NUMERICAL SIMULATIONS OF THE MAGNETIC RAYLEIGH–TAYLOR INSTABILITY IN THE KIPPENHAHN–SCHLÜTER PROMINENCE MODEL. II. RECONNECTION-TRIGGERED DOWNFLOWS

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ABSTRACT

The launch of the Hinode satellite has allowed high-resolution observations of supersonic bright downflows in quiescent prominences, known as prominence knots. We present observations in the Ca II H spectral line using the Solar Optical Telescope on board the Hinode satellite of a descending plasma knot of size ~900 km. The knot initially undergoes ballistic motion before undergoing impulsive accelerations at the same time as experiencing increases in intensity. We also present a subset of our three-dimensional magnetohydrodynamic simulations, performed to investigate the nonlinear stability of the Kippenhahn–Shl¨uter prominence model to the magnetic Rayleigh–Taylor instability in which interchange reconnection occurs. The interchange reconnection in the model breaks the force balance along the field lines which initiates the downflows. The downflows propagate with a downward fluid velocity of ~15 km s−1 and a characteristic size of ~700 km. We conclude that the observed plasma blob and the simulated downflow are driven by the breaking of the force balance along the magnetic field as a result of a change in magnetic topology caused by reconnection of the magnetic field.

Key words: instabilities – magnetic reconnection – magnetohydrodynamics (MHD) – methods: numerical – Sun: filaments, prominences

Online-only material: animation, color figures

1. INTRODUCTION

Quiescent prominences are large structures that exist in quiet regions of the solar corona of ~104 K (Tandberg-Hanssen 1995), ~1011 cm−3 (Hirayama 1986) plasma, predominantly found at high heliographic latitudes. The characteristic gas pressure of 0.6 dyn cm−2 (Hirayama 1986) and characteristic magnetic field strength of 3–30 G (Leroy 1989) give a plasma β ~ 0.01–1. Quiescent prominences often exist in the corona for weeks, implying a global stability.

In contrast to the global stability, quiescent prominences are known to be locally highly dynamic phenomena. Early observations of quiescent prominences found downflows (Engvold 1981), vortices of approximately 105 km x 105 km in size (Liggett and Zirin 1984), counterstreaming flows (Zirker et al. 1998), and large-scale instabilities of unknown origin (Stellmacher & Wiehr 1973). A full description of the current understanding of the structure and dynamics of quiescent prominences can be found in many reviews of the subject (see, for example, Tandberg-Hanssen 1995; Labrosse et al. 2010; Mackay et al. 2010).

Observations by the Solar Optical Telescope (SOT; Tsuneta et al. 2008) on the Hinode satellite (Kosugi et al. 2007) have provided a timely reminder that on a small-scale, quiescent prominences are highly dynamic and unstable phenomena. Berger et al. (2008) and de Toma et al. (2008) independently re-discovered the large-scale instabilities of Stellmacher & Wiehr (1973) using SOT and the Mauna Loa Solar Observatory, respectively. Both observations reveal large-scale dark “bubble-like” upflows that propagate from the chromosphere into, and in some cases through, the overlying prominences. Berger et al. (2010) analyzed the instability characteristics of the bubbles to show that they can break into small upflow plumes of approximately 2–6 Mm in width that propagate 10–15 Mm in height from their origin points. Measurements using the Atmospheric Imaging Assembly instrument on the Solar Dynamics Observatory show that the temperature of the plumes, and thus by inference the source region bubbles, was at least 250,000 K and more likely as high as 1.2 MK (Berger et al. 2011).

Hillier et al. (2011b) presented observations of plasma blob ejections from the top of a quiescent prominence that are impulsively accelerated to Alfvén velocity and then follow ballistic motion. Zapi´or & Rudawy (2007) described the three-dimensional (3D) velocity of “blobs” ejected from active region prominences and found that there was a significant line-of-sight velocity component associated with the ejections. Bright descending knots that form inside prominence threads and travel at an average speed of 16 km s−1 were observed by Chae (2010). The observed knots were impulsively accelerated throughout their lifetimes. An interchange reconnection process allowing material to be passed between prominence dips, based on the model of Petrie & Low (2005) that allowed the plasma to fall through the prominence, was proposed to explain this process. Haerendel & Berger (2011) described a “droplet” model for prominence downflow formation where the packets of plasma drop through a predominantly horizontal field. The velocity of the falling plasma is kept constant by the formation of Alfvén waves in the horizontal magnetic field. Low et al. (2012) show that “catastrophic radiative cooling” can lead to the formation of dense plasma concentrations that break the frozen-in flux condition to form discrete downflow blobs or droplets.

