THE DYNAMIC FORMATION OF PROMINENCE CONDENSATIONS

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ABSTRACT

We present simulations of a model for the formation of a prominence condensation in a coronal loop. The key idea behind the model is that the spatial localization of loop heating near the chromosphere leads to a catastrophic cooling in the corona. Using a new adaptive grid code, we simulate the complete growth of a condensation and find that after ~5000 s it reaches a quasi-steady state. We show that the size and growth time of the condensation are in good agreement with data and discuss the implications of the model for coronal heating and for observations of prominences and the surrounding corona.

Subject headings: Sun: chromosphere — Sun: corona — Sun: prominences

1. INTRODUCTION

Solar prominences and filaments are beautiful phenomena both in their appearance and in their physics. They are observed in Hα as large masses (~10ⁱ⁶ g) of dense (>10¹¹ cm⁻³) chromospheric plasma that seem to float high up in the tenuous (~10⁶ cm⁻³) solar atmosphere, the million-degree corona. They can have a myriad of shapes but often they resemble a long, arched wall and have widths ranging from 10⁸–10⁹ cm, heights ranging from 10¹⁰ cm down to the chromosphere itself, and lengths that can be of the order of the Sun’s circumference for high-latitude polar crown filaments. Typical lifetimes of prominences can range from as short as hours for active region filaments to as long as months for quiescent prominences. Prominences usually disappear by eruption. Large coronal mass eruptions and two-ribbon flares are almost always accompanied by prominence eruption. Consequently, the structure of prominences, especially their magnetic structure, is vital for understanding the major solar drivers of geomagnetic disturbances.

There are three central issues concerning the physics of prominences: how does mass condense in the corona, how is it supported, and why does it erupt? The coronal magnetic field undoubtedly plays a vital role in all three issues (e.g., Tandberg-Hanssen 1995; Priest 1989). We believe, however, that these questions are essentially decoupled in that while the eruption mechanism depends on the global topology of the field (Antiochos 1998; Antiochos, DeVore, & Klimchuk 1999), and the support depends on the local geometry of a sheared three-dimensional arcade (Antiochos, Dahlburg, & Klimchuk 1994), the condensation formation is due to the properties of the coronal heating process and the cooling of coronal plasma by radiation and conduction (Antiochos & Klimchuk 1991; Dahlburg, Antiochos, & Klimchuk 1998). In this paper we focus on the third question of condensation formation and, in particular, address the dynamics of this process.

When viewed with high spatial resolution, it can be seen that the mass of a prominence is not one continuous structure but is actually a collection of small condensations or knots that form on timescales of tens of minutes and have size scales of the order of 1000 km (Engvold 1976; Zirker 1989). Given typical coronal gas pressures ~10⁻¹⁴ ergs cm⁻³ and magnetic field strengths of 10 G, the plasma is low beta, and we expect that each condensation is isolated from the others by the magnetic field. Therefore, the canonical picture of a prominence is that of an arcade of flux tubes, or coronal loops, each of which has somewhere along its length a plug of cool material that is observed as an Hα knot, and the ensemble of all the knots forms the global structure that we call a prominence.

An important point is that prominence knots cannot be considered simply as the condensation of coronal plasma, because they generally contain more mass than the coronal portion of the flux tubes they occupy. Material must be brought up from the chromosphere in order to produce the large mass of a prominence, which has led many authors to consider the so-called siphon flow model for condensation formation (e.g., Pikel’ner 1971; Engvold & Jensen 1977; Priest & Smith 1979; Ribes & Unno 1980; Poland & Mariska 1986; An, Bao, & Wu 1988; Mok et al. 1990). It has proved to be difficult, however, for such models to produce condensations of sufficient mass (Poland & Mariska 1986). The observation that the knot contains a large mass implies a large amount of chromospheric evaporation (e.g., Antiochos & Sturrock 1978) and, hence, a large coronal heating rate in the loop. But a high heating rate implies high corona-like temperatures rather than the cool chromospheric temperatures needed for a prominence.

