Impact of Dynamic State on the Mass Condensation Rate of Solar Prominences

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Abstract

The interiors of quiescent prominences are filled with turbulent flows. The evolution of upflow plumes, descending pillars, and vortex motions has been clearly detected in high-resolution observations. The Rayleigh–Taylor instability is thought to be a driver of such internal flows. Descending pillars are related to the mass budgets of prominences. There is a hypothesis of dynamic equilibrium where the mass drainage via descending pillars and the mass supply via radiative condensation are balanced to maintain the prominence mass; however, the background physics connecting the two different processes is poorly understood. In this study, we reproduced the dynamic interior of a prominence via radiative condensation and the mechanism similar to the Rayleigh–Taylor instability using a three-dimensional magnetohydrodynamic simulation including optically thin radiative cooling and nonlinear anisotropic thermal conduction. The process to prominence formation in the simulation follows the reconnection–condensation model, where topological change in the magnetic field caused by reconnection leads to radiative condensation. Reconnection is driven by converging motion at the footpoints of the coronal arcade fields. In contrast to the previous model, by randomly changing the speed of the footpoint motion along a polarity inversion line, the dynamic interior of prominence is successfully reproduced. We find that the mass condensation rate of the prominence is enhanced in the case with dynamic state. Our results support the observational hypothesis that the condensation rate is balanced with the mass drainage rate and suggest that a self-induced mass maintenance mechanism exists.

Key words: instabilities – Sun: corona – Sun: filaments, prominences

Supporting material: animation

1. Introduction

Solar prominences or filaments are cool dense plasma clouds suspended in coronal magnetic fields. In limb observations, quiescent prominences appear to be composed of fine vertical threads (Engvold 1976, 1981; Berger et al. 2008, 2010; Chae et al. 2008; Chae 2010). Both upflows and downflows are present in prominences (Zirker et al. 1994, 1998) and show turbulent characteristics (Freed et al. 2016; Hillier et al. 2017). Rising dark plumes and descending pillars associated with internal flows have been clearly detected via high-resolution observations of the Solar Optical Telescope on the Hinode satellite (Berger et al. 2008, 2010, 2011, 2017; Chae et al. 2008; Chae 2010).

The downflows and upflows are related to the mass budget of the prominences. Prominences drain their mass via descending pillars (Hillier et al. 2012b; Liu et al. 2012) and obtain mass from upflow plumes (Berger et al. 2010), chromospheric jets (Chae 2003), flux emergence (Okamoto et al. 2008, 2009, 2010), and radiative condensation in the corona (Berger et al. 2012; Liu et al. 2012). Though it is difficult to estimate the mass supply or drainage rate of each process in most prominences, the condensation rate and the mass drainage rate have been reported to be comparable and temporally correlated with each other in one in situ prominence formation event (Liu et al. 2012). According to estimates, the timescale of mass cycle (the total mass divided by the mass drainage rate) is less than one hour, which is much shorter than the typical lifetime of prominences (a few hours to months; Mackay et al. 2010). This result suggests a dynamic equilibrium where the prominence mass is maintained by the complete balance between continuous and substantial mass supply and drainage. To reveal the mass maintenance of prominences, the physical interactions between the mass drainage and the condensation, as well as the detailed mechanisms of each process, need to be investigated.

In most previous studies, the mass drainage and the condensation have been separately investigated. For example, the Rayleigh–Taylor instability is thought to be a mechanism driving upflow plumes and descending pillars (Berger et al. 2010; Ryutova et al. 2010). Previous theoretical studies of the Rayleigh–Taylor instability using ideal magnetohydrodynamic (MHD) simulations succeeded in reproducing upflows and downflows with speeds consistent with observational values (Hillier et al. 2011, 2012a, 2012b; Keppens et al. 2015; Xia & Keppens 2016a); however, these simulations neglected the mass supply to the prominence via radiative condensation due to their use of the adiabatic assumption.

