KINETIC SIMULATION OF SLOW MAGNETOSONIC WAVES AND QUASI-PERIODIC UPFLOWS IN THE SOLAR CORONA

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ABSTRACT

Quasi-periodic disturbances of emission-line parameters are frequently observed in the corona. These disturbances propagate upward along the magnetic field with speeds of ~100 km s−1. This phenomenon has been interpreted as evidence of the propagation of slow magnetosonic waves or has been argued to be a signature of intermittent outflows superposed on the background plasmas. Here we aim to present a new “wave + flow” model to interpret these observations. In our scenario, the oscillatory motion is a slow-mode wave, and the flow is associated with a beam created by the wave–particle interaction owing to Landau resonance. With the help of a kinetic model, we simulate the propagation of slow-mode waves and the generation of beam flows. We find that weak periodic beam flows can be generated by Landau resonance in the solar corona, and the phase with the strongest blueward asymmetry is ahead of that with the strongest blueshift by about 1/4 period. We also find that the slow wave damps to the level of 1/e after the transit time of two wave periods, owing to Landau damping and Coulomb collisions in our simulation. This damping timescale is similar to that resulting from thermal conduction in the MHD regime. The beam flow is weakened/attenuated with increasing wave period and decreasing wave amplitude since Coulomb collisions become more and more dominant over the wave action. We suggest that this “wave + flow” kinetic model provides an alternative explanation for the observed quasi-periodic propagating perturbations in various parameters in the solar corona.

Key words: Sun: corona – Sun: oscillations – waves

1. INTRODUCTION

Quasi-periodic upward propagating intensity disturbances (PIDs) are frequently observed along the magnetic field structures in the solar corona. These disturbances were observed in polar coronal plumes (Ofman et al. 1997; DeForest & Gurman 1998; Banerjee et al. 2000) and in the coronal fan loops at the edges of active regions (Berghmans & Clette 1999; De Moortel et al. 2000). The parameters of the PIDs are summarized by De Moortel (2009) as follows: an oscillation period of 145–550 s, a propagation speed of 45–205 km s−1, a relative intensity amplitude of 0.7%–14.6%, and detected wavelengths of 2.9–23.2 Mm. These PIDs are often accompanied by Doppler shifts in the spectral lines. The phenomenon of PIDs had been almost universally interpreted as the propagation of slow magnetosonic waves (Nakariakov et al. 2000). The slow-mode waves may play an important role in the heating of the chromosphere, the generation of solar spicules, and the development of coronal loops (Hollweg et al. 1982; Shibata et al. 1982; Porter et al. 1994; Kumar et al. 2006; Yang et al. 2013; Liu et al. 2014). The slow-mode waves in the corona may also be connected with (driven by) the spicule/jet flows in the chromosphere (Jiao et al. 2015; Samanta et al. 2015). Another important piece of evidence supporting the slow-mode wave scenario comes from the good correlation between intensity and Doppler-shift variations as derived from the spectroscopic observations by the Extreme-ultraviolet Imaging Spectrometer onboard Hinode (Wang et al. 2009).

However, this interpretation has been challenged by some authors (De Pontieu & McIntosh 2010; Tian et al. 2011), who claimed that the quasi-periodic “red minus blue” (R–B) asymmetries found in the spectral lines of the intensity disturbance region are signatures of quasi-periodic upflows. Hence, the debate on whether the slow wave or the intermittent outflow corresponds to the real nature of the disturbances was initiated and continues. The slow wave and the flow are thought to be related to different dynamic processes. The intermittent outflows inferred at the edges of the active regions are thought to be the possible source of the solar wind (Sakao et al. 2007; Harra et al. 2008; Harra et al. 2008; McIntosh & De Pontieu 2009; He et al. 2010; Warren et al. 2011). Ascertain the nature of the intensity disturbances is crucial for clarifying the physical processes in the dynamic solar atmosphere.