The model that we use in this work is the Kippenhahn–Schluter model (Kippenhahn & Schluter 1957; Priest 1982). The Kippenhahn–Schluger model for the solar prominence uses the Lorentz force of a curved magnetic field to support plasma against gravity, where vertically the magnetic tension balances gravity and horizontally the gas pressure balances the magnetic pressure. This model describes the local structure of the prominence, without including a corona, and is uniform in the vertical direction. This model has been shown to be linearly stable to ideal magnetohydrodynamic (MHD) perturbations...
The intensity profile along the be the center of the plasma blob. The newly determined central to those observed by Berger at al. (2010). (2012), where the instability was found to create upflows similar

Figure 1. Quiescent prominence observed in the Ca II H-line 396.8 nm spectral line on 2007 October 3 01:19 UT. Disk position 41°N 84°W. The pixel size is 0.108 arcsec pixel\(^{-1}\). The arrow denotes the position of the downward propagating blob. Intensity masking of the disk and the off-limb spicules is used to show the off-limb prominence.

(Kippenhahn & Schlüter 1957; Anzer 1969). The nonlinear stability of this model to the magnetic Rayleigh–Taylor instability was investigated by Hillier et al. (2011a) and Hillier et al. (2012), where the instability was found to create upflows similar to those observed by Berger et al. (2010).

In this paper, we present a study of how current sheets created by the nonlinear stage of the magnetic Rayleigh–Taylor instability lead to reconnection that triggers supersonic downflows in the Kippenhahn–Schlüter model that are similar to the downflows presented in Chae (2010). First, the observations of a fast downward blob will be presented in Section 2. In Section 3, we will describe the numerical method, and the results are then presented and explained in Section 4 with a summary and discussion of both the observations and simulations given in Section 5.

2. OBSERVATIONS OF DESCENDING PLASMA BLOB

Figure 1 shows a quiescent prominence seen on the NW solar limb on 2007 October 3 observed by the SOT with the Ca II H filter at a cadence of 30 s. The time series of this observation was between 01:16 UT and 04:59 UT. This prominence presents many interesting dynamic features, for example, at the start of this observation (01:16 UT), a large bubble has formed inside the prominence similar to those described in Berger et al. (2010). There are also a number of bright threads and downwardly propagating knots that occur during the duration of the observations as well as the upwardly ejected plasma blobs presented in Hillier et al. (2011b).

This section focuses on the occurrence of a falling plasma blob. The initiation of this bright blob is highlighted by the arrow in Figure 1. To determine the physical parameters of the plasma blob, the position is first determined by eye detection, and this position is then refined by employing a brightest pixel detection routine in a 15 × 15 pixel square. This point is considered to be the center of the plasma blob. The newly determined central point of the plasma blob is used to perform a Gaussian fit of the intensity profile along the X-direction through the center of the plasma blob. The full width at half-maximum of the Gaussian curve is considered to be the size of the plasma blob. This size was found to be \( \sim 900 \) km.

Figure 2 shows the temporal evolution of the plasma blob, where the arrows highlight the position of the blob. The evolution can be described as follows. The boundary between the prominence and the large bubble begins to undulate, from which upflows and downflows propagate. The downflow presented in this paper follows free-fall motion, then reaches a terminal velocity. Once this state has been achieved, there is occasional impulsive downward acceleration of the plasma blob. These dynamics can be seen clearly in Figure 3.