Radiative condensation has been investigated as a mass supply mechanism for prominences. Several models with different drivers have been proposed. One is the evaporation–condensation model in which chromospheric evaporation injects mass into the coronal loops, leading to radiative condensation. Early studies using one-dimensional hydrodynamic simulations, where the magnetic fields were assumed to be fixed, were not able to discuss interactions with the Rayleigh–Taylor instability (Antiochos & Klimchuk 1991; Antiochos et al. 1999; Karpen et al. 2001, 2003, 2005, 2006; Karpen & Antiochos 2008; Xia et al. 2011; Luna et al. 2012). Recently, this model was investigated using multidimensional MHD simulations including radiative cooling and thermal conduction. In two-dimensional simulations, prominences were reproduced in the form of a stationary slab without a dynamic interior (Xia et al. 2012; Keppens & Xia 2014). Using a three-dimensional simulation, Xia & Keppens (2016b) found that a
dynamic fragmented prominence is reproduced with this model. Moreover, in their simulations, the total mass of the simulated prominence is maintained via the balance between the mass input from chromospheric evaporation and the mass drainage from prominences, showing a solution achieving dynamic equilibrium. The mass drainage in their simulation was driven by overdense clusters of fragmented condensations that were not sustained by the weak coronal magnetic fields. The role of the Rayleigh–Taylor instability in this simulation was not clear. In addition, even though parameterized heating is an important factor to determine the mass flux of the evaporated flow in this model, neither observations nor theories have yet to guarantee its presence. It is still unclear whether the dynamic equilibrium is self-induced or contingent on the parameter settings.

Another condensation model, the reconnection–condensation model, has also been proposed in previous theoretical studies (Neuman 1983; Kaneko & Yokoyama 2015, 2017). This model was recently demonstrated using three-dimensional MHD simulations including thermal conduction and optically thin radiative cooling in a study by Kaneko & Yokoyama (2017). In their simulations, a flux rope is formed in the sheared arcade fields by reconnection via converging motion at the footpoints of the arcade fields. Radiative cooling is enhanced due to dense coronal plasmas trapped inside the flux rope; however, thermal conduction along the long coronal loops after reconnection cannot compensate for the radiative losses, leading to radiative condensation. The simulated prominence ended up in a stationary slab without the dynamic features seen in observations. The convergence of magnetic patches of opposite polarities has been detected below prominences in observations of the photospheric magnetic field (Rondi et al. 2007; Schmieder et al. 2014; Yang et al. 2016). The issue remaining for this model, therefore, is to reproduce the internal dynamics associated with the mass budget of prominences.

In this study, we show that random footpoint motion leads to the formation of prominences with dynamic fine structures by radiative condensation and the mechanism similar to the Rayleigh–Taylor instability in the framework of the reconnection–condensation model. Moreover, we show that the dynamic state has an impact on the growth of radiative condensation by comparing simulations with and without the dynamic state. The numerical settings are described in Section 2. The results of the simulations are shown in Section 3. We discuss the results in Section 4, and the conclusions are given in Section 5.

2. Numerical Settings

The simulation domain is a rectangular box whose Cartesian coordinates \((x, y, z)\) extend to \(-12 \text{Mm} < x < 12 \text{Mm}, 0 < y < 50 \text{Mm}, \) and \(0 < z < 24 \text{Mm},\) where the \(y\)-direction corresponds to the height and the \(xz\)-plane is parallel to the horizontal plane. The initial corona is under hydrostatic stratification with a uniform temperature \((T_{\text{cor}} = 1 \text{MK})\) and uniform gravity \((g_{\text{cor}} = 270 \text{ m s}^{-2})\). The initial density profile is given as

\[
n = n_{\text{cor}} \exp \left[ \frac{-y}{L_z} \right] \tag{1}
\]

where \(n\) is number density, \(n_{\text{cor}} = 10^{9} \text{ cm}^{-3}\) is the number density at the coronal bottom, and \(L_z = k_B T_{\text{cor}}/(mg_{\text{cor}}) = 30 \text{ Mm}\) is the coronal scale height, where \(k_B\) is the Boltzmann constant and \(m\) is the mean molecular mass. The mean molecular mass of a prominence depends on the helium abundance and the ionization degree. An accurate treatment of the ionization degree requires non-LTE (local thermodynamic equilibrium) modeling. For simplicity, we set \(m = m_p\), where \(m_p\) is the proton mass.

