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Weakly damped Alfvén waves as drivers for spicules

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Abstract. We present an analytical model for the damping of Alfvén waves in the partially ionized chromosphere. The damping is due to collisions between ions and neutrals. The ion-neutral collision frequency in this environment is such that the ion and neutral populations are almost perfectly collisionally coupled, leading the Alfvén wave to behave as if it acts on the *whole* plasma (i.e. including neutrals). The small but finite coupling time between ions and neutrals leads to damping of the Alfvén waves. We find that this type of damping of upward traveling Alfvén waves with frequencies between 0.2 and 0.6 Hz, can cause not only significant heating but also upward motion of the upper chromospheric plasma. In addition the upward force and heating associated with this type of damping can sustain, both dynamically and thermodynamically, an already formed chromospheric spicule. The energy flux carried by the Alfvén waves needed for this type of support of a spicule does not seem to be in contradiction with observational and theoretical evidence for the presence of Alfvén waves in the chromosphere.

Key words: Sun: chromosphere – Sun: transition region – MHD – waves

1. Introduction

The solar chromosphere has a very inhomogeneous structure, determined mostly by the magnetic field. The transition from the chromosphere to the corona is dominated by spicules, very dynamic structures which are observed at the solar limb. Their probable counterparts on the solar disk, the so-called dark mottles, appear above the magnetic field concentrations at the edges of the supergranular cells.

Spicules are jets of chromospheric nature, that protrude into the corona. With their diameters of 300 to 1500 km (0.4” to 2”), spicules are observationally challenging. However, many spicule properties now seem well established (Beckers 1972). During their lifetime of five to ten minutes, they reach heights of 5000 to $10\,000$ km above the photosphere, i.e., the spicular plasma flows with upward velocities of 20 to 30 km/s. The motion of the plasma in a spicule appears to be different from a purely ballistic one. Some of the spicules fade out of view after having reached maximum height, whereas some fall back to the chromosphere. Spicules carry a mass flux into the corona that is one hundred times larger than the coronal mass loss due to the solar wind. Both the temperature (5000 to $15\,000$ K) and the density ($3\,10^{10}$ – $1.5\,10^{11}$ cm^{-3}) seem to show little variation along the length of a spicule. Given their location and the propagation speed of velocity disturbances (Alfvénic in nature), it seems certain the driving mechanism is linked to the magnetic field (Beckers 1968, Beckers 1972).

Spicules have received considerable attention from theorists over the years. However, most theoretical models and the numerical simulations based on them have difficulties explaining or reproducing all of the observed properties of spicules (Hollweg, 1992).

In this paper we will elaborate on an idea for a spicular driving mechanism proposed by Haerendel (1992a). This driving mechanism is based on the presence of upward traveling Alfvén waves in the partially ionized chromosphere. As we will show below, the damping of these Alfvén waves due to the collisions between the ions and the neutral particles can not only heat the chromospheric plasma, it can also provide a driving force to accelerate the matter upwards. The heating effect through ion-neutral damping of Alfvén waves was first proposed by (Piddington 1956) and is also studied in tokamak studies (Amagashi & Tanaka 1993). Osterbrock (1961) analyzed ion-neutral damping of Alfvén waves traveling through the chromosphere. Hartmann & MacGregor (1980) studied damping of Alfvén waves in stellar winds. Haerendel (1992a) was the first who drew attention to the momentum this type of damping can impart on the *solar* plasma along the magnetic field lines, so that the magnetic force can even balance the gravitational force.

Several spicular models based on Alfvén waves have been studied in the past. Hollweg et al. (1982) developed a numerical model which is based on an upward propagating Alfvénic pulse that steepens into a fast shock by non-linear coupling of the transverse wave into motions parallel to the background magnetic field. This fast shock interacts with the transition region and drives it upward, thus forming a spicular structure. Mariska & Hollweg (1985) and Hollweg (1992) discovered that the resulting structures are too low and too cold to be spicules, despite shock heating by the two acoustic-gravity pulses caused by the initial passage of the Alfvénic pulse. Related work by Ster-

Send offprint requests to: B. De Pontieu ling & Hollweg (1984) comes to the conclusion that already formed spicules may act as a resonant cavity for Alfvén waves, which are consequently dissipated via a turbulent cascade of their energy to higher wavenumbers and which could thus heat the spicular matter.