We have described elsewhere how one can reconcile the seemingly contradictory requirements of high coronal heating and low coronal temperatures (Antiochos & Klimchuk 1991; Dahlburg et al. 1998). The key idea is that if the heating in a coronal loop has a strong spatial dependence, then it is possible to have a high heating rate and still develop a massive condensation near the loop center.
that if the heating is uniform per unit volume or uniform per unit mass, which are the usual assumptions in coronal loop models (e.g., Rosner, Tucker, & Vaiana 1978; Vesecky, Antiochos, & Underwood 1979), then the loop can only have a hot structure with coronal temperatures throughout. This is true even if the loop has a dipped geometry suitable for supporting prominence material (Antiochos & Klimchuk 1991). If the heating is sufficiently localized near the loop base, however, the plasma structure can change to one in which there are chromospheric temperatures in the dipped portion of the loop.

In our previous paper (Antiochos & Klimchuk 1991), we verified this hypothesis with numerical simulation. We began the simulation with a coronal loop that had a dipped geometry and spatially uniform heating. This loop developed a typical hot coronal equilibrium solution with a temperature maximum greater than $10^6$ K at the loop midpoint and a thin transition region at the base. We then increased the heating, but only in a small region near the base. At first, this heating increase caused the temperature to rise throughout the loop, so that the central dipped portion became even hotter. Chromospheric evaporation was driven by the temperature rise, which caused the density and, consequently, the radiative losses to increase throughout the loop, even in the upper section where the heating was unchanged. Eventually, the increase in radiative losses produced a catastrophic cooling at the loop center and the formation of a chromospheric condensation.

It is worth noting that this temperature collapse is physically due to a loss of equilibrium rather than to a thermal instability. For small increases of the heating near the base, it is possible to find static equilibrium solutions that have a hot corona throughout the loop. These solutions are characterized by a slight increase in the temperature at the loop midpoint as the heating increases. But if the base heating is increased sufficiently, then no static solution with a hot midpoint is possible, because that would imply too high a radiative loss rate there. The only static solutions are those in which the temperature in the middle section of the loop is very low, $T \ll 10^5$ K, so that the radiative losses there are reduced (Dahlburg et al. 1998).

Although our previous simulation verified the basic idea of the model, there was a problem with that calculation that prevented us from comparing the model with observations, especially the timescale for prominence formation. The code could not follow the condensation significantly into its growth phase. As the condensation formed, a transition region with very large temperature gradients formed on each side of the condensation, exactly like the transition region at the footpoints. This transition region moved down the loop as the condensation grew, but since our code used a fixed grid, it was not possible to resolve a moving transition region and, consequently, the simulation broke down almost immediately after the condensation first appeared (Antiochos & Klimchuk 1991). In this paper we present simulations using a new code that has a fully adaptive grid (MacNeice et al. 1998) and, therefore, can follow a condensation throughout its evolution, allowing us to compare the growth times with data. These new calculations also include more realistic forms for the loop geometry and the heating.

2. THE NUMERICAL MODEL

Since the magnetic field in a prominence dominates the plasma, we make the usual assumption of a one-dimensional coronal loop model in which the field determines the loop geometry and all plasma quantities vary only along the field. Our code solves the standard set of transport equations for a one-dimensional plasma:

$$\frac{\partial}{\partial t} \rho + \frac{\partial}{\partial s} (\rho V) = 0 ,$$

$$\frac{\partial}{\partial t} \rho V + \frac{\partial}{\partial s} (\rho V^2 + P) = - \rho g_s (s) ,$$

and

$$\frac{\partial}{\partial t} U + \frac{\partial}{\partial s} \left( UV - 10^{-6} T^{5/2} \frac{\partial}{\partial s} T \right) = - P \frac{\partial}{\partial s} V + E(t, s) - n^2 \Lambda(T) .$$

In equations (1)–(3), $s$ is distance along the loop, i.e., along the magnetic flux tube; $\rho$ is the plasma mass density; $V$ is the velocity along the loop; $P$ is gas pressure; $g_s$ is the component of gravity along the field; $U = 3P/2$ is the internal energy; $10^{-6} T^{5/2}$ is the Spitzer (1962) coefficient of thermal conduction; $E(t, s)$ is the assumed form for the coronal heating rate; $n$ is the electron number density; and $\Lambda(T)$ is the radiative loss coefficient for optically thin emission (e.g., Cook et al. 1989). Note that we assume a fully ionized hydrogen plasma, so that $\rho = 1.67 \times 10^{-24} n$ and the equation of state is simply $P = 2nkT$.

