DYNAMICS OF CORONAL RAIN AND DESCENDING PLASMA BLOBS IN SOLAR PROMINENCES. II. PARTIALLY IONIZED CASE

R. Oliver1,2, R. Soler1,2, J. Terradas1,2, and T. V. Zaqqarashvili3,4

1 Departament de Física, Universitat de les Illes Balears, E-07122 Palma de Mallorca, Spain; ramon.oliver@uib.es
2 Institute of Applied Computing & Community Code (IAC3), UIB, Spain
3 Institute of Physics, IGAM, University of Graz, Universitätsplatz 5, 8010, Graz, Austria
4 Abastumani Astrophysical Observatory, Ilia State University, 0162 Tbilisi, Georgia

Received 2015 October 28; accepted 2016 January 5; published 2016 February 16

ABSTRACT

Coronal rain clumps and prominence knots are dense condensations with chromospheric to transition region temperatures that fall down in the much hotter corona. Their typical speeds are in the range 30–150 km s\(^{-1}\) and of the order of 10–30 km s\(^{-1}\), respectively, i.e., they are considerably smaller than free-fall velocities. These cold blobs contain a mixture of ionized and neutral material that must be dynamically coupled in order to fall together, as observed. We investigate this coupling by means of hydrodynamic simulations in which the coupling arises from the friction between ions and neutrals. The numerical simulations presented here are an extension of those of Oliver et al. to the partially ionized case. We find that, although the relative drift speed between the two species is smaller than 1 m s\(^{-1}\) at the blob center, it is sufficient to produce the forces required to strongly couple charged particles and neutrals. The ionization degree has no discernible effect on the main results of our previous work for a fully ionized plasma: the condensation has an initial acceleration phase followed by a period with roughly constant velocity, and, in addition, the maximum descending speed is clearly correlated with the ratio of initial blob to environment density.

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

1. INTRODUCTION

Several relevant facts converge in the physics of coronal rain; some of them are outlined as follows.

1. Its genesis. Coronal rain blobs are presumably formed when heating at the feet of active region loops triggers catastrophic cooling events (Müller et al. 2004, 2005; Antolin et al. 2010; Fang et al. 2015; Moschou et al. 2015). Plasma at coronal heights then cools from coronal to chromospheric temperatures (Antolin et al. 2015; Vashalomidze et al. 2015) and starts to fall under the action of gravity. The study of coronal rain has the potential to improve understanding of coronal heating processes and their catastrophic cooling (see Antolin et al. 2010).

2. Its dynamics. During their fall, rain blobs usually reach smaller than free-fall velocities and rather constant accelerations (Wik et al. 1996; Schrijver 2001; De Groof et al. 2004, 2005; Zhang & Li 2009; Antolin et al. 2010, 2012; Antolin & Verwichte 2011; Antolin & Rouppe van der Voort 2012). The simple numerical simulations of Oliver et al. (2014, hereafter Paper I) indicate that the presence of the condensation in the coronal loop results in a rearrangement of the hot plasma and that a large pressure gradient builds up. This provides an upward force that allows the blob to fall at a reduced, smaller-than-gravity acceleration. Antolin et al. (2015) give another explanation, by which the condensing material drags and compresses the magnetic field. As a consequence, the magnetic pressure downstream of the blob increases, and so an upward force that counteracts gravity is established.

3. The potential of coronal rain as a magnetic field tracer. Antolin & Verwichte (2011), Antolin & Rouppe van der Voort (2012), Harra et al. (2014), and Scullion et al. (2014) have used coronal rain blobs to reveal the multistranded magnetic structure of the magnetic tubes in which the cool clumps move. These magnetic tubes are often part of the diffuse (and hard to resolve) component of the corona, which in a single active region contains of the order of 100,000 strands (Klimchuk 2015). Furthermore, Antolin & Rouppe van der Voort (2012) found that coronal rain condensations often occur simultaneously in neighboring strands, which implies that these strands have a coherent cooling. This evidence points to a heating mechanism that affects adjacent strands at the same time and in a similar manner; this brings us back to the first issue, namely, the coronal rain formation.

4. Its role in the chromosphere–corona mass cycle. Antolin & Rouppe van der Voort (2012) and Antolin et al. (2015) estimated the downward mass flux per loop to be of the order of (1–5) \times 10^9 g s\(^{-1}\). This is a lower limit to the true mass drainage by coronal rain since the smallest mass clumps have widths below the observational resolution (Antolin & Rouppe van der Voort 2012; Scullion et al. 2014; Antolin et al. 2015); the same conclusion has been reached by Fang et al. (2013, 2015) and Moschou et al. (2015) on the basis of coronal rain numerical simulations. According to Antolin & Rouppe van der Voort (2012), who used the mass flux estimates of Beckers (1972) and Pneuman & Kopp (1978), the mass supplied by spicules to the solar corona amounts to 1.5 \times 10^{10} g s\(^{-1}\). In addition, a condensation/drainage rate of 10^{10} g s\(^{-1}\) during the formation of a prominence was estimated by Liu et al. (2012). It then seems that coronal rain contributes significantly to the mass interchange between chromosphere and corona.
5. Coronal rain encompasses transition region to chromospheric temperatures and displays neutral and ionized species falling in unison. This co-spatiality of neutral and ionized material has been observed in Hα and He ι λ304 (De Groof et al. 2004, 2005), in Hα and Ca ii λ8542 (Ahn et al. 2014), and in Hα and Ca ii H (Chae 2010). The last observation corresponds to cool emission features in a hedgerow prominence. More observations, with even better spatial resolution, are desirable to have more information on the dynamics of neutrals and ions in falling plasmas.

In this work we propose to use coronal rain and falling prominence knots as a test bed for exploring the interaction of the ionized and neutral plasma fractions. The numerical simulations of Paper I show that the descending cool clumps tend to achieve a more or less constant speed whose value depends, among other factors, on the ratio of blob to environment (i.e., loop) density. In the absence of ion–neutral interactions, the neutral fraction of a blob falls in a loop evacuated of neutrals, so that the density ratio is very large and the falling speed becomes much larger than that of the ionized fraction of the blob. In other words, very quickly the neutral part of the cool clump would move away from the ionized part, which contradicts the observations. Therefore, ion–neutral friction must be large enough to force a common dynamics of the two species. To investigate the ion–neutral coupling, a two-fluid model is adopted here (Section 2). The coronal rain path is assumed vertical since loop curvature does not induce qualitative changes in the blob dynamics (Paper I). A relevant assumption in our model is that once the cool blob is formed, the cooling process triggered by thermal instability ceases to operate. This may contradict the observations of Antolin et al. (2015), who find that coronal rain appears intermittent and clumpy at coronal heights, although it becomes more continuous at the chromospheric level. These authors suggest that this is caused by the smearing effect of a continued cooling of the rain even after it has formed. The results and discussion are presented in Sections 3 and 4, respectively. We note that our conclusions can also be relevant for the dynamics of descending prominence knots (Engvold 1976; Liu et al. 2012).

