariant along \( y \), the tube’s axial direction, and is given by \( \mathbf{B} = \mathbf{V} \times \mathbf{A} \), where \( \mathbf{A} = A_r \mathbf{e}_r \), and

\[
A_r(x, z) = B_d \frac{z - z_d}{r_1^2}, \tag{15}
\]

with \( r_1 = \sqrt{x^2 + (z - z_d)^2} \) being the distance from the source. We choose \( z_d \) to be \(-100 L_0\) so that the initial sub-surface flux tube is far from the source of the dipole field. To cover various dipole strengths, we perform three simulations, each with a different value of \( B_d \). Simulations SD (strong dipole), MD (medium dipole), and WD (weak dipole) have values \( B_d = [10, 7.5, 5] \times 10^3 B_{d0} \), where \( B_{d0} = B_0 L_0^{-2} = 3.76 \times 10^8 \text{Tm}^{-2} \), respectively. This gives a maximum magnetic field strength at the surface \((z = 0)\) of \([2.6, 1.95, 1.3] \times 10^{-3} \text{T} \), respectively. These choices of dipole strength allow for a range in the plasma \( \beta \) profile, as shown in Figure 1, where \( \beta = \rho J_0(\mathbf{B}^2/\mu_0) \). These profiles are consistent with the models of \( \beta \) in the solar atmosphere developed by Gary & Alexander (1999) and Gary (2001). Simulations no dipole (ND) and ND1 have no pre-existing dipole field.

A right-hand twisted magnetic flux tube is inserted at \( x = 0 \), \( z = z_t = -12 L_0 \), aligned along the \( y \) axis, and is given by

\[
B_y = B_0 e^{-r_1^2/R_1^4}, \tag{16}
\]

\[
B_0 = q r B_y, \tag{17}
\]
where $r = \sqrt{(x^2 + (z - z_1)^2)}$. The width of the tube is $R = 2.5L_0$ and the strength at $r = 0$ is $B_t = 5B_0$. The twist parameter is $q = 1/R$. Figure 2 shows some selected magnetic field lines from the initial configuration. The superposition of the flux tube and the dipole field is shown in Figure 2 (panel (b)). The conventional wisdom of active region formation is that large-scale Ω-shaped flux tubes, which are anchored well below the visible surface, extend through the surface and into the corona. Hence, the initial flux tube used here, which is initially horizontal and line tied at the side boundaries, is not a very realistic initial condition. However, by perturbing the density in the initial flux tube in a certain way, an Ω-shaped tube can be created. The convection zone flux tube is made buoyant at the center $y = 0$ and is neutrally buoyant at its ends at the $y$ boundaries. This is done by perturbing the background density $\rho_0(z)$ and background specific energy density $\epsilon_0(z)$ to

$$
\rho(r,z) = \rho_0(z) \left(1 + \frac{p_1(r)}{\rho_0(z)} e^{-z^2}\right) \quad \text{and} \quad \epsilon(r,z) = \frac{(\rho_0(z) + p_1(r))}{\rho(y - 1)},
$$

where $\lambda = 10L_0$, $\rho_0(z)$ is the original pressure profile, and $p_1(r)$ is determined by solving $\nabla p_1(r) = j \times B(r)$ for the flux tube’s field. As will be shown later, this creates sub-photospheric “legs” of the emerging tube that have a significant vertical component. To optimize the confining effect of the dipole field, the direction of the dipole field is chosen so that it is aligned with the field lines in the top edge of the flux tube, i.e., $B_{\omega,\text{dip}} > 0$.

It is worth making a point here regarding the use of the phrases “flux tube” and “flux rope.” In previous studies of flux emergence, “flux tube” has been used to describe the sub-surface initial magnetic field configuration and “flux rope” has been used to describe the presence of a collection of field lines wrapped around a central field line we designate as the axis. In that sense, the original sub-surface flux tube is also a flux rope. In this paper, we adopt the previously accepted practice of calling the original sub-surface field configuration a “flux tube” and the twisted coronal structure that is formed during the emergence process a “flux rope.”