Figure 3 shows the vertical, i.e., against gravity, (dashed line) and horizontal (dash and dot line) velocities as well as the intensity (solid line) of the plasma blob. The velocity at time \( t_n \) is calculated using the slope given by a linear fitting of the position at time \( t_{n-1} \), \( t_n \), and \( t_{n+1} \). The dash and three dot line is the reference slope for free fall under gravitational acceleration at the solar surface of \( 2.7 \times 10^5 \) m s\(^{-2}\). Initially, the blob descends at almost free-fall velocity, with downward acceleration of \( 2.7 \times 10^5 \) m s\(^{-2}\). Then, from \( \sim 200–350 \) s, the blob maintains an almost constant downward velocity of \( \sim 28 \) km s\(^{-1}\). After \( 350 \) s, the blob undergoes impulsive accelerations followed by decelerations. During the periods of acceleration, the blob is accelerated at approximately free fall under gravitational acceleration. These peaks in downward acceleration coincide with peaks in the intensity of the blob; this was also noted in Chae (2010).

3. SIMULATION MODEL FOR PLASMA BLOBS

To simulate the observed bright plasma downflows, we use the 3D conservative ideal MHD equations. Constant gravitational acceleration is assumed, but viscosity, heat conduction, and radiative cooling terms are neglected. Though this paper focuses on results that are the consequence of reconnection, the effect of diffusion is neglected. This includes ambipolar diffusion, which has been shown to be able to alter the prominence geometry (Hillier et al. 2010), but the timescales of interest here are much shorter than those expected for ambipolar diffusion, and so have been neglected. The reconnection in this simulation takes place due to numerical diffusion. As a conservative scheme is used, the magnetic energy dissipated is converted to internal energy. We assume the medium to be an ideal gas.

The equations are non-dimensionalized using the sound speed \( C_s = 13.2 \) km s\(^{-1}\), the pressure scale height \( \Lambda = C_s/(\gamma g) = RgT/(\mu g) = 6.1 \times 10^7 \) cm, the density at the center of the prominence \( \rho(x = 0) = 10^{-13} \) g cm\(^{-3}\), and temperature \( T = 10^4 \) K giving a characteristic timescale of \( \tau = \Lambda/C_s = 47 \) s. We take \( \gamma = 1.05 \) and \( \beta = 0.5 \), which gives an Alfvén velocity of \( V_A = C_s\sqrt{2/\gamma \beta} = 25.8 \) km s\(^{-1}\). As these simulations are being compared directly with observations, all the lengths, times, and velocities will be given in physical units, and the above values are given so that it can be understood how the physical values are calculated.

The initial model is as follows:

\[
B_x(x) = B_{x,0}
\]

\[
B_z(x) = B_{z,\infty} \operatorname{tanh} \left( \frac{B_{z,\infty} x}{2B_{x,0} \Lambda} \right)
\]

\[
\rho(x) = \frac{B^2_{x,\infty}}{8\pi} \cosh^{-2} \left( \frac{B_{x,\infty} x}{2B_{x,0} \Lambda} \right)
\]

\[
\rho(x) = \frac{1}{g} \frac{B^2_{x,\infty}}{8\pi} \cosh^{-2} \left( \frac{B_{x,\infty} x}{2B_{x,0} \Lambda} \right)
\]

\[
\rho(x) = \frac{1}{g} \frac{B^2_{x,\infty}}{8\pi} \cosh^{-2} \left( \frac{B_{x,\infty} x}{2B_{x,0} \Lambda} \right)
\]
where \( B_{x0} \) is the value of the horizontal field at \( x = 0 \) and \( B_{x\infty} \) is the value of the vertical field as \( x \to \infty \). This model is the Kippenhahn–Schlüter model as presented in Priest (1982).