2. MODEL

2.1. Governing Equations

To describe the temporal evolution of a partially ionized blob falling under the influence of gravity together with the pressure and friction forces, we consider the two-fluid equations for the charged (ions plus electrons) and neutral fractions (e.g., Draine et al. 1983; Smith & Sakai 2008; Meier 2011; Zaqareshvili et al. 2011b; Meier & Shumlak 2012; Khomenko et al. 2014a; Leake et al. 2014). We thus have the mass continuity equations for charged particles and neutrals,

$$\frac{\partial \rho_i}{\partial t} + \nabla \cdot (\rho_i v_i) = 0,$$

$$\frac{\partial \rho_n}{\partial t} + \nabla \cdot (\rho_n v_n) = 0,$$

where $\rho_i$ and $\rho_n$ are their respective densities (the electron mass is neglected) and $v_i$ and $v_n$ are their respective velocities.

We also have the momentum equations of charged particles and neutrals,

$$\rho_i \left[ \frac{\partial v_i}{\partial t} + (v_i \cdot \nabla) v_i \right] = -\nabla \rho_{ie} + \rho_i g - \alpha_{in} (v_i - v_n),$$

$$\rho_n \left[ \frac{\partial v_n}{\partial t} + (v_n \cdot \nabla) v_n \right] = -\nabla \rho_n + \rho_n g + \alpha_{in} (v_i - v_n),$$

with $g$ the acceleration of gravity, $\rho_{ie}$ the combined pressure of ions and electrons, and $\rho_n$ the neutrals’ pressure. Furthermore, the friction between the two plasma fractions causes their coupling through a force per unit volume proportional to $v_i - v_n$. To derive the proportionality constant, $\alpha_{in}$, we assume that the ion–neutral momentum transfer cross section ($\sigma_{in}$) is independent from the relative velocity of the colliding particles. This assumption is valid because charge exchange interactions (for which $\sigma_{in} \sim $ const.) generally contribute most of the momentum exchange compared to elastic scattering (for which $\sigma_{in}$ depends on the relative velocity of colliding particles). We thus use Equation (3.10) of Draine (1986), where a factor $\tilde{\sigma}$ is missing, or Equations (6), (7), (14), and (16) of Pinto & Galli (2008). Using the notation of Draine (1986), we can write

$$\alpha_{in} = \frac{I_1(s)}{2 \frac{m_i}{m_n}} \sigma_{in} \sqrt{\frac{8 k_B}{\pi}} \left( \frac{T_i}{m_i} + \frac{T_n}{m_n} \right) \rho_i \rho_n,$$

where $k_B$ is Boltzmann’s constant, $m_i$ and $m_n$ are the respective atomic masses of ions and neutrals, $T_i$ and $T_n$ are their respective temperatures, and $\sigma_{in}$ is the collision cross section between ions and neutrals (see, e.g., Braginskii 1965; Chapman & Cowling 1970, for its definition). Furthermore, $I_1(s)$ is given by Equation (3.11) of Draine (1986), with the parameter $s$ equal to

$$s = \frac{|v_i - v_n|}{c_r}, \quad c_r = \left[ 2 k_B \left( \frac{T_i}{m_i} + \frac{T_n}{m_n} \right) \right]^{1/2}.$$

Since our numerical simulations give very small values of $|v_i - v_n|$ compared to $c_r$, it is reasonable to consider $s \ll 1$. Taking into account that

$$\lim_{s \to 0} I_1(s) = \frac{8}{3},$$

we finally have

$$\alpha_{in} \approx \frac{4}{3} \frac{\sigma_{in}}{m_i + m_n} \sqrt{\frac{8 k_B}{\pi}} \left( \frac{T_i}{m_i} + \frac{T_n}{m_n} \right) \rho_i \rho_n.$$ 

Hence, Equation (8) is equivalent to Equations (7.4) and (7.5) of Braginskii (1965), who assumed an equal temperature of ions and neutrals. Note that since $\alpha_{in}$ depends on the density and temperature of ions and neutrals, it implicitly depends on position and time.

---

5 The prevalence of charge exchange over elastic scattering in proton-H momentum transfer can be seen in Figure 1(b) of Kestic & Schultz (1999), which shows that the line labeled “mt,sym” (which includes the symmetric process of resonant charge exchange) gives much higher values than that labeled “mt” (in which this process is not included). Appendix A3.2.3 and Figure A3.5 of Gerhardt (2004) lead to the same conclusion, namely, “the traditional result that charge transfer is the dominant process for momentum loss in ion–atom interactions.”
Here we assume an \( H \) plasma and so \( m_n \approx m_i = m_p \) (with \( m_p \) the proton mass). Regarding the proton–hydrogen collision cross section, slightly different values for \( \sigma_n \) are given in the literature (see Soler et al. 2015, for a discussion of this topic). To compute this parameter, Braginskii (1965) and Zaqarashvili et al. (2011a) use the hard sphere model for collision cross sections (see Chapman & Cowling 1970) and obtain \( \sigma_n = 4.70 \times 10^{-20} \text{ m}^2 \). Leake et al. (2012) use the value \( \sigma_n = 1.41 \times 10^{-19} \text{ m}^2 \), which is exactly three times the one used by the previous authors. On the other hand, Vranjes & Krstic (2013) use quantum mechanical theory to compute cross sections of several processes. From their results the value \( \sigma_n \approx 1.6 \times 10^{-18} \text{ m}^2 \) can be inferred for elastic scattering and charge transfer collisions and a proton energy around 1 eV. In this work we take \( \sigma_n \approx 0.75 \times 10^{-18} \text{ m}^2 \) and note that increasing \( \sigma_n \) would couple ions and neutrals more strongly, so that our results would not change. Our value is in close agreement with the one used by Leake et al. (2013) (based on the work of Draine et al. 1983), namely, \( \sigma_n = 1.16 \times 10^{-18} \text{ m}^2 \).

