Magnetohydrodynamic modeling of prominence formation within a helmet streamer

J. A. Linker, R. Lionello, and Z. Mikić
Science Applications International Corporation, San Diego, California, USA

T. Amari
Centre de Physique Théorique, École Polytechnique, Palaiseau, France

Abstract. We present a 2.5-D axisymmetric MHD model to self-consistently describe the formation of a stable prominence that supports cool, dense material in the lower corona. The upper chromosphere and transition region are included in the calculation. Reducing the magnetic flux along the neutral line of a sheared coronal arcade forms a magnetic field configuration with a flux rope topology. The prominence forms when dense chromospheric material is brought up and condenses in the corona. The prominence sits at the base of a helmet streamer structure. The dense material is supported against gravity in the dips of the magnetic field lines in the flux rope. Further reduction in magnetic flux leads to an eruption of the prominence, ejecting material into the solar wind.

1. Introduction

Prominences (called filaments when observed on the solar disk) support cool, dense chromospheric material ($\sim 10^4$ K and $10^{10} - 10^{11}$ cm$^{-3}$) against the solar gravity in the surrounding hot, tenuous corona ($\sim 10^6$ K and $10^7 - 10^9$ cm$^{-3}$). They are observed to lie above magnetic neutral lines in the photosphere and near the base of helmet streamers (regions of closed magnetic field that have confined the coronal plasma). The magnetic field in the prominence often exhibits “inverse polarity,” meaning that when the coronal magnetic fields embedded in the prominence cross over the neutral line they point in the direction opposite to that indicated by the photospheric magnetic field polarity. The prominence magnetic field is itself nearly aligned with the filament channel [Martin et al., 1994; Martin and Echols, 1994], indicating a highly sheared (and therefore magnetically energized) configuration.

Prominences have been studied for many years, yet the means by which these structures form and are maintained is still not understood, nor is their violent eruption. Three main difficulties confront any prospective theory attempting to describe the formation and evolution of prominences: (1) finding a magnetic configuration with “dips” (concave upward portions of flux tubes) that can gravitationally support the dense material; (2) understanding the mechanism by which chromospheric material is trapped in the dipped field lines and maintained there to form a condensation; and (3) elucidating the process that leads to the release of magnetic energy and the disruption of these structures. Because of the great complexity of the entire problem, these three individual aspects have typically been approached separately.

Models of magnetic field configurations for prominence support that develop the required dips and inverse polarity usually compute force-free magnetic fields and assume that the prominence material provides only a small perturbation to the magnetic structure [e.g., van Ballegooijen and Martens, 1989, 1990; Antiochos et al., 1994; Aulanier and Demoulin, 1998; Amari et al., 1999]. Given a favorable structure for supporting the filament, computation of the complicated dynamics and thermodynamics of condensations has been performed by assuming that the plasma flows along fixed magnetic flux tubes, which reduces the problem to one-dimensional hydrodynamics with energy transport [Poland and Mariska, 1986; Mok et al., 1990; Antiochos and Klimchuk, 1991; Antiochos et al., 1999a].

Condensation by radiative instability in sheared magnetic fields has also been studied in two-dimensional geometry [Sparks et al., 1990; Choe and Lee, 1992].

Models of prominence eruption typically start from configurations favorable for prominence support [e.g., van Ballegooijen and Martens, 1989; Priest and Forbes, 1990; Isenberg et al., 1993] and are closely related to the problem of coronal mass ejection (CME) initiation [Forbes and Priest, 1985; Linker and Mikić, 1995; Low, 1997; Mikić and Linker, 1997; Wu and Guo, 1997; Antiochos et al., 1999b; Lin and Forbes, 2000], as these phenomena are linked observationally [Handhausen, 1997]
and require the release of stored magnetic field energy and the opening of previously closed magnetic field regions [e.g. Aly, 1984; Sturrock, 1991; Forbes, 1992; Mikić and Linker, 1994; Antiochos, 1998]. In recent simulations [Mikić et al., 1999; Amari et al., 1999, 2000] we have found that magnetic flux emergence and cancelation in the photosphere can lead to the formation of magnetic flux ropes in sheared or twisted arcade configurations. When the flux cancelation reaches a critical threshold, the entire configuration erupts with a considerable release of magnetic energy.

