EMERGENCE OF A FLUX TUBE THROUGH A PARTIALLY IONIZED SOLAR ATMOSPHERE

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Received 2007 April 13; accepted 2007 May 19

ABSTRACT

For a magnetic flux tube, or indeed any flux, to emerge into the solar corona from the convection zone it must pass through the partially ionized layers of the lower atmosphere: the photosphere and the chromosphere. In such regions, the ion-neutral collisions lead to an increased resistivity for currents flowing across magnetic field lines. This Cowling resistivity can exceed the Spitzer resistivity by orders of magnitude and, in 2.5-dimensional (2.5D) simulations, has been shown to be sufficient to remove all cross field current from emerging flux. Here we extend this modeling into three dimensions (3D). Once again it is found that the Cowling resistivity removes perpendicular current. However, the presence of 3D structure prevents the simple comparison possible in 2.5D simulations. With a fully ionized atmosphere, the flux emergence leads to an unphysically low temperature region in the overlying corona, lifting of chromospheric material, and the subsequent onset of the Rayleigh-Taylor instability. Including neutrals removes the low-temperature region, lifts less chromospheric matter, and shows no signs of the Rayleigh-Taylor instability. Simulations of flux emergence therefore should include such a neutral layer in order to obtain the correct perpendicular current, remove the Rayleigh-Taylor instability, and get the correct temperature profile. In situations when the temperature is not important, i.e., when no simulated spectral emission is required, a simple model for the neutral layer is demonstrated to adequately reproduce the results of fully consistent simulations.

Subject headings: MHD — Sun: chromosphere — Sun: magnetic fields

1. INTRODUCTION

The emergence of new magnetic flux into the solar corona is responsible for the formation of active regions. The accepted view is that the emergence of Ω-shaped flux tubes through the photosphere is responsible for the formation of sunspots. It is clear therefore that the movement of magnetic flux from the convection zone up into the corona is one of the most significant drivers determining the structure of the corona. Furthermore, the input of new flux into the preformed corona is often the trigger mechanism for dynamic coronal activity such as prominence eruptions, flares, and CMEs. Any attempt at a unified model of solar activity must couple the magnetic field of the convection zone with that of the corona, and hence, a full understanding of the flux emergence process which connects them is essential.

The problem with studying flux emergence is that it must couple subphotospheric plasma with coronal plasma. In traversing this region, the magnetic field moves from regions which are convectively unstable to convectively stable, through orders-of-magnitude changes in equilibrium density and a rapid increase in temperature. The physics of each of these regions is therefore often dominated by different processes, and analytical treatment of the whole emergence process is therefore limited. As a result, this subject is now largely investigated by numerical simulations. A typical flux tube, assuming such a well-defined structure exists in the convection zone, must have sufficient twist to survive its transit of the convection zone (Moreno-Insertis & Emonet 1996; Dorch & Nordlund 1998). It will then reach the photosphere where the buoyancy instability becomes active, allowing the flux to escape into the corona (Matsumoto & Shibata 1992; Murray et al. 2006). Once the flux reaches the region of the chromosphere and corona, where the density drops by 6 orders of magnitude over a couple of megameters in height, the flux tube expands (Matsumoto et al. 1993; Magara & Longcope 2001; Fan 2001). Below the photosphere, the plasma-β β ≫ 1 and the twist of the flux tube give rise to a j × B force which is easily balanced by small changes in the much larger kinetic pressure terms. In the corona, the plasma is characterized by β ≪ 1, and any emerging j × B force cannot be balanced by kinetic pressure, and thus, the flux tube will expand rapidly into a configuration in which any residual j × B force is of the order of β; i.e., the flux expands until the coronal field is force free. The initial twist in the emerging flux tube affects the emergence process and how close the coronal field is to force free when it first reaches the corona (Abbet & Fisher 2003; Murray et al. 2006). This assumes that there is no overlying field with which the emerging flux can interact. Often such a field does exist, and the interaction of this new flux with existing magnetic structures can lead to complex, dynamic behavior (Archontis et al. 2004; Galsgaard et al. 2005). The emerging magnetic flux, since it may have a nonzero j × B force, is also capable of lifting chromospheric material up into the corona. As the field expands, reducing the magnetic forces, this heavier material may trigger the Rayleigh-Taylor instability (Isobe et al. 2005).