The code that we have developed (P. J. MacNeice & D. S. Spicer 1998, private communication) to solve these equations has two major improvements over our previous code. First, in order to advance the solution, it uses a second-order Godunov scheme with a MUSCL limiter (Van Leer 1979) applied to the characteristic variables. It is widely believed that higher order Godunov schemes, which solve the nonlinear Riemann problem at each cell face, are the most robust methods for numerically solving one-dimensional fluid flow problems like the one that we expect for a coronal loop plasma—a nonturbulent compressible flow in which steep gradients and/or discontinuities are likely to develop (Hirsch 1988; Colella 1990).

Second, the code employs a fully adaptive grid by using the parallel adaptive mesh (PARAMESH) refinement package (MacNeice et al. 1998). PARAMESH builds a hierarchy of subgrids to cover the computational domain, with spatial resolution varying to satisfy the demands of the application. These subgrid blocks form the nodes of a tree data structure. Each grid block has a logically Cartesian mesh, and the index ranges are the same for every block. In this case we use one-dimensional subgrid blocks with 20 grid cells and two guard cells at each end of the subgrid. When a subgrid block needs to be refined, it spawns two child blocks. These two children have their own 20 cell subgrids and together cover the same coordinate range as their parent, but now with mesh spacing one-half that of the parent’s mesh spacing. Any or all of these children can themselves be refined in the same manner. If consistent with resolution requirements, then PARAMESH also will “derefine” two adjacent blocks, i.e., merge them into a single block with twice the mesh spacing. PARAMESH provides routines that manage the necessary communication between subgrid blocks.

An adaptive grid is essential for our simulation because we expect a thin transition region to form somewhere in the
loop and to move with time. Unless this region is resolved, the plasma energetics and dynamics cannot be calculated accurately. Often a tricky issue is the selection of a criterion for refining and derefining the numerical grid. We find that density is the best quantity to use in determining the refinement because it exhibits a strong spatial variation, both in the transition region, where pressure is approximately constant, and in the chromosphere, where temperature is approximately constant. In the simulations presented here we use six levels of refinement, which corresponds to a minimum grid spacing of \(5 \text{ km}\). This does an adequate job of resolving the transition region throughout the evolution.

For these calculations we established a very simple refinement/derefinement criterion. At each grid cell \(i\) we computed the error measure

\[
\epsilon_i = \left| \rho_i - \rho_{i-1} \right| / \min (\rho_i, \rho_{i-1}), \tag{4}
\]

and for each subgrid found the maximum value

\[
\epsilon_{\text{max}} = \max (\epsilon_i, \epsilon_1, \ldots, \epsilon_{2^l}). \tag{5}
\]

Notice that we include one guard cell at each end of the subgrid (i.e., \(i = 0, l = 21\)) to enable refinement in advance of an arriving feature of the solution. If \(\epsilon_{\text{max}}\) exceeds 0.25, meaning that somewhere in the subgrid block the density varies by at least 25\% between grid cells, then that block is refined. If \(\epsilon_{\text{max}}\) is less than 0.05 on a pair of sibling subgrid blocks, then the block is derefined, provided that the lower resolution representation of its parent does not exceed the refinement threshold.

The refinement process was permitted to range freely between refinement levels 6 and 12. In this context, level 1 indicates that the entire computational domain is covered by one subgrid block. A level 1 subgrid would have a cell size of \(\Delta x = L/20\), where \(L\) is the total length of the computational domain. A level \(n\) subgrid therefore has a cell size of \(2^{-n} L/20\). The initial grid was uniformly refined at level 6, i.e., \(\Delta x = 2^{-5} L/20 = 0.0015625 L\). The subgrids refined to level 12 have a cell size that is a factor of \(2^6 = 64\) finer.

The geometry that we assume for the loop is shown in Figure 1. For simplicity, we took the loop to be symmetric about its midpoint. In a one-dimensional model, the loop (i.e., magnetic field) geometry influences the plasma evolution solely through its effect on gravity. The only significant restriction on the magnetic geometry is that the flux tube have a dipped section, as shown in Figure 1, which can support material against gravity. The geometry of Figure 1 would be appropriate for any of the usual prominence magnetic field models, such as the Kippenhahn-Schluter (1957) or Kuperus-Radu (1974) models, or even our recent three-dimensional model (Antiochos et al. 1994).