The excess line width enhancement beyond the pure-wave model may be explained by the superposition of waves on the background uncoupled plasmas, although the resultant frequency of the line width oscillation may be twice the original one (Verwichte et al. 2010; Wang et al. 2012). A dual MHD scenario/model has also been proposed to launch the slow mode waves as excited by the quasi-periodic flows at the footpoints of the magnetic flux tubes (Nishizuka & Harra 2011; Ofman et al. 2012; Wang et al. 2013). In this dual model, slow-mode waves are gradually decoupled from the upflows and propagate at a larger speed to higher altitudes than the latter. The PIDs in the coronal strands with extended length seem to be the signature of slow-mode waves rather than quasi-periodic flows in their model.
It seems that these two interpretations are incompatible in the MHD regime. Nevertheless, it is known that Landau damping of the slow wave can generate a beam and associated plasma flow. Consequently, the wave and flow may contemporaneously exist self-consistently, and the two diverse interpretations of the observed intensity disturbance can be compatible. To reconcile these two interpretations, we present a new scenario and call it the “wave + flow” kinetic scenario, which involves both the slow wave and the beam created by Landau resonance with the waves. We suggest that the observed R–B asymmetries in the spectral line may be the signal of beam related flows. We use kinetic simulation to test this scenario, and we reproduce the weak beam component in the ion velocity space. Consequently, the wave and their kinetic effects on the R–B asymmetry are at one footpoint of the loop. In their model, the propagating slow-mode waves and the flows are excited by a flare at one footpoint of the loop. In their model, the slow waves propagate in the loop, reflected by the footpoints of the loop, and interact with the flows. What is similar to our “wave + flow” scenario is that both waves and flows make contribution to the observed PID in their model. The difference between their model and this work lies in that the flows in their model are bulk jet flows rather than beam flow components, which contribute directly to the heat flux and thermal conduction in this work.

The kinetic simulation model is described in Section 2. In Section 3, the simulation setup and results are presented, including the introduction of slow-mode oscillation at the bottom of the simulation region and the corresponding response of the proton velocity distribution function (VDF) in terms of different moments (density, bulk velocity, and R–B asymmetry). The variations of the parameters of the slow mode wave and their kinetic effects on the R–B asymmetry of the proton VDF are investigated. The damping mechanism of the slow mode wave in our simulation is discussed as well at the end of this section. We conclude our paper with a discussion of the applicability of our “wave + flow” model in Section 4. The limitations of our work are also discussed in this section.

2. SIMULATION MODEL

In this section, we briefly describe the model we used in our kinetic simulation. This model was first introduced by Vocks & Marsch (2001). It is based on the Boltzmann equation,

\[ \frac{\partial f}{\partial t} + (\mathbf{v} \cdot \nabla)f + \left[ \mathbf{g} + \frac{q}{m} (\mathbf{E} + \mathbf{v} \times \mathbf{B}) \right] \cdot \nabla_{\mathbf{v}} f = \left( \frac{\partial f}{\partial t} \right)_{\text{Coul}}, \]

(1)

where \( \mathbf{B} \) is the magnetic field, \( \mathbf{E} \) is the electric field, \( q \) is the particle’s charge and \( m \) its mass, and \( g \) is the gravitational acceleration. The term on the right-hand side is the Fokker–Planck collision term to describe the Coulomb collisions. The form of the Fokker–Planck term is given by Equation (2):

\[ \left( \frac{\partial f}{\partial t} \right)_{\text{Coul}} = -\Gamma \cdot \frac{\partial f}{\partial \mathbf{v}} \cdot \mathbf{H} + \frac{1}{2} \Gamma \cdot \left( \frac{\partial^2 f}{\partial \mathbf{v} \cdot \mathbf{v}} \right). \]

(2)

For the concrete expressions of \( \Gamma, H, \) and \( G, \) see Boyd & Sanderson (2003, p. 544). The first term in Equation (2) produces a deceleration of a beam of particles, and the second term accounts for the spreading of a unidirectional beam in the velocity space (Boyd & Sanderson 2003, p. 544). In statistics, the Coulomb collisions tend to shape the distributions into the Maxwellian distribution. The VDF \( f(r, \mathbf{v}, t) \) depends on three velocity and three spatial coordinates, and on time. The computational cost is high to solve the Boltzmann equation in the complete six-dimensional phase space. To simplify, it is necessary to reduce the dimensions of \( f. \) Since the ion gyroradii is short compared to other characteristic timescales, it is reasonable to assume gyrotropy (Vocks & Marsch 2001). Therefore, the number of velocity coordinates can be reduced from three to two: \( \mathbf{v} \rightarrow (v_{\parallel}, v_{\perp}) \) (Vocks 2002). In this kinetic model, the ion velocity component parallel to the background magnetic field is of primary interest, and only waves propagating in the parallel direction are considered. Accordingly, a “reduced VDF” (Marsch 1998; Vocks & Marsch 2001; Vocks 2002; Vocks & Marsch 2002) is introduced in this model and obtained by integrating over the velocity component perpendicular to the background magnetic field:

\[ F_k(v_{\perp}) = 2\pi \int_{v_{\perp}}^{\infty} 2^{k+1} f(v_{\parallel}, v_{\perp}) dv_{\perp} \quad k = 0, 1, 2, \ldots . \]