All of the works cited above have studied flux emergence by using the magnetohydrodynamic (MHD) equations appropriate for a fully ionized plasma. They make the further assumption that the parallel and perpendicular resistivities are equal; whereas these differ by a factor ≈ 2 for a fully ionized plasma. This factor of 2 is routinely ignored in MHD simulations, as the resistivity used, for numerical reasons, exceeds the physical value by many orders of magnitude. For simulations of the coronal or convective regions of the Sun’s atmosphere, this is a perfectly valid approximation. However, the photosphere and chromosphere are not fully ionized plasma due to their low temperature. It is well known that ion-neutral collisions add an effective anisotropic resistivity into the single-fluid equations (Cowling 1957; Braginskii 1965). This additional effect, the Cowling resistivity, acts only on perpendicular currents, i.e., those flowing across the magnetic field, and can be many orders of magnitude larger than the parallel Spitzer resistivity in the chromosphere (Khodachenko et al. 2004). This dissipation of perpendicular currents by Cowling
resistivity has been used to study the damping of MHD waves in the chromosphere (Goodman 2000; Leake et al. 2005) and flux emergence in 2.5 dimensions (2.5D; Leake & Arber 2006). In Leake & Arber (2006), it was shown that the Cowling resistivity was sufficient to destroy all perpendicular currents in 2.5D flux emergence simulations. This is significant as it forces emerged flux to be force free as it traverses the chromosphere. The restricted resistivity has been used to study the damping of MHD waves in close fluid and include the effects of partial ionization through the standard Saha equation (Brown 1973) which can be solved for the steady state ionization equation based on the local temperature with the radiation temperature fixed at the photospheric value (Thomas & Athay 1961). We have

\[
\frac{n_i^2}{n_a} = f(T),
\]

\[
f(T) = \left(\frac{2\pi m_i k_B T}{h^3}\right)^{3/2} \exp\left(-\frac{X_i}{k_B T}\right),
\]

\[
b(T) = T \frac{w T_R}{w T_R} \exp\left[\frac{X_i}{4k_B T} \left(\frac{T}{T_R} - 1\right)\right],
\]

where \(T_R\) is the temperature of the photospheric radiation field and \(w = 0.5\) is its dilution factor. Using this equation, the ratio of the number density of neutrals to ions is given by

\[
r = \frac{n_a}{n_i} = \frac{1}{2} \left[\frac{1}{1 + \sqrt{\left(1 + \frac{4\rho/m_i}{n_i/n_a}\right)}}\right],
\]

and the neutral fraction \(\xi_n = n_a/n_i\) is

\[
\xi_n = \frac{r}{1 + r}.
\]

The dominant effect of the neutral atoms is how they modify Ohm’s law and consequently lead to an anisotropic dissipation of current. The full derivation of the resistive terms in equations (3) and (4) (Cowling 1957; Braginskii 1965) shows that it is the collisions between ions and neutrals which are most important. For a hydrogen plasma and in the limit when classical resistivity \(\eta\) can be ignored, the Cowling resistivity is given by

\[
\eta_{C} = \frac{\xi_n^2 B^2}{\alpha_n},
\]

with

\[
\alpha_n = 1 + \frac{1}{2} (1 - \xi_n) \frac{\rho^2}{m_i \sqrt{\pi m_i \Sigma_{in}}},
\]

where \(\Sigma_{in} = 5 \times 10^{-19}\) m\(^2\) is the ion-neutral collision cross section. When these formulae are applied to model chromospheres (Khodachenko et al. 2004), it is found that the Cowling resistivity \(\eta_{C}\) can exceed the classical resistivity \(\eta\) by many orders of magnitude. From 2.5D simulations (Leake & Arber 2006), it has been shown that the equations above include the dominant corrections to Ohm’s law and that the neglect of the Hall term is valid for these flux emergence studies. Furthermore, it was also shown that in the upper chromosphere the dominant contribution to Ohm’s law is through the Cowling resistivity, not the advective term.