**Figure 2.** Initial 3D configuration for simulation SD. Panel (a) shows the initial sub-surface flux tube and the dipole field, represented by the black field lines that originate from a line along $y = 0$ on the bottom boundary. Panel (b) shows a magnified slice of the horizontal field $B_t$ in the $y = 0$ plane as a color shading and a projection onto the $y = 0$ plane of the field lines of the dipole. (A color version of this figure is available in the online journal.)

3. RESULTS

3.1. Partial Emergence of a Sub-surface Flux Tube

The partial emergence of a sub-surface flux tube into the solar atmosphere has been studied and commented upon in a number of previous studies (Fan 2001; Archontis et al. 2004; Manchester et al. 2004; Murray & Hood 2008; MacTaggart & Hood 2009) and we direct the reader to those studies for a more detailed description. The salient points are these: the flux tube rises buoyantly until it reaches the convectively stable photosphere/chromosphere, where it temporarily halts and undergoes a large amount of horizontal expansion. Then, the upper portions of the deformed tube emerge via the magnetic buoyancy instability (Acheson 1979) through the photosphere/chromosphere, transition region, and into the corona.

The emergence through the surface of the rising flux tube in simulation SD is shown in Figure 3. Note that the boundary conditions employed here allow the same field lines to be tracked throughout the simulation to a good approximation by using the same seed point on the side ($y = \pm \max(y)$) boundary. Unless stated otherwise, the same field lines are drawn for each panel in a given figure. The black line in Figure 3 is the field line that intersects the side ($y = \pm \max(y)$) boundaries at the location of the original convection zone tube axis (note that at this early stage this is one single field line).

As shown in Figure 3, panel (a), the upper field lines of the flux tube at $t = 35 t_0$ have penetrated the surface ($z = 0$) from below. These field lines have a high tilt relative to the axis of the tube ($y$ axis) and they create a bipolar structure on the surface with a neutral line parallel to the axis of the tube. As time progresses, field lines emerge with less tilt, i.e., more aligned with the axis of the flux tube. This creates an apparent shearing of the bipole, as shown in Figure 3, panels (e) and (f), and the two polarity regions drift apart. This behavior is representative of the observed evolution of emerging active regions (Luoni et al. 2011). The emerging field pushes the pre-existing dipole field both vertically and horizontally. These upper field lines of the flux tube are nearly parallel to the pre-existing dipole field and so this minimizes the amount of reconnection between the two flux systems.
Figure 3. Early emergence of a convection zone magnetic flux tube into the solar atmosphere for simulation SD at times 35\(t_0\) (panels (a) and (d)), 45\(t_0\) (panels (b) and (e)), and 50\(t_0\) (panels (c) and (f)). The black line originates from the location of the original flux tube axis on the side (\(y = \pm \text{max } y\)) boundaries. The colored lines originate from a circle on both the side boundaries centered on this axis. The gray lines originate at the lower boundary and belong to the dipole field. The color shading slice shows the vertical magnetic field \(B_z\) at the surface (\(z = 0\)). This figure illustrates that the apparent shearing of the bipolar structure is associated with the emergence of the flux tube’s axis and the magnetic field that is more aligned with this axis.

(A color version of this figure is available in the online journal.)