The loop is comprised of three distinct sections. First there is a straight vertical tube of height \(s_1\); this is roughly the chromospheric section. The exact position of the top of the chromosphere changes during the calculation. This section of the loop \((s \leq s_1)\) is given by

\[
z(s) = s, \quad \text{hence} \quad g_\parallel (s) = g_\odot, \quad \tag{6}
\]

where \(z(s)\) is height at position \(s\), and \(g_\odot = 2.7 \times 10^4 \text{ cm s}^{-2}\) is the solar gravity. Next is a quarter-circle of arc length \(s_2 - s_1\), and, therefore, radius \(R = (s_2 - s_1)/\pi\). Note that the maximum height of the loop is \(s_1 + R\). This section \((s_1 < s \leq s_2)\) is given by

\[
z(s) = s_1 + R \sin \left( \frac{s - s_1}{R} \right),
\]

hence

\[
g_\parallel (s) = g_\odot \cos \left( \frac{\pi}{2} - \frac{s - s_1}{s_2 - s_1} \right). \tag{7}
\]

The final section is a quasi-cosine curve that has a maximum height at \(s_2\) to match the quarter-circle, a minimum height at the loop midpoint, \(s = L\), and a total depth \(D\). This part \((s_2 < s \leq L)\) is given by

\[
z(s) = s_1 + R - \frac{D}{2} \left[ 1 - \cos \left( \frac{\pi}{2} - \frac{s - s_2}{L - s_2} \right) \right],
\]

hence

\[
g_\parallel (s) = -g_\odot \frac{\pi D}{2(L - s_2)} \sin \left( \frac{\pi}{2} - \frac{s - s_2}{L - s_2} \right). \tag{8}
\]

From the form above, we note that the gravity term in the momentum equation is negative for \(s \leq s_2\) and positive for \(s > s_2\). The gravity term is continuous but does not have continuous derivatives.

The geometry parameters that we must specify (Fig. 1) are the position of the top of the chromosphere \(s_1 = 1 \times 10^9 \text{ cm}\) and the position of the loop apex \(s_2 = s_1 + 2.5 \times 10^6 \text{ cm}\) (therefore the maximum loop height \(s_1 + R = s_1 + 5 \times 10^9 / \pi \text{ cm}\)), the total loop half-length, \(L = 11 \times 10^9 \text{ cm}\), and the depth of the coronal dip, \(D = 5 \times 10^8 \text{ cm}\). These values are much more appropriate for quiescent prominence loops than the values used in our previous work, which are suitable only for low-lying active region prominences (Dahlburg et al. 1998).

For the radiative losses we take a simple analytic form. Since the radiative losses depend sensitively on elemental abundances that vary substantially on the Sun (Cook et al. 1989), there is no longer a “standard” curve to use. The important point is that our model does not depend on the details of the radiative loss curve; it is sensitive only to the spatial dependence of the heating. In the upper transition region and corona \((T \geq 10^5 \text{ K})\), we take \(\Lambda(T) = 1.0 \times 10^{-17}/T\). In the lower transition region \((10^5 > T \geq 30,000 \text{ K})\), we take \(\Lambda(T) = 1.0 \times 10^{-37}T^3\). A temperature at the base of the model of 30,000 K is assumed. We would like the chromospheric region of the loop to remain at this temperature as much as possible; consequently, we take the radiative losses to vanish below 29,500 K. For 30,000 > \(T > 29,500 \text{ K}\), \(\Lambda(T) = 2.7 \times 10^{-24} [(T - 29,500)/500]\), and \(\Lambda(T) = 0.0\) for \(T \leq 29,500 \text{ K}\). In addition, we drop the density dependence of the radiative losses in the chromospheric region, where the density can become extremely large owing to the exponential increase per scale height. In the Sun’s chromosphere the radiative losses are limited by radiative transfer, which
is not included in our model. As with $g_j(s)$, the assumed form for $A(T)$ is continuous but without continuous derivatives.

Finally, the form of the heating needs to be specified. We set $E = E_0 + E_s(s, t)$, where $E_0 = 1.5 \times 10^{-5}$ ergs cm$^{-3}$ s$^{-1}$ is the uniform background heating, which stays on throughout the simulation. The spatially dependent heating $E_s$ is ramped up over 1000 s only after the loop has settled into a static equilibrium with the uniform heating. We set $E_s$ to be constant in the chromosphere, $E_s = 10^{-3}$ ergs cm$^{-3}$ s$^{-1}$ for $s \leq s_1$, and we have an exponential drop-off in the corona,

$$E_s = 10^{-3} \exp \left[ -\left( s - s_1 \right) / \lambda \right] \quad \text{for} \quad s > s_1,$$

where the damping length $\lambda$ is chosen to be 10,000 km. This form for the spatial dependence of the heating is much more physical than that used previously, in which we assumed that the heating was concentrated in a very narrow region centered at the highest point in the loop (Antiochos & Klimchuk 1991). Here there is no special relation between the heating and the loop geometry. We are simply assuming that the heating propagates into the loop from below and is deposited over some spatial scale small compared to the loop length.