The theoretical investigations described above couple the salient processes and focus on modeling individual aspects of the problem. This approach is useful for revealing the basic underlying physics. However, a complete picture of prominence formation, evolution, and eruption ultimately requires a comprehensive model of all of the processes together. This is particularly true now that different models can more or less equally describe the basic features of the observations, albeit somewhat superficially. Eventually, models will need to produce more detailed predictions that can be tested directly by observations (for example, by producing simulated emission that can be compared with images from spacecraft such as SOHO or TRACE). This challenging goal requires that the complex thermodynamic processes of the upper chromosphere and transition region be incorporated into multidimensional magnetohydrodynamic (MHD) computations.

The purpose of this paper is to demonstrate that we can now begin to study the problem of prominence formation and eruption with a more comprehensive approach. We show that when energy transport processes are included into calculations similar to our recent models of prominence support and eruption, chromospheric material can be trapped on helical field lines and lifted into a stable configuration in the lower corona as a result of flux cancelation at the neutral line. When further flux cancelation occurs, the entire configuration erupts into the corona.

The plan for the rest of the paper is as follows: In section 2 we describe our computational methodology. Section 3 describes results for both polytropic and full thermodynamic simulations. Section 4 discusses the implications of our work and directions for future progress.

2. Methodology

Lionello et al. [1999a; 2001] describe the use of our time-dependent three-dimensional (3-D) MHD model to compute helmet streamer configurations in the solar corona that include the lower solar atmosphere (upper chromosphere and transition region). Those calculations form the basis for the studies discussed in this paper; here we note some important details relevant to the present calculations.

In this paper we confine ourselves to azimuthally symmetric, two-dimensional solutions (with three components of magnetic field and velocity). We solve the following set of equations in spherical coordinates:

\[ \nabla \times \mathbf{B} = \frac{4\pi}{c} \mathbf{J}, \]

\[ \frac{1}{c} \frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E}, \]

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

\[ \frac{1}{\gamma - 1} \left( \frac{\partial T}{\partial t} + \mathbf{v} \cdot \nabla T \right) = -T \mathbf{v} \cdot \nabla \mathbf{v} - \frac{m}{k\rho} (\nabla \cdot \mathbf{q} + n_e n_p Q(T) - H_{ch} - H_n - H_v), \]

\[ \rho \left( \frac{\partial \mathbf{v}}{\partial t} + \mathbf{v} \cdot \nabla \mathbf{v} \right) = \frac{1}{c} \mathbf{J} \times \mathbf{B} - \nabla p + \rho g + \nabla \cdot (\rho \nu \nabla \mathbf{v}), \]

where \( \mathbf{B} \) is the magnetic field; \( \mathbf{J} \) is the electric current density; \( \mathbf{E} \) is the electric field; \( \rho, \mathbf{v}, p, T \) are the plasma mass density, velocity, pressure, and temperature, respectively; \( \mathbf{q} = -g_{0t} R_s^2 / r^2 \) is the gravitational acceleration (with \( R_s \) the solar radius); \( \eta \) is the resistivity; and \( \nu \) is the kinematic viscosity. In the energy equation, equation (5), \( Q(T) \) is the (optically thin) radiation loss function [as given by Athay 1986]; \( n_e \) and \( n_p \) are the electron and proton number density (which are equal for a hydrogen plasma); \( \gamma = 5/3 \) is the polytropic index; \( H_{ch} \) is the coronal heating term (a parameterized function); \( H_n = \eta J^2 \) and \( H_v = \nu \rho \nabla \cdot \nabla v \) are the ohmic and viscous heating terms (neglected in these simulations); and \( q \) is the heat flux.

