Shock–Cloud Interaction in the Solar Corona

Takuya Takahashi\textsuperscript{1,2}

\textsuperscript{1} Department of Astronomy, Kyoto University, Sakyo, Kyoto, 606-8502, Japan; takahashi@kwasan.kyoto-u.ac.jp
\textsuperscript{2} Kwasan and Hida Observatories, Kyoto University, Yamashina, Kyoto 607-8471, Japan

Received 2016 May 26; revised 2017 January 21; accepted 2017 January 23; published 2017 February 17

Abstract

Flare-associated coronal shock waves sometimes interact with solar prominences, leading to large-amplitude prominence oscillations (LAPOs). Such prominence activation gives us a unique opportunity to track the time evolution of shock–cloud interaction in cosmic plasmas. Although the dynamics of interstellar shock–cloud interaction has been extensively studied, coronal shock–solar prominence interaction is rarely studied in the context of shock–cloud interaction. Associated with the X5.4 class solar flare that occurred on 2012 March 7, a globally propagated coronal shock wave interacted with a polar prominence, leading to LAPO. In this paper, we studied bulk acceleration and excitation of the internal flow of the shocked prominence using three-dimensional magnetohydrodynamic (MHD) simulations. We studied eight MHD simulation runs, each with different mass density structure of the prominence, and one hydrodynamic simulation run, and we compared the result. In order to compare the observed activation of prominence with the corresponding simulation, we also studied prominence activation by injection of a triangular-shaped coronal shock. We found that the prominence is first accelerated mainly by magnetic tension force as well as direct transmission of the shock, and later decelerated mainly by magnetic tension force. The internal flow, on the other hand, is excited during the shock front sweeps through the prominence and damps almost exponentially. We construct a phenomenological model of bulk momentum transfer from the shock to the prominence, which agreed quantitatively with all the simulation results. Based on the phenomenological prominence activation model, we diagnosed physical parameters of the coronal shock wave. The estimated energy of the coronal shock is several percent of the total energy released during the X5.4 flare.

Key words: magnetohydrodynamics (MHD) – shock waves – Sun: corona – Sun: filaments, prominences

1. Introduction

The interaction between interstellar clouds and shock waves associated with, for example, supernova remnants, H\textsc{ii} regions, or stellar winds has been studied as one of the most fundamental problems of interstellar gas dynamics. Shock–molecular cloud (MC) interaction is especially important as a process that dynamically drives star formation. Studies using hydrodynamics or magnetohydrodynamics (MHD) simulations revealed that strong interstellar shock waves injected into the molecular cloud crush and destroy it through shock injection and ensuing hydrodynamical instabilities such as Kelvin–Helmholtz (KH), Rayleigh–Taylor (RT), and Richtmyer–Meshkov instabilities, although inclusion of the magnetic field strongly suppress those instabilities (Woodward 1976; Nittmann et al. 1982; Klein et al. 1994; Mac Low \\& Zahnle 1994). Magnetic field orientation and mass density structure within clouds also affect significantly the later phase of the dynamics of molecular clouds impacted by shock waves (Poludnenko et al. 2002; Patnaude \\& Fesen 2005; Shin et al. 2008). Recent observations and numerical simulations reveal that multiphase, multiscale dynamics of filamentary cloud–shock interaction within a cloud is important for the evolution of star-forming molecular clouds with the help of thermal or self-gravitational instabilities (Inoue \\& Inutsuka 2012; Dobashi et al. 2014; Matsumoto et al. 2015).

On the other hand, shock waves are frequently observed in the corona of the Sun (Moreton 1960; Uchida 1968; Thompson et al. 2000; Warmuth et al. 2001; Vršnak et al. 2002; Liu et al. 2010). In the corona, magnetic field reconnection allows magnetic field energy to be released in a catastrophic way, resulting in the largest explosion in the solar system. These are called solar flares. During solar flares, typically \(10^{29}–10^{33}\) erg of magnetic field energy stored in the corona is converted to thermal, kinetic, radiation, and high-energy particle kinetic energy in a short (several minute) timescale (Shibata \\& Magara 2011).

As a result of sudden energy release in solar flares, part of the energy propagates globally in the corona as a form of nonlinear fast-mode MHD wave (or shock). Recently, high-cadence extreme-ultraviolet (EUV) observation of the solar corona by the Atmospheric Imaging Assembly (AIA; Title \\& AIA team 2006; Lemen et al. 2012) on board the Solar Dynamics Observatory (SDO; Pesnell et al. 2012) has allowed detailed imaging observation of such global shock waves in the corona (Asai \\& Ishii 2011). These flare-associated shock waves in the lower solar corona are thought to be weak, fast-mode MHD shocks (Ma et al. 2011; Gopalswamy et al. 2012). Plasma ejections (coronal mass ejections [CMEs]) associated with solar flares propagate in the interplanetary space, sometimes with clear shock fronts at their noses (Wang et al. 2001). These propagating shock waves in the upper corona or interplanetary space are observed in coronagraph images or in radio dynamic spectrum observations (Kai 1970). They are also observed as a sudden change in plasma parameters of solar wind velocity, temperature, density, and magnetic field strength observed in situ in front of Earth (Wang et al. 2001).

We also have “clouds” in the solar corona, the “prominences.” They are cool and dense plasma floating within a hot and rarefied corona. They are supported by a magnetic tension force against gravity (Labrosse et al. 2010; Mackay et al. 2010). Recent high-resolution and high-sensitivity observations by the Solar Optical Telescope (SOT; Tsuneta et al. 2008) on board the Hinode satellite (Kosugi et al. 2007) revealed the
highly dynamic nature of solar prominence, with continuous oscillation and turbulent flow (Berger et al. 2008, 2011). They give us a rare opportunity for studying dynamics of partially ionized plasma in detail, whose plasma parameters are very difficult to reach in the laboratory or in other space objects. The excitation mechanism of such chaotic flows within prominence material has also been discussed recently. Hillier et al. (2013) discussed photospheric motion as one possible mechanism that drives long-frequency small-amplitude prominence oscillations. Nonlinear MHD waves propagating upward toward the prominence foot are studied by Ofman et al. (2015). The magnetic RT instability invoked by interchange reconnection between magnetic field lines supporting prominence plasmas is discussed as an excitation mechanism for multiple plume-like upflows observed with Hinode/SOT, which help mix up prominence plasmas (Hillier et al. 2012). Magnetic KH instability excited near the absorption layer of transversely oscillating prominence in the corona helps cascade energy and heat the prominence plasma through turbulence excitation and current sheet formation (Antolin et al. 2015; Okamoto et al. 2015).

Sometimes, coronal shock waves associated with flares hit solar prominences and lead to excitement of large-amplitude prominence oscillations (LAPOs). These oscillations give us information on physical properties of prominences such as magnetic field strength, density, and eruptive stability and have been studied widely (Isobe et al. 2007; Gilbert et al. 2008).

In contrast to interstellar shock–molecular cloud interaction, which has been studied widely, the excitation process of large-amplitude solar prominences or prominence activations by coronal shock injection has not been studied in detail in the context of shock–cloud interaction. In contrast to the situation in the interstellar medium, where strong shock waves often interact with molecular clouds, the prominence excitation is the interaction between prominence and weak fast-mode MHD shock waves with fast-mode Mach number being between 1.1 and 1.5 (Narukage et al. 2002; Grechnev et al. 2011; Takahashi et al. 2015). The timescale of shock–prominence interaction is several minutes, which offers us a unique opportunity to study in detail the time evolution of the shock–cloud interaction process in cosmic plasmas.

In this paper, we analyze observational data of shock–prominence interaction obtained by SDO/AIA and compare them with numerical MHD simulations. Especially, we focus quantitatively on how wave momentum is transferred to cloud material through shock injection, flow drag, and magnetic tension acceleration processes with 3D MHD simulations. We also discuss the effect of internal structure such as volume filling factor on chaotic flow excitation and its damping. We compare the MHD simulation results with hydrodynamic simulations and discuss the role of magnetic field in bulk prominence acceleration, as well as the excitation and damping of internal flows. We make a phenomenological model that describes momentum transfer from the coronal shock to the prominence. We validate the phenomenological model by comparison with simulations and apply the phenomenological model of prominence activation in the context of diagnosing coronal shock properties. Lastly, we compare the shock–prominence interaction with the interstellar shock–molecular cloud interaction in the context of the MHD shock–cloud interaction.

2. EUV Observation of a Prominence Activation by the Coronal Shock Wave

On 2012 March 7, an X5.4 class flare occurred at NOAA active region (AR) 11429 located at the northeast quadrant of the solar disk. The soft X-ray light curve obtained by GOES peaked at 00:24 UT on March 7 (Figure 1(a)). This flare is the second-largest one in the current solar cycle (cycle 24). The flare was associated with a very fast CME whose velocity was about 2684 km s\(^{-1}\). Figure 1(b) is the composite of coronagraph images obtained by SOHO/LASCO C2 and EUV images obtained by SDO/AIA 193 Å band, both taken at 00:36 UT. We can see the shock front surrounding the CME ejecta in the SOHO/LASCO C2 images (Figure 1(b)). Figure 1(d) shows the SDO/AIA 193 Å difference image at 00:18 UT. We see a dome-like bright structure expanding above AR 11429 in Figure 1(d).

Takahashi et al. (2015) estimated the propagation speed of the leading shock front ahead of the CME as 672 km s\(^{-1}\) based on the analysis of the dynamic spectrum obtained with the Hiraiso Radio Spectrograph (HRAS; Kondo et al. 1995). It seems strange that the estimated speed of the leading shock front ahead of the CME is much slower than the CME speed estimated with SOHO/LASCO. We looked at the coronagraph observation data by STEREO-B/COR1 and radio dynamic spectrum once again to check the consistency. The height of the leading edge of the CME ejecta seen in the STEREO-B/COR1 image taken at 00:26UT is larger than 2R\(\text{e}\) measured from the photosphere, where R\(\text{e}\) is the solar radius. On the other hand, the radio dynamic spectrum during the flare period is complicated and composed of Type II, Type IV, and possibly Type III bursts. We noticed that Takahashi et al. (2015) mistook the Type IV burst for a Type II burst, resulting in a CME speed estimation inconsistent with coronagraph observations. From 00:17:10 UT (indicated as P\(_1\)) in Figure 2) to 00:17:38 UT (indicated as P\(_2\)) in Figure 2), we see a clear linear structure in the radio dynamic spectrum showing the characteristic signature of Type II bursts, whose frequency drifted from 112 to 88 MHz during the period. Assuming that the Type II burst signature corresponds to the first harmonics of the plasma oscillation at the upstream of the leading shock front, we got the propagation speed of 1.9 \(\times\) 10\(^3\) km s\(^{-1}\) based on Newkirk (1961) and Mann et al. (1999), which seems to be consistent with coronagraph observations.

