Heating of a Quiet Region of the Solar Chromosphere by Ion and Neutral Acoustic Waves

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Received 2019 January 30; revised 2019 April 3; accepted 2019 April 20; published 2019 June 17

Abstract

Using high-resolution numerical simulations we investigate the plasma heating driven by periodic two-fluid acoustic waves that originate at the bottom of the photosphere and propagate into the gravitationally stratified and partially ionized solar atmosphere. We consider ions + electrons and neutrals as separate fluids that interact between themselves via collision forces. The latter play an important role in the chromosphere, leading to significant damping of short-period waves. Long-period waves do not essentially alter the photospheric temperatures, but they exhibit the capability of depositing a part of their energy in the chromosphere. This results in up about a five times increase of ion temperature that takes place there on a timescale of a few minutes. The most effective heating corresponds to waveperiods within the range of about 30–200 s with a peak value located at 80 s. However, we conclude that for the amplitude of the driver chosen to be equal to 0.1 km s^{-1}, this heating is too low to balance the radiative losses in the chromosphere.

Key words: methods: numerical – Sun: activity – Sun: chromosphere – Sun: transition region

1. Introduction

One of the main open questions of solar physics is energy transport from the photosphere to the chromosphere and further up to the corona. It is well known that chromospheric plasma radiates more than that of the corona, thus an additional source up to the corona. It is well known that chromospheric plasma transport from the photosphere to the chromosphere and further to the corona. Cuntz et al. (1996) showed that magnetoacoustic waves, while propagating through the gravitationally stratified medium of the chromosphere, lead to alterations in plasma temperature.

We aim to contribute to the investigations mentioned above by performing novel and high-resolution two-fluid simulations of acoustic waves propagating in the solar atmosphere. Using two-fluid numerical simulations, Maneva et al. (2017) showed that magnetooacoustic waves, while propagating through the gravitationally stratified medium of the chromosphere, lead to alterations in plasma temperature.

We aim to contribute to the investigations mentioned above by performing novel and high-resolution two-fluid simulations of acoustic waves propagating in the solar atmosphere and to quantify the amount of heating occurring in the chromosphere. This paper is organized as follows. In Section 2, we describe the physical model of the solar atmosphere. Numerical simulations are presented in Section 3. A summary and conclusions are presented in the last section.

2. Physical Model

Here we describe the physical model of the solar atmosphere we employed. In particular, we consider a gravitationally stratified and partially ionized solar atmosphere. With the use of a realistic height-dependent temperature profile of the semi-empirical model of Avrett & Loeser (2008), we determine uniquely the equilibrium mass density and gas pressure profiles that fall off with height. To describe chromospheric plasma we adopt the set of equations for two fluids, mainly for ions + electrons and neutrals, with the contribution of both species depending on local ionization level. These two-fluid equations were derived by a number of authors (e.g., Smith & Sakai 2008; Zaqarashvili et al. 2011; Meier & Shumlak 2012; Soler et al. 2013; Balster et al. 2018, and references therein). They were implemented into the JOANNA code (Wójcik 2017), which was recently used by Kuźma et al. (2017) and Srivastava et al. (2018) to study two-fluid plasma jets and by Wójcik et al. (2018) to simulate two-fluid acoustic cutoff periods. In our model we assume initial thermal balance between ion and neutral components of plasma, \( T_i = T_n = T_0 \) (Oliver et al. 2016), and neglect all electron components due to the small mass of electrons in relation to ions and neutrals. Note that the subscripts \( i \) and \( n \) correspond to ions (protons) and neutrals (hydrogen atoms), respectively. See also Section 3 of Wójcik...
et al. (2018) for an extended discussion on two-fluid hydrostatic equilibrium.