(3)

A Boltzmann equation for the reduced VDFs can be derived by integrating over \( v_{\perp} \) in the same way as in the definition of \( F_k \) (Equation (3)):

\[ \frac{\partial F_k}{\partial t} + v_{\parallel} \frac{\partial F_k}{\partial s} + \left( \frac{q}{m} E_\parallel - g \cos \psi \right) \frac{\partial F_k}{\partial v_{\parallel}} + 1 \frac{1}{2} \frac{\partial A}{\partial s} \left( \frac{\partial F_{k+1}}{\partial v_{\parallel}} + 2v_{\parallel}(k + 1)F_k \right) - \frac{\partial F_k}{\partial t} \]

(4)

where \( E_\parallel(s) \) is the electric field component parallel to the magnetic field, \( A(s) \) is the cross-sectional area of the magnetic flux tube, and \( \psi(s) \) is the angle between the magnetic field and the direction normal to the solar surface. For the Fokker–Planck collision term on the right-hand side, the formulation of Ljepojevic & Burgess (1990) is used. Only the spatial coordinate \( s \) parallel to the magnetic field is considered in this model. Actually, there is no coupling between the reduced VDFs of different orders on the left-hand side of Equation (4), when \( A(s) \) is constant. However, there is still coupling between \( F_k \) and \( F_{k+1} \) in the Fokker–Planck term on the right-hand side of Equation (4). For the details of the coupling, see Equations (A16)–(A27) of Vocks (2002). The Boltzmann equation for the reduced VDFs \( F_k \) depends on \( F_{k+1} \), and therefore a cut-off is needed. The assumption

\[ F_k(v_{\parallel}) = k!(2v_{\parallel, \perp})^{k-1} F_0(v_{\parallel}), \]

(5)

is applied in this model. This assumption is exact for a Maxwellian in \( v_{\perp}, \) and it is easily satisfied in the collisional plasma.

Equation (4) for \( k = 0, 1 \) is used to determine the governing equations in this model. The particle density \( N, \) drift velocity \( v_{\parallel}, \) parallel temperature \( T_\parallel, \) and heat flux \( q_{\parallel} \) can be obtained from \( F_0. \) The perpendicular temperature \( T_\perp \) and the heat flux \( q_{\perp} \) can
be obtained from $F_1$. The Boltzmann equations for the reduced VDFs are only used to describe the ion kinetics. Electrons are dealt with in the fluid approximation. We note that the fluid description for electrons is just a rough approximation, since electrons are highly mobile. Electrons described with the VDFs are only used to describe the ion kinetics. Electrons are highly mobile. Electrons described with the equation for electrons is just a rough approximation, since the electron mass is quite small. Under the assumption that the electron mass is quite small. Under the approximation for electrons, we can deduce the form of $E_{\parallel}$:

$$E_{\parallel} = -\frac{1}{e \sum_j q_j N_j} \frac{\partial}{\partial s} \left( k_B T_e \sum_j q_j N_j \right).$$

3. SIMULATION SETUP AND RESULTS

3.1. Introducing the Slow-Mode Waves

The simulation results on the evolution and kinetic effect of the slow-mode waves propagating in a magnetic flux tube are presented in this section. We assume that the lower boundary of the magnetic flux tube is located in the low corona, and the flux tube is normal to the surface of the Sun and has a constant cross-sectional area. The model plasmas in the flux tube consist of protons and electrons. The gravity is considered in our simulation, with the gravitational acceleration $g = -274 \text{ m s}^{-2}$, and thus a gradient exists in the profile of the proton’s number density. At the beginning of the simulation, the number density of proton $N_0$ is set at the bottom boundary as $N_0 (h = 0 \text{ Mm}) = 1 \times 10^9 \text{ cm}^{-3}$ and at the top boundary as $N_0 (h = 100 \text{ Mm}) = 1.9 \times 10^7 \text{ cm}^{-3}$. The initial temperature of the plasma in the whole magnetic flux tube is $T_0 = 1 \text{ MK}$. The plasma is initially in an equilibrium state.