The footprints of the dome-like shock front that propagated to the north in the lower corona were especially bright in AIA 193 Å images. That shock front propagated further to the north and hit a prominence located at the north pole. The interaction between the shock wave and the polar prominence resulted in LAPO. We call the excitation process of LAPO “prominence activation” hereafter. Figure 3 shows the time evolution of the prominence activation seen in AIA 193 and 304 Å images. The FOV of Figure 3 is shown as the black rectangle in Figure 1(c).

The polar prominence was seen bright in AIA 304 Å images, while it was seen dark in AIA 193 images. Figures 3(a) and (b) show the prominence just before it is activated by the shock. Figures 3(c) and (d) show the prominence just as it is hit by the coronal shock front. The shocked part of the prominence seen in the AIA 304 Å image became about twice as bright as its original brightness. The bright part propagated in the direction of shock propagation, and at the same time the prominence

---

3 See http://cdaw.gsfc.nasa.gov/CME_list/UNIVERSAL/2012_03/univ2012_03.html.
started to move in the shock propagation direction. Figures 3(i) and (j) show time–distance diagrams along the cut AB shown in Figure 3(a) between 00:20 UT and 01:30 UT. In Figure 3(i), the coronal shock front is seen as a bright propagating structure. The propagation speed of the coronal shock front in the plane of the sky is measured to be 380 km s\(^{-1}\) from Figure 3(i). When the coronal shock front reached the dark prominence, the prominence was abruptly accelerated. The initial activated prominence speed was 48 km s\(^{-1}\) measured from Figure 3(i). We can clearly see the sudden brightening of the prominence during its activation in Figure 3(j). In Figure 3(j), we can see the somewhat chaotic movement of prominence threads during LAPO. We note here that we have neglected the line-of-sight component of the speed of shock propagation and activated prominence, so the estimated shock speed and prominence velocity should be regarded as lower limits. Figures 3(g) and (h) show the prominence when its displacement from the original position was largest. Note that the timescale of prominence activation is several minutes in this event, while the period of LAPO is longer than an hour. The white rectangle in Figure 3(i) corresponds to the prominence activation process that is also shown in Figure 20(a).

3. 3D MHD Simulation of Prominence Activation

In order to study in detail the physics of prominence activation, we carried out three-dimensional MHD simulations.

3.1. Numerical Methods

We numerically solved the following resistive MHD equations:

\[
\frac{\partial \rho}{\partial t} = -\nabla \cdot (\rho V)
\]  

**Figure 1.** (a) GOES soft X-ray light curve. The X5.4 class flare occurred at 00:04 UT and peaked at 00:24 UT. (b) Composite of coronagraph image obtained with \textit{SOHO}/LASCO C2 and 193 Å passband image by \textit{SDO}/AIA. We can see the shock front surrounding the CME ejecta. (c) \textit{SDO}/AIA 193 Å band image showing the emission from 1 MK coronal plasma at 00:18:19 UT. The pixels around flaring active region AR 11429 are saturated. The black rectangle \textquotedblleft R1\textquotedblright shows the field of view (FOV) of Figures 3(a)–(h). (d) Difference image made from two successive snapshots of the AIA 193 Å passband taken at 00:18:19 UT and 00:18:07 UT. A dome-like disturbance expanding above AR 11429 is clearly seen.
\[ \rho \frac{\partial \mathbf{V}}{\partial t} = -\rho \mathbf{V} \cdot \nabla\mathbf{V} - \nabla p + \mathbf{J} \times \mathbf{B} \]  
\[ \frac{\partial p}{\partial t} = -\nabla \cdot (\rho \mathbf{V} p) - (\gamma - 1) (\rho \nabla \cdot \mathbf{V} - \eta_0 \mathbf{J}^2) \]  
\[ \frac{\partial \mathbf{B}}{\partial t} - \nabla \times (\mathbf{V} \times \mathbf{B} - \eta_0 \mathbf{J}) = -\nabla \psi \]  
\[ \frac{\partial \psi}{\partial t} + \frac{c_s^2}{c_p^2} \nabla \cdot \mathbf{B} + \frac{c_s^2}{c_p^2} \psi = 0, \]  

where \( \rho, p, \mathbf{B}, \) and \( \mathbf{V} \) are mass density, gas pressure, magnetic field, and velocity, respectively. \( \mathbf{J} = \nabla \times \mathbf{B} \) is the electrical current density. \( \eta_0 \) is uniform electrical resistivity, and \( \gamma = 5/3 \) is the specific heat ratio. In the induction equation, the additional variable \( \psi \) is introduced in order to remove the numerical \( \nabla \cdot \mathbf{B} \) as proposed by Dedner et al. (2002). The numerical scheme we used is the Harten–Lax–van Leer–Discontinuities approximate Riemann solver (Miyoshi & Kusano 2005) with a second-order total variation diminishing Monotonic Upstream-Centered Scheme for Conservation Laws and second-order Runge–Kutta time integration.

3.2. Initial Conditions

We studied eight simulation cases of coronal shock–prominence interaction, with different density structures of the prominence for each run. In runs \( A_1, A_2, A_3, \) and \( A_4, \) the prominence is a uniformly dense spherical plasma of radius \( R_p. \) Initially, the center of the spherical prominence was at \((x, y, z) = (-R_p, 0, 0). \) In this model, \( R_p = 3 \) is used. The whole prominence region is where \( r_p = \sqrt{(x + \bar{R}_p)^2 + y^2 + z^2} < R_p. \) We modeled the prominence as a sphere of plasma instead of a cylindrical plasma in this paper in order to also study the effect of magnetic tension force induced by shock injection, which is revealed to be important in prominence activation, as discussed in Section 3.4. When the timescale of prominence activation is much smaller than that of LAPO (as in the case of the event discussed in Section 2), the global magnetic field structure will not play a significant role in the dynamics of the prominence activation process. In that case we can separate the physics that govern prominence activation from that governing the ensuing LAPO. In order to focus only on the physics of prominence activation, we think of such a situation in this study. We set the initial magnetic field to be uniform, neglecting the effect of global loop curvature and gravity. As a result, gas pressure also becomes uniform owing to total pressure balance between corona and prominence. The coronal shock front was initially at \( x = -2R_p. \) The fast-mode shock waves propagate in the positive \( x \)-direction with density compression ratio \( r_{sh, cor} = 1.37. \) The initial magnetic field is in the \( z \)-direction, so the injected fast-mode shock is a perpendicular one. We note that, in reality, oblique components may play roles in prominence activation, especially in the plane perpendicular to the shock propagation direction. We do not take into account such oblique dynamics induced by oblique shocks and focus only on the dynamics of activated prominence in the direction of shock propagation, and we discuss the effect of prominence mass on the activation dynamics.

In runs \( B_1, B_2, B_3, \) and \( B_4, \) the prominence consists of 200 randomly sized spherical clumps of spherical shape put inside the whole sphere of radius \( R_p. \) (Figures 4(a) and (b)). From the comparison of runs \( B_i \) with runs \( A_i, \) we study the role of the internal density structure of the prominence in its bulk acceleration and excitation of internal chaotic flow, and we also see how well uniform prominence approximation works. In reality, we only solved the numerical box with positive \( y \) and \( z \) coordinates in order to reduce the numerical cost (i.e., the motions of only \( \sim 50 \) clumps are simulated). From here, we express the density distribution of prominence as \( \rho_p \) for convenience.

Figure 2. Radio burst associated with the X5.4 flare on 2012 March 7. We can see a clear linear structure showing the characteristic signature of a Type II burst that drifted from \( f = 112 \) MHz at 00:17:10 UT (indicated as \( P_1 \)) to \( f = 88 \) MHz at 00:17:38 UT (indicated as \( P_2 \)). A Type IV burst is also seen, as indicated in the figure.
Figure 3. Time evolution of prominence activation by coronal shock waves seen in the AIA 193 and 304 Å bands. Panels (a) and (b) show the prominence just before the coronal shock front reached the prominence. After the coronal shock waves arrive at the prominence (panels (c) and (d)), they propagate further, accelerating the prominence in the direction of propagation (panels (e) and (f)). During acceleration, the prominence strongly brightens in the AIA 304 Å images. Panels (g) and (h) show the prominence with its maximum displacement from the original position seen in (a) and (b). (i) and (j) show time–distance plots along cut AB shown in (a) from 00:20 UT to 01:30 UT. The propagating shock front is seen as a bright linear feature. When the shock front hits the prominence (which is seen as the dark trajectory in panel (i)), it is suddenly accelerated. After the shock has passed, the prominence threads move in a somewhat disordered manner, as seen in panel (j).
In runs $B_3$, clumps are distributed so that the average mass density within the prominence region ($r_p < R_p$) becomes $\rho_{p,t} = \rho_{\text{clump}} + (1 - f_i) \rho_{\text{cor}}$, where $\rho_{\text{cor}}$ is the mass density of the background (shock upstream) corona and $\rho_{\text{clump}} = 100 \rho_{\text{cor}}$ is the mass density of each clump. The volume filling factors $f_i$ are set to be $(f_1, f_2, f_3, f_4) = (0.05, 0.1, 0.3, 1)$. In runs $A_3$, the mass density of the prominence is uniformly set to be $\rho_p = f_i \rho_{\text{clump}} + (1 - f_i) \rho_{\text{cor}}$ i.e., both $A_3$ and $B_3$ prominences have the same average density over the entire volume; the only difference is that $A_3$ prominences are uniform but those in $B_3$ are made of clumps). The equational form of the initial condition is as follows:

$$
\rho = \begin{cases} 
\rho_{\text{sh,cor}} \rho_{\text{cor}} & (x < -2R_p) \\
\rho_{\text{cor}} & (-2R_p < x, r_p > R_p) \\
\rho_p & (-2R_p < x, r_p < R_p) 
\end{cases}
$$

$$
p = \begin{cases} 
R_{\text{sh,cor}} p_0 & (x < -2R_p) \\
p_0 & (-2R_p < x) 
\end{cases}
$$

$$
B_z = \begin{cases} 
R_{\text{sh,cor}} B_0 & (x < -2R_p) \\
B_0 & (-2R_p < x) 
\end{cases}
$$

$$
V_x = \begin{cases} 
V_{\text{sh,cor}} & (x < -2R_p) \\
0 & (-2R_p < x) 
\end{cases}
$$

(6) (7) (8) (9) (10) (11)

with $\rho_{\text{cor}} = 1$, $p_0 = 1$, $B_0 = \sqrt{8 \pi p_0 / \beta}$. The variables $r_{\text{sh,cor}}$, $R_{\text{sh,cor}}$, and $V_{\text{sh,cor}}$ in the above equations are density jump (=compression ratio), pressure jump, and plasma velocity of coronal shock waves, respectively. Plasma $\beta$ is assumed to be $\beta = 0.2$ to model the low-beta corona, where the Lorentz force dominates the gas pressure gradient force in accelerating the prominence. The electrical resistivity $\eta_0$ is set to be $3.0 \times 10^{-4}$ in all cases, in order to prevent numerical instability.