We finally have the energy equations of charged particles and neutrals, written in terms of their pressures,

\[
\begin{align*}
\frac{\partial p_{ie}}{\partial t} + (v_i \cdot \nabla) p_{ie} &= \gamma p_{ie} \nabla \cdot v_i = (\gamma - 1) Q_{ie}^m, \quad (9) \\
\frac{\partial p_n}{\partial t} + (v_n \cdot \nabla) p_n &= \gamma p_n \nabla \cdot v_n = (\gamma - 1) Q_n^m. \quad (10)
\end{align*}
\]

Here \( \gamma \) is the ratio of specific heats and \( Q_{ie}^m \) and \( Q_n^m \) represent the heat generation due to ion–neutral collisions and thermal transfer and are given by

\[
\begin{align*}
Q_{ie}^m &= \alpha_{ie} \left[ \frac{1}{2} |v_i - v_n|^2 + \frac{3}{2} k_B \left( T_n - T_i \right) \right], \quad (11) \\
Q_n^m &= \alpha_{in} \left[ \frac{1}{2} |v_i - v_n|^2 + \frac{3}{2} k_B \left( T_i - T_n \right) \right]. \quad (12)
\end{align*}
\]

The second term in these two formulae can be derived from Equation (3.13) of Draine (1986), together with the assumption \( \epsilon \ll 1 \), as before. We note that although some authors (e.g., Smith & Sakai 2008; Zaqarashvili et al. 2011b) have used different expressions for these quantities, the ones provided here are correct. Leake et al. (2014) have almost identical expressions, but their Equation (16) misses the factor 1/2 before the temperature difference. It is also worth mentioning that Terradas et al. (2015) solved Equations (1)–(12) to study the support of neutrals against gravity in prominences. Although they state that \( Q_{ie}^m = -Q_{in}^m \) (or \( W_{ie} = -W_{in} \) in their notation), this is just an error in the text and their solutions are completely valid.

Since our aim is to concentrate on the dynamics of a partially ionized blob falling in the solar atmosphere, a number of assumptions have been made when writing the above equations. Effects such as ionization/recombination have been ignored. Furthermore, collisions between electrons and neutrals usually have a negligible effect, and for this reason we have considered a vanishing electron–neutral friction coefficient (\( \alpha_{en} = 0 \)) in the momentum and energy equations. Joule heating and the divergence of the heat flux in the energy equations have also been neglected. In addition, in the following only motions in the vertical direction are taken into account (see Paper I for a discussion of the curvature of the blob path and its effect on the blob dynamics). Moreover, the effect of magnetic fields is ignored.

A Cartesian coordinate system with the \( z \)-axis pointing in the vertical direction is considered. Next, we insert \( \rho_i = \rho_i(z, t), \rho_n = \rho_n(z, t), p_{ie} = p_{ie}(z, t), p_n = p_n(z, t), v_i = v_i(z, t) \hat{e}_i, \) and \( v_n = v_n(z, t) \hat{e}_n \) into Equations (1)–(4), (9), and (10) and substitute \( \mathbf{g} = -g \hat{e}_z \). The horizontal components of the momentum equations are identically zero, and so we end up with six nonlinear partial differential equations (PDEs) for the six unknowns,

\[
\begin{align*}
\frac{\partial \rho_i}{\partial t} &= -v_i \frac{\partial \rho_i}{\partial z} - \rho_i \frac{\partial v_i}{\partial z}, \quad (13) \\
\frac{\partial \rho_n}{\partial t} &= -v_n \frac{\partial \rho_n}{\partial z} - \rho_n \frac{\partial v_n}{\partial z}, \quad (14) \\
\rho_i \frac{\partial v_i}{\partial t} &= -\rho_i v_i \frac{\partial v_i}{\partial z} - \rho_i \frac{\partial p_{ie}}{\partial z} - g \rho_i - \alpha_{in}(v_i - v_n), \quad (15) \\
\rho_n \frac{\partial v_n}{\partial t} &= -\rho_n v_n \frac{\partial v_n}{\partial z} - \rho_n \frac{\partial p_n}{\partial z} - g \rho_n + \alpha_{in}(v_i - v_n), \quad (16) \\
\frac{\partial p_{ie}}{\partial t} &= -v_i \frac{\partial p_{ie}}{\partial z} - \gamma p_{ie} \frac{\partial v_i}{\partial z} + (\gamma - 1) Q_{ie}^m, \quad (17) \\
\frac{\partial p_n}{\partial t} &= -v_n \frac{\partial p_n}{\partial z} - \gamma p_n \frac{\partial v_n}{\partial z} + (\gamma - 1) Q_n^m. \quad (18)
\end{align*}
\]

These expressions are supplemented with the ideal gas law for charged particles and neutrals,

\[
p_{ie} = 2\rho_i R T_i, \quad p_n = \rho_n R T_n, \quad (19)
\]

with \( R = k_B/m_p \) the ideal gas constant. The reason for the presence of the factor 2 in Equation (19) for charged particles is that the partial pressures of ions and electrons are summed up under the conditions of same temperature and same number density.

2.2. Static Equilibrium

In Paper I we considered a fully ionized blob falling in a vertically stratified atmosphere representing the coronal part of the plasma along the blob path. Coronal \( H \) is fully ionized, and for this reason the neutral fraction of the blob moves in an environment devoid of neutrals. This causes numerical problems when solving the two-fluid system of PDEs presented above, and for this reason we here assume that the ionized and neutral fractions of the blob move in an atmosphere in which both gas fractions are present. The amount of neutrals is kept as small as possible, and so ions are predominant in the coronal part of the plasma. For a homogeneous equilibrium temperature, Equations (13)–(18) have the following static solution:

\[
\begin{align*}
\rho_i(z, t = 0) &= \rho_{i0} e^{-z/H_i}, \rho_i(z, t = 0) = \rho_{i0} e^{-z/H_i}, \quad (20) \\
\rho_n(z, t = 0) &= \rho_{n0} e^{-z/H_n}, \rho_n(z, t = 0) = \rho_{n0} e^{-z/H_n}, \quad (21)
\end{align*}
\]

where the density and pressure at the coronal base \( (z = 0) \) satisfy the ideal gas law for the charged and neutral fractions of the gas,

\[
p_{i0} = 2\rho_{i0} R T_0, \quad p_{n0} = \rho_{n0} R T_0. \quad (22)
\]

\( T_0 \) is the homogeneous equilibrium temperature, assumed equal for all species in the initial state. Furthermore, the charged particles’ and neutrals’ vertical scale heights \( (H_i \text{ and } H_n) \)
depend on the equilibrium temperature as follows:

\[ H_i = \frac{2RT_0}{g}, \quad H_n = \frac{RT_0}{g}. \quad (23) \]

Then, the isothermal equilibrium assumption results in \( H_i = 2H_n \).