The initial magnetic field is a linear force-free arcade given as

\[
B_x = -\left( \frac{2L_a}{\pi a} \right) B_a \cos \left( \frac{\pi x}{2L_a} \right) \exp \left[ -\frac{y}{a} \right], \tag{2}
\]
\[
B_y = B_a \sin \left( \frac{\pi x}{2L_a} \right) \exp \left[ -\frac{y}{a} \right], \tag{3}
\]
\[
B_z = -\sqrt{1 - \left( \frac{2L_a}{\pi a} \right)^2} B_a \cos \left( \frac{\pi x}{2L_a} \right) \times \exp \left[ -\frac{y}{a} \right]. \tag{4}
\]

where \(B_a = 3 \text{ G}, L_a = 12 \text{ Mm}, \) and \(a = 30 \text{ Mm}\). The polarity inversion line (PIL) on the surface \(y = 0\) is located along \(x = 0\).

The three-dimensional MHD equations including nonlinear anisotropic thermal conduction and optically thin radiative cooling are as follows:

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \tag{5}
\]
\[
\frac{\partial (\rho \mathbf{v})}{\partial t} + \nabla \cdot \left( \rho \mathbf{v} \mathbf{v} + \rho I - \frac{BB}{4\pi} + \frac{B^2}{8\pi} \mathbf{I} \right) - \rho g = 0, \tag{6}
\]
\[
\frac{\partial}{\partial t} \left[ e_{\text{th}} + \frac{1}{2} \rho v^2 + \frac{B^2}{8\pi} \right] + \nabla \cdot \left[ \left( e_{\text{th}} + p + \frac{1}{2} \rho v^2 \right) \mathbf{v} + \frac{c}{4\pi} \mathbf{E} \times \mathbf{B} \right] \nonumber
\]
\[
= \rho g \cdot \mathbf{v} + \nabla \cdot \left( \kappa T^{5/2} \mathbf{b} \cdot \nabla T \right) - n^2 \Lambda(T) + H, \tag{7}
\]
\[
e_{\text{th}} = \frac{p}{\gamma - 1}, \tag{8}
\]
\[
T = \frac{m p}{k_B \rho}, \tag{9}
\]
\[
\frac{\partial \mathbf{B}}{\partial t} = -\mathbf{c} \nabla \times \mathbf{E}, \tag{10}
\]
\[
\mathbf{E} = -\frac{1}{c} \mathbf{v} \times \mathbf{B} + \frac{4\pi \eta}{c^2} \mathbf{J}, \tag{11}
\]
\[
\mathbf{J} = \frac{c}{4\pi} \nabla \times \mathbf{B}, \tag{12}
\]

where \(\kappa = 2 \times 10^{-6} \text{ erg cm}^{-1} \text{s}^{-1} \text{K}^{-7/2}\) is the coefficient of thermal conduction, \(\mathbf{b}\) is a unit vector along the magnetic field, \(\Lambda(T)\) is the radiative loss function of an optically thin plasma, \(H\) is the background heating rate per volume, and \(\eta\) is the magnetic diffusion rate. We use the same radiative loss function as that used in Kaneko & Yokoyama (2017). The radiative loss function under \(10^6 \text{ K}\) is assumed to have a dependence of \(T^3\). The background coronal heating is taken to be proportional to the magnetic energy density and is given as

\[
H = \alpha_H B^2, \tag{13}
\]
where $B$ is the magnetic field strength and $\alpha_H = 4.8 \times 10^{-7}$ s$^{-1}$ is a constant coefficient. The value of $\alpha_H$ is computed to satisfy the condition for thermal equilibrium for a uniform temperature:

$$n^2 \Lambda(T_{\text{eq}}) = H,$$

(14)

where $a = L_\alpha$ (see also Equations (13)–(15) in Kaneko & Yokoyama 2017 for the equilibrium condition).