3. RESULTS

In order to solve equations (1)–(3), appropriate boundary and initial conditions must be specified. As boundary conditions on the loop we assumed a rigid, constant temperature base at $s = 0$ and symmetry at the top, $s = L$. To derive the initial equilibrium, we simply started with a discontinuous temperature profile in which $T = 10^6$ K in the corona $s \geq s_1$ and $T = 30,000$ K in the chromosphere. The density was constant in the corona, $n = 0.7 \times 10^9$ cm$^{-3}$ for $s \geq s_1$, and increased exponentially with depth in the chromosphere at the appropriate scale height. Hence, the plasma was in approximate force balance to start with, but not in thermal equilibrium. We then let the system evolve until an equilibrium was reached.

After $\sim 35,000$ s the loop had settled into a static equilibrium with negligible residual motions, $V_{\text{max}} < 0.1$ km s$^{-1}$.

Fig. 2.—Temperature as a function of position along the loop for four times in the evolution. The temperature is measured in units of $10^6$ K and the distance in units of 10,000 km. The solid curve corresponds to $t = 0$, the time at which the localized heating is turned on. The dashed curve corresponds to $t = 10,000$ s, the dotted curve to $t = 30,000$ s, and the dash-dotted curve to $t = 50,000$ s.

We took this state to be the initial equilibrium, $t \equiv 0$ s. Figures 2, 3, and 4 show the temperature, density, and velocity profiles, respectively, at this time. Even though the loop has a dip geometry, the temperature and density profiles are exactly what one would expect for a standard coronal loop (Rosner, Tucker, & Vaiana 1978; Vesecky et al. 1979). The maximum temperature occurs at the loop midpoint and has a value of $1.2 \times 10^6$ K. Note also that the top of the chromosphere has moved only slightly upward from its starting position at $s_1 = 10,000$ km.

We then turned on the spatially dependent heating, $E_s$, ramping it up linearly over 1000 s. The plasma responded as expected by first heating and evaporating chromospheric material into the loop before a condensation appeared at the midpoint. This evaporative phase of the evolution lasted a considerable time, approximately 60,000 s. Figures 2, 3, and 4 show the plasma profiles at three times: 10,000, 30,000, and 50,000 s after the spatial heating turn-on. The temperature at the loop midpoint reached a maximum value greater than $1.7 \times 10^6$ K at around 10,000 s, but then started to decline as the radiation losses at the midpoint began to exceed the heating there. The position of the top of the chromosphere moved downward over 1000 km as a result of the temperature increase and accompanying pressure increase, but then it moved back to nearly its original position as the temperature dropped.

Fig. 3.—As in Fig. 2, but for number density in units of $10^6$ cm$^{-3}$. Note that the chromospheric and transition region portions of the loop are not shown, so that the coronal density structure can be seen clearly.

Fig. 4.—As in Fig. 2, but for velocity in units of 10 km s$^{-1}$.
We note from Figure 3 that the density throughout the loop increased steadily during the evaporative phase except near the very end \( (t > 50,000 \text{ s}) \), when the condensation started to form at the midpoint. Owing to the temperature collapse there, the pressure decreased rapidly at the midpoint, which tended to evacuate the loop leg. This effect can be seen in the velocity profiles. At early times \( (\sim 10,000 \text{ s}) \) the velocity is driven by evaporation, i.e., an overpressure at the top of the chromosphere, and consequently the velocity peaks \( (\sim 5 \text{ km s}^{-1}) \) near the loop base (Fig. 3). The evaporative velocity decreased, however, as the loop temperature decreased. Later in the evolution the velocity started to increase again, but now peaking nearer the midpoint, which indicates that it was driven by the underpressure in the condensation.

Once the condensation appeared at the loop midpoint, the subsequent evolution was dominated by its rapid growth. The temperature at the loop midpoint first dropped below the chromospheric value of 30,000 K at \( t = 54,000 \text{ s} \). Figures 5 and 6 show the temperature and density profiles over the last 10,000 km of our half-loop at this time and at four later times in the evolution, 500, 1000, 10,000, and 30,000 s after the condensation appearance. The corresponding velocity profiles over the whole loop are shown in Figure 7.