The final term in equation (4) is designed to model all of the missing acoustic shock heating, radiative transport, and thermal conduction terms. These terms would act to restore the equilibrium photosphere and chromosphere, but are too computationally expensive, or simply unknown, and cannot be explicitly included.

2. EQUATIONS AND INITIAL CONDITIONS

2.1. Equations

The standard MHD equations, in Lagrangian form, are modified to include the effects of anisotropic current dissipation. For simplicity, it is assumed that the atmosphere is composed entirely of hydrogen. The resulting set of equations apply to a single fluid and include the effects of partial ionization through the neutral fraction \(\xi_n = n_a/(n_i + n_a)\), where \(n_a\) is the neutral number density and \(n_i\) is the ion number density. We have

\[
\frac{Dp}{Dt} = -\rho \nabla \cdot v,
\]

\[
\frac{DIV}{D} = -\frac{1}{\rho} \nabla P + \frac{1}{\rho} J \times B + \frac{1}{\rho} \nabla \cdot S,
\]

\[
\frac{DB}{Dt} = \frac{1}{\rho} \nabla \cdot v - B(\nabla \cdot v) - \nabla \cdot \left(\frac{\eta_{f} J}{\rho} - \nabla \cdot (\eta_{L} J_{\perp})\right),
\]

\[
\frac{Dv}{Dt} = -\frac{F}{\rho} J + \frac{1}{\rho} \eta_{f} J_{\parallel}^2 + \frac{1}{\rho} \eta_{L} J_{\perp}^2 + \frac{1}{\rho} \frac{\eta_{r} S_{ij}}{\rho} - \frac{\epsilon - eo(\rho)}{\tau} v,
\]

where the parallel and perpendicular current vectors, \(J_{\parallel}\) and \(J_{\perp}\), respectively, are defined as

\[
J_{\parallel} = \frac{(J \cdot B) B}{||B||^2},
\]

\[
J_{\perp} = B \times (J \times B),
\]

\[
\rho = \text{mass density}, P = \text{gas pressure}, \epsilon = \text{internal specific energy density}, v = \text{center-of-mass velocity of the fluid}, B = \text{magnetic field}, g = \text{gravitational acceleration}, S = \text{stress tensor which has components} S_{ij} = \nu (\partial v_{ij} - \frac{1}{2} \delta_{ij} \nabla \cdot v),\]

where \(\eta_{f}\) is the magnetic field, \(\eta_{L}\) is the ion number density. We have \(\eta_{r} = \frac{1}{2} \left[(\partial v_{ij} \partial v_{kl} + (\partial v_{ij} \partial v_{kl}))\right].\)

Since the plasma is not fully ionized, the total pressure \(P\) and specific internal energy \(\epsilon\) include the neutral fraction \(\xi_n\) through

\[
P = \rho k_B T / \mu_m,
\]

\[
\epsilon = \frac{k_B T}{\mu_m (\gamma - 1)} (1 - \xi_n) X_i / \mu_i,
\]