The unit of speed in our simulation is $\sqrt{p_0 / \rho_0} = 1$. The corresponding sonic, Alfvénic, and fast-mode wave phase speeds in the corona are respectively expressed as

$$
C_{s,c} = \sqrt{\frac{\rho_0}{\rho_0}} = 1.29
$$

$$
C_{A,c} = \frac{B_0}{\sqrt{4 \pi p_0}} = \sqrt{\frac{2}{\gamma \beta}} = 3.16
$$

$$
C_{f,c} = \sqrt{C_{s,c}^2 + C_{A,c}^2} = \sqrt{1 + \frac{2}{\gamma \beta}} = 3.42.
$$

(12) (13) (14) (15) (16)

From MHD Rankine–Hugoniot relations for perpendicular shocks, sonic and fast-mode Mach numbers $M_{s,c}$ and $M_{f,c}$ are
respectively expressed by \( r_{sh,cor} \), \( \gamma \), \( \beta \) as
\[
M_{s,c} = \sqrt{\frac{2r_{sh,cor}((2-\gamma)r_{sh,cor} + \gamma r_{sh,cor}(\beta+1))}{\beta\gamma((\gamma+1)-r_{sh,cor}(\gamma-1))}} = 3.39
\]
(17)
\[
M_{f,c} = \sqrt{\frac{2r_{sh,cor}((2-\gamma)r_{sh,cor} + \gamma(\beta+1))}{((\gamma+1)-r_{sh,cor}(\gamma-1))(\gamma\beta+2)}} = 1.28.
\]
(18)

From MHD Rankine–Hugoniot relations, the pressure jump \( R_{sh,cor} \) and the plasma velocity of the shocked corona \( V_{sh,cor} \) are expressed respectively as follows:
\[
R_{sh,cor} = \gamma M_{s,c}^2 \left( 1 - \frac{1}{r_{sh,cor}} \right) - \frac{2 r_{sh,cor} - 1}{\beta} + 1 = 1.80
\]
(19)
\[
V_{sh,cor} = (1 - 1/r_{sh,cor})M_{f,c}C_{f,c} = 1.18.
\]
(20)

The simulation box is \( x \in [-60, 60], y \in [0, 60], z \in [0, 60] \) discretized with nonuniformly arranged \( n_x \times n_y \times n_z = 800 \times 400 \times 400 \) grid points. Especially, the inner region of \( x \in [-7.5, 7.5], y \in [0, 4.5], z \in [0, 4.5] \) is discretized with uniformly set \( n_x \times n_y \times n_z = 400 \times 200 \times 200 \) grid points. We applied reflective boundary conditions on \( y = 0 \) and \( z = 0 \) planes, while for the other boundaries we applied free boundary conditions. We use sparse grids in outer space so that we can neglect unwanted numerical effects on prominence dynamics from outer free boundaries.

3.3. Momentum Transport from a Coronal Shock Wave to a Prominence

Figures 4(a) and (c) show the density distribution of two different times \( (t = 0 \) and \( t = 10, \) respectively) in simulation run \( A_3 \). The prominence hit by the coronal fast-mode shock is compressed and starts to move in the direction of shock propagation (positive x-direction). The fast-mode shock front

![Figure 5: Density, vorticity, and total pressures in runs A3 at t = 10. Arrows represent velocity fields in the plane of the plot. Logarithmic density distributions in the XY- or XZ-plane are shown with color contours in panels (a) and (b), respectively. In panels (c) and (d), vorticity components perpendicular to the XY- or XZ-plane are shown with color contours, respectively. In panel (e), total plasma pressure (magnetic pressure + gas pressure) in the XY-plane is shown. Panel (f) is the schematic figure showing how Kelvin–Helmholtz instability evolves along the velocity shear layer in the XY-plane. The resultant formation of a low-pressure region behind the cloud is also indicated.](image-url)
transmitted into the prominence material experiences multiple reflections at the prominence–corona boundary. Figure 5(c) shows the distribution of the \(z\)-component of vorticity \((\nabla \times V)_z\) in the \(XY\)-plane. In Figure 5(e), we find sharp velocity shear at the prominence–corona boundary. The velocity shear results in the development of KH instability in the \(XY\)-plane (Figures 5(a) and (f)). The magnetic tension force suppresses the velocity shear in the \(XZ\)-plane (Figure 5(d)). Associated with the KH vortex behind the cloud, the low-pressure region is formed behind the cloud, helping the prominence to be accelerated in the \(x\)-direction (Figures 5(e) and (f)). Magnetic tension forces also help stretch the cloud in the \(x\)-direction. In Figure 4, the time evolution of the density structure in simulation runs A3 and B3 is shown with color contours in logarithmic scale, together with magnetic field vectors in the plane of the plots shown by white arrows. The cloud in run A3 is deformed by KH instability that has developed behind the cloud. In run B3, the time needed for each clump to be deformed by shear flows is much shorter than in run A3 because of the small length scale of each clumps. In runs B3, the shocked clumps interact with each other through the flow field around them, making the overall flow and density structure more complicated compared with runs A3 in a short timescale. The time evolution of density structure in runs A1 and B3 are compared in Figure 6 together with the magnetic field.

Figure 7 shows the time evolution of the center-of-mass velocity \(V_p\) of the prominence in runs A1 and B1. The center-of-mass velocity of prominence \(V_p\) is defined as follows:

\[
V_p = \frac{\int_{\rho > \rho_{\text{threshold}}} \rho V_x \, dx \, dy \, dz}{\int_{\rho > \rho_{\text{threshold}}} \rho \, dx \, dy \, dz},
\]

where \(\rho_{\text{threshold}}\) is the threshold mass density and set to be 2 in all simulation runs. We regard the plasma with density \(\rho > \rho_{\text{threshold}}\) as prominence in this analysis. The volume integral is done over the region where the mass density is larger than \(\rho_{\text{threshold}}\). Practically, we first flag the region where \(\rho\) is larger than \(\rho_{\text{threshold}}\) in the whole computational box, and then we sum up the quantities within the flagged region. We tried various values of \(\rho_{\text{threshold}}\) and found no significant change in the analysis results. It takes more time in runs A1 and A2 for \(V_p\) to approach \(V_{\text{sh,cor}}\) than in runs B1 and B2, although the time evolution of \(V_p\) in runs A3 and A4 is very similar to that in runs B3 and B4.

### 3.4. Phenomenological Model of Prominence Activation

Here, we focus on the mechanism of the momentum transfer from the coronal shock wave to the prominence. Takahashi et al. (2015) discussed the momentum transfer mechanism of shock injection into the prominence and estimated the resultant prominence velocity \(V_{p,ss}\) with the help of one-dimensional linear theory as

\[
V_{p,ss} = \frac{2}{1 + \sqrt{\chi}} V_{\text{sh,cor}},
\]

where \(\chi = \rho_p/\rho_{\text{cor}}\) is the ratio of the mass density between corona and prominence.

Multidimensionality and nonlinearity effects on prominence activation that are neglected in the 1D model are important, especially in quantitative discussion. In this section, we make a phenomenological model of prominence activation taking into consideration the effect of fluid drag force and magnetic tension force, and we compare it with the result of 3D MHD simulations of prominence activation. The schematic figure of the phenomenological model and corresponding simulation snapshots are shown in Figure 9. We first construct the phenomenological model in this section and then compare the predicted prominence center-of-mass motion with the numerical simulation results. In our case, the center of mass of the prominence moves parallel to the direction of shock propagation (\(x\)-direction) because the coronal shock that activates the spherical prominences is a perpendicular one. In this section, we focus on how the prominence center of mass is accelerated in the \(x\)-direction as a result of the interaction with coronal perpendicular shock waves. We define the timescale in which the injected shock front sweeps the prominence material as the “shock-sweeping” timescale, \(\tau_{ss} \approx 2R_p/C_{f,p}\), where \(C_{f,p} = C_{f,c}/\sqrt{\chi}\) is the fast-mode phase speed in the prominence. This is similar to the “cloud-crushing” timescale, \(\tau_{cc}\), discussed in Klein et al. (1994), during which a molecular cloud is crushed by strong shock waves in the interstellar medium. This shock-sweeping mechanism will accelerate the prominence to have a velocity of \(V_{p,ss}\) according to 1D linear theory. During the shock-sweeping process (\(t < \tau_{ss}\)), the mean acceleration of the prominence through shock-sweeping \(\alpha_{ss}\) is approximated as

\[
\alpha_{ss} \approx \frac{V_{p,ss}}{\tau_{ss}} \approx \frac{A_{ss} V_{\text{sh,cor}} C_{f,c}}{\sqrt{\chi (1 + \sqrt{\chi})} R_p}.
\]

\(A_{ss}\) is an ad hoc factor of order unity introduced to take into account multidimensional effects such as shock refraction and is set to be \(A_{ss} = 0.5\) throughout the paper. The difference between the velocity of shocked coronal plasma \(V_{\text{sh,cor}}\) and the prominence itself \(V_p(t)\) causes the pressure difference between the front and back side of the prominence, which accelerates the prominence to the direction of shock propagation (fluid drag force). The prominence acceleration by the fluid drag force \(\alpha_{\text{drag}}(t)\) is approximated as

\[
\alpha_{\text{drag}}(t) \approx -\frac{C_d S_{\text{sh,cor}} \rho_{\text{cor}} S_p (V_p(t))}{M_p} - \frac{V_{\text{sh,cor}}}{V_p(t) - V_{\text{sh,cor}}},
\]

where \(S_p \approx \pi R_p^2\) and \(M_p \approx \frac{\pi}{3} \rho_p R_p^3\) are the cross-sectional area and total mass of the prominence assuming the spherical shape of the prominence, respectively. \(C_d\) is a drag coefficient. In the following discussion, we approximate the drag coefficient by unity, i.e., \(C_d \approx 1\). It is revealed in the following discussion that the drag acceleration \(\alpha_{\text{drag}}\) is much smaller than the other two acceleration mechanisms in weak shock acceleration, which is the case in prominence activation by coronal shocks. Substituting these into Equation (24), we get

\[
\alpha_{\text{drag}}(t) \approx -\frac{3S_{\text{sh,cor}}}{4\sqrt{\chi}} \tilde{V}_p(t)[\tilde{V}_p(t)],
\]

where \(\tilde{V}_p(t) = V_p(t) - V_{\text{sh,cor}}\) is the prominence center-of-mass speed relative to the ambient shocked corona. Also, the velocity difference between the prominence and ambient corona distorts the magnetic field lines penetrating the prominence. This curved magnetic field accelerates the prominence by the magnetic tension force. The equation of motion of the prominence accelerated by the \(x\)-component of
the magnetic tension force is approximately written as
\[ M_p \alpha_{\text{tension}}(t) = \frac{2}{4\pi} S_p'(t) \tan \theta, \]
with \( \alpha_{\text{tension}} \) being the acceleration by the magnetic tension mechanism. Here, \( S_p'(t) \) is the effective cross-sectional area of the part of the prominence in the \( XY \)-plane through which shocked coronal magnetic field lines penetrate, \( \tan \theta = B_x/B_z \) is the inclination of distorted magnetic field lines (Figure 9).