Figure 4. Partial emergence of the convection zone flux tube for simulation SD at times 50\(t_0\) (panels (a) and (d)), 55\(t_0\) (panels (b) and (e)), and 60\(t_0\) (panels (c) and (f)). The field lines shown are the same as in Figure 3. The color shading of the vertical magnetic field at the surface has been saturated to highlight the neutral line (which is located in white regions between red and blue regions). Above the flux tube axis (black line), the field lines are concave down and are able to drain mass and continue to rise. Below the axis, the field lines are concave up. These sections of the field lines carry mass that cannot be drained and are

Figure 4 shows the same field lines at later times, but with the color shading of \(B_z\) at the surface now saturated to highlight the neutral line (which appears white between red and blue). The sections of the field lines that cross above the axis of the flux tube (black line) are concave down. These sections of field lines are able to rise further as they drain mass. Beneath the emerged axis field line, the sections of the field lines are concave up. These sections carry mass that cannot be drained and are
Figure 5. Rotation of sunspots and the formation of a coronal flux rope at times $t = 80, 100, 120,$ and $140 t_0$. The color shading shows the vertical component of the magnetic field in the $z = 0$ plane. The arrows represent horizontal velocities perpendicular to the magnetic field in the $z = 0$ plane and are scaled by magnitude. The black and purple field lines originate at $\pm \max y$, respectively, at the intersection of the original convection zone tube axis and the side boundaries. The yellow lines originate at the lower boundary and belong to the dipole field. The velocity vectors show a strong shearing component, but also suggest a rotational motion near the center of each polarity region. A new flux rope axis can be defined at the location of the O-point in the $y = 0$ plane when the black and purple field lines separate. This new axis field line is shown as the green field line.

(A color version of this figure is available in the online journal.)

therefore unable to rise further into the atmosphere, as originally found in the simulations of Fan (2001) and Manchester et al. (2004). As a result, the original flux tube emerges only partially: the sections of the field that are concave down can expand into the corona, while the sections beneath the axis that are concave up remain trapped near the surface. At $t = 60 t_0$, the original flux tube axis has emerged to $3 L_0$ above the surface and there is an O-point above the surface in the $y = 0$ plane that this axis goes through. Previous authors have reported on the location of the original flux tube axis and found that for similar flux tube parameters as those used in this paper, the flux tube axis remains close to or below the surface, typically less than $3 L_0$ above the surface (Magara 2001; Murray et al. 2006; Fan 2001). However, these simulations do not explore the later evolution of the flux tube axis. As we shall show, the original axis of the convection zone flux tube splits into two new field lines and these new field lines twist around a new coronal flux rope axis.

3.2. Formation of a Coronal Flux Rope

Figures 5 and 6 show the active region for simulation SD at times $[80, 100, 120, 140] t_0$. The color shading shows the vertical component of the magnetic field in the $z = 0$ plane. Figure 5 also shows the horizontal velocities perpendicular to the magnetic field on the $z = 0$ plane as vectors, which excludes flows caused merely by plasma draining along field lines. Also shown in Figures 5 and 6 are selected field lines. The yellow field lines are the dipole field and originate at the lower boundary. The black and purple field lines are line tied at and originate from, the $y = \pm \max y$ boundaries, respectively, at the location of the original flux tube axis. These two field lines are coincident early in the simulation and pass through the O-point located in the $y = 0$ plane just above the surface as the original flux tube partially emerges. As can be seen in Figures 5 and 6, these two field lines, which were once the same, separate (perhaps due to magnetic diffusion) and appear to twist about each other as time progresses. They also rise higher into the corona as they do so.

It should be noted that due to a non-zero value of resistivity, the tracking of field lines cannot be exact (even with zero resistivity, there is some numerical diffusion to the scheme used to solve the induction equation). We have performed an additional simulation with ideal $\eta = 0$ and found that this splitting of the original axis field line also occurs, which suggests it may be independent of the choice of resistivity used.
Figure 6. Same as Figure 5 but from a different viewpoint. In addition, blue iso-surfaces of $j \geq 0.035 \mu_0 \beta$ are plotted for $z \geq 5 L_0$.

(A color version of this figure is available in the online journal.)