where \(k_B\) is Boltzmann’s constant, \(\gamma\) is the ratio of specific heats, \(\mu_m = m_i / (2 - \xi_n)\) is the reduced mass, and \(X_i\) is the ionization energy of hydrogen. Equation (4) can be used to numerically advance \(\epsilon\), but \(\xi_n\) is a function of temperature \(T\) so equation (8) must be solved implicitly for \(T\), which can then be used to specify \(P\) through equation (7). For direct comparison with Leake & Arber (2006) and all previous 3D flux emergence simulations, here we simply set \(\mu_m = m_i / 2\) and ignore the \(X_i\) term in equation (8). As discussed in Leake & Arber (2006), this is unlikely to affect the emergence process through the chromosphere. The chromosphere is not in LTE, and the radiation temperature and thermodynamic temperature cannot be assumed to be the same. The complete calculation of the ionization state of hydrogen therefore requires the solution of the 3D radiative transfer and ionization equations. To save time for the 3D emergence problem, a simplified reduced model for \(\xi_n\) is used. This is based on a modified Saha equation (Brown 1973) which can be solved for the steady state ionization equation based on the local temperature with the radiation temperature fixed at the photospheric value (Thomas & Athay 1961).
As a first attempt at modeling the detailed heating and cooling terms omitted from equation (4), we add a Newton cooling term \(-\epsilon / \epsilon_0 / \tau\), where \(\tau\) is the timescale of the relaxation. The equilibrium specific energy density \(\epsilon_0\) is chosen to be a function of the density \(\rho\). The reasoning for this is related to the nature of these simulations. The buoyancy force drives magnetic field in the convection zone upward into the photosphere, where the field then expands into the atmosphere above. Thus, as a parcel of plasma from the convection zone of density \(\rho\) is moved upward into the photosphere, its temperature should be relaxed to its own initial temperature, rather than the local plasma temperature, which is of a different density.

A form for the timescale of this relaxation is required. For this the approach of Gudiksen & Nordlund (2005) is adopted. In simulating coronal heating, they chose \(\tau\) to depend on some power of the density,

\[
\tau = 0.1 (\rho / \rho_{ph})^{-1.7},
\]

so that at the relatively dense photosphere \((\rho = \rho_{ph})\) the timescale is about 0.1 s and is large enough that the effect becomes negligible in the sparse corona.

2.2. Initial Conditions

The modified MHD equations are normalized by division of the SI variables by photospheric values. The basic units are

\[
L_{ph} = 150 \text{ km}, \quad \nu_{ph} = 6.5 \text{ km s}^{-1},
\]

\[
\rho_{ph} = 2.7 \times 10^{-4} \text{ kg m}^{-3}, \quad B_{ph} = 1200 \text{ G},
\]

which gives the derived units

\[
t_{ph} = 23 \text{ s}, \quad T_{ph} = 6420 \text{ K},
\]

\[
P_{ph} = 1.2 \times 10^{4} \text{ Pa}.
\]

From here on, unless stated, all quoted values are internal code variables and should be multiplied by the above values to recover the SI variables. The differential equations (1) – (5) are advanced in time numerically using the Lagrangian remap code Lare3d (Arber et al. 2001). The physical domain simulated extends vertically from \(-20\) (3000 km below the surface) to \(130\) (19,500 km above). The horizontal extent is \(75\) (11,250 km) about the center of the domain, i.e., \(-75 \geq x \geq 75\) and \(-75 \geq y \geq 75\). The \(z\)-axis is vertical, the \(y\)-axis is across the tube, and the \(x\)-axis is aligned with the initial tube axis. Simulations have been run on 128\(^3\), 320\(^3\), and 512\(^3\) grids to check convergence. The computational grid was always uniform so that the minimum grid spacing used was \(\Delta x = 0.3\), i.e., \(\approx 44\) km.

The anisotropy in the resistivity prevents the induction equation from being cast in simple diffusive form. In order to estimate the relative magnitudes of the resistive and the implicit numerical diffusion contributions to equation (3), we consider the one-dimensional (1D) model equation \(\partial_t B + \nu \partial_x B = \eta \partial_x^2 B\). For the second-order accurate scheme employed here, the leading-order error term introduced in this model equation is of order \(\nu \Delta x^2 \partial_x \eta \partial_x B\) so that if a typical scale length in the dynamic evolution is \(L\), this gives an effective numerical resistivity of \(\nu \Delta x^2 / L\). In the upper chromosphere, where the Cowling resistivity is dominant, in normalized units the maximum Alfvén speed is 0.6. The worst case for numerical resolution corresponds to \(L \approx 1\), as this would place three grid points across a slightly diffuse shock. This gives a normalized implicit numerical resistivity of about 0.06. Note that this estimate is based on the fastest phase speed and shortest gradient scale length, and thus, a value of 0.06 represents an absolute upper limit on the numerical resistivity. The typical normalized \(\eta_c\) found in the simulations is of order 10, corresponding to a real magnetic diffusivity of \(9.75 \times 10^{6} \text{ m}^2 \text{ s}^{-1}\), and is therefore larger than numerical resistivity. Note that this value of diffusivity corresponds to a magnetic Lundquist number of order 0.1. This is larger than that found for estimates based on the quiet chromosphere in Leake & Arber (2006), because chromospheric material expands as it is lifted due to flux emergence and the associated adiabatic cooling increases the Cowling resistivity.