Assuming that the inclination is determined by the balance between displacement of magnetic field lines by velocity difference \( -V_p(t)/\delta t \) and its relaxation by Alfvén wave propagation \( C_A \delta t \), \( \tan \theta \) is approximated as
\[ \tan \theta \approx -\frac{V_p(t)}{C_A}, \]
where \( C_A \approx n_{\text{sh,cor}} B_0 \sqrt{4\pi n_{\text{sh,cor}} m_p} \) is the Alfvén wave phase speed in the shocked coronal plasma. Here, we approximate \( S_p'(t) \) as follows:
\[
S_p'(t) \approx \begin{cases} 
S_p \frac{t}{\tau_{\text{sp}}} & (0 < t < \tau_{\text{sp}}) \\
S_p & (\tau_{\text{sp}} < t)
\end{cases}
\]
where the shock passage timescale \( \tau_{\text{sp}} \approx 2R_p/C_{F,c} = \tau_{\text{sh}}/\sqrt{\beta} \) is the timescale in which the shock front passes over the prominence. Then \( \alpha_{\text{tension}} \) is finally approximated as
\[ \alpha_{\text{tension}}(t) \approx \begin{cases} 
\frac{3x^{3/2} V_p(t) C_{A,\infty} t}{2 \gamma_{\text{sp}}} & (0 < t < \tau_{\text{sp}}) \\
\frac{3x^{3/2} V_p(t) C_{A,\infty}}{2 \gamma_{\text{sp}}} & (\tau_{\text{sp}} < t)
\end{cases}
\]

The order-of-magnitude ratio of maximum accelerations for each mechanism is roughly \( \alpha_{\text{drag}} : \alpha_{\text{tension}} : \alpha_{\text{drag}} \sim 1 : 1 : (1 - 1/\beta_{\text{sh,cor}}) M_{F,c} \), if \( \beta \) is much smaller than unity and the shock is not strong. When the coronal shock is weak, \( \alpha_{\text{drag}} \) is negligible compared with the other two. After the shock front
has swept the prominence, the main acceleration mechanism is magnetic tension force acceleration.

Summarizing the above discussion, the prominence acceleration in the phenomenological model \( \alpha_{\text{ph}}(t) \) is written as

\[
\alpha_{\text{ph}}(t) \approx \begin{cases} 
\alpha_{\text{ss}} + \alpha_{\text{drag}}(t) + \alpha_{\text{tension}}(t) & (0 < t < \tau_{\text{ss}}) \\
\alpha_{\text{drag}}(t) + \alpha_{\text{tension}}(t) & (\tau_{\text{ss}} < t),
\end{cases}
\]

or explicitly, \( \alpha(t) \) is expressed as follows:

\[
\alpha_{\text{ph}}(t) \approx \begin{cases} 
A_{\text{ss}} \frac{V_{\text{sh,cor}} \nabla \cdot \mathbf{E}_{\text{cr}}}{4 \xi_{p} r_{p}} - \frac{3A_{\text{ss}} V_{\text{sh,cor}} \nabla \cdot \mathbf{E}_{\text{cor},l}}{4 \xi_{p} r_{p}} & (0 < t < \tau_{\text{ss}}) \\
A_{\text{ss}} \frac{V_{\text{sh,cor}} \nabla \cdot \mathbf{E}_{\text{cr},l}}{4 \xi_{p} r_{p}} + \frac{3A_{\text{ss}} V_{\text{sh,cor}} \nabla \cdot \mathbf{E}_{\text{cor},l}}{4 \xi_{p} r_{p}} & (\tau_{\text{ss}} < t < \tau_{\text{sh}}).
\end{cases}
\]

The time evolution of the prominence velocity in the phenomenological model \( V_{\text{ph}}(t) \) is obtained by solving the prominence equation of motion,

\[
\frac{dV_{\text{ph}}(t)}{dt} = \alpha_{\text{ph}}(t).
\]

Substituting the simulation parameters for run \( A_{i} \) into the above equation of motion, we get the resultant time evolution of phenomenological prominence velocity. The time evolutions of prominence velocity in runs \( A_{i} \) and models \( A_{i} \) are compared in Figure 8.

On the other hand, the fluid equation of motion is written as

\[
\rho \frac{dV}{dt} = -\nabla p + \frac{B^2}{8\pi} - \frac{(\mathbf{B} \cdot \nabla) \mathbf{B}}{4\pi}.
\]

If we assume that the prominence mass \( M_{p} = \int_{V > V_{\text{threshold}}} \rho \, dV \) is constant, the volume integral of the above equation of motion over the prominence volume leads to the following relations:

\[
\alpha(t) = \alpha_{\text{pp}}(t) + \alpha_{\text{mt}}(t),
\]

with \( \alpha(t) \), \( \alpha_{\text{pp}}(t) \), and \( \alpha_{\text{mt}}(t) \) being the acceleration of the center of mass of the prominence in the x-direction, prominence acceleration by total pressure gradient force, and prominence acceleration of magnetic tension force, respectively. \( \alpha(t) \), \( \alpha_{\text{pp}}(t) \), and \( \alpha_{\text{mt}}(t) \) are respectively written (and calculated from simulation results) as follows:

\[
\alpha(t) = \frac{1}{M_{p}} \int_{V > V_{\text{threshold}}} \frac{dV}{dt} \, dxdydz,
\]

\[
\alpha_{\text{pp}}(t) = -\frac{1}{M_{p}} \int_{V > V_{\text{threshold}}} \frac{\partial}{\partial t} \left( p + \frac{B^2}{8\pi} \right) \, dxdydz,
\]

\[
\alpha_{\text{mt}}(t) = -\frac{1}{M_{p}} \int_{V > V_{\text{threshold}}} \left( \frac{\mathbf{B} \cdot \nabla \mathbf{B}}{4\pi} \right) \, dxdydz.
\]

In Figures 10(a)–(h), we show the total pressure gradient acceleration \( \alpha_{\text{pp}} \) and magnetic tension acceleration \( \alpha_{\text{mt}} \) with dashed and solid lines, respectively, in MHD simulation runs \( A_{i} \) and \( B_{i} \). Figures 10(i)–(l) show \( \alpha_{\text{ss}} + \alpha_{\text{drag}}(t) \) and \( \alpha_{\text{tension}} \) in models \( A_{i} \). A sudden drop in total pressure gradient force acceleration in, for example, run \( A_{2} \) (Figure 10(c)) corresponds to the first reflection of the wave that has swept the whole prominence (Figure 9(e)). The oscillations of accelerations in runs \( A_{i} \) are the results of multiple wave reflections at the prominence–corona boundary (Figure 9(f)), and the oscillation period is approximately \( \tau_{\text{ss}} \) for each run. In runs \( B_{1} \) and \( B_{2} \), the total pressure gradient force acceleration is smaller than that in runs \( A_{1} \) and \( A_{2} \). That is one reason why it took more time for the prominence velocity \( V_{p} \) in runs \( A_{1} \) and \( A_{2} \) to approach the velocity of shocked corona \( V_{\text{sh,cor}} \) than in runs \( B_{1} \) and \( B_{2} \). After the shock-sweeping time \( \tau_{\text{sh}} \), \( \alpha_{\text{ss}} + \alpha_{\text{drag}}(t) \) in the phenomenological models \( A_{i} \) are smaller than pressure gradient force acceleration \( \alpha_{\text{pp}} \) in runs \( A_{i} \). Also, the magnetic tension force acceleration \( \alpha_{\text{tension}}(t) \) in the phenomenological models \( A_{3} \) and
A_4 underestimates the actual simulation results \( \alpha_{\text{mt}} \) in runs \( A_3 \) and \( A_4 \). Roughly, the time evolution of \( \alpha_{\text{op}} \) and \( \alpha_{\text{mt}} \) in runs \( A_i \) resembles that of \( \alpha_{\text{op}} + \alpha_{\text{drag}}(t) \) and \( \alpha_{\text{tension}} \) in models \( A_i \), respectively. This is natural because \( \alpha_{\text{op}} \) in models \( A_i \) represents the acceleration by fast-mode shock transmission into the prominence, which is mainly due to the total pressure gradient. Also, \( \alpha_{\text{drag}} \) in models \( A_i \) is mainly due to a pressure gradient that lasts longer than \( \tau_{\text{ss}} \) resulting from velocity difference between the corona and prominence. Due to velocity shear at the prominence–corona boundary in the \( XY \)-plane, KH instability develops and a low-pressure region is formed behind the cloud associated with induced vortices (Figures 5(a), (c), and (e)). Also, we see a high-pressure region ahead of the cloud due to flow collision. This type of flow structure results in the large-scale total pressure gradient that accelerates the prominence in the \( x \)-direction as a drag force. Formation of induced vortices and the associated low-density region behind the cloud is clearer in the hydrodynamic case (simulation run \( C_3 \)) discussed in Section 3.6 (Figure 13).