Figure 7 shows the $(x, y)$ locations of the intersection of the two former axial field lines with the surface ($z = 0$) at different times in simulation SD. The black diamonds (purple stars) represent locations for the black (purple) field line in Figure 5. The locations are taken relative to the center of the polarity region, which we define as the $(x, y)$ location where the field line that goes through the central O-point in the corona (the green line in Figure 5) intersects the surface. The results from both polarity regions are superimposed onto one single plot. Figure 7 shows a partial rotational motion of these locations around the center of each polarity region, with the same sign of rotation for each polarity region. A rotational motion is also suggested from the vectors of horizontal velocity perpendicular to the magnetic field in Figure 5. These motions reflect the transport of twist from the convection zone into the corona (see below). Since the green field line in Figure 5 passes through the O-point without exhibiting significant writhe, it can be considered as a good approximation of the axis of the successively forming coronal flux rope.

Previous simulations have suggested two different mechanisms for the formation of a coronal flux rope during magnetic flux emergence. Magara (2006) and Fan (2009) suggested that rotational motions, brought about by an equilibration of twist along emerging field lines, can twist up the coronal sections of field lines to create a new flux rope. On the other hand, Manchester et al. (2004), Archontis & Török (2008), and Archontis & Hood (2012) suggested that the reconnection of emerged sections of sheared field lines can create twisted field lines, resulting in a flux rope structure in the corona. We now briefly discuss these two mechanisms.

Figure 8 shows simulation SD at times $t = 50 t_0$, $t = 100 t_0$, and $t = 200 t_0$. To give a sense of the local twist per unit length, the field lines are colored with the quantity $\alpha L_0 = \mu_0 \mathbf{L}_0 \mathbf{J} |\mathbf{B}|^2$. As the upper part of the flux tube emerges into the atmosphere, the field lines expand into the low-$\beta$ atmosphere, increasing their length. As $\alpha$ is related to the twist per unit length and the tube expands faster than the twist propagates upward, this creates a gradient in $\alpha$ along the expanding field line. Such a gradient was also observed in the simulations of Fan (2009) and discussed in Longcope & Welsch (2000). A gradient in twist along a section of a flux tube will drive torsional Alfvén waves that equilibrate this twist. Figure 9 shows the quantity $\alpha$ at times $t = 100 t_0$ and $t = 200 t_0$ as a function of height along a portion of the purple field line from Figures 5 and 6 as it penetrates the surface and passes into the corona. The magnitude of the gradient of $\alpha$ around $z = 0$ clearly decreases with time, indicating that the twist is equilibrating along this section of the flux tube. Magara (2006) and Fan (2009) suggested that the torsional motions brought about by this process are capable of causing sunspot rotation that twists up the magnetic field in the corona. This idea is also supported by recent observations of the formation of active regions that suggest that sunspot rotation can be attributed to the emergence of twisted magnetic fields (Kumar et al. 2013).

It has also been suggested that magnetic reconnection is responsible for the formation of coronal flux ropes, by a process similar to what has been suggested based on observations of photospheric flux cancellation (van Ballegooijen & Martens 1989). In flux emergence simulations, the reconnection is driven by a combination of shearing flows, caused by Lorentz forces in the expanding field and inflows caused by pressure gradients (Manchester et al. 2004; Archontis & Török 2008; Archontis & Hood 2012). However, we see no direct evidence of magnetic reconnection, such as an X-point, outflow jets, or curved reconnecting field, underneath the flux rope axis in the simulations described in this paper. This may be due to the combined effects of (1) the limited expansion of the emerging field in the corona due to the presence of the dipole field (which may suppress the amplification of reconnection below the flux rope to a level at which it does not produce noticeable outflow.
Figure 8. Twist along field lines at times $t = 50 t_0$, $t = 100 t_0$, and $t = 200 t_0$ for simulation SD. Red–white–blue field lines are colored with $\alpha L_0 = \mu_0 L_0 | \mathbf{J} \cdot \mathbf{B} | / |\mathbf{B}|^2$ and originate from both side ($y = \pm \min y$) boundaries. The solid yellow field lines originate at the base of the domain and belong to the magnetic dipole field. This figure demonstrates how the emergence of field into the corona causes a gradient in twist along a field line as it goes from the convection zone into the corona. (A color version of this figure is available in the online journal.)