As the Kippenhahn–Schlüter model is linearly stable to ideal MHD perturbations, a nonlinear perturbation is necessary. The perturbation considered here is a high-temperature, low-density tube placed in the center of the Kippenhahn–Schlüter model. To create a system that allows the excitement of the interchange of the magnetic field, the effect of a low-density (high-temperature) tube placed inside the prominence of the form

\[
\rho'(x, z) = -0.25 \rho_{nd} \times \rho(x) \times \left[ \tanh \left( \frac{|z| - H_z/2}{0.3\Lambda} \right) + 1.0 \right] \\
\times \left[ \tanh \left( \frac{|x| - W_x/2}{0.3\Lambda} \right) + 1.0 \right]
\]

is initially considered. Where \( \rho_{nd} \) is the normalized density difference, \( H_z \) is the height of the bubble and \( W_x \) is the width of the bubble. The density of the tube at \( x = z = 0 \) is \( 0.3 \rho(0) \) with width \( 2\Lambda = 1.22 \text{ Mm} \) and height \( 8\Lambda = 4.88 \text{ Mm} \). Therefore, the total density at the start of the simulation is given by \( \rho(x, z) = \rho(x) + \rho'(x, z) \). The pressure distribution is left unchanged, therefore the temperature distribution is given by \( T(x, z) = \gamma p(x)/(\rho(x) + \rho'(x, z)) \). A full description of the initial conditions and the meaning of all parameters can be found in Hillier et al. (2012).

Figure 4 shows a visual representation of the initial conditions. The color contour represents the mass density; the white lines represent selected magnetic field lines. To excite the interchange mode, a velocity perturbation in \( v_y \) was imposed, where \( v_y \) was given as a sum of sinusoidal curves of different wavelength. This velocity perturbation takes the form \( v_y(x, y) = \Delta \Sigma_{l=x}(l)) \sin(2.0\pi y/(l + \text{Rand}(l, x, y))) \cos(2.0\pi x/(l + \text{Rand}(l, x, y))) \). Where the maximum value of the perturbation \( |v_y(x, y)| < 0.01 C_s \), and \( \text{Rand}(l, x, y) \) is a random number where \( |\text{Rand}(l, x, y)| \leq 0.5 \).

To reduce computational time, we assume a reflective symmetry boundary at \( x = 0 \). Due to the nature of the magnetic field at the top and bottom (z) boundary and at \( x = L_x \), the choice for boundary is very limited. A free boundary is assumed at \( x = L_x \) with a damping zone (damping time \( \tau_D = 1.0/4.4\tau = 10.5 \text{ s} \))
for the hydromagnetic variables and $B_z$ (to maintain the angle of the magnetic field at the boundary). For the top and bottom boundary, a periodic boundary is assumed. A periodic boundary is also used at $y = 0, Z$.

The scheme used is a two-step Lax–Wendroff scheme based on the scheme presented in Ugai (2008), using the artificial viscosity and smoothing also presented in this paper. The grid size is uniform in the $y$-direction, and in the $x$–$z$ plane we take a grid of $75 \times 400$ grid points, where the total area of the calculation domain is $3.5\Lambda \times 85\Lambda \sim 2.14 \times 51.9$ Mm. A fine mesh is assigned to an area of $40 \times 320$ grid points, of actual size $1.2\Lambda \times 30\Lambda \sim 0.73 \times 18.3$ Mm, around the upper contact discontinuity allowing the plumes to be resolved. The non-uniform grid is made through a continuous construct using hyperbolic tangents. The equations are as follows:

$$dx(i) = 0.03 + 0.5 \times 0.1 \times (\tanh((i - 70)/30) + 1)$$

$$dz(k) = 0.1 + 0.5 \times 0.9(\tanh((k - 320)/10) + 1)$$

$$- 0.5 \times 0.9(\tanh((k - 370)/10) + 1),$$

where $i$ and $k$ are the $i$th and $k$th grid points in the $x$- and $z$-directions, respectively. In the $y$-direction, a total of 150 grid points were used with $dy = 0.05\Lambda \sim 0.03$ Mm.

4. SIMULATION RESULTS OF RECONNECTION-TRIGGERED DOWNFLOWS

Figure 5 shows the evolution of the upflows. Upflows of size $\sim 1.2$–$1.8$ Mm ($2\Lambda$–$3\Lambda$) in width with velocity $\sim 6$ km $s^{-1}$ ($\sim 0.45C_s$) can be clearly seen in panel (d) of the figure. A full

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**Figure 4.** Contour plots of the initial density distribution for the standard model for (a) the $x$–$z$ plane at $y = 0$ (with magnetic field lines) and (b) the $y$–$z$ plane at $x = 0$ as well as the 3D isosurface of the density distribution showing selected field lines. All physical quantities are initially constant in the $y$-direction. The initial velocity perturbation is applied to the upper contact discontinuity in the $y$-direction. The 3D density isosurface only shows the distribution for $z > 0$, so only one high-density region is displayed.