The two-fluid equations are taken from Zaqarashvili et al. (2011), with the exception for the heat production terms in the total energy equations that result from ion–neutral collisions. These terms are written here in similar (but generalized here for the solar mean atomic masses) form as in Oliver et al. (2016). See also Wójcik et al. (2018). Thus, ion and neutral components of plasma are governed by the following set of equations:

\[
\frac{\partial \varphi_i}{\partial t} + \nabla \cdot (\varphi_i V_i) = 0, \tag{1}
\]

\[
\frac{\partial \varphi_n}{\partial t} + \nabla \cdot (\varphi_n V_n) = 0, \tag{2}
\]

\[
\varphi_i \frac{\partial V_i}{\partial t} + (V_i \cdot \nabla) V_i = -\nabla p_{\kappa e} + \varphi_i g - \alpha_{in}(V_i - V_n), \tag{3}
\]

\[
\varphi_n \frac{\partial V_n}{\partial t} + (V_n \cdot \nabla) V_n = -\nabla p_n + \varphi_n g + \alpha_{in}(V_i - V_n), \tag{4}
\]

\[
\frac{\partial E_i}{\partial t} + \nabla \cdot ((E_i + p_{\kappa e}) V_i) = -\alpha_{in} V_i(V_i - V_n) + Q_{in}^i + \varphi_i g \cdot V_i, \tag{5}
\]

\[
\frac{\partial E_n}{\partial t} + \nabla \cdot ((E_n + p_n) V_n) = \alpha_{in} V_n(V_i - V_n) + Q_{in}^n + \varphi_n g \cdot V_n, \tag{6}
\]

\[
Q_{in}^i = \alpha_i \left[ \frac{1}{2} |V_i - V_n|^2 + \frac{3k_B}{m_i \mu_i + \mu_n}(T_i - T_n) \right], \tag{7}
\]

\[
Q_{in}^n = \alpha_n \left[ \frac{1}{2} |V_i - V_n|^2 + \frac{3k_B}{m_i \mu_i + \mu_n}(T_i - T_n) \right]. \tag{8}
\]

These equations are supplemented by the ideal gas laws as

\[
p_n = \frac{k_B}{m_i \mu_n} \varphi_n T_n, \quad p_i = \frac{2k_B}{m_i \mu_i} \varphi_i T_i. \tag{9}
\]

Here \( E_i \) and \( E_n \) are, respectively, ion and neutral energy densities:

\[
E_i = \frac{p_{\kappa e}}{\gamma - 1} + \frac{1}{2} \varphi_i V_i^2, \quad E_n = \frac{p_n}{\gamma - 1} + \frac{1}{2} \varphi_n V_n^2. \tag{10}
\]

where \( \gamma = 5/3 \) is the ratio of specific heats, \( \alpha_{in} \) and \( \alpha_{ni} \) are friction coefficients between ions and neutrals, \( \mu_i \) and \( \mu_n \) are mean atomic masses of, respectively, ions and neutrals, which are taken from the solar abundance model (e.g., Vögler et al. 2005), \( t \) is time, \( k_B \) is the Boltzmann constant, \( m_H \) is hydrogen mass, and \( T_i \) and \( T_n \) are temperatures of, respectively, ions and neutrals, \( V_i \) and \( V_n \) are their corresponding velocities, and \( p_{\kappa e} \) and \( p_n \) are gas pressures. A constant solar gravity magnitude \( g = 274.78 \text{ m s}^{-1} \) points in the negative \( y \)-direction. To estimate \( \alpha_{in} \), we use the formula provided by Braginskii (1965), with the collisional cross section taken as the quantum-mechanical from Vranjes & Krstic (2013), who showed that the classical hard-sphere model may lead to underestimation of the cross-section values; they also derived from the quantum-mechanical model the integral cross section \( \sigma_{in} \) for ions (protons) collisions with neutrals (neutral hydrogen atoms). For typical chromospheric plasma temperature in the range of \( (6–10) \times 10^4 \text{ K} \) this cross section is equal to \( 1.89 \times 10^{-18} \text{ m}^2 \), that is, about three orders of magnitude larger than that in the hard-sphere model. Following Zaqarashvili et al. (2011), we assumed \( \alpha_{in} = \alpha_{ni} \). However, the ion–neutral collision frequency differs from the neutral–ion collision frequency (Ballester et al. 2018).