Coronal heating is specified as a nonuniform profile that varies exponentially with radial distance, and both the heat deposition length scale (\( \lambda \)) and flux vary with latitude. The length scale varies from \( \lambda = 0.7 R_S \) at the poles to \( \lambda = 0.1 R_S \) at the equator, and the heat flux at \( r = R_S \) is \( 10^5 \text{ ergs cm}^{-2} \text{s}^{-1} \) at the poles and \( 5 \times 10^5 \text{ ergs cm}^{-2} \text{s}^{-1} \) at the equator. These parameters have been chosen to yield the strongly concentrated heating and higher densities seen in active regions near the equator of the simulation while also providing the more distributed heating present in coronal holes in the open field portions of the simulation. Lionello et al. [2001] describe further details of the coronal heating parameters and the use of both collisional (Spitzer’s law) and collisionless [Hollweg, 1978] thermal conduction as a function of radial distance from the Sun. For the results shown in this paper, we have modified the Spitzer law to reduce the steepness of the temperature and density gradients in the lowest part of the transition region (see Appendix A for a more detailed discussion). We have found that this procedure allows coarser meshes to be used and yields solutions that are qualitatively similar to those obtained with the full Spitzer thermal conductivity.
The calculation described here has been performed on a $201 \times 201$ nonuniform $(r, \theta)$ grid, with the mesh points highly concentrated in the equatorial region near the lower boundary. A mesh with $\Delta r \approx 6 \times 10^{-4} R_s$ in the upper chromosphere (near the lower radial boundary) and $\Delta \theta \approx 0.2^\circ$ near the equator is used. Other calculations with $121 \times 101$ mesh points, with coarser mesh resolution, were also performed prior to this calculation in order to scope out the parameters and find a regime in which a prominence was formed. The results with the higher-resolution mesh allowed us to use a larger Lundquist number, and we report only these results here. This simulation used 20 hours of CPU time on a Cray T90 supercomputer, and we did not deem it necessary to pursue calculations with higher mesh resolution considering the limitations on our computing resources. We believe that increasing the resolution beyond that used in the present calculation would not change the results significantly. However, when we begin modeling more realistic prominences in the future we will verify that a further increase in the spatial resolution does not change the results.

The simulation domain extends out to $30 R_s$. At this upper boundary the flow is supersonic and super-Alfvénic, and we implement boundary conditions that utilize the characteristics to allow only outgoing waves there [Linke and Mikić, 1997]. A uniform resistivity $\eta$ has been used, corresponding to a resistive diffusion time $\tau_R = 4\pi R_s^2/(\eta c^2) = 4 \times 10^3$ hours (for a length scale of $R_s$). Underneath the relaxed helmed streamer, the Alfvén speed ($V_A$) near the equator at $10,000 - 20,000$ km altitude is about $1800$ km s$^{-1}$, so the Alfvén travel time ($\tau_A = R_s/V_A$) is $6.3$ min and the Lundquist number $\tau_R/\tau_A \approx 3.8 \times 10^5$. A uniform viscosity $\nu$ is also used, corresponding to a viscous diffusion time $\tau_v = R_s^2/\nu$ such that $\tau_v/\tau_A = 3.8 \times 10^5$. Typically, we find that much higher values of $\tau_R$ can be specified for our algorithm if $\tau_v$ is kept at a relatively smaller value.

The method of solution of (1)-(6) has been described previously [Mikić and Linke, 1994; Linke and Mikić, 1997; Lionello et al., 1998; Linker et al., 1999; Mikić et al., 1999]. For the parameter regime of the calculations we describe here, stiffness of the equations is introduced by the combination of both the high Alfvén speed and thermal conduction in the transition region and lower corona, and the use of small mesh cells in this region to capture the steep gradients in density and temperature there. The stiffness introduced in the time integration of the equations by the Alfvén speed is treated efficiently using a semi-implicit method [Mikić and Linke, 1994]. The accuracy of the semi-implicit method has been studied previously [Schnack et al., 1987; Mikić et al., 1988], and it is known that when the time step falls below the Courant-Friedrichs-Lewy (CFL) limit, the semi-implicit method is identical to an explicit method. The stiffness introduced by the parabolic equations resulting from the thermal conduction, resistivity, and viscosity are treated using standard fully implicit methods.

3. Results

Amari et al. [1999, 2000] have described how magnetic flux changes at the photosphere, when introduced into a three-dimensional sheared arcade field, can yield solutions with stable magnetic flux ropes suitable for prominence support and how further flux changes can lead to magnetic energy release and eruption. These calculations were performed in the "zero beta" limit of equations (1)-(6): equations (1)-(3) and (6) are solved with $P = 0$ and a fixed profile assumed for $\rho$ (equilibria found in this manner are force-free solutions). Mikić and Linke [1999] have shown that when the same procedure is applied to polytropic MHD solutions with helmet streamer configurations (solutions of equations (1)-(6) with the energy source terms in (5) set to zero and $\gamma = 1.05$), the streamer disrupts and material is ejected out into the solar wind. Figure 1 shows the eruption of a flux rope in a three-dimensional calculation of this kind. A detailed description of these results is presently in preparation. Here we describe how the inclusion of processes in the lower solar atmosphere in these calculations (full solutions of equations (1)-(6)) leads to the formation of a prominence-like structure and its subsequent eruption.