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

**Figure 5.** Temporal evolution of the density distribution showing the downflows for $t = 709, 1568, 2103$, and $2403$ s ($t = 15.3, 35.8, 45.5$, and $52.0$), respectively, taken in the $y$–$z$ plane at $x = 0$. The black diamond highlights the position of the test particle, where the velocity of this test particle is shown in Figure 6. The vertical line in panel (a) marks the position where the forces taken in the $x$–$z$ plane we

(A color version of this figure is available in the online journal.)
description of the evolution of the magnetic Rayleigh–Taylor instability in the Kippenhahn–Schlüter prominence model can be found in Hillier et al. (2011a) and Hillier et al. (2012). The main point of interest in the magnetic Rayleigh–Taylor instability is that the nonlinear inverse cascade, driven by interaction between flows created by the instability, creates the dynamics found. Therefore, the sizes of the structures that are found are larger than those of the most unstable wavenumber of the magnetic Rayleigh–Taylor instability. What makes the two cases presented here of particular interest is that reconnection occurs, triggering fast downflows. The characteristics of these downflows will now be explained.

The first impulsive downflow is excited at an approximate time of 2103 s (t = 45.5r). This time (t = 2103 s = 45.5r) is shown in Figure 5(c). This downflow (and three others that occur in the same prominence thread) has a velocity \( v_z \approx 15 \text{ km s}^{-1} \) (1.2\( C_s \)) and a characteristic size of \( \sim 700 \text{ km} \) (1.2\( \Lambda \)). From these snapshots of the dynamics, it is possible to see that the downflow of interest is significantly faster than any downflow created by the nonlinear ideal processes of the magnetic Rayleigh–Taylor instability in this model.

Figure 6 shows the \( v_x \) and \( v_y \) components of the velocity of a test particle and the normalized density \( (\rho/\rho_0) \) at the position of the test particle in the simulation domain as dash-dotted, dashed, and solid lines, respectively. The particle was inserted in the \( x = 0 \) plane, so the symmetry of the simulation means that the particle does not move from this plane. The white diamond in Figure 5 follows the position of a test particle inserted into the \( y-z \) plane. The figure clearly shows that this particle follows the evolution of the fast downflow. The approximate time when the dense downflows that interact with the particle appear in the \( x = 0 \) plane are shown by the dash-triple-dot vertical lines. The velocity plot shows that the particle is strongly accelerated downward to about 13 \text{ km s}^{-1}, also gaining a horizontal velocity of approximately 3 \text{ km s}^{-1}. At the same time as this strong acceleration, the density at the position of the test particle jumps significantly. A similar impulsive acceleration is found for the second reconnection event. Though not shown in the velocity of the test particle, there is a maximum velocity of 16 km s\(^{-1}\) found in the descending blob, therefore supersonic velocities \( (>C_s) \) are formed associated with the formation of high-density regions.

Figure 7 shows the forces acting in the \( z \)-direction at \( x = 0 \text{ Mm}, y = 1.9 \text{ Mm} \) (3.1\( \Lambda \)) between \( z = 1.2–7.3 \text{ Mm} \) (2\( \Lambda \)–12\( \Lambda \)) at time \( t = 2008, 2103, \) and 2222 s (\( t = 43.5, 45.5, \) and 48.1\( r \)). At first, a density increase gives a net downward force, i.e., the plasma blob is accelerated downward by gravity. As the flow propagates downward, stretching of the magnetic field results in an increase in tension, decelerating the dense blob.

Figure 8 shows selected field lines plotted with a density isosurface to display the 3D structure of the downflow. Due to the high velocity of the downflow, the magnetic field is dragged down, creating a tension force that works against the downflow. The stretching of the magnetic field that results in the increased tension can be seen. The dynamics of the downward propagating plasma blob can be clearly seen in the online movie of Figure 8.