There were two distinct phases to the condensation evolution: a transient-motion phase characterized by large upward and downward flows, and then a quasi-steady state characterized by slow upflows throughout the loop. At first, the condensation grew very rapidly as the middle section of the loop dropped down to the base temperature of 30,000 K. This temperature collapse and the resulting pressure decrease were so fast that a weak rebound shock wave was launched from the loop midpoint toward the base (Fig. 7). It can be seen from the figure that the wave speed was 160 km s\(^{-1}\), which agreed well with the sound speed in the corona. Weak shocks, like ordinary acoustic waves, are expected to travel at the sound speed (e.g., Zeldovich & Raizer 1968). The maximum upward velocity of the coronal plasma was less than 60 km s\(^{-1}\), however, so the bulk motion was clearly subsonic and no strong shocks developed. The wave hit the chromosphere and then bounced back up the loop with a reduced amplitude. As shown in Figure 7, the loop oscillated back and forth for several bounces with an oscillation period \( \sim 1000 \text{ s} \), but after \( \sim 5000 \text{ s} \) the loop settled into a quasi-steady state with a maximum upward velocity \( V < 5 \text{ km s}^{-1} \) and with a coronal gas pressure \( P \sim 10^{-1} \text{ ergs cm}^{-3} \). This flow implied a substantial enthalpy flux \( 5PV/2 = 10^5 \text{ ergs cm}^{-3} \text{ s}^{-1} \), which was sufficient to balance the excess radiative losses in the dipped portion of the loop, so that a steady state became possible.

We let the simulation run for over 30,000 s after the condensation first appeared in order to verify that a true steady state was achieved. There was no significant change in the velocity profile after the first \( \sim 5000 \text{ s} \). The evolution consisted of a very slow growth of the condensation. In principle, the condensation should expand until it encompasses most of the dipped portion of the loop in Figure 1, but this would require a timescale much longer than typical prominence lifetimes. The key point is that the total mass \( M \) in the condensation must increase as \( \exp \left( z/H_g \right) \), where \( z \) is the vertical depth of the condensation and \( H_g \) is the gravitational scale height \( \sim 200 \text{ km} \) for typical prominence temperatures. Hence, for a fixed velocity, the condensation's expansion rate slows down exponentially rapidly with time,
so we do not expect condensations to become significantly larger than their scale heights.

The evolution calculated by our simulation agrees very well with prominence observations. From the time that the midpoint temperature first dropped below 30,000 K, the condensation required 500 s to reach a half-width of 1000 km and approximately 1000 s to reach a half-width of 2000 km, as shown in Figures 5 and 6. These numbers are typical of the timescales quoted for the formation time of prominence condensations (Engvold 1976; Zirker 1989). Furthermore, the condensation growth slowed down abruptly after it reached a half-width of order 1000 km; the half-width was only 5000 km after 30,000 s.

4. DISCUSSION

There are two important geometrical features that we have not included in the calculation above. First is the effect of asymmetry, either in the loop geometry or in the heating rates at the two base regions. In contrast to the simulation above, condensation will not first appear at the bottom of the dip in a realistic flux tube that has no symmetry; it appears at some arbitrary point along the dipped portion of the loop and then accelerates toward the bottom. We noted this fact in our previous work, in which we simulated a loop with unequal heating at the two legs (Antiochos & Klimchuk 1991), but could not at that time calculate the growth or motion of the condensation. Furthermore, we expect from our three-dimensional models for prominence magnetic fields (Antiochos et al. 1994; Antiochos 1996) that the loops will have long, stretched-out, concave-up sections with quite shallow dips, even shallower than in Figure 1. If so, a condensation could move a considerable distance along this dipped section before settling down to its equilibrium position. We propose that this is the origin of the disconnected condensations along the body of the prominence (Zirker, Engvold, & Martin 1997; S. F. Martin 1998, private communication). Simulations of condensation formation in asymmetric loops will clearly be very informative.