To drive reconnection, the footpoint velocities perpendicular and parallel to the PIL are given in $y < 0$ as

$$v_x = -v_0(t) \sin \left( \frac{\pi x}{2L_\alpha} \right) f(z),$$

(15)

$$v_y = 0,$$

(16)

$$v_z = v_x,$$

(17)

where $t$ represents time, $v_0(t)$ is the speed depending on time, and $f(z)$ represents a random number with an amplitude of $0.5 \leq f(z) \leq 1.5$. The velocity component $v_x$ represents the converging motion used to drive reconnection. The component $v_z$ represents the anti-shearing motion necessary to trigger radiative condensation and to prevent the eruption of a flux rope (Kaneko & Yokoyama 2017). The speed $v_0(t)$ is given as

$$v_0(t) = v_{00}, \ (0 < t < t_1)$$

(18)

$$v_0(t) = v_{00} \frac{t_2 - t}{t_2 - t_1}, \ (t_1 < t \leq t_2)$$

(19)

$$v_0(t) = 0, \ (t \geq t_2),$$

(20)

where $v_{00} = 12$ km s$^{-1}$, $t_1 = 1200$ s, and $t_2 = 1440$ s. For three grids below $y = 0$, the magnetic fields are numerically calculated with the given velocities (Equations (15)–(20)) based on the induction equation and a free boundary condition at the bottom boundary. Via this manipulation, the given footpoint velocity is smoothly adopted onto the domain solving the full MHD equations. The gas pressure and density below $y = 0$ are assumed to be in hydrostatic equilibrium at a uniform temperature of 1 MK. The free boundary condition is applied to

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**Figure 1.** Time evolution of the emission measure along the $x$-axis for case R. The animation runs from $t = 0$ to 104 minutes and is annotated with the flow arrows also presented in Figure 2. (An animation of this figure is available.)

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all the variables at the top boundary. The anti-symmetric boundary condition is applied to $v_x$, $v_y$, $B_y$, and $B_z$, and the symmetric boundary condition is applied to the other variables at the boundaries in the $x$-direction. Periodic boundary conditions are applied to all variables at the boundaries in the $z$-direction under the assumption that a flux rope sustaining a prominence is sufficiently long.

For fast magnetic reconnection, we adopt the following form of the anomalous resistivity (e.g., Yokoyama & Shibata 1994):

$$\eta = 0, \ (J < J_c)$$

$$\eta = \eta_0 (J/J_c - 1)^2, \ (J \geq J_c),$$

where $\eta_0 = 3.6 \times 10^{13}$ cm$^2$ s$^{-1}$ and $J_c = 25$ erg$^{1/2}$ cm$^{-3/2}$ s$^{-1}$. We restrict $\eta \leq \eta_{\text{max}} = 1.8 \times 10^{14}$ cm$^2$ s$^{-1}$.

For comparison, we performed an additional simulation for footpoint motions without random components by fixing $f(z_b) = 1$. This simulation corresponds to a 2.5-dimensional simulation in the $xy$-plane where the translational symmetry is constrained in the $z$-direction (along the PIL). Hereafter, we call the case with random footpoint motion case R, and the case with footpoint motion without random components case N. For case N, the basic parameters and boundary conditions, except for the velocities at the footpoints, are the same as those for case R.

### 3. Results

Figure 1 shows snapshots of the emission measure, which is defined as

$$\text{EM}(y, z) = \int n(x, y, z)^2 \, dx$$

when observed with the line of sight in the $x$-direction, for case R. The process leading to radiative condensation is the same as that in previous studies (Kaneko & Yokoyama 2015, 2017). A flux rope is created by the reconnection of the sheared arcade fields via converging footpoint motion. Radiative cooling inside the flux rope is enhanced because the dense plasmas in the lower corona are trapped. Due to the topology change of the magnetic fields, the relaxation effect via thermal conduction becomes ineffective at compensating for radiative loss inside the flux rope, leading to radiative condensation. As shown in Figure 1(a), the density of the condensation is nonuniform in the $z$-direction because random components of the footpoint motion add fluctuation to the interior of the flux rope prior to condensation (note that the footpoint motion is stopped prior to condensation). As the mass of the prominence increases, multiple spikes grow (Figures 1(b) and (c)). Thin vertical threads are formed after the spikes reach the bottom boundary (Figure 1(d)).