Magnetic field fluctuations, in particular in the longitudinal direction, will be introduced when the slow-mode wave is propagating oblique to the background magnetic field. There will be a good anti-correlation between the oscillations of magnetic field strength and the density fluctuations, in addition to the correlation between the density and the longitudinal velocity fluctuations, for the quasi-perpendicular propagating slow-mode waves (see He et al. 2015, and references therein). In this work, we just consider the parallel propagation of slow-mode waves, which do not have magnetic field perturbations. It would be helpful to diagnose the propagation direction of the wavefront if the PIDs observed in the non-uniform magnetized plasma environment if the wavefront can be recorded and tracked. In a traditional stratified corona, plasma beta reduces with height due to the increase of the Alfvén speed, the large spatial gradient of which may cause a partial reflection of low-frequency upward propagating Alfvén waves. The terminology of the slow-mode wave for the wave under investigation holds throughout the simulation domain of the lower solar corona, where plasma beta is always less than unity.

We introduce the slow wave into the flux tube by setting the parameters of the plasma at the lower boundary. Here we set:

$$N_b(t) = [1 + \epsilon \cos (2\pi t / \tau)] N_0(0),$$

$$\nu_b(t) = \epsilon \cos (2\pi t / \tau) \nu,$$

$$T_b(t) = [1 + (\gamma - 1) \cos (2\pi t / \tau)] T_b(0),$$

where $N_b(t)$, $\nu_b(t)$, and $T_b(t)$ are the number density, drift velocity, and temperature, respectively, at the bottom boundary at simulation time $t$. The plasma at this boundary is assumed to be in a state of thermal equilibrium, the velocity distribution of which is Maxwellian. The constant $\epsilon$ is the relative amplitude of the slow wave, $\tau$ is its period, and $\gamma = 5/3$ is the adiabatic exponent. As an example, the amplitude is chosen as $\epsilon = 0.1$ and the period is chosen as $\tau = 120 \text{ s}$. The acoustic speed is set to $c_s = 166 \text{ km s}^{-1}$, which is calculated theoretically from isotropic MHD assuming the processes are adiabatic. In this paper, a positive velocity means that it is in the upward direction.

Two simulations with and without Coulomb collisions are run to study the effect of Coulomb collisions. For the first simulation, the governing equation is Equation (4), the Boltzmann equation, in which the Coulomb collision is considered. For the second simulation, the governing equation is the Vlasov equation, the form of which is similar to Equation (4) but without taking into account the Fokker–Planck collision term. The simulation results are given in Figure 1 for comparison. Figure 1 displays the height profiles of the number density, drift velocity, parallel temperature, and perpendicular temperature at $t = 300 \text{ s}$ and $t = 360 \text{ s}$. The simulation results of the collisional case are shown in the left panel, and the results of the collisionless case are shown in the right panels. In Figure 1, $\delta N$, $\delta v$, and $\delta T$ denote the fluctuations away from the initial quantity, i.e., $\delta N(h) = N_p(h) - N_0(h)$, where $N_p(h)$ is the number density at $t = 300 \text{ s}$ or $t = 360 \text{ s}$, and $N_0(h)$ is the number density at $t = 0 \text{ s}$. $T_p(h)$ is the height profile of the temperature at $t = 0 \text{ s}$.

In Figure 1, periodic variations can be found in the profiles of number density, drift velocity, and temperature. The third peak of the drift velocity arrives at the height $h = 4.7 \text{ Mm}$ at $t = 300 \text{ s}$ and the height $h = 14.7 \text{ Mm}$ at $t = 360 \text{ s}$. The propagation speed of velocity perturbation is thus about $167 \text{ km s}^{-1}$, which is approximate to the acoustic speed of the plasma ($c_s$) as calculated theoretically under the MHD regime. Accordingly, we consider that the slow wave is successfully introduced into the magnetic flux tube.

The temperature of the plasma in the left panel of Figure 1 is anisotropic, the amplitude of $\delta T_{\perp}$ is larger than the amplitude of $\delta T_{\parallel}$, and the phase of $\delta T_{\parallel}$ is ahead of the phase of $\delta T_{\perp}$. $T_{\parallel}$ and $T_{\perp}$ are coupled by the Coulomb collisions here. As a result, a perturbation can be found in the profile of $T_{\perp}$. However, the Coulomb friction is not so strong as to establish an isotropic temperature distribution, so the plasma remains anisotropic in temperature. As can be seen in the right panels of Figure 1 for the case without Coulomb collisions, there is no coupling...
between the parallel temperature and the perpendicular temperature, i.e., no oscillation of the perpendicular temperature associated with the parallel temperature oscillation. What should be emphasized is that the waves, which can be excited from the thermal anisotropy instability and thereby coupling the parallel and perpendicular temperatures, are not considered in both the collisional and collisionless cases.