### 2.3. Mass Condensation

At \( t = 0 \) a dense blob composed of charged and neutral material is added to the previous static atmosphere, and its temporal evolution is then investigated. Since the condensation process is ignored, a fully condensed blob is initially at rest at a certain height and is left to evolve under the action of gravity, the pressure gradient, and the friction force between ions and neutrals according to Equations (13)–(18). At \( t = 0 \) the ionized and neutral parts of the blob have densities described by

\[
\rho_i(z, t = 0) = \rho_{i0}\exp\left(-\frac{z - z_0}{\Delta}\right), \quad (24)
\]

\[
\rho_n(z, t = 0) = \rho_{n0}\exp\left(-\frac{z - z_0}{\Delta}\right), \quad (25)
\]

where \( \rho_{i0}, \rho_{n0} \) are the maximum blob densities of ions and neutrals at the initial time, \( z_0 \) is its initial position, and its length in the vertical direction is of the order of \( 2\Delta \). The maximum density value, \( \rho_{i0} + \rho_{n0} \), used in this work is in agreement with the rain core density estimations of Antolin et al. (2015), which are in the range \( 2 \times 10^{10}–2.5 \times 10^{11} \, \text{cm}^{-3} \). Similar values have been found in numerical simulations by Müller et al. (2004, 2005), Antolin et al. (2010), and Fang et al. (2015).

### 3. RESULTS

The initial density distribution is the sum of \( \rho_i \) in Equation (20) and \( \rho_n \) in Equation (24) for ions and the sum of \( \rho_i \) in Equation (21) and \( \rho_n \) in Equation (25) for neutrals. The initial pressure of the two gas components is described by Equations (20) and (21), so that at \( t = 0 \) the blob is not in mechanical equilibrium and starts to fall. The initial velocity is zero for both ions and neutrals. The temporal evolution of the system is described by Equations (13)–(18); details of the imposed boundary conditions and the numerical method used to solve this system of PDEs can be found in Paper I.

Unless stated otherwise, we use the following parameter values at \( t = 0 \): the ion and neutral densities at the coronal base are \( \rho_{i0} = 4.5 \times 10^{-12} \, \text{kg m}^{-3} \) and \( \rho_{n0} = 0.5 \times 10^{-12} \, \text{kg m}^{-3} \). The temperature is \( T_0 = 2 \times 10^6 \, \text{K} \), which results in the scale heights of ions and neutrals \( H_i \approx 120 \, \text{Mm} \) and \( H_n \approx 60 \, \text{Mm} \). Furthermore, the blob has an initial length given by \( \Delta = 0.5 \, \text{Mm} \) and is released at the height \( z_0 = 50 \, \text{Mm} \). The only two parameters that remain to be fixed are the maximum blob density of ions and neutrals: \( \rho_{i0} \) and \( \rho_{n0} \). The ionization degree of a coronal rain blob or a falling prominence knot is not well known, and for this reason we select different values of \( \rho_{i0} \) and \( \rho_{n0} \). The results shown in this paper correspond to two cases: (1) a blob made of equal quantities of ions and neutrals (\( \rho_{i0} = \rho_{n0} \), Sections 3.1–3.4), and (2) a blob containing 10% ions and 90% neutrals (\( \rho_{i0} = 0.9\rho_{n0} \), Section 3.5).

### 3.1. Blob Dynamics in the Absence of Ion–Neutral Friction

We start considering the decoupled dynamics of neutrals and charged particles by imposing \( \alpha_{in} = 0 \). Then our PDEs split into the system of Equations (13), (15), and (17) for charged particles and the independent system of Equations (14), (16), and (18) for neutrals. Each of these systems of PDEs is identical to that solved in Paper I; thus, the dynamics of the neutral and charged parts of the falling blob is already known. In fact, although the blob is initially composed of a mixture of the two species, when \( \alpha_{in} = 0 \), they will in general behave as two independent blobs, as we explain now. The main conclusion of Paper I is that for a fixed blob length (i.e., a fixed \( \Delta \)), the blob dynamics is to a great extent determined by the density contrast of the blob with respect to the ambient plasma. For equal blob density (\( \rho_{i0} = \rho_{n0} \)), the density contrast of neutrals is larger than that of ions for two reasons: the smaller neutral density at the coronal base and the smaller neutral scale height. For this reason, neutrals fall much faster than the ionized material (Figure 1(a)). As can be seen in this figure, the neutrals’ density also displays two side effects of a large density contrast, namely, a strong increase with time of the blob density and a proportional reduction in its length so as to maintain the blob mass almost constant during its fall (these features were also found in Paper I). Figure 1(b) shows the pressure of the ionized and neutral fractions. We see that \( p_g \) develops a gradient that remains constant and moves with the blob. This gradient is roughly equal to the gravity force, thus allowing the ionized blob to achieve a more or less constant downward speed (see Paper I). On the other hand, because of the larger density contrast, the neutral blob requires a higher pressure gradient to counteract the gravity force. Figure 1(b) shows how this pressure gradient builds up during the simulation.

### 3.2. Blob Dynamics with Ion–Neutral Friction: Temperature and Ion–Neutral Collision Frequency

Now we repeat the study of Section 3.1 but use the values of the friction coefficients, \( \alpha_{in} \), given by Equation (8). The time variations of the density and pressure for selected times are presented in Figures 1(c) and (d). These figures are explained in Section 3.3 because to understand the spatial structure of the pressure, we first need to pay attention to the blob temperature, which is determined from the ideal gas law for neutrals and charged particles. At \( t = 0 \) (Figure 2(a)) the temperatures of charged particles and neutrals are everywhere equal to \( 2 \times 10^6 \, \text{K} \), except at the blob, where they smoothly decrease to their minimum values \( T_i \approx 57,700 \, \text{K} \) and \( T_n \approx 4350 \, \text{K} \). For \( t > 0 \) (Figures 2(b)–(d)) the two temperatures very quickly become identical, in about 4 ms, and take practically equal values (around 40,000 K) during the whole simulation. This fast temperature balance is achieved by the thermal energy transfer induced by the difference \( T_i - T_n \) in the last term of the quantities \( Q_{in}^\alpha \) and \( Q_{in}^\beta \) (Equations (11) and (12)).