Figure 2 shows the velocity fields inside the prominence, which is defined as

$$V_{y}^p(y, z) = \frac{1}{M_e} \int_{T < T_{\text{K}}} \rho_v \, dx,$$

$$V_{z}^p(y, z) = \frac{1}{M_e} \int_{T < T_{\text{K}}} \rho_z \, dx,$$

$$M_e = \int_{T < T_{\text{K}}} \rho \, dx.$$  \hfill (24)

$$M_e = \int_{T < T_{\text{K}}} \rho \, dx.$$  \hfill (25)

$$M_e = \int_{T < T_{\text{K}}} \rho \, dx. $$\hfill (26)

The strong downflows are concentrated in the descending spikes (Figure 2). The downward speed of the spikes is approximately 12 km s$^{-1}$, which is consistent with observational values of 10–15 km s$^{-1}$ (Berger et al. 2008; Chae 2010; Hillier et al. 2012b). The spikes are reflected at the bottom boundary and upflows or vortex motions are created (see around $(y, z) = (4 \text{ Mm}, 5 \text{ Mm})$ in Figure 2(b)). The spikes are squeezed via the interactions of the flows. Eventually, thin vertical threads form along the paths of the descending spikes. The widths of the threads in our simulation are approximately 1000 km, comparable to the observed width of 600 km (Chae 2010).
Figure 3 shows three-dimensional snapshots of the simulation result from different angles. The prominence is located along the dips of the flux rope. The magnetic fields maintain a coherent flux rope structure even though the local density and velocities evolve in a highly nonuniform manner. As shown in Figure 3(b), the vertical threads are not manifestations of the vertical magnetic fields. Figure 4 shows the density and temperature distribution in the $z = 12$ Mm cross section at the same time. A low-density and high-temperature cavity forms around the prominence due to the mass depletion after condensation.

Figure 5 shows the distribution of the vertical flows, location of prominence, and magnetic field during condensation. The strong downflows along the coronal magnetic field connected to the prominence are the condensation flows (panels (b), (c), and (d)). Both downflows and upflows are found along the interface between the prominence and the lower corona (panels (a) and (d)), which is indicative of ongoing Rayleigh–Taylor instability. The growth rates of both flows, which are represented by dashed lines 1 and 2 in Figure 5(a), were compared with the analytical linear growth rate of the Rayleigh–Taylor instability. First, the analytical growth rate was estimated. The linear growth rate of the magnetic Rayleigh–Taylor instability in a uniform magnetic field is given by

$$\sigma_i^2 = g\kappa k = \frac{k \cdot B_{\perp}}{2\pi (\rho_+ + \rho_-)}.$$  

where $k$ is the wave vector, and $\Lambda = (\rho_+ - \rho_-)/(\rho_+ + \rho_-)$ is the Atwood number, $\rho_+$ and $\rho_-$ are the densities of heavy and light fluids, respectively (e.g., Hillier 2018). As shown in Figure 5(b), the wave vectors of the perturbation are virtually perpendicular to the magnetic field. Hence, we focus on the growth rate of the interchange mode $(k \cdot B = 0)$. The wavelength of the perturbation is approximately $\lambda = 2\pi/k = 7$ Mm, as shown by the black arrow representing the interval between the upflows in Figure 5(a). Figure 6 shows the profile of the density and the Atwood number along the dashed lines in Figure 5(a), where the Atwood number in panel (b) is computed by substituting the density at $y = 8$ Mm with $\rho_-$ and the local density with $\rho_+$. The Atwood number of the prominence varies from 0.8 to 0.95. By substituting $A = 0.95$ and $\lambda = 2\pi/k = 7$ Mm for the first term of the right side of Equation (27), the analytical linear growth rate of the interchange mode $\sigma_i = 1.5 \times 10^{-2}$ s$^{-1}$. The growth rate of the flows was then measured in the simulation. Figure 7 depicts the time evolution of both the flows along the dashed lines in Figure 5(a). At $t = 72$ minutes, the upflow begins to grow, and the finite speed of the downflow at $t = 72$ minutes is due to the condensation flows. Figure 8(a) depicts the time evolution of the maximum speed of the downflow $V_{\text{down}}$ and that of the upflow $V_{\text{up}}$ in Figure 7. The growth rate $\sigma$ of both the flows can be estimated by