Only the simulation results of the collisional model are discussed further below, since the Coulomb collision should be considered in our work.

### 3.2. The Kinetic Effect of the Slow-mode Wave

The proton kinetics, i.e., characteristics of the proton velocity distribution, in association with the slow wave is investigated in this sub-section. The proton VDF \( F_0(h, v) \) along the magnetic flux tube is displayed in Figure 2(a). The perturbation is more obvious in the half of the phase space where the parallel velocity \( v_\parallel > 0 \) and less obvious in the other half where \( v_\parallel < 0 \). This asymmetric pattern in \( F_0(h, v) \) is determined by the nature of the slow wave. The perturbations of drift velocity, number density and temperature are nearly in the same phase for the slow-mode wave. When the bulk drift velocity becomes larger, the number density becomes larger and the temperature becomes higher. As a result, the velocity distribution is wider when the bulk drift velocity \( v_\parallel > 0 \) and narrower when \( v_\parallel < 0 \). So the perturbation is more obvious in the velocity domain with \( v_\parallel > 0 \).

To evaluate the deviation of VDF from the Maxwell distribution and reveal the beam component flow, a function \( \delta F_{0, \text{non-Maxw}}(h, v) \) is defined as a measure of the non-thermal state:

\[
\delta F_{0, \text{non-Maxw}}(h, v) = F_0(h, v) - F_{0, \text{fit}}(h, v),
\]

where \( F_{0, \text{fit}}(h, v) \) is the Maxwellian fitting of \( F_0(h, v) \). The values of \( \delta F_{0, \text{non-Maxw}}(h, v) \) at \( t = 300 \) s are displayed in Figure 2(b). The VDF \( F_0(h, v) \) of protons deviates from the Maxwellian distribution periodically. In the velocity regions of \(-300 \) km s\(^{-1}\) < \( v_\parallel \) < \(-100 \) km s\(^{-1}\) and 100 km s\(^{-1}\) < \( v_\parallel \) < 300 km s\(^{-1}\), the deviation from a Maxwellian distribution is more obvious, for the reason that the frequency of collisions between local protons there and the major population of protons with small \( |v_\parallel| \) is lower, while the number density of protons is not very small there. Stronger beam component flows can be found in the region of 100 km s\(^{-1}\) < \( v_\parallel \) < 300 km s\(^{-1}\), since the particles in this velocity range move nearly in phase with the propagating wave electric field \( E \). These particles can be involved in Landau resonance with the slow-mode wave.

Figure 2(e) illustrates the VDFs at four different heights with four different phases: (e1) \( h = 5 \) Mm, \( v_d = 0 \), \( \partial u_\parallel / \partial h < 0 \); (e2) \( h = 10 \) Mm, \( v_d = 0 \), \( \partial u_\parallel / \partial h < 0 \); (e3) \( h = 15 \) Mm, \( v_\parallel = 0 \) is at the minimum; and (e4) \( h = 20 \) Mm, \( v_d = 0 \), \( \partial u_\parallel / \partial h > 0 \). The Coulomb friction cannot be neglected, thus the VDFs (blue solid lines) in Figure 2(b) are not far from the Maxwellian ion distributions (red dashed lines). Nevertheless, we can find a weak asymmetry in these VDFs whenever we perform a R–B profile asymmetry analysis on them. The non-thermal components as deviating from the Maxwellian distribution are plotted with green dashed lines in Figure 2(e). We perform the R–B analysis in a similar way as Verwichte et al. (2010) did. We treat the part with \( v < v_d \) as the red wing and the part with...
\( v > v_d \) as the blue wing in our work. The results of the R–B analysis at 120–300 km s\(^{-1} \) off the velocity center \( (v_d) \) are shown in Figure 2(b). The R–B values above/below 0 indicate redward/blueward asymmetries. Among these four phases of the slow wave, the R–B estimate with maximum absolute value is \(-3.5\% \) (panel (e2) of Figure 2).

Height profiles of the electric field and the R–B values are shown in Figures 2(c) and (d), respectively. The peaks of the blueward asymmetry (the minimum of R–B) profile are located near the peaks of electric field \( E \). At the heights where the blueward asymmetry is peaked, a balance is achieved between the formation of beam flows caused by Landau resonance and the destruction of beam flows dominated by Coulomb friction. The traditional particle trapping scenario in Landau resonance is not applicable here, as the mean free path of the proton is far smaller than the wave length.