3.5. Excitation and Damping of Internal Flow

While the coronal shock wave gives its bulk momentum to the prominence, internal flow is also excited within the prominence. Internal flow excitation during interstellar shock–cloud interaction has been studied widely as a measure of internal turbulence. The internal flow is discussed using velocity dispersion in \( x, y \), and \( z \) components,

\[
\delta V_x = \sqrt{\langle V_x^2 \rangle - \langle V_x \rangle^2} \quad (38)
\]
\[
\delta V_y = \sqrt{\langle V_y^2 \rangle - \langle V_y \rangle^2} \quad (39)
\]
\[
\delta V_z = \sqrt{\langle V_z^2 \rangle - \langle V_z \rangle^2} \quad (40)
\]

where the mass-weighted average of physical quantity \( f \) is calculated as

\[
\langle f \rangle = \frac{\int_{\rho > \rho_{\text{threshold}}} \rho f \, dx dy dz}{\int_{\rho > \rho_{\text{threshold}}} \rho \, dx dy dz}. \quad (41)
\]
Figure 10. Time evolution of prominence accelerations. In the top two rows, total pressure gradient force acceleration $a_{tp}$ and magnetic tension force acceleration $a_{mt}$ in simulation runs $A_i$ and $B_i$ are shown as dashed and solid lines, respectively. In the bottom row, the summation of shock-sweeping acceleration and fluid drag acceleration $a_{ss} + a_{drag}$ and magnetic tension force acceleration $a_{tension}$ in models $A_i$ are shown as dashed and solid lines, respectively.

Figure 11. Time evolution of velocity dispersion $\delta V_x$ (solid lines), $\delta V_y$ (dashed lines), and $\delta V_z$ (dot-dashed lines) in runs $A_i$ and $B_i$. Exponential damping timescales are estimated for the three variables in each run by least-squares fitting during the time between $t = 10$ and $t = 20$ in runs $A_i$ and $B_i$ ($i$ is from 1 to 3). The resultant slopes for exponential damping are shown as thick lines for each variable. Estimated exponential damping timescales $\tau_{dx}$, $\tau_{dy}$, and $\tau_{dz}$ are listed in Table 1.
We assumed that $\langle V_1 \rangle$ and $\langle V_2 \rangle$ should be 0 because of the symmetry around the x-axis.

Figure 11 shows the time evolution of velocity dispersion $\delta V_x$, $\delta V_y$, and $\delta V_z$ in runs $A_1$ and $B_1$ in semi-log graphs. They show the excitation of the internal flow through the passage of shock and its damping roughly in an exponential manner. In runs $A_1$, we also find oscillatory behavior of $\delta V_x$ with their period of roughly 0.5$\tau_{ss}$. We estimated the damping times (e-folding times) $\tau_d$ for $\delta V_x$, $\delta V_y$, and $\delta V_z$ in runs $A_1$ and $B_1$ based on least-squares fitting by the bisector method (Isobe et al. 1990) between $t = 10$ and $t = 20$. The resultant slopes for exponential damping are shown as thick lines for each variable in Figure 11. We did not estimate them in runs $A_4$ and $B_4$ because the simulation time span did not cover their damping phase. The estimated damping times are shown in Table 1. The estimated damping times fall in the range between $\tau_{ss}$ and 2$\tau_{ss}$ in almost all the cases. $\tau_{d,xy}$ is roughly twice as large as $\tau_{d,vy}$ and $\tau_{d,xz}$.

We show the time evolution of the center-of-mass position of the prominence $\langle x \rangle$ in Figure 12. Uniform mesh is used between $X = \pm 7.5$, shown as dashed lines in Figure 12. In all eight simulation runs, the prominence acceleration phase is mostly solved within the uniform mesh region. In Figure 12, we have no significant change of damping behavior when the cloud center of mass crossed the $X = 7.5$ plane.

### 3.6. Comparison with Hydrodynamic Simulation

In this section, we study hydrodynamic simulation of shock–cloud interaction (run $C_3$) and compare it with MHD simulation run $A_3$. By directly comparing run $C_3$ with run $A_3$, we discuss hydrodynamic effects that were suppressed in MHD run $A_3$ by the presence of magnetic field. Such hydrodynamic effects are expected to be more significant in the interstellar shock–cloud interaction. In practice, we modified the initial plasma beta in run $A_3$ to $1.0 \times 10^{10}$ and use it as an initial condition of the MHD simulation. The pressure jump $R_{sh,cor}$, plasma velocity of shocked corona $V_{sh,cor}$, and shock Mach number $M_{sc}$ are $R_{sh,cor} = 1.70$, $V_{sh,cor} = 0.44$, and $M_{sc} = 1.25$, respectively. The minimum plasma beta that appeared throughout the simulation was $2.7 \times 10^4$, which is still much larger than unity. The magnitude of plasma beta is of the order of the ratio of the pressure gradient force term to the Lorentz force term in the MHD equation of motion. Because plasma beta is very large ($>10^5$) throughout the simulation, the effect of magnetic field on the dynamics in simulation run $C_3$ is negligible. So, we call run $C_3$ the “hydrodynamic” simulation hereafter.

| Run  | $f_c^a$ | $\bar{p}_p^b$ | $\tau_{ss}^c$ | $\tau_{d,xy}^d$ | $\tau_{d,vy}^e$ | $\tau_{d,xz}^f$ | $\tau_{d,xy}/\tau_{ss}^g$ | $\tau_{d,vy}/\tau_{ss}^h$ | $\tau_{d,xz}/\tau_{ss}^i$ |
|------|---------|--------------|---------------|----------------|----------------|----------------|--------------------------|--------------------------|--------------------------|
| Run $A_1$ | 0.05 | 5.95 | 4.28 | 4.28 | 4.57 | 9.14 | 1.00 | 1.07 | 2.14 |
| Run $A_2$ | 0.10 | 10.9 | 5.79 | 8.06 | 7.19 | 7.71 | 1.39 | 1.24 | 1.33 |
| Run $A_3$ | 0.30 | 30.7 | 9.72 | 12.4 | 14.8 | 26.6 | 1.28 | 1.52 | 2.73 |
| Run $A_4$ | 1.00 | 100. | 17.5 | ... | ... | ... | ... | ... | ... |
| Run $B_1$ | 0.05 | 5.95 | 4.28 | 5.31 | 3.30 | 8.67 | 1.24 | 0.77 | 2.03 |
| Run $B_2$ | 0.10 | 10.9 | 5.79 | 6.60 | 5.32 | 12.0 | 1.14 | 0.92 | 2.07 |
| Run $B_3$ | 0.30 | 30.7 | 9.72 | 16.0 | 10.0 | 12.9 | 1.65 | 1.02 | 1.32 |
| Run $B_4$ | 1.00 | 100. | 17.5 | ... | ... | ... | ... | ... | ... |
| Run $D_1$ | 0.30 | 30.7 | 9.72 | 9.94 | 9.94 | 8.54 | 1.02 | 0.88 | 1.06 |

Notes.

$^a$ The volume filling factor of the prominence.

$^b$ Volume-averaged density of the prominence.

$^c$ Shock-sweeping timescale.

$^d$ Damping timescale for $\delta V_x$.

$^e$ Damping timescale for $\delta V_y$.

$^f$ Damping timescale for $\delta V_z$.

$^g$ Damping timescale for $\delta V_x$ in units of shock-sweeping timescale.

$^h$ Damping timescale for $\delta V_y$ in units of shock-sweeping timescale.

$^i$ Damping timescale for $\delta V_z$ in units of shock-sweeping timescale.

![Figure 12. Time evolution of x-coordinate of prominence center of mass in runs $A_i$ (solid lines) and runs $B_i$ (dot-dashed lines). The boundary between the uniform and nonuniform numerical grid is located at $X = \pm 7.5$ and expressed as dashed lines.](image-url)
Figure 13. (a) Volume rendering of density distribution in simulation run $C_3$ at $t = 60$. The values in the color bars correspond to logarithm of density. Panels (b) and (c) show the density (in logarithmic scale) and pressure distributions in the $y = 0$ plane in simulation run $C_3$ at $t = 60$, respectively. The arrows in panels (b) and (c) show the velocity field in the $y = 0$ plane.

Figure 13(a) shows the volume rendering of the density distribution in run $C_3$ at $t = 60 (=2.3\tau_{ss})$. After being crushed by the injected shock wave, the cloud in run $C_3$ continues to be deformed. This is quite different from the time evolution of the cloud shape in MHD run $A_3$. In run $A_3$, magnetic field suppresses the cloud deformation and damps the internal flow rapidly in the timescale of $\sim \tau_{ss}$. Without such a magnetic suppression of cloud deformation, the time evolution of the cloud and the ambient flow structure in run $C_3$ are very different from those in run $A_3$. One characteristic flow structure in run $C_3$ is the formation of a large coherent vortex flow behind the cloud. The ambient flow converging toward the cloud along the $x$-axis and diverging in the $YZ$-plane associated with the vortex helps flatten and stretch the cloud further in run $C_3$ (Figures 13(b) and (c)).

In Figures 14(a)–(d), snapshots of density and pressure distribution in the $XZ$-plane at $t = 0.31\tau_{ss}$ in simulation runs $A_3$ and $C_3$ are shown, respectively. The velocity amplitudes of transmitted shock front $V_{sf1}$ and that of the shock front from behind the cloud $V_{sf2}$ in simulation runs $A_3$ and $C_3$ are shown in Figures 14(e) and (f), respectively. Shock front 2 appears when the injected shock waves have passed over the cloud at $t \approx \tau_{sp}$, where $\tau_{sp} = 2R_p/C_{T,c}$ is the shock passage timescale. The ratio of $V_{sf2}$ to $V_{sf1}$ in MHD cases is much smaller than unity. For example, $V_{sf2}/V_{sf1}$ is 0.29 at $t = 3 = 0.31\tau_{ss} = 1.7\tau_{sp}$ in run $A_3$ (Figure 14(e)). In the hydrodynamic case, on the other hand, $V_{sf2}$ is similar to $V_{sf1}$. $V_{sf2}/V_{sf1}$ is 0.70 at $t = 8 (=0.31\tau_{ss} = 1.7\tau_{sp})$ in run $C_3$ (Figure 14(f)). In the hydrodynamic case, the shock-sweeping acceleration mechanism that is driven by the propagation of shock front 1 in Figure 14(f) is almost canceled out by propagation of shock front 2 after $t = \tau_{sp}$.