$$\sigma = \frac{1}{V} \frac{dV}{dt}.$$  

where $V = V_{\text{up}}$ is substituted for the upflow, and $V = V_{\text{down}} - V_c$ is substituted for the downflow, where $V_c$ is the speed of the condensation flow. We assume that $V_c = V_{\text{down}}(t = 72$ minutes) does not change in the short duration. Figure 8(b) shows the measured growth rates of both the flows, which are comparable with the analytical growth rate. This suggests that a mechanism similar to the Rayleigh–Taylor instability (RTI-like mechanism) facilitates the corrugation of the interface between the prominence and the lower corona. The situation reproduced in our simulation is not the rigorous form of the magnetic Rayleigh–Taylor instability discussed in the previous studies because the condensation flows coexist in the interface, indicating that the system is not in a state of complete thermal and mechanical equilibrium. As shown in Figure 5(d), the upflows are located between the condensation downflows (see around $(x, y) = (1$ Mm, $10$ Mm)), and are gradually canceled. Therefore, the
upflows do not evolve into the dark plumes observed in prominences (Berger et al. 2008, 2010). Note that the acceleration of the downflow in our simulation is estimated to be \(-10 \text{ m s}^{-2}\) from the inclination of \(V_{\text{down}}\) in Figure 8(a), which is significantly smaller than the gravitational acceleration \(g = -270 \text{ m s}^{-2}\).

To investigate the impact of the dynamic state on the condensation rate, we compare the cases with and without dynamic state. Figure 9 shows snapshots of the simulation results for case N in which the footpoint motion does not contain random components. Due to the constraints of the 2.5-dimensional assumption, all variables are uniform in the \(z\)-direction. The temporal evolution of the total mass and the mass growth rate of the prominences in cases R and N are shown in Figure 10. The mass growth rate of the prominence is computed as

\[
M_{\text{prom}} = \frac{d}{dt} \int_{T<10^5 \text{K}} \rho(x, y, z) dx dy dz. \tag{29}
\]

Until \(t = 50\) minutes, the mass growth rates are the same for these two cases. After \(t = 50\) minutes, the mass growth rate in case R becomes larger than that in case N. As shown in Figure 10(b), the mass growth rate in case R is enhanced around \(t = 80\) minutes just after the spikes begin to evolve (see also Figure 1(b)). The mass drainage rate is also enhanced with
Figure 5. Panels (a), (b), (c), and (d) show snapshots at $t = 76$ minutes in the $x = 0.7$ Mm, $y = 9.6$ Mm, $z = 6.4$ Mm, and $z = 10.3$ Mm planes, respectively. The colors represent the vertical velocity $v_y$. The locations of panels (c) and (d) correspond to those of dashed lines 1 and 2 in panel (a), respectively. The solid lines are density isocontour of $n = 6 \times 10^9$ cm$^{-3}$ indicating the location of prominence. The arrows in panel (b) represent the magnetic field vector $(B_y, B_z)$, and those in panels (c), and (d) represent $(B_x, B_y)$.

Figure 6. Profiles of density and the Atwood number along the dash lines in Figure 5(a). The solid and dashed lines represent the values of lines 1 and 2, respectively.
the evolution of the spikes. The dashed–dotted line in Figure 10(b) shows the temporal evolution of the mass drainage rate computed as

$$M_{\text{drain}} = \int_{T<10^5 \text{K}} \rho v_y(x, y_c, z) dx dz,$$

(30)

where $y_c = 5 \text{ Mm}$ is selected. The mass drainage rate is comparable to and temporally correlated with the condensation rate. Our simulation results support the observational findings of Liu et al. (2012) and give insights into the self-induced mass maintenance mechanism via the coupling of the radiative condensation and the Rayleigh–Taylor instability.