Zaqarashvili et al. (2011b) show that the relative velocity between ions and neutrals decreases exponentially in time, as \( \exp(-t/\tau) \). Soler et al. (2013) find an identical behavior for the vorticity component parallel to the equilibrium magnetic field. We next check whether the energy interchange between ions and neutrals found here also has the same exponential
timescale, which is given by

$$\tau = \frac{1}{n_{\text{ni}} + n_{\text{in}}}$$

where $n_{\text{ni}}$ is the collision frequency between neutrals and ions,

$$n_{\text{ni}} = \frac{\alpha_{\text{in}}}{\rho_n} = 72.5 \left( \frac{T_i + T_n}{10^4 \text{ K}} \right)^{1/2} \frac{n_i}{10^{10} \text{ m}^{-3}}$$

$$= 433 \left( \frac{T_i + T_n}{10^4 \text{ K}} \right)^{1/2} \frac{\rho_i}{10^{-10} \text{ kg m}^{-3}}.$$  

The collision frequency between ions and neutrals, $n_{\text{in}}$, follows from this formula with the subscripts $i$ and $n$ interchanged and with $\alpha_{\text{in}} = \alpha_{\text{ni}}$. Plugging in this expression the characteristic values of the blob core at $t = 0$ ($T_i \simeq 57,700 \text{ K}$, $T_n \simeq 4350 \text{ K}$, and $\rho_i \simeq \rho_n \simeq 10^{-10} \text{ kg m}^{-3}$), we obtain $n_{\text{in}} \simeq n_{\text{ni}} \simeq 1080 \text{ Hz}$, i.e., a timescale $\tau \simeq 0.5 \text{ ms}$. (We note that using the temperature values $T_i \simeq T_n \simeq 40,000 \text{ K}$ achieved at $t = 4 \text{ ms}$, we obtain a very similar exponential timescale, namely, $\tau \simeq 0.4 \text{ ms}$.) The exponential temperature transition from initial to final values is plotted with dashed lines in Figure 2 (d), and it shows an excellent agreement with the behavior of the numerical simulation, in line with Zaqarashvili et al. (2011b) and Soler et al. (2013).

### 3.3. Blob Dynamics with Ion–Neutral Friction: Density, Pressure, and Velocity Drift

From Figure 1(c) it is obvious that ion–neutral friction provides a coupling strong enough so as to force the neutral and charged fractions of the blob to move as one. The full blob now falls as a single entity, at an intermediate speed between that of the decoupled ionized and neutral blobs of Figure 1(a). Visual inspection of Figure 1(c) points out that the falling speed is more or less constant; this will be discussed later. The numerical results are stored for values of $z$ whose separation is 20 km. For all times the ionized and neutral blobs have their maximum density position exactly on the same grid point. For this reason, from now on we will use “maximum density
position” to refer to the common maximum density height of both ions and neutrals.

We next turn our attention to the pressure structure (see Figure 1(d)). Again, a pressure gradient develops at the blob position, although its spatial structure is more complex than that of the decoupled case. As we explain next, this is a consequence of the neutral and charged temperature equalization just described. If the thermal exchange terms in the energy equation are set to zero, then the two species are dynamically coupled through the terms proportional to $v_i v_n$ in the momentum equations and the condensation falls exactly as in Figure 1(c). Neutrals and charged particles are also coupled by means of the term proportional to $(v_i - v_n)^2$ in the energy equations, but the omission of the quantities proportional to $T_i - T_n$ makes their temperatures independent from one another. The neutral and ionized parts of the blob fall together, and pressure gradients similar to those in Figure 1(b) develop.

We recall that these pressure gradients result from the presence of the blob at $t = 0$ and its lack of mechanical equilibrium in the vertically stratified atmosphere. On the other hand, when the terms proportional to $T_i - T_n$ in the energy equation are retained, the ion temperature quickly varies at the beginning of the simulation from $\approx 57,700$ K to $\approx 40,000$ K, and this induces a pressure decrease on top of the pressure gradient at the blob position, hence the shape of solid lines in Figure 1(d). In the case of neutrals, the temperature readjustment is much larger (from $\approx 4350$ K to $\approx 40,000$ K), and so the pressure increase at the blob position prevails over the pressure gradient established at $t = 0$ (dashed lines in Figure 1(d)).

This vertical structure of $p_n$ and $p_i$ points out a complex force balance inside the condensation. The pressure gradient at the bottom (top) part of the neutral blob, i.e., left (right) of the vertical dotted line in Figures 1(c) and (d), points downward (upward). The opposite applies to the ionized part of the condensation, and so the pressure gradient alone tends to expand the neutral fraction of the blob and to compress the ionized fraction. Nevertheless, Figure 1(c) does not display changes in the blob shape, the reason being that the friction force acts in the opposite direction to the pressure gradient. In the case of neutrals the friction force (per unit volume) equals $\alpha_{in} (v_i - v_n)$. The vertical distribution of the velocity drift, $v_i - v_n$, for the times of Figures 1(c) and (d) is shown in Figure 3. One can see that it provides an upward force on neutrals at the bottom part of the blob and a downward pull at its top. These forces balance $-\nabla p_n$ and prevent the neutral blob from expanding vertically, and similarly, for the ionized part of

---

**Figure 2.** Vertical distribution of the temperature for (a) $t = 0$, (b) $t = 0.4$ ms, (c) $t = 2$ ms. (d) Temporal variation of the temperature at the maximum density position. Red and black solid lines, respectively, correspond to ions and neutrals, while dashed lines give a transition from the initial to the final temperature of the form $(T_{\text{initial}} - T_{\text{final}}) \exp(-t/\tau) + T_{\text{final}}$ with $\tau = 0.5$ ms. The results of this figure correspond to the numerical simulation of Figures 1(c) and (d).
the blob, these combined forces mean that it is not compressed vertically.

It may then seem that the friction force is important in keeping the blob structure, but this is only true because of our choice of initial conditions. As we have mentioned, although the initial temperature difference of neutrals and charged particles inside the blob is removed in a time $t \approx 2 \text{ ms}$, it gives rise to a drift $|v_i - v_n| \approx 200 \text{ m s}^{-1}$ at the blob edges. This drift is absent when the terms proportional to $T_i - T_n$ in the energy equations are ignored and also when the initial blob densities of neutrals and ions are selected so as to have equal initial temperatures at the blob center. Hence, the ion–neutral velocity drifts of Figure 3(a) and the resulting complex pressure structure of Figure 1(d) are unrealistic features of our simulations, and they can be ignored because they do not modify the blob dynamics, whose study is the main aim of this paper.