Figure 9. $\alpha L_0$ as a function of $z$ along a field line at two different times in the simulation SD, $t = 100 t_0$ and $t = 200 t_0$. The field line is the same for both times and originates at the location of the original convection zone flux tube axis on the $y = \min y$ boundary (the purple field line in Figures 5 and 6). After $t = 80 t_0$, this field line splits from the axis of the flux tube and expands into the corona. A decrease in the gradient in $\alpha$ between $z = -5 L_0$ and $z = 0$ can be seen from time $100 t_0$ to $200 t_0$. The original value of $\alpha$ along this field line at $t = 0$ is $0.76/L_0$.

velocities) and (2) the relatively high resistivity used here, which may suppress the build up of a steep current layer. We also see no direct evidence for reconnection in the simulation with $\eta = 0$, which suggests that the confinement by the dipole field in our simulations, rather than the relatively high resistivity, is the reason that magnetic reconnection beneath the flux rope axis is suppressed. We conclude that the formation process in our simulations is primarily due to the rotation of the polarity regions and the twisting of the field.

Figure 6 also shows an iso-surface of current density above $0.03 j_0$ in the region above $z = 5 L_0$ to highlight the current distribution below the flux rope. At $t = 80 t_0$, the current density is larger above the two regions of concentrated opposite polarity vertical magnetic field than above the center of the bipolar region. After $t = 80 t_0$, there is an increase in current density in the center. The predominant shape of the current sheet when viewed from above is of two distorted J-shapes that merge later to form one S-shape, a process that has been reported in previous flux emergence simulations (Fan 2009; Archontis & Hood 2012). Recent extreme ultraviolet observations of active regions have shown that high temperature (6 MK), J-shaped loops exist before the formation of coronal flux ropes (Liu et al. 2010) and that these J-structures combine to form a single S-shaped structure when the flux rope is formed (e.g., McKenzie & Canfield 2008; Aulanier et al. 2010). Such S-shaped sigmoid structures have been observed as precursors to CMEs (Sterling 2000).

Figure 10 shows the later evolution of simulation SD. The rotation of the two opposite polarity regions decreases after $t = 180 t_0$, but there is still significant twisting of the field lines that extend into the corona. The field lines that defined the original convection zone flux tube’s axis (black and purple) both wrap around the new flux rope axis in the corona. They also have a pinched U-shape at the center of the active region, which creates the strong current sheet structure.

(A color version of this figure is available in the online journal.)
3.3. Confinement by the Overlying Field

Figure 11 compares the simulation with ND and the simulation with the strongest dipole (SD) at a late stage in the flux rope formation process \( t = 180 t_0 \) as the envelope of the flux rope expands further into the corona. In simulation SD, the dipole field, which was chosen to be aligned so as to minimize reconnection with this envelope field, constrains the expansion (both vertically and horizontally).

Figure 12 shows the height of the axis of the coronal flux rope and the height of the envelope field as a function of time. The height of the axis of the coronal flux rope is found by locating the point along the \( z \)-axis at which the horizontal field \( B_y \) is zero. This point is approximately the location of an O-point, i.e., where \( \sqrt{B_x^2 + B_y^2} = 0 \), and the field line that goes through this O-point appears to have very little writhe, as shown in Figure 10. The axis of the new flux rope is therefore well represented by this field line. We define the height of the envelope field by the intersection of the \( z \)-axis and the contour of \( B_z = 1.0 B_{\text{min}} \). Ideally, we would use the separatrix between the dipole field and the expanding flux rope field to measure the height of the envelope field, but no such separatrix exists in simulation ND. The value of \( 0.1 B_{\text{min}} \) was chosen because the intersection of the \( z \)-axis and this contour is where the separatrix between the dipole field and the expanding flux rope field is located in simulations SD, MD, and WD for times \( t > 50 t_0 \).