Figure 11 shows the differential emission measures (DEMs) during condensation for cases R and N defined as

$$\text{DEM}(T) = n_e^2 \frac{dx}{dT},$$

(31)

where the line of site is in the $x$-direction. The DEMs are averaged over $0 < y < 14 \text{ Mm}$ and $0 < z < 24 \text{ Mm}$. We compared the DEM in case R at $t = 88.0 \text{ minutes}$ to that in case N at $t = 97.6 \text{ minutes}$ because the total mass of the plasmas cooler than 0.8 MK are the same for those times. In case R, the DEM in the lower temperature region ($0.1 \text{ MK} < T < 0.65 \text{ MK}$) is larger than that in case N. Figure 12 shows the temperature in the $x = 0.3 \text{ Mm}$ plane in each case. In case R, the prominence–corona transition region (PCTR) is deformed by the Rayleigh–Taylor instability and the area of the lower temperature region (the area represented by the green to yellow colors in Figure 12) is broader than that in case N. The difference in the DEM profiles is therefore due to the deformation via the Rayleigh–Taylor instability in the nonlinear regime. The radiative cooling rate of the plasmas in the temperature range of $0.3 \text{ MK} < T < 0.5 \text{ MK}$ is higher than those for other temperatures. Because a larger mass is distributed in this temperature region, condensation proceeds more efficiently in case R than in case N.
4. Discussion

We reproduced a prominence with vertical threads and internal flows within the framework of the reconnection–condensation model. In the model, reconnection via footpoint conversing motion creates a flux rope structure, and the radiative condensation inside the flux rope leads to formation of prominence. In contrast to the homogeneous footpoint motion in the previous simulations (Kaneko & Yokoyama 2015, 2017), the speed of the footpoint motion varies randomly along the PIL in the present simulation. Because of the density and velocity fluctuations inside the flux rope given by the random footpoint speed, radiative condensation heterogeneously proceeds. As a result, the prominence with highly dynamic interior was reproduced. We confirmed that the RTI-like mechanism facilitates the corrugation of the interface between the prominence and the lower corona. The descending spikes reflected at the bottom boundary generates the upflows. Due to the collision of the downflows and the upflows, the spikes evolves into the thin vertical threads. The downward speeds of the spikes and the widths of the vertical threads are consistent with the typical values in observations.

We found that the mass condensation rate is enhanced in a dynamic state. The condensation rate was comparable to the mass drainage rate of the downflows in our simulation. Significant mass drainage via downflows has actually been found in observations of prominences (Zirker et al. 1994; Liu et al. 2012). A study by Liu et al. (2012) suggested that the condensation rate is comparable to the mass drainage rate to maintain the total mass of the prominence. Our results support this suggestion from the observations. In a recent simulation by Xia & Keppens (2016b), the mass circulation between the chromosphere and the corona via a prominence was reproduced based on the evaporation–condensation model. However, the presence of the Rayleigh–Taylor instability was unclear in their fragmented condensations. It is also likely that the condensation rate in their simulation could be enhanced because the fragmentation extends the total volume of PCTRs.

In our simulation, the thin vertical threads form after the spikes are reflected at the bottom boundary, where the velocities are numerically fixed to zero. Previous MHD simulation including the chromosphere showed that the materials in falling spikes of the Rayleigh–Taylor instability are reflected at the top of the chromosphere and returned to
The dashed solid and dashed lines represent the growth rates in cases R and N, respectively. Panel coronal heights (Keppens et al. 2015). It is possible to assume the artificial boundary in our simulation as the interface between the chromosphere and the corona. In observations, the complicated flow pattern in the prominences suggests that the collisions between internal flows are likely to occur; however, the reflection of downflows at the chromosphere has not been clearly detected.

The observed dark plumes (Berger et al. 2008, 2010) are not reproduced in our simulation. In observations (Berger et al. 2011; Dudik et al. 2012; Berger et al. 2017), the dark plumes (corresponding to the bubbles of the Rayleigh–Taylor instability) originated at the interface between the prominence and the buoyant cavity, which is a semicircular dark region between the prominence and the chromosphere. The buoyant cavities (also called prominence bubbles) were conjectured to be emerging fluxes. To explain the origin of dark plumes inside prominence bodies, it might be necessary to consider additional emerging fluxes interacting with the flux ropes.