Figure 9 shows (a) a closeup of the evolution of three dense blobs with the arrows showing the projected velocity vector in the \( x = 0 \) plane and (b) the change in pressure from the original distribution with the projected velocity vector and selected projected magnetic field lines in the \( y = 2.74 \text{ Mm} \) (4.5\( \Lambda \)) plane. Panels (i)–(iv) show times \( t = 2103, 2222, 2347, \) and 2450 s (\( t = 45.5, 48.1, 50.8, \) and 53.1\( r \)). The position of the \( y = 2.75 \text{ Mm} \) plane is marked by the white dashed line in row (a) panel (iv). The long arrows on (a)(i) (black vertical) and (b)(i) (white oblique) show the position of an oblique plane

![Figure 6. Velocity and normalized density at the position of a test particle inserted into the y–z plane. Vertical lines show approximate times at which the reconnection downflows interact with the test particle.](image-url)

![Figure 7. Distribution of the force acting in the vertical direction at x = 0 Mm, y = 1.9 Mm (3.1\( \Lambda \)) between z = 1.2–7.3 Mm (2\( \Lambda \)–12\( \Lambda \)) at time t = 2008, 2103, and 2222 s (\( t = 43.5, 45.5, \) and 48.1\( r \)). The line along which these forces are calculated is shown in panel (a) of Figure 5.](image-url)
Figure 8. 3D structure of fast downflows with selected magnetic field lines at times $t = 2008$ and $2222$ s ($t = 43.5$ and $48.1\tau$).

(An animation and a color version of this figure are available in the online journal.)

Figure 9. Row (a) shows the evolution of the density with arrows showing velocity in the $x = 0$ plane and row (b) shows the evolution of the pressure change from initial distribution with arrows showing velocity and selected magnetic field lines in the $y = 2.74$ Mm (4.5\text{\textcircled{A}}) plane. Panels (i)-(iv) show times $t = 2103, 2222, 2347,$ and $2450$ s ($t = 45.5, 48.1, 50.8, \text{ and } 53.1\tau$). Bright, fast downflows can be seen in panels (i), (iii), and (iv) of row (a). Row (b) shows that the downflow seen in (a)(iv) is the result of velocities created along the magnetic field. The arrows on row (a) and (b) panel (i) show the position of the oblique plane used in Figure 10. A shock has formed at coordinates [1.0, 2.4] in panel (b)(iii).

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

Row (a) of Figure 9 shows that the dense blobs create vortex flows as they propagate downward. In panel (ii), the role-up created by the Kelvin–Helmholtz instability of the downward flow can be seen. Row (b) panel (iii) shows fast flows that have developed along the magnetic field. The region where the fast downflow is initiated is at the same ($y, z$) coordinates as a vortex created by the downward propagation of a fast downflow. A shock is formed where these fast flows fall down the field
lines. This is similar to the accretion shock created at the foot points of magnetic loops formed through the Parker instability (Matsumoto et al. 1988; Shibata et al. 1989). Panel (iv) shows that these inflows become fast (>Cₜ) downflows at the center of the prominence.

Figure 10 shows (a) the velocity perpendicular to the oblique plane indicated in Figure 9(i) with arrows showing the in-plane component and (b) the absolute value of the current density with selected projected magnetic field lines. The vertical unit L refers to the length along the cut. The long arrows on (a)(i) (black vertical) and (b)(i) (white oblique) of Figure 9 show the position of the oblique plane used. The direction of the arrows refers to the direction of increasing L. The four columns labeled (i), (ii), (iii), and (iv) show times τ = 2103, 2222, 2347, and 2450 s (τ = 45.5, 48.1, 50.8, and 53.1τ). Figure 10 row (a) panel (iii) shows a region of fast flow perpendicular to the plane at [2.7, 1.0] ([4.43A, 1.64A]) from row (b) panel (iii), it can be seen that the region where the flow perpendicular to the plane occurs is cospatial with a region where reconnection has created an o-point. Other o-points can be seen, but these are not created by reconnection. They are created by convergence of the component of the magnetic field in the plane. It can be seen that the o-point formed in row (b) panel (iii) is cospatial with a region of high velocity. This velocity is almost anti-parallel to the direction of the magnetic field (as can be seen in Figure 9 row (b) panel (iii)), therefore only hydrodynamic forces can accelerate the plasma.