From \( t = 200 t_0 \) to \( t = 450 t_0 \), there is a slow rise of the flux rope, which appears to tend to a stable position. The vertical velocities at the envelope field fall from a typical value of \( 1.5 v_0 \) at \( t = 200 t_0 \) to \( 0.15 v_0 \) at \( t = 450 t_0 \). As can be seen in Figure 12, the height of the axis of the coronal flux rope at time \( t = 450 t_0 \) is smaller for larger dipole field strength, as expected. The simulation with no dipole, simulation ND, exhibits the strongest expansion, which continues until the envelope field of the flux rope approaches the damping region near the top boundary. From Figure 12, it appears that the flux rope in simulations SD, MD, WD, and ND are ultimately stable, but given that the envelope field is so close to the damping region in simulation ND, the effect of the boundary conditions on the stability cannot be ruled out. To investigate this, we perform an additional simulation, ND1, where the top boundary is extended further out, as is the damping region near this boundary (as described in Section 2). We find that envelope field does not extend past \( \pm 80 L_0 \) in the \( x \)- or \( y \)-directions and so the side boundaries and the damping region at these side boundaries do not play a role in the stability of the flux rope. Therefore, we do not change these in simulation ND1. Figure 12 shows only a small difference in the curves between simulations ND and ND1. In both cases, the height of the envelope field appears to saturate at \( 180 L_0 \), which is well below the height of the top damping region for simulation ND1. We conclude that the confinement of the flux rope in the case of an initially field free corona is not a consequence of the boundary conditions, but of the self-stabilization of the flux rope by its own envelope field.

Previous simulations by Archontis & Hood (2012) with the same initial tube strength and twist as the simulations in this paper and without any pre-existing coronal field also suggest that the flux rope is ultimately stable. However, the flux rope axis in their simulations reaches a lower height of \( 62.3 L_0 \) above the surface compared with heights of \( 108 L_0 \) and \( 110 L_0 \) for the flux ropes in simulations ND and ND1 in this paper, respectively. This fact, together with the fact that in the simulations of Archontis & Hood (2012) the envelope of the flux rope reaches a height of \( 127 L_0 \) above the surface, at the boundary of the damping region between \( 127 L_0 \) to \( 130 L_0 \) above the surface, and thus very close to the top boundary at \( 130 L_0 \) above the surface, suggests that their boundary conditions are affecting the ultimate height of the flux rope. However, their conclusion, that the coronal flux rope is stabilized by its own envelope field, is supported by our simulations ND and ND1.
In the simulations of Manchester et al. (2004), which used a slightly thinner tube ($w = 2 \, L_0$ compared with $w = 2.5 \, L_0$ here) and were placed initially higher in the convection zone ($z = -10 \, L_0$ compared with $z = -12 \, L_0$ here), the O-point of the flux rope that formed in the corona rose to a height of $50 \, L_0$ by $t = 70 \, t_0$ at the end of the simulation. This is higher than the flux rope reaches by $t = 70 \, t_0$ in simulations ND and ND1, but the flux ropes in simulations ND and ND1 achieve heights well above $50 \, L_0$ later in time, when they become stable. While the flux rope rises quickly during the initial phase presented in Manchester et al. (2004), based on the findings in this paper, it seems likely that the flux rope formed in Manchester et al. (2004) would also ultimately be confined by its own envelope field, if its evolution was followed long enough.