In observations, a possible origin of the footpoint converging motion is the collisions of supergranular diverging flows, which have typical speeds of \(0.3 \text{ km s}^{-1}\) on the photosphere and lifetimes of one day (Rondi et al. 2007; Schmieder et al. 2014). In the present simulations, we set the footpoint motion with a speed of \(v_{00} = 12 \text{ km s}^{-1}\) and a duration of 1440 s. The footpoint speed in our simulations is much faster than the typical values in the observations; however, the migration distance of the magnetic components (the product of the speed and duration) is comparable. The footpoint speed in our simulation is still much slower than the sound speed and the Alfvén speed in the corona. It is likely that this result will not change significantly even if the typical observational speed is applied.

In our simulation, the descending spikes (and the perturbation at the interface) have a spatial scale of approximately 7 Mm. It was difficult to clarify the origin of this spatial scale from a single simulation. However, it is possible that the spatial scale is the wavelength of the interchange mode of the Rayleigh–Taylor instability, which depends on the numerical viscosity in the simulation. To estimate the effect of numerical viscosity, the simulation results must be compared at different resolutions. It is also important to note that the growth of the upflow bubbles and the descending spikes may be asymmetrically influenced by radiative condensation, whereby the condensation flows cancel out the upflows and amplify the downflows. To clarify the effect, a simpler model including an equilibrium interface is needed, and a parameter survey of the perturbation wavelength should be performed.

The development of MHD simulations that includes a self-consistent ionization effect is one of the important future steps for the realistic numerical modeling. Numerical study of the magnetic Rayleigh–Taylor instability including ambipolar diffusion have indicated that the growth rate and the flow speed are affected by the cross-field diffusion (Khomenko et al. 2014). The analytical studies by Low et al. (2012a, 2012b) also suggest that prominence condensation inevitably creates a discrete current, leading to the spontaneous cross-field mass transport with the presence of neutrals. The ionization effect must be considered to obtain a more accurate evaluation of mass flux associated with downflows.

Recent observational studies have revealed the turbulent properties in quiescent prominences (Leonardis et al. 2012; Freed et al. 2016; Hillier et al. 2017). A break in the scaling exponent has been found to exist at a spatial scale of 2000 km in the power spectrum or structure function. Our simulation was unable to reproduce such a turbulent nature probably due to the numerical viscosity. Simulations with higher resolution are required to understand the turbulent characteristics of solar prominences.

![Figure 10](image1.png)

Figure 10. Panel (a) shows the temporal evolution of the total mass of cool plasma \((T < 10^5 \text{ K})\). The solid and dashed lines represent cases R and N, respectively. Panel (b) shows the mass growth rate and the drainage rate. The solid and dashed lines represent the growth rates in cases R and N, respectively. The dashed–dotted line represents the mass drainage rate in case R.

![Figure 11](image2.png)

Figure 11. Averaged DEMs during condensation. The solid and dashed lines represent the DEMs of case R at \(t = 88.0\) minutes and that of case N at \(t = 97.6\) minutes, respectively.
Figure 12. Snapshots of simulation results in the $x = 0.3$ Mm plane. The colors represent the temperature. Panels (a) and (b) show the snapshots of case R at $t = 88.0$ minutes and that of case N at $t = 97.6$ minutes, respectively.

5. Conclusions

We reproduced a prominence with vertical threads and internal downflows within the framework of the reconnection–condensation model. In the present model, footpoint motion with a random speed along a PIL is given. The Rayleigh–Taylor instability is triggered during radiative condensation. The spikes eventually evolve into descending pillars and thin vertical threads. It was found that the mass condensation rate is enhanced to the same level as the mass drainage rate in the dynamic state. The extension of the PCTR at lower temperatures, where the radiative loss is the highest, leads to higher mass condensation rate. Our results reveal the impact of the dynamic state on the radiative condensation rate and support the observational hypothesis claiming a balance between the condensation rate and the mass drainage rate in prominences.

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