The true main role of the friction force, on the other hand, is to provide the necessary forces for the common descent of ions and neutrals in the condensation. One must bear in mind that when we set $\alpha_{in} = 0$, we found that neutrals fall much faster than charged particles. To see the net effect of the competing pressure gradient and friction force when friction is included, their sum is represented in Figure 3(b). Although the drift velocity is dominant and has opposite signs on the blob edges, when these two forces are combined together they give a net upward (positive) force acting on the whole blob that makes neutrals fall more slowly than in the uncoupled case. It is worth mentioning that $v_i - v_n \approx 0.07 \text{ m s}^{-1}$ at the blob center, much smaller than the value $|v_i - v_n| \approx 200 \text{ m s}^{-1}$ found at the blob edges. Note that Figure 3(b) corresponds to neutrals; when charges are considered, the sum $-\nabla p - \alpha_{in} (v_i - v_n)$ gives a downward (negative) force that increases the falling acceleration of charged particles. Then, the descent of neutrals is slowed down and that of ions is accelerated in the exact amount to make them fall as one.

3.4. Blob Kinematics with Ion–Neutral Friction

More insight into the blob kinematics can be gained by representing the height and velocity of the maximum density position, for both charges and neutrals; here we identify the behavior of the maximum density position with that of the blob. Figure 4 confirms previous results of Figure 1: the two blob fractions move at very different speeds throughout the simulation when $\alpha_{in} = 0$ (solid lines), but they move as a single entity when $\alpha_{in} \neq 0$ (symbols). Perhaps the most interesting result of Figure 4 is that the falling speed in the
presence of ion–neutral friction is not the average of the uncoupled values; it is actually much closer to the small velocity of the uncoupled ionized blob. The results presented here are in agreement with those of a fully ionized plasma: the velocity of the uncoupled ionized blob. The results presented so far the maximum density of ions and neutrals in the blob is the same, namely, \( \rho_{b0} = \rho_{b0} = 10^{-10} \text{kg m}^{-3} \). The relative abundance of the two species in coronal rain is controlled by ionization and recombination processes, which are ignored in this work. Therefore, our choice of \( \rho_{b0} \) and \( \rho_{b0} \) is arbitrary and, if incorrect, may lead to the wrong blob dynamics. Now we present the results for a blob with a very different ionization degree to ascertain the effect of this parameter. We use \( \rho_{b0} = 0.2 \times 10^{-10} \text{kg m}^{-3} \), \( \rho_{b0} = 1.8 \times 10^{-10} \text{kg m}^{-3} \), which implies that the total blob mass is equal to the one used before. With the new ionization fraction, the charged and neutral species also fall in unison because of the friction force. In addition, the blob height and velocity are indistinguishable from those obtained before (i.e., the red and black symbols in Figure 4). The same applies to the blob acceleration, which is shown in Figure 6 with an orange line. Other line colors refer to the various terms in the momentum equation. A comparison of the charged particles’ acceleration for the two ionization fractions (Figures 5(a) and 6(a)) shows that, although the total acceleration is the same, the acceleration caused by the pressure gradient and friction force are much larger in the blob with a smaller percentage of ions. This difference is much less pronounced in the case of neutrals; see Figures 5(b) and 6(b). The main conclusion of this section is that the blob ionization fraction is irrelevant in the blob dynamics.

### 3.6. Blob Maximum Falling Velocity

In Paper I we defined the density ratio as the ratio of maximum blob density to initial coronal density at the blob position, computed from Equation (20) with \( z \) substituted by the blob height at \( t = 0 \). Here the density ratio is defined, in a similar way, as the sum of the maximum neutral and ion density at \( t = 0 \) divided by the total environment density at the initial blob position, which is calculated from Equations (20)
In Paper I it was found that the density ratio determined the maximum descending speed, \(v_{\text{max}}\), which is the maximum unsigned value of \(v\); for example, in the simulation of Figure 4(b) (red and black symbols) we have \(v_{\text{max}} \approx 30 \text{ km s}^{-1}\). The variation of \(v_{\text{max}}\) with the density ratio for a fully and a partially ionized plasma is presented in Figure 7, where the values of the filled circles are taken from Paper I. The blue squares (orange triangles) correspond to numerical simulations in which the blob is 50% ionized and 50% neutral (10% ionized and 90% neutral) material. They follow the same trend of the fully ionized case. Moreover, we have seen before that the blob ionization degree has no effect on its dynamics, and so the positions of squares and triangles coincide in this plot. Therefore, the blob maximum speed is insensitive to the blob ionization degree, even in the limit of a fully ionized blob.

### 3.7. Velocity Drift

In Figure 8 we plot the drift velocity at the position of maximum density as a function of time. The various colors and line styles represent different initial blob density and ionization degree. We obtain drift velocities (of the order of 0.01–0.15 m s\(^{-1}\)) much smaller than the blob speed (of the order of 10–100 km s\(^{-1}\)). Then, over the course of a few hundred seconds the neutral and ionized fractions of the blob drift by less than 1 km, an insignificant amount. Figure 8 also shows that the drift velocity steadily increases in time until it reaches a more or less constant value. For the initial densities \(\rho_{\text{blob}} = \rho_{\text{hot}} = 10^{-10} \text{ kg m}^{-3}\) (red solid line), the time required for the drift velocity to reach this value is of the order of 400 s, after which the blob acceleration approximately vanishes, such as shown in Figures 4(b) and 5. When the initial blob peak density is increased (green and blue lines), the time needed for \(v_i - v_n\) to stabilize also increases because it takes longer to achieve a negligible blob acceleration; more details on this issue are given in Paper I. We also remark that the sign of \(v_i - v_n\) is positive and so neutrals fall faster than charged particles (in agreement with the uncoupled case of Figure 4). Hence, the former provide a downward friction force on the latter, as discussed in Section 3.3.

Finally, if the maximum blob density is kept but its ionization degree is varied (compare solid and dashed lines...
with the same color), the drift velocity suffers an increase that becomes larger for denser blobs. Such an increase in \( v_i - v_n \) helps provide an additional upward force on neutrals so that they can fall together with the ions present in the blob.