By including a dipole field in simulations SD, MD, and WD, we are also able to constrain the flux rope at lower heights than its own envelope field is able to hold it at. Thus, the dipole field is suppressing the rise of the coronal flux rope. The magnetic forces and plasma $\beta$ in the $y = 0$ plane are shown in Figure 13 for simulation SD at time $t = 280 \, t_0$. The magnitudes of the magnetic tension force and the magnetic pressure forces in the $y = 0$ plane are shown in panels (a) and (b), respectively. Above a height of $z = 10 \, L_0$, the two forces approximately cancel and the magnetic field associated with the flux rope and dipole field has a small Lorentz force relative to the magnitude of the magnetic pressure and tension forces. As can be seen from panel (d), $z = 10 \, L_0$ is the height at which the plasma $\beta$ undergoes the transition from above unity to below unity. Above $z = 10 \, L_0$, the magnetic field configuration has approximately $j \times B = 0$.

3.4. Distribution of Electric Currents

At present, there is much debate as to whether the electric currents in active regions are “neutralized,” in the sense that, for a single sunspot or active region polarity, the direct current (current aligned parallel to the axial magnetic field of the flux tube associated with that sunspot) is surrounded by a return current (aligned anti-parallel to the axial field), which cancels this direct current out. This is important for flare and CME modeling as some models use an initial magnetic field with a net current (i.e., return currents are either absent or not large enough to neutralize the direct currents), e.g., the TD model. Although in
the simulations presented here the initial sub-surface flux tube is current neutralized, i.e., it has no net current, there is significant distortion of the magnetic field by the emergence process and so it is not clear that the resulting coronal flux rope will also be current neutralized.

To investigate this issue, we plot in Figure 14 the electric currents at time $t = 240 \, t_0$ in simulation SD. Panels (a) and (b) show slices of vertical current $j_y$ at a height of $z = 10 \, L_0$, the top of the photosphere/chromosphere region in the model, where $\beta = 1$ for simulation SD. This height is chosen so as to eliminate any overshoot convective flows that distort the magnetic field. Panels (a) and (b) also show current field lines (field lines of $j_y$). The current field lines are colored by the sign of $j_y$ at their location of origin on the side boundary: red for direct current ($j_y > 0$) and blue for return current ($j_y < 0$). The field lines are located at regular values of radius at the convection zone flux tube on the side boundary. These radial values are $0.4 \, L_0 + (0.8 \, n \, L_0)$ for $n = 0, 7$ ($j_y$ changes sign at the radial value $2.4 \, L_0$). The number of field lines at each radial value $r$ is proportional to the total unsigned axial current in the annulus $2 \pi r \, dr$ centered on that radius $r$, where $dr = 0.8 \, L_0$. The total number of field lines is 30, so each field line represents $1/30$ of the total unsigned axial current in the entire flux rope. For a given current field line there are only two routes by which it can return to a side boundary. First, it can exit through the opposite boundary. Second, it can reverse direction and return to the same boundary in a region of opposite current. Figure 14 shows that it is mostly current field lines that originate in regions of direct ($j_y > 0$) current on the side boundary that enter the corona above $z = 10 \, L_0$. Figure 14, panel (c) shows that a strong central positive $j_z$ develops above $z = 10$ in the $y = 0$ plane. As predicted by the 2.5D model of Longcope & Welsch (2000), there is a return current that flows along the interface between the sub-photosphere and corona, although the simulations in this paper show that some return current extends into the corona.

Note that a single blue line emerges into the corona in Figure 14, panels (a) and (b). This current field line originates from a region of negative $j_y$ (but positive $B_z$) on the $y = \min \, y$ plane and so is considered return current. If this field line were to follow a simple $\Omega$ shaped path from one boundary to the other, it would intersect the $z = 10 \, L_0$ plane such that the current normal to that plane $j_z$ would be anti-parallel to the magnetic field normal to the plane $B_z$ and it would be considered return current in the corona. However, this current field line, which is representative of many others, performs a complicated circuit, first crossing underneath the active region before passing into the corona and back into the convection zone. This loop-like circuit results in the field line having $j_z$ parallel to the magnetic field $B_z$ on the $z = 10 \, L_0$ plane. In this sense, for the $z = 10 \, L_0$ plane, the blue field line is a direct current field line, even though for its seed location on the $y = \min \, y$ plane it is a return current field line. This changing of currents is due to the complicated current structure underneath $z = 10 \, L_0$, where the currents are far from force-free and plasma motions can dominate over magnetic forces.