Taking the above effect into account, the phenomenological model of cloud acceleration in the hydrodynamic case is
modified from model $A_i$ to be as follows:

$$
\alpha_{ph}(t) \approx \begin{cases} 
A_p \frac{v_{ih, cor} v_{ih, cor} - v_p(t)^2}{2 \sqrt{1 + v_p(t)^2}} + \frac{3v_{sh, cor} (v_{sh, cor} - v_p(t))^2}{4xR_p} & (0 < t < \tau_{sp}) \\
\frac{3v_{sh, cor} (v_{sh, cor} - v_p(t))^2}{4xR_p} & (\tau_{sp} < t). 
\end{cases}
$$

(42)

The time evolution of prominence velocity $V_p$ in simulation run $C_3$ is shown as a solid line in Figure 15(a), together with that of the phenomenological model $C_3$ discussed above, shown as a dashed line. Model $C_3$ captures abrupt acceleration characteristics before $t = \tau_{sp}$, but underestimates the acceleration in the later phase. In model $C_3$, the acceleration decreases because relative velocity $v_{sh, cor} - v_p$ becomes smaller with time, but in run $C_3$, the acceleration is almost constant after $t = \tau_{sp}$. This is partly because of the deformation (flattening and stretching) of the prominence in the hydrodynamic case. If we denote the cross-sectional area of the cloud in the $YZ$-plane as $S_{yc}(t)$, the magnitude of aerodynamic drag force that mainly accelerates the cloud in run $C_3$ is of the order of $\rho_{sh, cor} v_p(t)^2 S_{yc}(t)$. In model $C_3$, we assumed that the prominence has a spherical shape, although in run $C_3$, the cloud is flattened and stretched in the $YZ$-plane. We study $\langle r^2 \rangle = \langle x^2 + z^2 \rangle$ as an effective cross-sectional area of the prominence in the $YZ$-plane where aerodynamic drag works to accelerate the prominence in the $x$-direction. Figure 15(c) shows the time evolution of $\langle r^2 \rangle$. $\langle r^2 \rangle$ at $t = 90$ is almost 5 times that of initial $\langle r^2 \rangle$. We note that in run $C_3$ the drag acceleration mechanism works much more than in model $C_3$ because $S_{yc} \sim \langle r^2 \rangle$ evolves in time, as discussed above (Figure 15(c)).

The internal flow structure is also different in the hydrodynamic case compared with the MHD case. Figure 15(d) shows the time evolution of velocity dispersions $\delta V_x = \sqrt{\langle V_x^2 \rangle - \langle V_x \rangle^2}$ and $\delta V_y = \sqrt{\langle V_y^2 \rangle}$ as a proxy of internal flow within the prominence. In run $C_3$, $\delta V_x$ and $\delta V_y$ increase with time after the shock passage. This is different from the results in run $A_3$, where the velocity dispersions damp in an exponential manner owing to the presence of magnetic field.

### 3.7. Prominence Activation by Triangular Wave Packet

Coronal shock waves that activate prominences in reality are not blast waves, but wave packets with finite width. The motion of activated prominence depends not only on plasma velocity of the shock but also on wave packet width. In this section, we analyze 3D MHD simulation results that reproduced prominence activation by triangular wave packets and compare with the phenomenological model. The density profile of the prominence in this simulation (let us call it run $D_3$ hereafter) is the same as that in simulation run $A_3$. In simulation run $D_3$, we have a triangular wave packet with velocity amplitude $V_{sh, cor}$ and wave packet width $w$ that activate prominence. The initial plasma velocity distribution for simulation run $D_3$ is as follows (solid line in Figure 16(a)):

$$
V_x = \begin{cases} 
2V_{sh, cor} \frac{x + (2R_p + 2w)}{2w} & (-2R_p - 2w < x < -2R_p) \\
0 & (x < -2R_p - 2w, -2R_p < x). 
\end{cases}
$$

(43)

with the wave packet width being $w = 10$ (Figure 16(a)). The density, pressure, and magnetic field in the corona are all uniform, initially (solid lines in Figures 16(b)–(d)).
The initial condition is a “superposition” of two triangular fast-mode wave packets propagating in opposite directions to each other with velocity amplitude $V_{sh, cor}$ and wave packet width $w$. If the wave packets were linear ones (i.e., $v_C = f_C$), the wave propagating in the positive $x$-direction would keep its velocity amplitude and wave packet width unchanged without any interaction with the oppositely directed wave packet. In simulation run $D_3$, the wave packet propagating in the positive $x$-direction interacts with the prominence.

We checked by nonlinear 1D MHD numerical simulation (without prominence) how initially superposed wave packets (shown as solid lines in Figure 16) evolve in time. Dashed and dot-dashed lines in Figure 16 denote plasma parameter distribution at times $t = 2$ and $t = 6$, respectively, in the 1D simulation (Figure 16). We see in Figure 16(a) that the initial single peak of superposed wave packets in plasma velocity split into two oppositely directed wave packets, as expected from linear theory. Because of the nonlinearity, however, the velocity amplitude decays slowly and the wave packet broadens as they propagate.

We make the phenomenological model that describes the prominence center-of-mass motion in simulation run $D_3$. We call it model $D_3$. Model $D_3$ can be obtained simply by replacing $V_{sh, cor}$ and $V_p(t) = V_p(t) - V_{sh, cor}$ in model $A_3$ with $V_{wp, cor}(t)$ and $V_p(t) = V_p(t) - V_{wp, cor}(t)$, respectively. $V_{wp, cor}(t)$ is the plasma velocity in the corona around the prominence shocked by the triangular wave packet and is approximated as follows:

$$V_{wp, cor}(t) = \begin{cases} V_{sh, cor} \frac{2\tau_w - t}{2\tau_w} & (0 < t < 2\tau_w) \\ 0 & (2\tau_w < t) \end{cases}$$

with $\tau_w = w/C_{f,c}$ being the “wave packet passage” timescale. $V_{wp}$ is an approximation based on the fact that the injected coronal shock is weak and the activated prominence moves much slower than the coronal fast-mode phase speed. The solid and dashed lines in Figure 17(a) show prominence center-of-mass acceleration by magnetic tension force $\alpha_{mt}$ and that by total pressure gradient force $\alpha_{tp}$ in simulation run $D_3$, respectively. The solid and dashed lines in Figure 17(b) show prominence center-of-mass acceleration by magnetic tension force $\alpha_{mt}$ and that by total pressure gradient force $\alpha_{tp}$ in simulation run $D_3$, respectively.
mechanism $\alpha_{\text{tension}}$ and that by both shock-sweeping and fluid drag mechanisms $\alpha_{\text{ss}} + \alpha_{\text{drag}}$, respectively. We find from the plot in Figure 17(a) that the prominence is mainly accelerated by magnetic tension force first and then decelerated also by magnetic tension force. Model $D_3$ captures such a characteristic response of the prominence to wave packet injection. We find some oscillations in both $\alpha_{\text{ms}}$ and $\alpha_{\text{ss}}$ in simulation run $D_3$. They result from multiple reflections of wave packets within the prominence. The effect of multiple reflections is not included in model $D_3$.

Solid and dashed lines in Figure 18(a) show time evolution of prominence center-of-mass position $X_p$ in simulation run $D_3$ and that expected from model $D_3$, respectively. Solid and dashed lines in Figure 18(b) show the prominence center-of-mass speed $V_p$ in simulation run $D_3$ and that expected from model $D_3$, respectively. From Figure 18, we see that the center-of-mass motion of activated prominence expected with model $D_3$ quantitatively agreed with those in simulation run $D_3$. We call the time interval during which the prominence is accelerated to its maximum speed by the shock wave the “acceleration phase” of the prominence activation. The acceleration phase continued until $t \approx 4$ in run $D_3$ (Figure 18(b)). Solid, dashed, and dot-dashed lines in Figure 19 show time evolution of velocity dispersions $\delta V_x$, $\delta V_y$, and $\delta V_z$ in simulation run $D_3$, respectively. Exponential damping time-scales are estimated for the three velocity dispersions by least-squares fitting by the bisector method during the time between $t = 10$ and $t = 20$. The resultant slopes for exponential damping are shown as thick lines for each component. Damping times $\tau_{dx}$, $\tau_{dy}$, and $\tau_{dz}$, are listed in the last row in Table 1. We find that the damping times for all three components are similar to the shock-sweeping timescale $\tau_{\text{ss}}$. We note that $\tau_{dz}$ in simulation run $D_3$ is much smaller than that in run $A_3$.

4. Coronal Shock and Prominence Diagnostics Using Prominence Activation

In this section, we try to diagnose coronal shock properties and prominence properties using prominence activation with the help of the phenomenological model discussed above. Figure 20(a) is a time–distance diagram of prominence activation made from AIA 193 Å passband images. This corresponds to the white rectangle in Figure 3(i). The white crosses denote prominence positions during the prominence activation estimated by eye. We denote the prominence displacement in the plane of the sky at time $t = t_i$ ($i = 1$–14) as $L_{dx,i}$. On the other hand, the position of the prominence activated by the triangular wave packet can be predicted by the phenomenological model discussed in the previous section.

Here, we try to fit the observed time evolution of the activated prominence position with the phenomenological model expectation. I used Equation (31) modified with $V_{wp,cor}$ as described in Section 3.7 for the fitting. We assume the coronal temperature to be $T = 10^6$ K. This leads to the coronal
sound speed being \( c_{t,c} \approx 1.8 \times 10^2 \text{km} \text{s}^{-1} \), assuming the specific heat ratio to be \( \gamma = 5/3 \). The local density gap between corona and prominence is assumed to be \( \chi = 100 \). We think of two different values for the angle \( \phi \) between the line-of-sight and shock propagation direction, which are \( \phi = 0^\circ \) and \( 45^\circ \). The thickness of the prominence core seen as the dark structure in AIA 193 Å images is about 10\(^9\), which corresponds to the estimated prominence radius of \( R_p \approx 3.6 \times 10^3 \text{ km} \). The fast-mode wave speed in the corona \( C_{f,c} \) is estimated to be \( C_{f,c} \approx C_{w}\cos \phi \), with \( C_{w} = 380 \text{ km} \text{s}^{-1} \) being the wave propagation speed in the plane of the sky. The coronal fast-mode wave speeds in the \( \phi = 0^\circ \) and \( 45^\circ \) cases are \( 3.8 \times 10^2 \text{ km} \text{s}^{-1} \) and \( 5.4 \times 10^2 \text{ km} \text{s}^{-1} \), respectively. The plasma beta is obtained with \( \beta = (2/\gamma)(C_{t,c}/C_{A,c})^2 \), where \( C_{A,c} = \sqrt{C_{f,c}^2 - C_{t,c}^2} \) is the coronal Alfvén speed in the perpendicular propagating case. In the \( \phi = 0^\circ \) and \( \phi = 45^\circ \) cases, plasma beta is calculated to be \( \beta = 0.33 \) and \( \beta = 0.15 \), respectively. Assuming that the shock propagation direction and the center-of-mass velocity of activated prominence are parallel (which is correct in the perpendicular shock case), the displacements of the activated prominence at time \( t - t_F \) are estimated to be \( X_{p,i} = L_{p,i} \cos \phi \). Then, the number of remaining free parameters of the phenomenological model is three: prominence volume filling factor \( f_V \), compression ratio of coronal shock wave \( \gamma \), and wave packet width \( w \). We denote the time evolution of the activated prominence position expected from the phenomenological model as \( X_{p,\text{model}}(f_V, \gamma, w; t) \). We searched for best-fit values for \( f_V, \gamma \), and \( w \) with which the sum of squared residuals \( \text{SSR}(f_V, \gamma, w) = \Sigma(X_{p,\text{model}}(f_V, \gamma, w; t) - X_{p,i})^2 \) is minimized in the parameter space \( f_V \in [0.01, 1.0], \gamma = [1.07, 1.9], \) and \( w = [0.02R_p, 0.4R_p] \), with \( R_p \) being a solar radius. The fitting results for the \( \phi = 0^\circ \) and \( \phi = 45^\circ \) cases are shown in Figures 20(b) and (c).