At the bottom of the photosphere, given by \( y = 0 \), we set the periodic driver in the vertical component of ion and neutral velocities, i.e.,

\[
V_{i,n}(y, t) = V_0 \sin \left( \frac{2\pi}{P_d} \right), \tag{11}
\]

where \( V_0 \) is the amplitude of the driver and \( P_d \) its period. This driver excites upwardly propagating ion and neutral acoustic waves (e.g., Zaqarashvili et al. 2011). We set \( V_0 = 0.1 \text{ km s}^{-1} \), but allow \( P_d \) to vary within the range of \( 5 \text{ s} \leq P_d \leq 300 \text{ s} \). These values fit to the typical characteristics of flow associated with the solar granulation (e.g., Musielak et al. 1994; Hirzberger 2003). Among others, Tu & Song (2013) and Arber et al. (2016) determined the amplitude of Alfvén waves with the use of a piecewise power law. Thus, more realistic ways of implementing the amplitude \( V_0 \) are possible.

3. Numerical Results

We solve the two-fluid system of equations numerically using the JOANNA code (Wójcik 2017). In our problem, we set the Courant–Friedrichs–Lewy number (Courant et al. 1928) equal to 0.8 and adopt WENO3 with HLLD Riemann solver (Miyoshi et al. 2010). Typically, our one-dimensional numerical box covers the region between the bottom of the photosphere \( (y = 0 \text{ Mm}) \) and the low corona \( (y = 2.5 \text{ Mm}) \), and by default it is represented by \( 25 \times 10^3 \) identical numerical cells. This results in the uniform spatial grid of size \( \Delta y = 100 \text{ m} \). Above this region, namely for \( 2.5 \text{ Mm} < y < 30 \text{ Mm} \), we stretch the grid, dividing it into 128 cells who size increase with height. As short-waveperiod waves require very a fine spatial resolution, we refine our spatial grid for shorter \( P_d \).

We investigate the impact of acoustic waves propagating through the photosphere and chromosphere on plasma heating. Figure 1 shows the vertical profile of \( V_i \) (solid line) and \( V_n \) (dashed line) drawn at \( t = 10^4 \text{ s} \) for \( P_d = 5 \text{ s} \) (top) and the zoomed in regions in the photosphere (panel (a)) and upper chromosphere (panel (b)). In the photosphere, where ions and neutrals are strongly coupled, both \( V_i \) and \( V_n \) overlap each other (panel (a)). Higher up, in the chromosphere and upper chromosphere, where ions and neutrals start to decouple, the difference between ion and neutral wavefront positions is clearly seen. We infer that the amplitudes of the excited ion and neutral acoustic waves are essentially not affected by ion–neutral collisions in the photosphere. However, higher up, the amplitude grows \( e \)-times with the pressure-scale height, while waves steepen into shocks.

From Equations (7) and (8) we infer that the velocity difference between ions and neutrals, \( \delta V = V_i - V_n \), plays a role in plasma heating associated with propagating acoustic waves. Figure 2 illustrates \( \delta V(y) \) for the driving period \( P_d = 5 \text{ s} \) at \( t = 10^4 \text{ s} \). At the bottom of the photosphere both
ions and neutrals remain strongly coupled, therefore they propagate with essentially the same velocity. The difference between ion and neutral velocities grows with height, reaching a magnitude of about $14 \text{ m s}^{-1}$ at the transition region, which is located at $y \approx 2.1 \text{ Mm}$. Because at higher altitudes, the acoustic wave profiles steepen into shocks, these differences are present at these shocks and they result from the structured nature of shocks in a two-fluid regime (Hillier et al. 2016), which is clearly seen in Figure 1 (top and panel (b)).

Figure 3 shows the temporal and spatial evolution of relatively perturbed temperatures of ions, $\delta T_i/T_i = (T_i - T_0)/T_0$, where $T_0(y)$ is the hydrostatic temperature (Avrett & Loeser 2008). The low-period waves, mainly of $P_d = 5 \text{ s}$, are significantly damped with height (top left), and this damping results from ion–neutral collisions. It is expected, as oscillation periods are larger than ion–neutral collision times, that the damping efficiency will decrease and e-times growth of the wave amplitude with a pressure-scale height takes a leading role. During the damping process the energy carried by shocking acoustic waves is dissipated in the upper photosphere and lower chromosphere. For $P_d = 5 \text{ s}$, the increase of the initial temperature is up to 16% on a timescale of $10^3 \text{ s}$, and the heating occurs mainly from the level $y = 0.6 \text{ Mm}$ up to $y = 0.8 \text{ Mm}$ and from $y = 1.1 \text{ Mm}$ up to $y = 1.4 \text{ Mm}$.