4. DISCUSSION

In this work we have studied the dynamics of a partially ionized coronal rain blob or falling prominence knot. We have extended the model of Paper I (for the fully ionized case) and have included separate mass continuity, momentum, and energy equations for the charged and neutral fractions of the plasma. These equations contain the frictional interaction between neutrals and ions, and so the coupling between the two species is only dynamic (through the friction coefficient \( \alpha_{\text{in}} \)), while ionization/recombination has been neglected. This means that the relative density of ions and neutrals in the blob is arbitrarily imposed. For this reason we have considered different ionization degrees of the blob. Our results show that this parameter is irrelevant and that the mass of the cold lump determines its dynamics. Hence, any physical mechanism that does not change the lump’s mass as it falls but varies its ionization degree is not expected to modify the falling blob behavior.

In our simple model we have also set aside some important physics. We have excluded plasma cooling and heating, together with thermal conduction and viscosity. Hence, our model does not self-consistently lead to the formation of cold coronal rain condensations through a catastrophic cooling process, and so we are forced to include a plasma clump as a localized density enhancement at \( t = 0 \). This clump then finds itself in a rather quiescent environment compared to that in which a catastrophic cooling event has just taken place. Moreover, coronal rain blobs often increase their length during their fall, which is not observed in our simulations. On the contrary, both in Paper I and in this work falling clumps tend to become shorter in time, an effect that is more perceptible for denser clumps. This can perhaps be the consequence of omitting some physical ingredients in the energy equations. Finally, since the model is one-dimensional, several fundamental plasma instabilities such as Rayleigh–Taylor and Kelvin–Helmholtz are absent. The extension of the model to two and three dimensions is crucial to understand the effects of the previous instabilities on the plasma blobs. Interestingly, partial ionization can influence the growth rates of the instabilities, as has been recently studied by Soler et al. (2012), Díaz et al. (2012, 2014), Khomenko et al. (2014b), and Martínez-Gómez et al. (2015).

The ion–neutral friction force is large enough to make the ionized and neutral fractions of the cool clump fall together. This friction force is proportional to the ion–neutral drift velocity, which for typical coronal rain and prominence knot conditions is of the order of 0.01–0.15 m s\(^{-1}\) at the center of the condensation. This is an insignificant value compared to the typical descending velocities of 10–100 km s\(^{-1}\). The main forces acting on the condensation are gravity, the pressure gradient (caused by the rearrangement of the pressure in the blob surroundings), and the friction force (which is proportional to the drift velocity). In the case of charged particles, both the pressure gradient and the friction force point upward. In the case of neutrals, however, the latter has the same value but points downward. The interplay of forces is such that both ions and neutrals experience the same acceleration regardless of their respective densities. It is worth pointing out that, in their study on the support of neutrals against gravity in solar prominences, Terradas et al. (2015) showed that ions remain static while neutrals slip down across the almost horizontal magnetic field with a small downward velocity. The drift velocities found by Terradas et al. (2015) are of the order of a few meters per second, i.e., 10 times larger than those obtained here, but rather small nonetheless. Analogous values to those in Terradas et al. (2015) were found by Gilbert et al. (2002) for the diffusion of neutral atoms in a hydrogen–helium static prominence plasma. Our derivation of Equations (8), (11), and (12) relies on the assumption that drift velocities are small compared to \( c_r \), or equivalently, \( s \ll 1 \) (see Equation (6)). Under some circumstances (Figure 3) our simulations give \( |v_i - v_n| \) of the order of a few hundred meters per second, and so it may be necessary to consider finite \( s \). The charge exchange effects for \( s > 0 \) are presented in, e.g., Meier (2011) and Meier & Shumlak (2012), and an application to solar prominences can be found in Terradas et al. (2015). An additional outcome of nonzero relative ion–neutral velocities may be the presence of the critical ionization velocity effect, by which the neutral gas can be ionized by collisions with ions; see Lai (2001).

The results obtained in Paper I (for a fully ionized condensation) are also found here. In particular, the falling material displays two separate phases: first, the blob is accelerated, and next, it maintains a practically constant speed (see red and black symbols in Figure 4), with the duration of the first phase lasting longer for a denser blob. In addition, the emission of small-amplitude sound waves at the start of the descent is present in a partially ionized clump and makes itself visible by means of the oscillations in some curves of Figures 5, 6, and 8. Finally, the important correlation between the blob maximum speed and the initial density ratio derived in Paper I remains unchanged once partial ionization is included.

We have also assessed the importance of heating by collisions and the energy interchange between charged particles and neutrals, given by the terms \( Q_i^{\text{in}} \) and \( Q_n^{\text{in}} \) in Equations (17) and (18). We have shown that an initial temperature imbalance between ions and neutrals in the blob is removed in a time of the order of 4 ms, and that despite the brevity of this interval, it leads to unrealistic pressure gradients and velocity drifts in the blob edges. Nevertheless, these unrealistic features are innocuous because they do not affect the blob dynamics. We have repeated some numerical simulations with \( Q_i^{\text{in}} = Q_n^{\text{in}} = 0 \) and have obtained identical results: the temporal variations of the blob height and velocity are the same, the accelerations caused by the forces in the momentum equation are the same, and the temporal variation of the drift velocity at the maximum density position is also unchanged. We conclude that the blob dynamics is insensitive both to the thermal structure of the blob and to the energy interplay between neutrals and ions, because both terms in the heat generation quantities \( Q_i^{\text{in}} \) and \( Q_n^{\text{in}} \) (see their definition in Equations (11) and (12)) are negligible. The first one, caused by heating because of friction between the two species, is proportional to \( (v_i - v_n)^2 \), and we have seen that the drift velocity is very small. Moreover, the second term, which accounts for the thermal energy transfer, is proportional to the temperature difference between the two species, and, as mentioned before, \( T_i - T_n \) vanishes very quickly and remains zero during the whole temporal evolution.
A final comment about the interaction between the charged and neutral portions of the condensation is needed. Here this interaction has been found to be strong enough to tie together ions and neutrals when the clump trajectory is vertical. Coronal rain blobs usually follow curved paths, so one may wonder whether the friction force provided by ions would be sufficient to drag neutrals along the curved magnetic field, or else neutrals would slip through the field lines and would separate from the ions (contradicting the observations). This question cannot be answered with the one-dimensional model of this work, and two-dimensional numerical simulations as in Terradas et al. (2015) are required to ascertain whether the physics of a partially ionized blob would be qualitatively the same in the case of a trajectory guided by curved magnetic fields.