As best-fit parameters, we estimate \( \gamma \) and \( w \) of the coronal shock wave to be \( 1.17R_p \) and \( 0.16R_p \) in the \( \phi = 0^\circ \) case and \( 1.17R_p \) and \( 0.22R_p \) in the \( \phi = 45^\circ \) case, with \( R_p \) being the solar radius. Best-fit parameters for the \( \phi = 0^\circ \) and \( \phi = 45^\circ \) cases are listed in Table 2. Fast-mode Mach numbers \( M_f \) of the coronal shock in the \( \phi = 0^\circ \) and \( \phi = 45^\circ \) cases are 1.2 and 1.13, respectively. From mass conservation at the shock front, the velocity amplitude of the injected triangular wave is expressed as \( V_{sh,\text{cor}} = M_fC_{f,c}(1 - 1/r) \). From this, the plasma velocity amplitude for the injected triangular wave is estimated as \( 62 \) and \( 88 \text{ km} \text{s}^{-1} \) in the \( \phi = 0^\circ \) and \( \phi = 45^\circ \) cases, respectively. The best-fit \( f_V \) in both cases were 0.01, which is the smallest value in the free parameter space.

Based on the phenomenological model (Equation (31)), on the other hand, when \( f_V \) is much smaller than unity, the prominence is accelerated to its maximum speed almost within the shock-sweeping timescale \( \tau_{sw} \), while the subsequent prominence deceleration occurs within the wave packet passage timescale \( \tau_{wp} \). With the estimated parameters of \( f_V = 0.01, \quad R_p = 3.6 \times 10^3 \text{ km}, \quad C_{f,c} \approx 500 \text{km} \text{s}^{-1}, \quad \text{and} \quad w = 0.2R_p \), the acceleration and deceleration timescales are roughly \( \tau_{sw} \approx 20 \text{ s} \) and \( \tau_{wp} \approx 300 \text{ s} \), respectively. In order to estimate \( f_V \) correctly, we have to time-resolve the acceleration phase whose timescale \( \sim \tau_{sw} \) reflects \( f_V \) directly. Compared with the estimated acceleration timescale of \( \sim 20 \text{ s} \), the AIA time cadence of \( 12 \text{ s} \) is not high enough to track the acceleration phase of the prominence activation in this event, although we can track the subsequent deceleration phase with sufficient time resolution. As seen in the time–distance plot of Figure 20(a), the prominence appears to be accelerated to its maximum speed right after the arrival of the shock. We think that the evaluated value of \( f_V \) in this analysis is not accurate enough, due to the lack of fully time-resolved observation of the acceleration phase of the prominence activation in this event.

Then, we estimate the energy of the coronal shock wave \( E_{sh,\text{cor}} \) associated with the X5.4 flare. The coronal emission measure near the prominence before the arrival of the shock was \( EM \approx 1 \times 10^{27} \text{ cm}^{-3} \). The emission measure is calculated at point A in Figure 3(a), based on a method proposed in Cheung et al. (2015). Assuming the line-of-sight distance \( d \) to be of order of the distance to the (solar) horizon as seen from a point above the photosphere by a pressure scale height \((h = 50 \text{ Mm})\), we get \( d \approx \sqrt{h(2R_p + h)} \approx 2 \times 10^4 \text{ Mm} \). The coronal proton number density \( n \) at point A is estimated from the relation \( EM \approx n^2d \) to be \( n \approx 2 \times 10^8 \text{ cm}^{-3} \). The corresponding coronal mass density is \( \rho \approx mn_p \approx 3 \times 10^{-16} \text{ g cm}^{-3} \), with \( m_p \approx 1.7 \times 10^{-24} \text{ g} \) being the proton mass. The energy flux of the shock at point A is \( E_{sh,\text{cor}} \approx \rho V_{sh,\text{cor}}^2 C_{f,c} \approx 1 \times 10^6 \text{ erg cm}^{-2} \text{s}^{-1} \). The timescale for the triangular wave to pass through the fixed point A is \( \tau \approx w/C_{f,c} \approx 3 \times 10^2 \text{ s} \). The surface area of the spherically...
expanding dome of the shock front in the corona is approximated as $S \sim 2\pi L^2$, with $L$ being the distance between flaring AR and point A. We approximate $L \sim R_c$ and get $S \sim 3 \times 10^{22} \text{cm}^2$. From above, the energy budget of the coronal shock wave is estimated as $E_{FS,\text{sh,cor}} \sim 10^{31} \text{erg}$. The total energy released in the X5.4 class flare that occurred on 2012 March 7 was associated with very fast CME, with its speeds of about 2700 km s$^{-1}$ estimated from coronagraph observation by SOHO/LASCO. A global shock front is formed around the expanding CME ejecta. The shock front had a dome-like form, especially with the bright structure propagating to the north at the foot of the dome. The northward disturbance hit a polar prominence, leading to the excitation of LAPO. During the prominence activation, the prominence was strongly brightened, receiving momentum in the direction of shock propagation.

In order to explain the observational signature of prominence activation, we have done a three-dimensional MHD simulation of the coronal fast-mode shock–prominence interaction. Especially, the momentum transfer mechanism from the shock to the prominence is studied in detail. The shock injection into the prominence material compressed and accelerated the prominence. The velocity shear at the corona–prominence boundary resulted in KH instability. KH instability was stabilized by magnetic tension force in the plane containing the initial magnetic field lines. By analyzing the simulation results and comparing them with phenomenological models, magnetic tension force acceleration was also found to be very important. The accelerated prominence velocity asymptotes to the value of coronal shocked plasma velocity when the shock is a blast wave after some timescale depending on different prominence density. When the volume filling factor is small like in runs $B_1$ and $B_2$, the acceleration timescale is longer compared with the case with uniformly distributed prominence density in runs $A_1$ and $A_2$. This may be because of the suppression of the shock-sweeping acceleration mechanism in “clumpy” clouds in runs $B_i$. Both the total surface area of clumps and the local density gap between the clump and corona in runs $B_i$ are larger than those in runs $A_i$. This makes it difficult for injected shock fronts in runs $B_i$ to penetrate deep into the cloud as a whole so that they could exchange momentum with cloud materials. When the volume filling factors are larger than 0.3, the resultant time evolution of the mean velocity in runs $B_3$ and $B_4$ is very similar to that of runs $A_3$ and $A_4$. 

5. Summary and Discussion

Recent high time and spatial resolution EUV observations of the solar corona by SDO/AIA enabled us to study in detail the time evolution of coronal shock waves associated with flares. We can now also study the interaction between coronal shock waves and prominences using AIA. In this paper, we studied the excitation process of LAPO through the interaction between the prominence and coronal shock wave, with the help of three-dimensional MHD simulation.
The Astrophysical Journal, 836:178 (21pp), 2017 February 20

Takahashi

Figure 20. Application of the phenomenological model to observations. Extracted prominence positions from the time–distance diagram of prominence activation are shown as white crosses in panel (a). Panels (b) and (c) show the fitting results by the phenomenological model of triangular wave–prominence interaction in the cases of \( \phi = 0^\circ \) and \( \phi = 45^\circ \), respectively. Black crosses in panels (b) and (c) are the extracted position from observation, and the solid lines are the best-fit curves from the phenomenological model. The fitting result is listed in Table 2.

We also studied the time evolution of the velocity dispersion of shocked prominence material in each \((x, y, \text{and} z)\) component. The velocity dispersion is excited during the shock sweeps through the cloud and then damps almost exponentially. The exponential damping timescales of velocity dispersions in each component \( \tau_{d,x}\), \( \tau_{d,y} \), and \( \tau_{d,z} \) are estimated and summarized in Table 1. In almost all the simulation runs, \( \tau_{d,x} \) and \( \tau_{d,y} \) are roughly comparable to the shock-sweeping timescale \( \tau_{st} \), while \( \tau_{d,z} \) is about twice as large as \( \tau_{st} \). A possible reason for the discrepancy of the damping times among the components is as follows. When the randomized flow is directed to the positive \( x \)-direction at a certain time and location, the magnetic field lines originally directed to the \( z \)-direction \( \left( B_z \right) \) will be distorted there in the \( AX \) plane, resulting in the electric current directed in the negative \( y \)-direction \( \left( -J_y \right) \). The Lorentz force \( \sim -J_y \times B \), acts in the negative \( x \)-direction, which pulls back the flow originally directed to the positive \( x \)-direction. This helps damp the velocity dispersion in the \( x \)-direction, making \( \tau_{d,x} \) small. The same damping mechanism works on the \( y \)-component of velocity dispersion as well, but not on the \( z \)-component.

In interplanetary shock–cloud interaction, both the shock-sweeping mechanism and fluid drag force accelerate the cloud. KH instability is also important in mixing MC materials, which will affect the star formation process taking place. In solar coronal shock–prominence interaction, magnetic tension force is more important in accelerating the prominence than fluid drag force, because in prominence activation, plasma beta is typically smaller than unity and the shock is not strong. When the prominence has internal density structure like in runs \( B_i \), plasma mixing might also occur in a short timescale.

When the plasma beta is much larger than unity (which is a reasonable assumption in some molecular clouds), the cloud acceleration and internal flow excitation show very different characteristics. The shock-sweeping acceleration mechanism is effective only during the shock passage time \( \tau_{sp} \), and the pressure gradient force due to the velocity difference between the cloud and the ambient plasma works as an accelerator. The cloud is flattened with the help of ambient flow that converges toward the cloud along the \( x \)-axis and diverges in the \( YZ \)-plane associated with a coherent vortex formed behind the cloud. The flattening effect increases the effective cross section of the cloud, helping the cloud acceleration by fluid drag. The excited internal flow does not decay in hydrodynamic simulation run \( C_3 \), though in MHD simulation runs \( A_3 \) the internal flow damps in an exponential manner, mainly due to the Lorentz force.