According to linear theory (e.g., Zaqarashvili et al. 2011), acoustic waves with longer periods, which thus have oscillation frequencies that are much lower than ion–neutral collision frequencies, are weakly damped and they possess the capability to propagate upward, reaching upper atmospheric regions if their periods are lower than cutoff periods (Wójcik et al. 2018).
Wave amplitude growth with height is a dominant factor over the wave amplitude damping resulting from ion–neutral collisions. As a result, the wave amplitude experiences a net growth with height and linear theory is too approximate.

The top right panel of Figure 3 corresponds to $P_d = 30$ s. Acoustic waves deposit their energy mainly in the chromosphere, at the height of $y > 0.5$ Mm. The waves propagate higher up. However, as the damping depends on the ion–neutral collision coefficient (Equations (1)–(2)), and thus effectively on the mass density and temperature, higher up, the amplitude of disturbance is not sufficient to heat the chromosphere significantly. As $P_d = 180$ s is lower than the acoustic cutoff period (Wójcik et al. 2018), according to the theory originally developed by Lamb (1909) the excited acoustic waves propagate freely across the chromosphere up to the transition region (bottom left panel). See Wójcik et al. (2018) for the corresponding numerical simulations. For $P_d = 300$ s cooling events discernible in Figure 4 are related to the fluctuations of perturbed ion temperature associated with acoustic waves; rarefactions that follow shock-fronts lead to pressure decrease (e.g., Kuźma et al. 2017) and consequently act contrary to plasma heating. Energy dissipated by these waves in the chromosphere is sufficient to heat up the plasma up to $10^4$ K on a timescale of $10^3$ s. We conclude here that shorter-waveperiod waves heat more lower atmospheric heights than higher-waveperiod waves. In all cases, plasma heating grows with height, up to $y = 1.5–1.8$ Mm. Higher up, $\delta T_i/T_i$ falls off with height.

Figure 3. Time–distance plots of relative perturbed temperature of ions, $\delta T_i/T_i$, in the case of the driving period $P_d = 5$ s, $P_d = 30$ s, $P_d = 180$ s, and $P_d = 300$ s (from top left to bottom right).

Figure 4. Averaged over time relative perturbed temperature of ions, $H$, vs. driving period $P_d$, at $y = 1.0$ Mm (squares), $y = 1.25$ Mm (diamonds), $y = 1.5$ Mm (triangles), and $y = 1.75$ Mm (circles).

$H(y) = \frac{1}{t_a} \int_0^{t_a} \frac{T_i(y, t) - T_0(y)}{T_0(y)} dt$, \hspace{1cm} (12)

versus $P_d$. Here $t_a = 5 \times 10^3$ s is the averaging time. The upper photosphere, at $y = 0.5$ Mm, is a bit heated solely by short-period waves (not shown), while the chromospheric plasma is effectively heated by waves with periods higher than 30 s and shorter than 120 s (diamonds), with a maximum close to $P_d = 80$ s. Note that the upper chromosphere is heated by a wider range of waves and more thermal energy is dissipated there (triangles and circles).
The total vertical energy flux transported through the medium can be estimated as

\[ F(y, t) \approx \varrho_i(y)c_s(y)V^2_0(y, t) + \varrho_n(y)c_s(y)V^2_n(y, t), \]  

(13)

where \( c_s(y) \) is the sound speed in the equilibrium, given by

\[ c_s(y) = \sqrt{\frac{\gamma (\varrho_i(y) + \varrho_n(y))}{\varrho_i(y) + \varrho_n(y)}}. \]  

(14)