3.1. Helmet Streamer Solution

In the first phase of the calculation, we generate a helmet streamer equilibrium [e.g., Linke and Mikić, 1995]. We start with a potential magnetic field in the corona that matches a specified distribution of radial magnetic field at the solar surface $B_0$. The distribution we choose is the sum of a weak dipole field ($B_{rodip} = 2.2$ G at the poles), and a stronger ($B_{rodip} = 9.5$ G) concentrated dipole near the neutral line. For simplicity we choose a configuration that is symmetric in latitude about the solar equator. Our intent in including the strong dipole is to model the large-scale effect of the field in an active region (where prominences frequently form). The equatorial position of the bipole at the equator is not very realistic for prominences, but is convenient for the illustrative purposes of this calculation. We place our lower boundary at the top of the chromosphere, and we impose a fixed temperature $T_0 = 20,000$ K and a plasma number density $n = n_p = n_e = 10^{11}$ cm$^{-3}$. We impose a spherically symmetric solar wind solution which includes the upper chromosphere and the transition region and is consistent with the chosen boundary values for temperature and density. This combination is not initially at equilibrium. We integrate the time-dependent MHD equations in time until the solution settles down to an equilibrium (for $380\tau_A$, where $\tau_A$ is the Alfvén time described in section 2). The final state models the coronal plasma and
Plate 1. The evolution of the plasma density and magnetic field in a thermodynamic MHD model (see text). The plasma density is depicted in color, and projections of the magnetic field lines are overlain on the density. (Corresponding plots of the temperature can be seen in Plate 2.) A close-up of the lower part of the simulation domain near the equator is shown. Plate 1a shows the helmet streamer configuration at the end of the shearing phase. Plate 1b shows the streamer after the magnetic flux has been reduced (reduction of 3.75% of the initial interior bipolar). A low-lying filament structure has formed and is just discernible near the lower boundary (see Plate 3 for a close-up view of this region). In Plate 1c, the filament is at a height of 140,000 km and is moving upward slowly. The enhanced density can be seen to lie near the bottom of the detached flux surfaces. Flux reduction (at 11.25%) is beyond the critical threshold for eruption (see text). Plate 1d shows the eruption proceeding, and dense material is carried upward into the corona, shown in Plates 1e and 1f.
Plate 2. The same as Plate 1, except that the evolution of the plasma temperature and magnetic field are depicted. The high density material shown in Plate 1 is also very cold.
Plate 3. A close-up view of the prominence-like structure when flux reduction is halted at a level of 3.75% ($t = 608 \tau_A$, Plates 1b and 2b) and the calculation is continued (see text), allowing a stable flux rope to form. The plasma density (Plate 3a) and temperature (Plate 3b) are shown together with projections of the magnetic field lines as in Plates 1 and 2. Tracings of the magnetic field lines, colored by temperature (Plate 3c) and density (Plate 3d) are shown. The helical field lines support cold ($T \sim 2 \times 10^4$ K) and dense ($n \sim 10^{10}$ cm$^{-3}$) material against gravity.
Figure 1. Simulated eruption of a magnetic flux rope on the Sun. Figure 1a shows a three-dimensional flux rope formed by reducing the magnetic flux in a sheared arcade, computed with the polytropic model (see text). The flux rope lies within a helmet streamer that is surrounded by open magnetic field lines along which the solar wind streams outward. When the amount of magnetic flux decrease is small enough, the flux rope is stable. Further reduction of the photospheric magnetic field leads to the eruption of the flux rope, as shown in Figures 1b-1d.

includes the transition region in the calculation. This procedure has been described by Lionello et al. [2001]. The final solution has a coronal streamer with closed field lines, surrounded by open field lines along which the solar wind flows outward. Figure 1 of Linker and Mikić [1995] shows an example of the magnetic configuration of a 2-D helmet streamer using a polytropic energy equation; Figure 1 and Figure 3 of Lionello et al. [2001] show examples of a helmet streamer configuration computed with thermodynamics.