The initial stratification is a simple 1D model of the temperature profile of the Sun, which includes the upper 3000 km of the convection zone, the photosphere/ chromosphere, the transition region, and the base of the corona. The temperature profile consists of a linear polytrope for the convection zone with a vertical gradient at the critical adiabatic value

\[
dT/dz = \gamma - 1 \quad T \quad dP/dz.
\]

The temperature in the photosphere and chromosphere is assumed to be constant at 1, as is the temperature in the corona at a temperature of 150. These two regions are connected by a transition region of width \(w_{tr} = 5\). We have

\[
T(z) = T_{ph} - \frac{q}{m + 1} z, \quad z < 0,
\]

\[
= T_{ph} + \left(\frac{t_{cor} - t_{ph}}{2}\right) \left[ \tanh \left(\frac{z - z_{cor}}{w_{tr}}\right) + 1 \right], \quad z > 0,
\]

where \(m = 1/(\gamma - 1)\) is the adiabatic index for a polytrope, \(z_{cor} = 25\) is the height of the corona, \(t_{ph}\) is the photospheric temperature, and \(t_{cor} = 150\). The density and pressure of the background atmosphere are found from solving the hydrostatic equation

\[
dP/dy = -\rho g.
\]

A magnetic tube is placed in the convection zone at \(z = -10\) with the profile

\[
B_z = B_0 \exp \left(\frac{-r^2}{a^2}\right),
\]

\[
B_\phi = q r B_z,
\]

where \(r\) is the radial distance from the tube center in the \((y, z)\)-plane. The strength of the field at the center of the tube, \(B_0\), is 5, and the radius of the tube, \(a\), is chosen to be 2; \(q\) is the amount of twist in the loop defined as

\[
q = B_z / r B_\phi.
\]

This is set to be the minimum required to avoid fragmentation during the rise through the convection zone and is defined as \(|q| = 1/\alpha\) (Moreno-Insertis & Emonet 1996).

A choice must be made as to how to initialize the rise of the flux tube in the convection zone. It is thought that flux tubes formed from the toroidal field in the tachocline remain connected to the large-scale field by their roots (Zwann 1978), while the apex of the tube rises to the surface. As a result, a flux tube which reaches the surface will be significantly “bent” into an \(\Omega\)-shape.
In order to force the tube into this shape in these simulations, the
center is made buoyant while the ends are left in mechanical equi-
librium. This is done by setting the pressure in the tube different
from the field-free atmosphere \( p_0(z) \) by \( p_1(r) \), where
\[
\frac{dp_1(r)}{dr} \hat{e}_r = j \wedge B, \tag{24}
\]
so that the pressure gradient matches the Lorentz force. The
density in the tube differs from the field-free density \( \rho_0(z) \) by
\( \rho_1(r) \), where
\[
\rho_1(r) = \alpha \frac{p_1}{p_0(z)} \rho_0(z) \exp \left( -\frac{x^2}{\lambda^2} \right), \tag{25}
\]
where \( \alpha \) is used to scale the initial perturbation. In this paper, un-
less stated otherwise, \( \alpha = 0.1 \). With this perturbation, the center
of the tube, at \( x = 0 \), is buoyant, while for \( x > \lambda \), the tube is in
mechanical equilibrium \( (\rho_1 = 0) \). The value of \( \lambda \) is chosen to be
20, as in Fan (2001).