To evaluate the influence of wave parameters on the periodic beam flow formation, nine simulations are run with different amplitudes and periods of the slow wave. The wave amplitudes \( \epsilon = 0.1, 0.2, 0.3 \), and the wave periods \( \tau = 60, 120, 240 \) s, are chosen. The VDFs of these nine simulations at the phase with the most blueward asymmetry, the profiles of \( \delta F_0(\nu)/F_{0,\text{max}} \), and the corresponding R–B values are illustrated in Figure 3, where \( F_{0,\text{max}} \) is the peak value of \( F_0 \). The amplitudes and periods are listed on the top of each plot. The damping timescale of the slow wave is determined by the wave period. We note that the waves with longer periods can propagate to higher positions, and Landau resonance can happen in a larger spatial range. As a result, the most obvious beams tend to appear at higher positions for the waves with longer wave periods. For a comparison between the simulation results with different wave periods \( \tau = 60, 120, 240 \) s, VDFs at a similar phase with maximum \( E_\parallel \) (e.g., \( E_\parallel = E_{\text{max}} \) at \( h = 5, 10, 20 \) Mm and \( t = 270, 300, 360 \) s for \( \tau = 60, 120, 240 \) s, respectively) are selected and shown in Figure 3. Obvious beam flows tend to appear in the simulations with decreasing wave periods from left to right in Figure 3 and increasing wave amplitudes from bottom to top. As the wave period shortens or the wave amplitudes, the plasma will undergo more intensive charge separation, and therefore the electric field will strengthen.

The time evolution of the number density, the drift velocity, and the R–B asymmetry at \( h = 10 \) Mm are displayed in Figure 4 for comparison with the observational results. The
strongest blueward asymmetries appear ahead of the phase with the strongest blueshift by about 30 s (1/4 of the wave period). It is different from the classical intermittent outflow scenario (De Pontieu & McIntosh 2010; Tian et al. 2011), in which the strongest blueward asymmetries appear at the time when the strongest blueshift appears. This difference arises because these authors assumed that the blueward asymmetry is caused by the superposition of intermittent fast flow on the background plasmas nearby, but is unresolved under the current spatial resolution of observations. In their scenario, the weak correlation between Doppler-shift and R–B asymmetry from observations is speculated to be caused by strong noise of the R–B asymmetry. We suggest an alternative explanation for the weak correlation: a phase difference between Doppler shift and R–B asymmetry as a result of slow wave kinetic effects.

3.3. Damping of the Slow-mode Wave

Here we discuss the damping mechanism. As shown in the left panels of Figure 1, the slow waves damp gradually after the propagation of several wavelengths. We suggest that the

Figure 3. VDFs of collisional simulations with gravity at the phase of $E_1 = E_{\text{max}}$ with different amplitudes ($\epsilon = 0.1, 0.2, 0.3$) and periods ($\tau = 60, 120, 240$ s). The blue solid lines denote $F_0(v)$, the red dashed lines denote the Maxwellian fitting of $F_0(v)$, and the green dashed lines denote the values of $\delta F_{0,\text{max}}(v)/F_0,\text{max}$, where $F_0,\text{max}$ is the maximum value of $F_0(v)$. Here the height and time corresponding to the phase of $E_1 = E_{\text{max}}$ are selected as $h = 5, 10, 20$ Mm and $t = 270, 300, 360$ s for different periods $\tau = 60, 120, 240$ s, respectively.

Figure 4. Time evolution of the number density, the drift velocity, and the R–B values at $h = 10$ Mm for the collisional simulation with gravity. The positive R–B values indicate redward asymmetries.
damping of the slow wave consists of the concurrent two-fold processes: (1) formation of heat flux contributed from the non-thermal velocity tail due to Landau resonance with the slow wave and (2) thermalization of non-thermal tail particles (dissipation of heat flux) by Coulomb collisions. This mechanism behaves similarly to the thermal conduction in MHD or fluid equations. This similarity of damping evolution can be demonstrated by comparing the height profile of the wave amplitude in the simulation with the theoretical height profile of amplitude as derived from the fluid equations.