3.2. An Energized Helmet Streamer

In the second phase of the simulation, we apply a shear flow near the neutral line that builds free magnetic energy into the streamer. This shear flow is not intended to model actual flows on the Sun. It is just a convenient mechanism for producing strongly sheared field lines that are nearly aligned with the neutral line, a frequently observed characteristic of filaments [Martin and Echols, 1994]. We use a shear profile that is like the one used by Mikić and Linker [1994], with a width $\Delta \theta_m = 8.5^\circ$. The shear is applied from $t = 380 \tau_A$ to $t = 570 \tau_A$, with a maximum velocity $v_0 = 0.005 V_{A0}$. Plates 1a and 2a show contours of the flux function (projections of the magnetic field lines) superimposed on the plasma density (Plate 1a) and plasma temperature (Plate 2a), at the end of the shearing phase. The maximum displacement of the foot points from their
original position is 0.8 \( R_s \), and the sheared field has a magnetic energy equal to 2.3\( W_{\text{pot}} \). The calculation is then continued in a relaxation phase to \( t = 589\tau_A \).

### 3.3. Creation of a Magnetic Flux Rope and Subsequent Eruption

In the final phase of the calculation, we change the magnetic flux at the boundary to generate a flux rope. On the Sun it is frequently observed that the magnetic fields in an active region tend to disperse days to weeks after its emergence. During this time, filaments are frequently observed to form along the neutral line. At times, these filaments disappear, presumably due to eruption, and may even reform in the same location later. This dispersal of magnetic flux is thought to occur on a small spatial scale by annihilation and submergence of magnetic dipole elements and has been modeled as a convective-diffusive process on a large scale [Wang et al., 1989; Wang and Sheeley, 1990]. The disappearance of photospheric magnetic flux has long been suspected of playing a role in filament formation and eruption [van Ballegooijen and Martens, 1989]. We model the effect of such changes by reducing the magnetic flux in the bipolar flux region. The electric fields that specify this change (described below) imply converging flows at the neutral line, as is believed to occur in flux cancelation. We find that the reduction in magnetic flux can create a filament and that further reduction in flux can make it erupt, as illustrated below.

The boundary conditions for flux reduction are essentially the same as those described by Amari et al. [2000]. The change in flux is applied by specifying the appropriate electric fields at the boundary. For example, when we seek steady state solutions of (1)-(6), we set the tangential component of the electric field at the boundary, \( \mathbf{E}_{\text{B}} \), to zero. This keeps \( B_{\text{r}} \) (the radial magnetic field at the solar boundary) fixed in time. In order to specify a desired change in the magnetic flux, we specify a nonzero \( \mathbf{E}_{\text{B}} \) that is consistent with the required \( \partial B_{\text{r}} / \partial t \). Reduction of \( B_{\text{r}} \) is equivalent to cancelation of flux at the neutral line and leads to the formation of a flux rope [Amari et al., 2000].

To compute the electric field required to drive a specific flux change, we note that, in general, \( \mathbf{E}_{\text{B}} \) can be expressed as \( \nabla \times \mathbf{\psi} + \nabla \phi \), where \( \mathbf{\psi} \) and \( \phi \) are arbitrary functions of \( \theta \) and \( \phi \) and \( \nabla \phi \) indicates transverse derivatives (in the \( \theta - \phi \) plane at \( r = R_s \)). The potential \( \phi \) changes \( \mathbf{E}_{\text{B}} \) without changing \( B_{\phi} \) and can be used to control the transverse magnetic field (i.e., the shear and the normal electric current), whereas the potential \( \psi \) changes the flux. For the simulation presented in this paper we used \( \phi = 0 \), which minimizes changes to \( B_{\phi} \). Then \( c
abla \phi \cdot \mathbf{\psi} = \partial B_{\text{r}} / \partial t \), which can be solved for \( \psi \) for the flux change specified by \( \partial B_{\text{r}} / \partial t \).

For the case presented here, we apply the appropriate electric field \( \mathbf{E}_{\text{B}} \) at the lower boundary, as calculated above, to reduce the dipole portion of the flux distribution \( B_{\text{odi}} \) by 15% (while leaving the dipole portion \( B_{\text{odi}} \) undisturbed), from \( t = 589\tau_A \) to 665\( \tau_A \). Plate 1 shows the resulting evolution of the plasma density and projected magnetic field lines (for a portion of the simulation); Plate 2 shows the corresponding evolution of the plasma temperature. At \( t = 608\tau_A \) when the flux has been reduced by 3.75% (Plates 1b and 2b), a flux rope containing cold dense chromospheric material (\( T = 2 \times 10^4 \) K and \( n > 1 \times 10^{10} \) cm\(^{-3} \)) has been lifted into the corona to form a prominence-like structure (barely visible near the lower boundary at the equator). In a detailed analysis of vector magnetograph and Hα observations of an emerging \( \delta \)-sunspot group, Lites et al. [1995] suggested that a twisted magnetic loop emerged from below the photosphere, carrying with it dense material that formed a filament. The filament-formation mechanism we describe is consistent with these observations. This is in contrast to the idea of siphon flows leading to the development of a condensation [e.g., Antiochos et al., 1999a], which cannot occur here because the flux rope field lines do not connect to the surface.