\[2.3. \text{Resistivity Models}\]

Three different models for the resistivity used in the flux
emergence have been studied. The first was the fully ionized
plasma model (labeled as FIP in later figures) in which ideal
MHD was used. This is the same model used by all previous 3D
flux emergence simulations and provides a benchmark against
which the effects of Cowling resistivity can be measured. The
second is the partially ionized plasma model (labeled as PIP in
later figures) which solves the equations including partial ioniza-
tion, the Cowling resistivity, and the Newton cooling as outlined
above. This is the same model that was used in Leake & Arber
(2006). The final model is based on a simple model for partial
ionization effects in which a time-independent perpendicular re-
sistivity profile is fixed in a layer of the upper chromosphere.

This model (labeled as \textit{Layer} in later figures) has \( \eta = 0 \) and the
Cowling resistivity fixed by
\[
\eta_C = 400B^2 \exp\left[ -(z - 10)^2 / 5 \right] \tag{26}
\]
in normalized units. This profile was chosen to closely match the
resistivity observed in the chromospheric layers during simula-
tions using the PIP model, but still retains the magnetic field
dependence.

\[3. \text{RESULTS}\]

The basic stages of the flux emergence process are the same as
in previous studies (e.g., Fan 2001; Archontis et al. 2004) with
the tube initially rising due to buoyancy. When the tube reaches
the photosphere, the buoyancy stops, the tube expands until suf-
ficient flux has built up for the magnetic buoyancy instability to
become important, and the field expands through the chromosphere.

\[\text{Fig. 1.} - \text{Field lines and flux for the partially ionized simulations at } t = 160. \text{ The shaded contour plot, at the photosphere } z = 0, \text{ shows vertical flux with positive flux shaded dark and negative flux shaded light. Dark field lines are those which connect to the initial equilibrium flux with } r = 2.0, \text{ while those shaded lighter correspond to field lines near the tube axis, i.e., } r = 0.5.\]

\[\text{Fig. 2.} - \text{Plot of the magnitude of } j_z \text{ as a function of height along the line } x = y = 0 \text{ for all three resistivity models at } t = 160.\]
The structure of the emerged field at $t = 160$ can be seen in Figure 1. In previous 2.5D studies of the effects of the partially ionized layers on flux emergence (Leake & Arber 2006), it was possible to quantify these effects by calculating the integrated perpendicular and parallel currents as functions of height. These calculations were conducted at the same time and only included flux inside the expanding envelope of flux. This simple measure of the effectiveness of the Cowling resistivity at removing perpendicular current is not possible in 3D due to the extra structure discussed below. Figure 2 shows the magnitude of $j_\perp$ as a function of height along the line $x = y = 0$ for all three resistivity models in the corona. While similar to the results in Leake & Arber (2006), there are a number of significant differences. First, the results in Figure 2 all show a peak in $j_\perp$ at the top of the emerged flux. This is the expanding shock between flux-free and emerging-flux regions and was used to define the region inside which the integrated flux of Leake & Arber (2006) was defined. More significantly, the larger $j_\perp$ in the FIP model simulations cannot now be exclusively attributed to $\eta_C$ removing this flux in the PIP simulations. The reason for this is that in 3D the chromospheric material raised into the corona in the FIP model is sufficient to trigger a Rayleigh-Taylor instability.

Figure 3 shows the structure of the perpendicular current density in a vertical slice through the center of the computational domain for the FIP and PIP models. The PIP model shows a slice through a symmetric, expanding shell of flux, while for the FIP model, there is a dropping central patch of enhanced $j_\perp$. This feature is due to the Rayleigh-Taylor instability (see discussion below), and as this leads to a bending and compression of field lines, it also contributes to the net $j_\perp$. Hence, it is not possible to directly attribute the reduction of $j_\perp$ in the PIP model, compared to the FIP model, shown in Figure 2 to the dissipation of $j_\perp$ in the partially ionized chromosphere, as some of the excess of $j_\perp$ in the FIP model is created in situ in the corona by the Rayleigh-Taylor instability.