As the first step, to simplify this comparison, gravity and gravitational stratification of the plasma are not considered in the test kinetic simulation and the fluid equations. A new test simulation without gravitational stratification was run. The background number density of protons is \( N_0 = 1 \times 10^8 \text{ cm}^{-3} \), and the background temperature is \( T_0 = 1 \text{ MK} \) in the simulation and the fluid equations. For the energy equation, the form in Owen et al. (2009) was used. The fluid equations read:

\[
\frac{\partial \rho_p}{\partial t} + \frac{\partial \rho_p v_p}{\partial h} = 0, \tag{12}
\]

\[
\frac{\partial \rho_p v_p}{\partial t} + v_p \frac{\partial \rho_p v_p}{\partial h} = -\frac{\partial P_p}{\partial h} + \rho_p E, \tag{13}
\]

\[
\frac{\partial \epsilon_p}{\partial t} + v_p \frac{\partial \epsilon_p}{\partial h} = - (\gamma - 1) \epsilon_p \frac{\partial v_p}{\partial h} + \frac{1}{\rho_p} \frac{\partial}{\partial h} \left( \kappa_p \frac{\partial T_p}{\partial h} \right), \tag{14}
\]

\[
\epsilon_p = \frac{P_p}{(\gamma - 1) \rho_p}, \tag{15}
\]

where \( \rho_p \) is the proton density, \( v_p \) the proton drift velocity, \( P_p \) the proton pressure, \( \rho_{p0} \) the proton charge density, \( \epsilon_p \) the specific internal energy of proton per unit mass, \( \gamma = 5/3 \) the ratio of specific heats, and \( T_p \) the proton temperature. The electric field \( E \) is the same as the form of \( E_i \) in Equation (7).

The coefficient of thermal conduction \( \kappa_p = 460 \text{ W m}^{-1} \text{ K}^{-1} \) is estimated from the ratio \( -Q_p(h)/(\partial T_p(h)/\partial h) \) in our simulation, where \( Q_p(h) \) is the heat flux of the protons calculated from Equation (16):

\[
Q_p = Q_{p,||} + Q_{p,\perp}
\]

\[
= \int \frac{1}{2} m_p (v_i - \bar{v}_i)^3 F_0(v_i) dv_i
\]

\[
+ \int \frac{1}{2} m_p (v_i - \bar{v}_i) F_1(v_i) dv_i, \tag{16}
\]

where \( \bar{v}_i \) is the bulk velocity of the plasma. The Coulomb friction between protons and electrons is not included in Equations (12)–(14) as electrons are not dealt with as kinetic particles in the simulation model. The comparison between the theoretical heat flux \(-\kappa_p (\partial T_p/\partial h)\) and the actual heat flux \( Q_p \) is given in Figure 5.

We linearize the Equations (12)–(14) and calculate the theoretical height profile of the velocity amplitude using the same method as Owen et al. (2009). The theoretical velocity amplitude is \( \delta v(h)/c_s = (\delta v(0)/c_s) \text{exp}(-k_i h) \), where \( k_i = 2.21 \times 10^{-8} \text{ m}^{-1} \). The theoretical drift velocity amplitude given by Equations (12)–(14) (black dashed line) and the drift velocity of the simulation at \( t = 390 \text{ s} \) (blue solid line) are shown in Figure 6(a). For the drift velocity of the simulation, a periodic average drift velocity is subtracted. The simulation result agrees well with the predicted one, which means that the damping of the slow wave is dominated by the thermal conduction if the gravity and gravitational stratification of the plasma are not considered. The drift velocities of the simulations with and without gravitational stratification at \( t = 390 \text{ s} \) are shown for comparison in Figure 6(b).

The slow wave becomes a gravity-magnetohydrodynamic wave (acoustic gravity wave for parallel propagation) as the gravity is taken into account. We calculate the dispersion relation and damping growth length based on linearizing the governing equations with gravity, thermal conduction, the hydrostatic equilibrium of the background, and a uniform background temperature \( T_0 \). The perturbations \( \delta v, \delta \rho, \) and \( \delta \epsilon \) have the form \( \delta v(t, h) = \delta v_0 e^{(kh - \omega t)} \). For the set of parameters \( (N_0 = 10^8 \text{ cm}^{-3}, \kappa_p = 460 \text{ W m}^{-1} \text{ K}^{-1} \) and \( \omega = 2\pi/120 \text{ rad s}^{-1} \), it is found that the imaginary part of the wavenumber \( (k = k_e + ik_i) \) is positive, indicating a damping of the gravity-magnetohydrodynamic waves, with \( k_i H \sim (1.2, 2.0) \) for normalized \( \rho \) decreasing from 1.0 at the bottom to 0.4 at a certain height, where \( H \) is the gravitational scale height and about 60 Mm for \( T_0 = 1 \text{ MK} \). This damping of the gravity-magnetohydrodynamic wave in the thermal-conductive atmosphere is consistent with the results in Owen et al. (2009). It is also found that \( k_i H \) remains about 1.3 for the thermal-conductive atmospheres of different density without gravity. Therefore, in our case with the parameters specified above, the gravity alters the damping of thermal conduction a little but not too much, which means that the wave amplitude still attenuates rather than amplifies in the atmosphere with gravitational stratification. The gravitational amplification effect on the wave amplitude would become more significant with decreasing thermal-conduction coefficient \( \kappa \) and wave frequency \( \omega \) as well as increasing density, according to our numerical calculation experiment and Equation (38) in Owen et al. (2009).
We note that whether the stratification or the thermal conduction dominates the propagation evolution of the slow wave is determined by the parameters such as plasma density and wave period. For example, in De Moortel & Hood (2004), the influence of stratification is stronger, as the wave period is longer and the plasma density is larger in their work. The other thing we need to emphasize is that the kinetic effects of electrons are not dealt with in our simulation. However, if we consider again the kinetic effects of electrons, the slow wave is expected to damp faster, as the thermal conduction of electrons is much larger than that of the protons.