At \( t = 646\tau_A \) shown in Plates 1c and 2c (flux reduction at 11.25%) the prominence has already begun to erupt. It has been lifted to a height of 0.2 \( R_s \) (140,000 km) and is slowly moving upward. By \( t = 684\tau_A \) (19\( \tau_A \) after flux reduction was halted) the flux rope and helmet streamer are erupting (Plates 1d and 2d), and the entire configuration is opened and material is carried upward (Plates 1e, 1f, 2c, and 2f).

There is a threshold of flux reduction for the eruption to occur. To demonstrate this behavior, we performed another case where flux reduction is halted at \( t = 608\tau_A \) (= 3.75% of the initial amount). When the calculation is continued from that point, no eruption occurs, and the flux rope containing chromospheric material relaxes to a stable state. Plate 3 shows a closeup of this configuration after it has relaxed for 19\( \tau_A \). On the left are shown projected field lines overlain on the plasma density (Plate 3a) and temperature (Plate 3b). Note the presence of high density and low temperature on the detached flux surfaces. On the right, field lines from the calculation are plotted with the color along the field line indicating the density (Plate 3c) and temperature (Plate 3d). Cool and dense material in the prominence is supported against gravity in the dips of the field lines of the flux rope. The height of the simulated prominence is about 15,000 km. As the calculation is continued further, the flux rope slowly diffuses because of the presence of finite resistivity.

If flux reduction is halted at a level of 7.5% of the initial amount and the calculation is continued, the flux rope and helmet streamer erupt in a manner similar to the case shown in Plates 1-2. We have found that the exact level of flux reduction necessary for eruption in a given configuration depends on the details of the surface
Figure 2. The prominence height as function of time for the eruptive phase of the calculation (Plates 1b-1f and 2b-2f).

flux distribution and the amount of shear introduced. Prior to reaching the critical level, the behavior of the system is quasi-static.

Figure 2 shows the prominence height as a function of time during the eruptive phase. The eruption that was produced in this case was not very fast (reaching ≈40 km s\(^{-1}\) at 2 \(R_\odot\)), although it is still accelerating at this stage. This is in contrast to some of the polytropic simulations we have run with this mechanism (speeds >500 km s\(^{-1}\)). A more concentrated shear and a larger initial magnetic field strength (enabling release of more magnetic energy) may be required to create a fast mass ejection containing a filament.

4. Discussion

The possibility that prominences are supported by magnetic flux ropes has been considered in recent years by a number of authors [e.g., van Ballegooijen and Martens, 1989, 1990; Rust and Kumar, 1994; Chen, 1996; Aulanier and Demoulin, 1998]. Our recent work [Amari et al., 1999, 2000] and the results we have presented here are similar in many respects to the force-free calculations of van Ballegooijen and Martens [1989], which showed how flux cancelation could lead to the formation of flux ropes. The computations we have described here show that when the full MHD equations are used, including an energy equation that takes into account energy transport in the upper chromosphere and transition region, chromospheric material can be lifted by the helical field lines and supported against gravity inside a helmet streamer. The flux rope can subsequently be ejected into the solar wind as a consequence of eruption.

The idealized calculations presented here are intended to illustrate that it is possible to perform self-consistent MHD calculations of prominences. Clearly, a 2-D axi-symmetric model has limited applicability to solar observations. In this geometry the helical field lines of the flux rope form a torus around the Sun. The connection of the magnetic fields embedded in the prominence to the photosphere, and the possible role of siphon flows in filling/draining the prominence, is an important question not addressed by this calculation. Preliminary one-dimensional solutions for energy flow along the magnetic field loops from a three-dimensional flux rope configuration indicate that condensations do indeed form, but a fully self-consistent 3-D simulation is required to investigate this question more completely. We have performed 3-D simulations of flux rope formation in polytropic helmet streamers (Figure 1), and we intend to extend the present results to 3-D as well.