The acceleration process for the fast downflows can be summarized as follows: the magnetic Rayleigh–Taylor instability drives the creation of current sheets. Reconnection occurs in these current sheets. If there is a significant difference in the distribution of field-aligned forces across the reconnecting field lines, then the change in topology allows these hydrodynamic forces to create field aligned flows. The flows toward the center of the prominence reach supersonic velocities that form a shock. The resulting increase in density at the center of the prominence breaks the vertical force balance accelerating the plasma downward as fast-moving, dense blobs.

To improve the clarity with which we can show this reconnection region, we devise a function to show the regions where x-points are created. This function can be written as follows:

\[ \text{Xpoint}(i, j) = \left\{ \begin{array}{ll} 0.5 \left( \frac{B_y(i, j + 1)}{|B_y(i, j + 1)|} - \frac{B_y(i, j - 1)}{|B_y(i, j - 1)|} \right) + 0.5 \left( \frac{B_l(i + 1, j)}{|B_l(i + 1, j)|} - \frac{B_l(i - 1, j)}{|B_l(i - 1, j)|} \right) & \text{if } B_y(i, j) > 0, \\ 0 & \text{otherwise} \end{array} \right. \]  

where i and j are the grid points in the oblique plane in the y- and l-directions, respectively, By is the y component of the magnetic field in the plane and Bl is the component of the magnetic field along the length of the oblique plane. This function gives a value of one when one of the components of the magnetic field reverses sign and a value of two when both the components of the magnetic field reverse sign at an x-point. When there is an o-point, the function returns a value of zero. For all other cases, a value of zero is returned. Therefore, using this method, the formation of x-points can be investigated. It should be noted that changing the minus sign to a plus sign would give something similar to a unit current.

Figure 11 shows the positions that this function takes on the value one (gray) and two (white) in association with the evolution of the magnetic field. This is shown at τ = 2103, 2158, and 2222 s (τ = 45.5, 46.7, and 48.1τ) in panels (i), (ii), and (iii), respectively, with panel (iv) showing an expanded view of
Figure 11. Positions of the x-points in the oblique plane at $t = 2103$, 2158, and 2222 s ($t = 45.5$, 46.7, and 48.1 τ) in panels (i), (ii), and (iii), respectively. The gray regions show where one component of the magnetic field changes sign and the white dots show where an x-point exists. Panel (iv) shows an expanded view of the area shown in the dashed box in panel (iii). The formation of the x-point (highlighted by the arrow in panel (iv)) at $\sim [2.35, 1.35]$ Mm is clearly associated with the formation of the o-point.

5. RELATION BETWEEN RECONNECTION-TRIGGERED DOWNFLOWS AND KNOTS

This paper presents 3D simulations of dense downflows triggered by reconnection in current sheets that are created by the nonlinear evolution of the magnetic Rayleigh–Taylor instability in the Kippenhan–Schlüter prominence model. Also presented are observations of a downwardly propagating plasma blob, excited by the falling of prominence material into a large bubble that has formed below the prominence, that shows intermittent acceleration associated with increases in the blob intensity.

The observations show a plasma blob of $\sim 900$ km in size being impulsively accelerated to $\sim 45$ km s$^{-1}$. The simulations show downflows of $\sim 700$ km in size impulsively accelerated to $\sim 15$ km s$^{-1}$. The dynamics of the impulsive acceleration and the formation could imply that they are the same phenomenon. The size is clearly well matched, but there is a factor of three difference in the velocity. There are three possible ways to explain this. First, there is a lower density in the large bubble than used in this simulation, this is suggested by the high temperature of the plumes found by Berger et al. (2011), which may allow the blob to flow through the cavity with greater ease. Second, there is a weaker magnetic field associated with the bright downflow, so greater stretching is required to create the tension force to balance gravity. Finally, there is a larger angle between the reconnection field components, resulting in a greater force imbalance (assuming approximately hydrostatic equilibrium along the field line). It can be concluded that the reconnection-triggered downflow model provides a good explanation for the observations as well as a good explanation of the results presented in Chae (2010).