In reality, the coronal shock wave that activate a prominence is not a blast wave as studied in runs \( A_3 \) or \( B_i \) but a wave packet with finite wave packet width. We studied the interaction between a coronal shock wave in the form of a triangular-shaped wave packet and a prominence in simulation run \( D_3 \). The shocked prominence is first accelerated and then decelerated by magnetic tension and total pressure gradient force in simulation run \( D_3 \). The phenomenological model \( D_3 \) well captures the characteristic dynamics of the prominence center of mass in both acceleration and deceleration phases, but slightly underestimates the impact of pressure gradient force. Especially, phenomenological model \( D_3 \) does not reproduce the prominence deceleration by total pressure gradient force that is present in run \( D_3 \). One of the possible reasons for the discrepancy is that the phenomenological model neglects the effect of multiple reflection of transmitted shock waves within the prominence. We compared the prominence center-of-mass position and velocity in simulation run \( D_3 \) with those expected by phenomenological model \( D_3 \) and found quantitative agreement between the simulation and the model.

| Table 2 |
|---|
| Estimated Coronal Shock Parameters |
| \( \beta^a \) | \( f_b^b \) | \( r^c \) | \( w/R_m^d \) | \( M_f^e \) | \( V_{\text{dis}, \text{cor}}^f \) |
| \( \phi = 0^\circ \) | 0.33 | 0.01 | 1.17 | 0.16 | 1.12 | 62 km s\(^{-1}\) |
| \( \phi = 45^\circ \) | 0.15 | 0.01 | 1.17 | 0.22 | 1.13 | 88 km s\(^{-1}\) |

Notes.

\(^a\) Plasma beta in the shock upstream corona.

\(^b\) Volume filling factor of the prominence.

\(^c\) Compression ratio of the coronal shock.

\(^d\) Wave packet width in the corona in units of a solar radius.

\(^e\) Fast-mode Mach number of the coronal shock wave.

\(^f\) Plasma velocity amplitude of the coronal shock wave.
We tracked the time evolution of the position of the activated prominence and fitted it with the phenomenological model. The best-fit curve for prominence movement agreed well with observation. As best-fit parameters, we obtained prominence volume filling factor $f_V$, coronal shock compression ratio $r$, and wave packet width $w$. The resultant compression ratio and fast-mode Mach numbers of the coronal shock were $1.17$ and $1.12$ in the $\phi = 0^\circ$ case and $1.17$ and $1.13$ in the $\phi = 45^\circ$ case, respectively. The estimated wave packet widths of the coronal shock were $0.16R_s$ and $0.22R_s$ in the cases of $\phi = 0^\circ$ and $\phi = 45^\circ$, respectively. They are comparable to typical EUV wave front widths of ~$100$ Mm, which are suggestive of coronal shock waves reported in Muhr et al. (2014). Both the estimated coronal shock propagation speed and the plasma velocity amplitude of $380$–$540$ km s$^{-1}$ and $68$–$88$ km s$^{-1}$ are reasonable values as a weak fast-mode coronal shock wave.

The shock wave that activated the prominence had likely been driven by the lateral expansion of CME ejecta in the lower corona (whose speed is much smaller than the radial ejection speed) and had propagated a considerable distance of ~$1R_s$ from the source active region. This is a possible reason why the coronal shock near the prominence was weak, although the associated CME was extremely fast, with its speed of almost ~$3000$ km s$^{-1}$ at a height of $2R_s$ estimated with SOHO/LASCO coronagraph observations. The best-fit value of the prominence filling factor, on the other hand, was $f_V = 0.01$, which we do not think is accurate partly because the observational data did not time-resolve the acceleration phase of activated prominence, which is vital for determining prominence $f_V$ based on our model. With the help of emission measure analysis, we estimated the energy of the coronal shock to be $E_{\text{shock}} \sim 10^{31}$ erg in this event. This was roughly several percent of the total released energy during the X5.4 flare.

The physical mechanism that mainly works in prominence activation differs with different wave packet widths. If the width of the wave packet is longer than $C_{f, e} \tau_{ms}$, magnetic tension force acceleration after Alfvén travel time $\tau_s = \frac{L}{c_s}$, where $L$ is the length of the magnetic loop. We note that $\tau_s$ is of order of the period of ensuing LAPO.

In this paper, we discussed the dynamics of the prominence–coronal shock interaction that leads to LAPO. The prominence activation event we studied was triggered by the arrival of coronal shock waves at a polar prominence that propagated from a faraway AR corona. Such globally propagating flare-associated shock waves in the corona or LAPOs are relatively rare phenomena. On the other hand, small-scale magnetic explosions (small flares and jets) always occur in the corona. We expect that the interactions between solar prominences and small-amplitude shock waves generated by such small-scale magnetic explosions are always occurring in the corona and might play a role in driving small-amplitude prominence dynamics such as small-amplitude oscillations and chaotic movements of plasma elements.

The SDO/AIA data are courtesy of NASA/SDO and the AIA science team. The simulation code used in this work is created with the help of the HPCI Strategic Program. Numerical computations were carried out on Cray XC30 at the Center for Computational Astrophysics, National Astronomical Observatory of Japan. The author is grateful to Dr. Kazunari Shibata and Dr. Ayumi Asai of Kyoto University, and Dr. Yuki Kubo of the National Institute of Information and Communications Technology for their helpful comments and discussions. The author is also grateful to the journal referee for his/her careful reading and lots of comments that significantly improved the paper. This work is financially supported by the Grant-in-Aid for JSPS Fellows 15J02548 and by JSPS KAKENHI Grant Numbers 16H03955.

References

Antolini, P., Okamoto, T. J., De Pontieu, B., et al. 2015, ApJ, 809, 72
Asai, A., & Ishii, T. 2011, ApJL, 745, L18
Berger, T., Testa, P., Hillier, A., et al. 2011, Natur, 472, 197
Berger, T. E., Shine, R. A., Slater, G. L., et al. 2008, ApJL, 676, L89
Chen, M. C. M., Boerner, P., Schnierj, C. J., et al. 2015, ApJ, 807, 143
Dedner, A., Kemm, F., Kröner, D., et al. 2002, JCoPh, 175, 645
Dobashi, K., Matsumoto, T., Shimoikura, T., et al. 2014, ApJ, 797, 58
Emslie, A. G., Dennis, B. R., Shih, A. Y., et al. 2012, ApJ, 759, 71
Gilbert, H. R., Doau, A. G., Young, D., Tripathi, D., & Alexander, D. 2008, ApJ, 685, 629
Gopalswamy, N., Niita, N., Akiyama, S., Mikela, P., & Yashiro, S. 2012, ApJ, 744, 72
Grechnev, V. V., Uralov, A. M., Chertok, I. M., et al. 2011, SoPh, 273, 433
Hillier, A., Isobe, H., Shibata, K., & Berger, T. 2012, ApJL, 756, 110
Hillier, A., Morton, R. J., & Erdelyi, R. 2013, ApJL, 779, L16
Inoue, T., & Inutsuka, S.-i. 2012, ApJ, 759, 35
Isobe, H., Tripathi, D., Asai, A., & Jain, R. 2007, SoPh, 246, 89
Isobe, T., Feigelson, E. D., Akritas, M. G., & Babu, G. J. 1990, ApJ, 364, 104
Kai, K. 1970, SoPh, 11, 310
Klein, R. I., McKee, C. F., & Colella, P. 1994, ApJ, 420, 213
Kondo, T., Isobe, T., Igi, S., Watari, S., & Tokumaru, M. 1995, J. Commun. Res. Lab., 42, 111
Kosugi, T., Matsuzaki, K., Sakao, T., et al. 2007, SoPh, 243, 3
Kretzschmar, M., de Wit, T. D., Schmutz, W., et al. 2010, NatPh, 6, 690
Labrosse, N., Heinzel, P., Vial, J.-C., et al. 2010, SSRV, 151, 243
Lemen, J. R., Title, A. M., Dlin, J. D., et al. 2012, SoPh, 275, 17
Liu, W., Nitta, N. V., Schnierj, C. J., Title, A. M., & Tarbell, T. D. 2010, ApJ, 723, L53
Ma, S., Raymond, J. C., Golub, L., et al. 2011, ApJ, 738, 160
Mac Low, M.-M., & Zahnle, K. 1994, ApJL, 434, L33
Mackay, D. H., Karpen, J. T., Ballester, J. L., Schmieder, B., & Aulanier, G. 2010, SSRV, 151, 333
Mann, G., Jansen, F., MacDowall, R. J., Kaiser, M. L., & Stone, R. G. 1999, A&A, 348, 614
Matsumoto, T., Dobashi, K., & Shimoikura, T. 2015, ApJ, 801, 77
Miyoshi, T., & Kusano, K. 2005, JCoPh, 208, 315
Moreton, G. E. 1960, AJ, 65, 494
Muhr, N., Veronig, A. M., Kienreich, I. W., et al. 2014, SoPh, 289, 456
Narukage, N., Hudson, H. S., Morimoto, T., et al. 2002, ApJL, 572, L109
Newkirk, G. A. 1961, ApJ, 133, 983
Nittmann, J., Falle, S. A. E. G., & Gaskell, P. H. 1982, MNRAAS, 201, 833
Ofman, L., Knizhnik, K., Kucera, T., & Schneider, B. 2015, ApJL, 813, 124
Okamoto, T. J., Antolini, P., De Pontieu, B., et al. 2015, ApJ, 809, 71
Pataudie, D. J., & Fresen, R. A. 2005, ApJL, 633, 240
Pesnell, W. D., Thompson, B. J., & Chamberlin, P. C. 2012, SoPh, 275, 3
Pholudenko, A. Y., Frank, A., & Blackman, E. G. 2002, ApJ, 576, 832
Shibata, K., & Magara, T. 2011, LRSIP, 8, 6
Shin, M.-S., Stone, J. M., & Snyder, G. F. 2008, ApJ, 680, 336
Takahashi, T., Asai, A., & Shibata, K. 2015, ApJL, 801, 37
Thompson, B. J., Reynolds, B., Aurass, H., et al. 2000, SoPh, 193, 161
Title, A. & AIA team 2006, BAAS, 38, 261
Tsuneta, S., Ichimoto, K., Katsukawa, Y., et al. 2008, SoPh, 249, 167
Uchida, Y. 1968, SoPh, 4, 30
Vrsnak, B., Warmuth, A., Brajša, R., & Hulsmeier, A. 2002, A&A, 394, 299
Wang, C., Richardson, J. D., & Burlaga, L. 2001, SoPh, 204, 413
Warmuth, A., Vrsnak, B., Aurass, H., & Hulsmeier, A. 2001, ApJL, 560, L105
Woodward, P. R. 1976, ApJ, 207, 484