4. DISCUSSION AND CONCLUSION

To interpret the observed quasi-periodic disturbances in emission intensity, Doppler shift, line width, and R–B asymmetry, we present a new “wave + flow” scenario. In our scenario, the oscillation is a slow-mode wave, and the flow is due to a beam component created by the kinetic effect of Landau resonance. We suggest that the quasi-periodic R–B asymmetries found in the spectral lines in intensity disturbance regions may be the signatures of flows. However, we do not know whether they can be generated in the low corona, where the plasma is semi-collisional. Therefore, we test our scenario by simulating the propagation of the slow wave in a magnetic flux tube with a kinetic model.

In our simulations, the plasma in the flux tube consists of proton particles and electron fluid. Weak periodic beam components are found in the velocity distributions of the protons. The formation of the beam is caused by Landau resonance between protons and slow waves, which is partially counteracted by Coulomb collisions between protons. The signatures of beam flows and R–B asymmetries are periodic as well. The R–B asymmetry with the strongest blue wing enhancement appears at the height where the electric field and the gradient of the proton number density are at their maxima. When we envision sitting at a constant height and watching the time variation, the strongest blueward asymmetry (minimum R–B) appears before the time of the strongest blueshift (maximum drift velocity). This phase relation between R–B asymmetry and drift velocity is different from the classical scenario of intermittent outflow (De Pontieu & McIntosh 2010), in which the strongest blueward asymmetry appears at the time of strongest blueshift. The phase difference between R–B asymmetry and Doppler shift in our model may be one of the reasons for the weak correlation found between them from observations. The other point of interest is that the slow waves damp fast. We suggest that this damping may result from concurrent two-fold processes: (1) formation of heat flux (non-thermal particles) by Landau resonance and (2) dissipation of the heat flux (thermalization of the non-thermal tail in the velocity distribution) by Coulomb collisions. This idea is corroborated by comparing the damping profiles from our kinetic model with those obtained from fluid equations with thermal conduction.

Our conclusions are as follows: (1) weak periodic beam components are generated due to Landau resonance in our simulation; (2) the strongest blueward asymmetry appears before the time of strongest blueshift by 1/4 of a wave period; (3) the main damping mechanism of the slow wave in our simulation is the kinetic process related to thermal conduction; and (4) our “wave + flow” scenario may be a viable explanation for the observed quasi-periodic intensity disturbances in the solar corona.

This paper concentrated on whether a beam can be generated in the normal environment of the low corona. In fact, quasi-periodic intensity disturbances were observed in the spectra of heavy ions like Fe xi, Fe xii, etc. To further test our scenario is...
to simulate the behaviors of such heavy ions, in addition to the behaviors of protons. However, the non-thermal turbulent fluctuations are important when studying the kinetic behaviors of heavy ions, as the non-thermal width of the spectral lines is larger than the thermal width of the heavy ion’s velocity distribution in the low corona. But it is difficult to introduce a turbulent fluctuation in the simulation, since the model we use is one-dimensional in velocity and in physical space, but the turbulent fluctuation is a two-dimensional or three-dimensional phenomenon in physical space.

For heavy ions, the thermal velocity is much smaller than the local sound speed, and thus few ions can resonate with the slow mode wave. However, after considering the turbulent fluctuations, the profiles of the heavy ion velocity distribution will become wider, and more ions can be in the region of resonance with slow waves. Whether beam flows can be generated or not in this environment needs to be tested with the help of higher dimensional kinetic simulation.

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