Wave energy is deposited in the form of thermal energy to overlying plasma by ion–neutral collisions. Figure 5 illustrates the collisional heating (first) part of the \( Q^{\text{in}} \) term (see Equations (7)–(8)) in the case of a driving period \( P_d = 80 \) s. In Table 1 we compare \( F(y, t = 2.5 \times 10^3 \) s\), evaluated at \( y = 0.5 \) Mm, \( y = 1.0 \) Mm, and \( y = 1.75 \) Mm (from Figure 3), that correspond to the lower, middle, and upper chromosphere, with the corresponding radiative energy losses as estimated by Withbroe & Noyes (1977). Note that the numerically obtained values are lower than the predictions. Thus, we conclude that acoustic waves with realistic amplitudes in the chromosphere do not carry a sufficient amount of energy to compensate radiative losses.

We have verified by inspection that for wave periods \( P_d = 180 \) s and \( P_d = 300 \) s consecutive heating shock-fronts are separated in space along the \( y \)-direction. The localized plasma pressure increase is compensated by rarefactions that follow the leading shock-front, before the consequent shock-front arrives at the same region, increasing the plasma temperature even more. Due to rarefactions, plasma temperature attains its quasi-stationary value. As a result, the temporally averaged relative perturbed temperature, \( H \), tends to be constant (0 in the case of \( P_d = 300 \) s) (see the bottom right panel of Figure 3 and the bottom panel of Figure 4). Thus, we infer that for very high period waves, plasma heating does not accumulate in the chromosphere, and the oscillations with these characteristic periods cannot effectively heat the plasma.

**4. Summary and Conclusions**

We performed numerical simulations of two-fluid acoustic waves propagating in the gravitationally stratified and partially ionized solar photosphere and chromosphere. We perturbed the initial hydrostatic equilibrium with the periodic driver in vertical components of both ion and neutral velocities, which operates at the bottom of the photosphere. We found that wave periods between about 30 and 200 s lead to significant heating of the chromosphere, while the upper photosphere remains hardly affected. The wave amplitude grows with height due to mass density falloff (e.g., Murawski et al. 2018), thus the amplitude growth can dominate over ion–neutral collisions damping. Such waves possess the capability to penetrate the upper layers of the solar atmosphere and dissipate their energy in the chromosphere. However, for low-amplitude waves, the deposited energy remains low and plasma heating is insignificant.

By performing our two-fluid simulations, we contributed to ongoing discussions of plasma heating by means of acoustic waves. Previous works such as Carlsson et al. (2007), Andic et al. (2008), and Sobotka et al. (2014) were based on observational estimations and single-fluid equations, and they reached the conclusion that energy carried by acoustic waves may compensate a substantial fraction of the radiative losses. Our results reveal that by implementing ion–neutral collisions, the plasma heating, which corresponds to the amplitude of the driver equal to 0.1 km s\(^{-1}\), is not sufficient to compensate the radiative losses.

The authors thank the referee for the stimulating comments, Dr. Nikola Vinas for discussions on solar abundance, and Dr. Iztvan Ballai for stimulating discussions. This work was done within the framework of the projects from the Polish Science Center (NCN) grants No. 2014/15/B/ST9/00106, 2017/25/B/ST9/0050, and 2017/27/N/ST9/01798. The JOANNA code was developed by Darek Wójcik at UMCS, Lublin, Poland.

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**Table 1**

Radiative Energy Losses and Averaged Energy Flux of Acoustic Waves in the Solar Chromosphere, \( F(y, t = 2.5 \times 10^3 \) s\), in the Case of the Driving Period \( P_d = 80 \) s

| Layer          | Radiative Losses (erg cm\(^{-2}\) s\(^{-1}\)) | \( F \) (erg cm\(^{-2}\) s\(^{-1}\)) |
|---------------|----------------------------------|------------------|
| Low chromosphere | \( 2 \times 10^6 \)                | \( 2 \times 10^8 \)       |
| Middle chromosphere | \( 2 \times 10^6 \)            | \( 2 \times 10^8 \)       |
| Upper chromosphere | \( 3 \times 10^5 \)             | \( 5 \times 10^7 \)       |