These results suggest that the coronal flux rope is not neutralized, in the sense that there is not a balance of direct and return current. Of course, there may be very diffuse return currents surrounding this flux rope and further, more rigorous, analysis is required to determine whether or not the flux rope is indeed un-neutralized. An in-depth analysis of the neutralization of active region currents in an analogous simulation is presented in Török et al. (2013).

4. DISCUSSION

The aim of this paper is to use 3D visco-resistive MHD simulations to investigate whether convection zone flux tube emergence could create coronal magnetic field configurations compatible with a flux rope model such as the TD model, where a net-current coronal flux rope is tethered by an overlying potential field. Consistent with previous simulations, we found that the initial convection zone flux tube partially emerged into the corona; only sections of field lines that were able to shed mass were able to emerge. The original flux tube axis first reached a height of 3 Mm above the surface. This is consistent with simulations by Magara (2001), Fan (2001), and Murray et al. (2006).

As a result of the transport of twist from the convection zone into the corona, torsional motions manifested themselves in corotation of the opposite-polarity regions and effectively twisted up the field in the corona, as originally shown by Fan (2009). The field line associated with the original convection zone flux tube axis separated into two field lines due to magnetic diffusion and became wrapped around a new flux rope axis in the corona. Two distinct J-shaped current layers beneath the new flux rope axis formed, which began to merge during the rotation.
of the sunspots. This process of emergence and equilibration of twist supports the conclusions from observations that sunspot rotation is driven by twisted flux tube emergence and that it can cause the formation of sigmoids prior to a solar flare (e.g., Min & Chae 2009; Kumar et al. 2013).

No obvious evidence of magnetic reconnection was seen at the location of the current layer below the new coronal flux rope axis, such as the evidence presented in Manchester et al. (2004), Archontis & Török (2008), and Archontis & Hood (2012). In those simulations, there was slow, steady reconnection at the location of the current sheet during the expansion of the emerged field in the corona and this reconnection amplified as the flux rope rose to successively larger heights. Because in the simulations in this paper the emerged field was constrained by the dipole field, we did not see this reconnection stage clearly. Since, however, we did see evidence of rotational motions in sunspots, as suggested by Fan (2009), we conclude that the flux rope formation process is predominantly due to these motions in our simulations.

By varying the height of the top boundary and the upper velocity damping region and finding that the ultimate height of the flux rope axis was unchanged, we removed the effect of the top boundary conditions on the stability of the flux rope and concluded that even without a dipole field in the corona, the flux rope was constrained by its own envelope field, which support the results by Archontis & Török (2008) and Archontis & Hood (2012), which were achieved for smaller simulation boxes.

By adding a dipole field, aligned so as to minimize reconnection with the emerging field in the corona, we were able to constrain the expansion of the active region into the corona. The stronger the dipole field, the lower the height of the newly formed coronal flux rope, as expected. Such a system of a coronal current-carrying flux rope (or, alternatively, a strongly sheared arcade) stabilized by an overlying potential field is a canonical configuration believed to produce solar eruptions. We found that the relatively simple, idealized initial conditions used in our simulations, with a twisted convection zone flux tube emerging into a dipole field representing a decaying active region, is able to robustly produce such a coronal configuration.

A simple analysis of the electric currents suggests that the majority of the return currents did not emerge into the corona and so a coronal flux rope with a non-neutralized current was created. Further analysis is presented in Török et al. (2013). The preliminary results presented here suggest that coronal flux rope models that consider only direct currents, such as the TD model, are compatible with the magnetic fields created by the emergence of a twisted magnetic flux tube.

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