The simulation results verify that the knot formation mechanism proposed by Chae (2010) can produce impulsively accelerated downflows. On top of this, the multiple knot formation
and supersonic speeds were found to match between simulations and observations. The absence of upflows due to this mechanism may imply that this type of interchange reconnection may not be a valid explanation for the upwardly ejected plasma blobs presented by Hillier et al. (2011b). The importance of this simulation is that it implies that the observed knots can be produced by reconnection, but do not rely on matter condensation to create them. Merely with the current sheets that are formed by subsonic flows, where many such flows have been observed in prominences, can reconnection be driven. These results predict that downward propagating knots can be formed where strong shear flows/vortices are excited creating current sheets.

When this reconnection mechanism is applied to a prominence, it could be suggested that the Joule heating from the current sheets may result in heating of the prominence material resulting in a reduction in the intensity instead of an increase as the material may no longer be in the correct temperature range to be visible in chromospheric lines. However, the reconnection presented here is the component reconnection in a high magnetic Reynolds number system, which occurs in the high-density material of a quiescent prominence. If we compare the terms in the energy equation to estimate the heating, then we find

$$\frac{R}{\mu(y-1)} \frac{\Delta T_{\text{PROM}}}{\Delta t} \sim \eta J^2,$$

where \(\Delta t\) is given as the Alfvén crossing time \(L/V_A\). This can be rearranged to give a value for the length scale associated with the heating,

$$L \sim \frac{R}{\mu(y-1)} \frac{\eta V_A}{\Delta T},$$

where \(V_A\) is the prominence Alfvén velocity. Taking a temperature change of \(10^4\), an Alfvén velocity of \(5 \times 10^6\) cm s\(^{-1}\), and a diffusion of \(10^7\) cm\(^2\) s\(^{-1}\), the length scale found is \(\sim 10^2\) cm. Therefore, the Joule heating that may occur from this reconnection would only be able to heat the prominence material so that it could no longer emit in chromospheric lines on scales that are much smaller than the observed phenomena. This would mean that Joule heating should not significantly affect the observed prominence flows.

It was shown that these fast downflows can begin to wind up the magnetic field. van Ballegooijen & Cranmer (2010) hypothesized that the nonlinear evolution of the Rayleigh–Taylor instability in a quiescent prominence would lead to the formation of threads of tangled magnetic field supporting the prominence material. The twisted field presented could represent the onset of the formation of this field.

This type of interchange/component reconnection has been observed in other systems. Isobe et al. (2006) showed how current sheets created by the nonlinear Rayleigh–Taylor instability acting on emerging flux led to reconnection and, as a result, formation of filamentary structure. Jiang et al. (2011) presented simulations that show how 3D Pestchek type reconnection can lead to outflows that propagate along the magnetic field. In this case, the plasma was initially accelerated by magnetic pressure, which is different from these simulations. Here, a different mechanism for acceleration is presented, but it is clear that this type of component reconnection can happen in many different solar phenomena, where flows can then be driven along the direction of the magnetic field. Further simulations looking at the long-term change in magnetic field distribution due to this reconnection mechanism for a global prominence model would be of great interest to show how the prominence structure changes with time.

To truly quantify the observed bright downflows in terms of this interchange reconnection model, we believe that further observation work is necessary. The first step would be to produce some statistics on the occurrence rate and lifetime of the bright downflows, and also investigate the initiation of the downflows in terms of the flows surrounding the initiation site as this could give some understanding of the formation of current sheets in prominences. The relationship between the acceleration, velocity, and the increase in intensity needs to be quantified to understand if slow-mode shocks can really explain the phenomenon. The simulated downflow shown in panel (ii) of Figure 9 shows a Kelvin–Helmholtz role-up of the plasma (see Figure 9 panel (a)(ii)); it needs to be understood if this is seen in the observations and under what conditions this may occur. The observed downflow in this paper first descends at almost free-fall velocity as it falls through the large bubble, and it would be very interesting to see if this can tell us something about the magnetic structure of the bubble. We aim to present such observations in a future paper.
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