The FIP model allows flux to emerge through the chromosphere with a larger $\mathbf{j} \times \mathbf{B}$ force than the PIP model in which the net $j_\perp$ is dissipated by Cowling resistivity. As a result, the FIP model lifts more chromospheric material into the corona and is susceptible to the Rayleigh-Taylor instability. This is shown by the isosurface of density in Figure 4 with the central density sheet dropping from the Rayleigh-Taylor instability clearly visible. Note that this density sheet is aligned with the magnetic field at the top of the emerging flux and, hence, as can be seen from Figure 1, is in a plane which crosses the tube axis. The $x$-direction is therefore
not an ignorable coordinate for this Rayleigh-Taylor mode, explaining its absence in the previous 2.5D work.

A common feature of all flux emergence simulations using a fully ionized MHD model is that the rapid expansion of the emerging flux, once it reaches the low-density corona, caused adiabatic cooling of the plasma to low temperatures. This can be seen for the FIP model in Figure 5 which compares the temperature as a function of height along a line at \( x = 10 \) and \( y = 0 \) where the \( x \)-coordinate is offset just enough to avoid the Rayleigh-Taylor instability–induced density sheet. With the FIP model, the temperature in the corona drops to 0.04 in normalized units, or \( \approx 250 \) K. Including the Newton cooling/heating term with a relaxation time specified in equation (16) only affects the temperature on the timescale of these simulations in a layer \( \approx 7 \) Mm above the photosphere and accounts for the ledge in temperature profile for both the Layer and PIP simulations between \( z = 0 \) and 7. In the corona, only the PIP simulations have a heating term due to the Cowling resistivity, and this maintains the temperature there to about 0.7 times the photospheric value.

4. CONCLUSIONS

Previous 2.5D simulations (Leake & Arber 2006) had shown that including the partially ionized layers of the solar atmosphere changed the dynamics of magnetic flux emergence by introducing a Cowling resistivity. This dissipated perpendicular current density and led to a force-free coronal field (inside the expanding region of emerged flux but ignoring the perpendicular current density at the interface of emerged flux and field-free corona). This simple picture cannot be supported in such a clear way for the full 3D simulations presented here. In 3D the most pronounced difference between partially ionized and fully ionized simulations is that the fully ionized simulations raise more chromospheric material. For the initial conditions used in this paper, this meant that the fully ionized model became unstable to the Rayleigh-Taylor instability, while the partially ionized model remained symmetric with no signs of instability. The presence of the Rayleigh-Taylor–induced perpendicular current density means that it is not possible to assess the effect of the Cowling resistivity on the amount of perpendicular current density emerging into the corona in isolation.

The primary result of these simulations is therefore that the inclusion of the Cowling resistivity affects the amount of chromospheric material uplifted into the corona. The process of removing \( j \perp \) allows most of the field to move through the chromospheric plasma rather than lifting it, thus reducing the amount of mass uplifted. Recent publications (Isobe et al. 2005) have suggested that the onset of the Rayleigh-Taylor instability during flux emergence may be the cause of coronal loops. This result may depend critically on the neutral hydrogen, but this was absent from all previous flux emergence simulations.

Often in coronal physics one is only concerned with the emergence of magnetic flux and its structure. In such situations, this paper has shown that a greatly simplified model of the Cowling resistivity is capable of reproducing the results of the full PIP simulations, except for the temperature profile. Since the present, and indeed all previous, flux emergence simulations have omitted a full treatment of thermal conduction, radiation effects, and coronal heating, it is unlikely that such simulations can achieve accurate temperature estimates for emergence. A practical conclusion from this work is therefore that the minimum physics required to obtain credible field structures in flux emergence is the Layer model presented in equation (26). This has the advantage of being easy to include into any code and yet accurately predicts the correct magnetic field structure and uplifting of chromospheric material.

This work was funded in part by the Particle Physics and Astronomy Research Council. The computational work was supported by resources made available through the UK MHD Consortium and Warwick University’s Centre for Scientific Computing.