Once the simulations become more realistic, it will be necessary to explain the many detailed features that have been observed in prominences [e.g., Martin et al., 1994]. For example, what are the implications of the observed relationship between the filament axial field and the skew of coronal arcades [Martin and McAllis-
ter, 1996], and the relationship between differential solar rotation and the axial field in polar crown filaments [van Ballegooijen et al., 1998]? Do these considerations imply that the axial field originates from below the photosphere [Priest et al., 1996]? Magnetic flux ropes are not the only candidates for explaining prominence support [Martín and Echols, 1994; Antiochos et al., 1994] and eruption [Antiochos et al., 1999b]. Understanding the pre-eruptive state of filaments, as well as their violent eruption, requires calculations that can predict observable quantities. We consider the calculations we have presented here as a first step of that process.

Appendix A: Modified Thermal Conduction

In the lower transition region, there is a close balance between thermally conducted heat ($\nabla \cdot \mathbf{q}$ in equation (5)) coming from the hot, tenuous corona and the energy radiated away ($n_e n_p Q(T)$ in equation (5)) by the ambient dense cold material. In this collisional regime, the Spitzer form of the thermal conductivity is appropriate:

$$\mathbf{q} = -\kappa || \mathbf{b} \cdot \nabla T,$$

where $\kappa || = \kappa_0 T^5/2$ ($\kappa_0 = 9 \times 10^{-7}$ in cgs units), $\mathbf{b}$ is the unit vector along the magnetic field, and $T$ is the temperature in K.

The steepness of the temperature in the transition region arises from the balance of conductive heat flux from the corona and radiative losses in the transition region. The temperature dependence of the Spitzer thermal conductivity ($\kappa || \sim T^{5/2}$) forms the steepest temperature gradient in the lower transition region. Lionello et al. [2001] used highly resolved meshes ($\Delta r \approx 8 \times 10^{-5} R_\odot$, or 56 km) to capture this gradient. It seems that an artificial broadening of this gradient should be possible without qualitatively changing the overall solution. With this goal in mind, we have formulated $\kappa ||$ in such a way as to preserve the essential physics of the Spitzer model for the upper transition region and corona, while reducing the steepness of the gradient in the lower part of the transition region:

$$\kappa || = \kappa_0 \left[ s T^{5/2} + (1-s) T^{\alpha} T_{\text{mod}}^{5/2-\alpha} \right],$$

where

$$s(T) = \frac{1}{2} \left( 1 + \tanh \left( \frac{T - T_{\text{mod}}}{\Delta T_{\text{mod}}} \right) \right).$$

The function $s$ varies smoothly between 0 and 1 and regulates the transition from the Spitzer regime of $\kappa ||$ (when $T > T_{\text{mod}}$) and the modified regime (when $T \leq T_{\text{mod}}$). In the modified regime, which applies to the lower transition region, $\kappa || \propto T^\alpha$, and the $\alpha$ power is $0 \leq \alpha \leq 5/2$. The transition between the two regimes occurs in a temperature interval depending on $\Delta T_{\text{mod}}$. For the simulation results presented here, we chose $\alpha = 0$ (corresponding to constant $\kappa ||$ for $T \leq T_{\text{mod}}$), $T_{\text{mod}} = 250,000$ K, and $\Delta T_{\text{mod}} = 20,000$ K.

This formulation makes $\kappa ||$ higher than the Spitzer value in the lower transition region, allowing thermal conduction to balance radiation loss with a smaller temperature gradient. We have found that this modified form of the thermal conduction yields solutions that are qualitatively similar to those obtained with the full Spitzer thermal conductivity, and we were able to increase the radial grid size in the lower transition region by almost an order of magnitude over that used by Lionello et al. [2001]. Quantitatively, the temperature of solutions on individual loops matches closely in the corona, although differences in the density can be as large as a factor of 2 different. We regard the modified thermal conduction as a useful tool for exploring the qualitative properties of solutions without using extremely fine meshes.

Acknowledgments. These results were presented at the International Conference on Solar Eruptive Events, March 6-8, 2000, at the Catholic University of America. This work was supported by the NASA Sun Earth Connections Theory Program, the NASA Supporting Research and Technology program, the NSF Space Weather program (jointly funded by NSF and AFSOR), and Boston University’s Integrated Space Weather Modeling project (funded by NSF). Computations were at the San Diego Supercomputer Center (SDSC) and the National Energy Research Supercomputer Center (NERSC).

Janet G. Luhrmann thanks Piet Martens and another referee for their assistance in evaluating this paper.

References