The second geometry feature missing from the present simulation is a variable cross-sectional area for the loop. The area factor is known to have a significant effect on the flow of mass and heat through a coronal loop (e.g., Antiochos & Sturrock 1976; Mariska & Boris 1983; Dowdy, Moore, & Emslie 1987). An interesting point is that the area factor of the dipped prominence loops should have a qualitatively different structure than that for all coronal flux tubes studied to date. Since the area of a flux tube is inversely proportional to field strength, the area of a usual flux tube such as a dipolar field increases monotonically with height or distance along the field and reaches a maximum at the apex. In the dipped flux tubes of our prominence magnetic models, however, the field strength has a local maximum near the dip; this is what causes the dips to form in the first place (Antiochos et al. 1994; Antiochos 1996). Consequently, the area reaches a local minimum near the bottom of the dip, implying that flow into the condensation will be accelerated owing to a nozzle effect. Further simulations are needed to determine whether the flows can become sufficiently large to have observable consequences.

The results of the simulation described in this paper have several interesting implications for coronal heating theories and for observations. The most important feature of our model is that the heating must be localized near the chromospheric footpoints of a loop; however, this localization is not severe, since we used a spatial scale of 10,000 km for the heating. On the other hand, our results do imply that any heating theory in which the energy input is predicted to occur uniformly, or to be localized near the loop top, would be incompatible with the observation of coronal condensations. Of course, even if the heating were uniform or localized near the top, transient condensations could still form if the heating underwent strong temporal variations, as in the case of postflare Hz loops.

Another implication of the model is that a prominence loop should be cooler and less dense (except in the condensation itself) than a standard, uniformly heated coronal loop with the same total heating. The reason is that in the standard coronal loop model the total energy input must be balanced by the total radiation emitted (e.g., Rosner et al. 1978; Vsecky et al. 1979), but in our prominence solution, some of the energy input goes into driving the mass flow, which lifts material against gravity and powers the extra radiation losses at the transition region of the condensation. The situation is physically similar to comparing a closed coronal loop in a quiet region with an open flux tube in a coronal hole. Even if the heating in each region has the same magnitude, the quiet-region static loop will have a substantially higher temperature and density than the coronal hole steady-flow flux tube and, hence, will appear much brighter in X-rays.

To verify this claim we ran a simulation in which the localized heating $E_1$ was spread out uniformly over the whole loop length, so that we would obtain a solution with no condensation. The resulting temperature and density profiles looked very similar to the first curves in Figures 2 and 3, except that the temperature maximum at the midpoint had a value over $2.6 \times 10^6$ K and the density was over $2.6 \times 10^9$ cm$^{-3}$. These numbers are to be compared to the coronal temperatures and densities of the quasi-steady state of the condensing loop above, $1.0 \times 10^6$ K and $1.5 \times 10^9$ cm$^{-3}$, respectively. Even though the total heating in both cases was identical, the "coronal loop" (uniform heating case) had a coronal density almost 2 times larger than the "prominence loop" (localized heating case) and, therefore, would have an X-ray brightness 4 times larger.

We conjecture that this effect may play a role in the origin of the dark cavities that are often seen to surround prominences. We expect the magnetic field in a prominence region to consist of long sheared flux tubes, with dips in which prominence condensations form, and overlying unsheared flux tubes that are relatively short (Antiochos et al. 1994; Antiochos & DeVore 1999). Because of this difference in length, the overlying tubes develop a usual hot static-loop plasma structure, while the inner sheared tubes develop condensations and, therefore, appear darker in X-rays. In this picture the region of dipped magnetic flux tubes threading a prominence can be thought of as forming a local coronal hole region in which plasma is undergoing a quasi-steady flow into the prominence mass. Even though the low density of the cavity would make such flows difficult to detect, it should be informative to search for them with SOHO and TRACE.

Our model predicts other flows that may be detectable by SOHO as well. The key point is that if coronal heating has the type of spatial variation assumed here, then condensations should form in sufficiently high loops even if they do
not have dips. If the loop does have a dip, then a quasi-steady state is possible in which the condensation settles to the bottom of the dip, but if the loop has no dip, we believe that no steady state is possible, and the loop must be constantly dynamic. We propose that this is the origin of the continuous dynamics observed in very high active region loops with the CDS instrument on SOHO (Brekke 1998; Kjeldseth-Moe & Brekke 1998). Further simulations with nondipped loop geometries are clearly needed.

The main conclusion of this paper is that our model successfully reproduces the observed timescales for condensation formation. Not only does the condensation in the simulation reach a size of a few thousand km in tens of minutes, but equally important, the condensation stops growing significantly once it reaches this size. Both results are necessary in order to agree with observations of prominence knots. We believe, therefore, that loss of thermal equilibrium driven by a spatially localized coronal heating is the answer to the question: How does mass condense in the corona?

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