J.T. acknowledges support from the Spanish Ministerio de Educación y Ciencia through a Ramón y Cajal grant. R.S. acknowledges support from MINECO and UIB through a “Ramón y Cajal” grant (RYC-2014-14970). R.O., R.S., and J.T. acknowledge support from MINECO and FEDER funds through project AYA2014-54485-P. The work of T.Z. was supported by the Austrian Fonds zur Förderung der Wissenschaftlichen Forschung (FWF) project P26181-N27, by the European FP7-PEOPLE-2010-IRSES-269299 project SOLSPANET, and by Shota Rustaveli National Science Foundation grant DI/14/6-310/12. The authors are also grateful to The Leverhulme Trust for funding under grant IN-2014-016 and thank the International Space Science Institute for the financial support and the facilities for two meetings on partially ionized plasmas in astrophysics and on coronal rain observations and modeling. The authors thank P. Antolin and D. Martinez-Gomez for useful comments that helped improve the manuscript. R.O. is indebted to D. W. Fanning for making available the Coyote Library of IDL programs (http://www.idlcoyote.com/). We are indebted to the referee for his/her detailed reading of our paper and for his/her very useful comments.

REFERENCES

Ahn, K., Chae, J., Cho, K.-S., et al. 2014, SoPh, 289, 4117
Antolin, P., & Rouppe van der Voort, L. 2012, ApJ, 745, 152
Antolin, P., Shibata, K., & Vissers, G. 2010, ApJ, 716, 154
Antolin, P., & Verwichte, E. 2011, ApJ, 736, 121
Antolin, P., Vissers, G., Pereira, T. M. D., Rouppe van der Voort, L., & Scullion, E. 2015, ApJ, 806, 81
Antolin, P., Vissers, G., & Rouppe van der Voort, L. 2012, SoPh, 280, 457
Beckers, J. M. 1972, ARA&A, 10, 73
Braginskii, S. I. 1965, RvPP, 1, 205
Chae, J. 2010, ApJ, 714, 618
Chapman, S., & Cowling, T. G. 1970, The Mathematical Theory of Non-uniform Gases. An Account of the Kinetic Theory of Viscosity, Thermal Conduction and Diffusion in Gases (3rd ed.; Cambridge: Cambridge Univ. Press)
De Groof, A., Bastiaensen, C., Müller, D. A. N., Berghmans, D., & Poedts, S. 2005, A&A, 443, 319
De Groof, A., Berghmans, D., van Driel-Gesztelyi, L., & Poedts, S. 2004, A&A, 415, 1141
Díaz, A. J., Khomenko, E., & Collados, M. 2014, A&A, 564, A97
Díaz, A. J., Soler, R., & Ballester, J. L. 2012, ApJ, 754, 41
Draine, B. T. 1986, MNras, 220, 133
Draine, B. T., Roberge, W. G., & Dalgarno, A. 1983, ApJ, 264, 485
Engold, V. 1976, SoPh, 49, 283
Fang, X., Xia, C., & Keppens, R. 2013, ApJL, 771, L29
Fang, X., Xia, C., Keppens, R., & van Doorsselaere, T. 2015, ApJ, 807, 142
Gerhardt, S. P. 2004, PhD thesis, Univ. Wisconsin–Madison
Gilbert, H. R., Hansteen, V. H., & Holzer, T. E. 2002, ApJ, 577, 464
Harra, L. K., Matthews, S. A., Long, D. M., Doschek, G. A., & De Pontieu, B. 2014, ApJ, 792, 93
Khomenko, E., Collados, M., Díaz, A., & Vitas, N. 2014a, PhP, 21, 092901
Khomenko, E., Díaz, A., de Vicente, A., Collados, M., & Luna, M. 2014b, A&A, 565, A45
Klimchuk, J. A. 2015, RSPTA, 373, 40256
Krstić, P. S., & Schultz, D. R. 1999, PhRvA, 60, 2118
Lai, S. T. 2001, RvGeo, 39, 471
Leake, J. E., Lukin, V. S., & Linton, M. G. 2013, PhPl, 20, 061202
Leake, J. E., Lukin, V. S., Linton, M. G., & Meier, E. T. 2012, ApJ, 760, 109
Leake, J. E., DeVore, C. R., Thayer, J. P., et al. 2014, SSRv, 184, 107
Liu, W., Berger, T. E., & Low, B. C. 2012, ApJL, 745, L21
Martínez-Gómez, D., Soler, R., & Terradas, J. 2015, A&A, 578, A104
Meier, E. T. 2011, PhD thesis, Univ. Washington
Meier, E. T., & Shumlak, U. 2012, PhPl, 19, 072508
Moschou, S. P., Keppens, R., Xia, C., & Fang, X. 2015, AdSpR, 56, 2738
Müller, D. A. N., De Groof, A., Hansteen, V. H., & Peter, H. 2005, A&A, 436, 1067
Müller, D. A. N., Peter, H., & Hansteen, V. H. 2004, A&A, 424, 289
Oliver, R., Soler, R., Terradas, J., Zaqarashvili, T. V., & Khodachenko, M. L. 2014, ApJ, 784, 21
Pinto, C., & Galli, D. 2008, A&A, 484, 17
Pneuman, G. W., & Kopp, R. A. 1978, SoPh, 57, 49
Schrőjer, C. J. 2001, SoPh, 198, 325
Scullion, P., Rouppe van der Voort, E., Wedemeyer, L., & Antolin, S. 2014, ApJ, 797, 36
Smith, P. D., & Sakai, J. I. 2008, A&A, 486, 569
Soler, R., Ballester, J. L., & Zaqarashvili, T. V. 2015, A&A, 573, A70
Soler, R., Carbonell, M., Ballester, J. L., & Terradas, J. 2013, ApJ, 767, 171
Soler, R., Díaz, A. J., Ballester, J. L., & Goossens, M. 2012, ApJ, 749, 163
Terradas, J., Soler, R., Oliver, R., & Ballester, J. L. 2015, ApJL, 802, L28
Vashalomidze, Z., Kukhiandze, V., Zaqarashvili, T. V., et al. 2015, A&A, 577, A136
Vranjes, J., & Krstić, P. S. 2013, A&A, 554, A22
Willk, J. E., Schmieder, B., Heinzel, P., & Roudier, T. 1996, SoPh, 166, 89
Zaqarashvili, T. V., Khodachenko, M. L., & Rucker, H. O. 2011a, A&A, 529, A82
Zaqarashvili, T. V., Khodachenko, M. L., & Rucker, H. O. 2011b, A&A, 534, A93
Zhang, J., & Li, L.-P. 2009, RA&A, 9, 1368