In the current work we employ a linearized and stationary approach and study a *train* of Alfvén waves with a frequency between 0.2 and 0.6 Hz, much higher than in the work referenced above. Whereas previous work did not take into account any type of damping of the Alfvén waves, we do include damping of the Alfvén waves due to collisions between ions and neutrals. We will show that upward-travelling, weakly-damped Alfvén waves with frequency f between 0.2 and 0.6 Hz provide enough upward momentum and heat to maintain an already existing spicule. A fully non-linear and time-dependent calculation will be studied in a future paper.

2. Analytical approach

First, we calculate the effect of collisions between ions and neutrals on the propagation of an Alfvén wave in a homogeneous plasma. We consider a plasma consisting of three species: electrons, ions and neutrals. Each of these species obeys a momentum conservation law of the form:

$$m_j n_j \left[\frac{\partial \boldsymbol{v_j}}{\partial t} + (\boldsymbol{v_j}.\boldsymbol{\nabla})\,\boldsymbol{v_j}\right] = e\,n_j\,(\boldsymbol{E} + \boldsymbol{v_j}\times\boldsymbol{B}) - \nabla p_j - m_j\,n_j\,(\boldsymbol{v_j} - \boldsymbol{v_k})\,\nu_{jk} + m_j\,n_j\,(\boldsymbol{v_j} - \boldsymbol{v_l})\,\nu_{jl}, \quad (1)$$

with j, k and l the species indices (= e, i or n), m the mass of a particle of that species, n the number density, $\boldsymbol{v}$ the velocity, e the elementary charge, $\boldsymbol{E}$ the electric field, $\boldsymbol{B}$ the magnetic field and ν_{jk} the collision frequency for elastic collisions between species j and k. The average momentum transfer collision frequency between species 1 and 2 is defined as:

$$\nu_{12} = \frac{m_2}{m_1+m_2}\, n_2 \left[\frac{8kT}{\pi m_{12}}\right]^{\frac{1}{2}} \sigma_{12}, \quad (2)$$

with $m_{12} = \frac{m_1\,m_2}{m_1+m_2}$, σ_{12} the collisional cross-section for collisions between the two species. It is assumed that the species have equal temperatures (Giraud & Petit 1978). From the momentum balance between both species, we know:

$$n_1 m_1 \nu_{12} = n_2 m_2 \nu_{21} \quad (3)$$

since $\sigma_{12} = \sigma_{21}$. Some text books define the collision frequency without the leading mass ratio $m_2/(m_1+m_2)$, with appropriate changes in the equations (Mitchner & Kruger 1973).

We now study the propagation of a linearly polarized Alfvén wave in a plane geometry, in the limit of small wave perturbations, i.e., we linearize all relevant equations:

$$\boldsymbol{E} = E_{1x}\,\boldsymbol{e_x}, \quad (4)$$

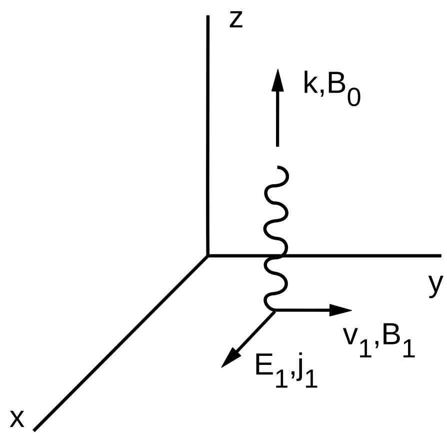

Fig. 1. Geometry of a linearly polarized Alfvén wave

$$\boldsymbol{B} = B_0\,\boldsymbol{e_z} + B_{1y}\,\boldsymbol{e_y}, \quad (5)$$

$$\boldsymbol{v_j} = v_{1jx}\,\boldsymbol{e_x} + v_{1jy}\,\boldsymbol{e_y}, \quad (6)$$

with j the species index, B_0 the background magnetic field, B_1 the perturbation magnetic field, E_1 the perturbation electric field and v_1 the perturbation velocity. For the geometry, see Fig. 1.

We further assume that the plasma is cold, since thermal motions are not important for Alfvén waves ($T_{i,e,n} = 0$). Since we are interested in waves with frequencies for which the slippage between ions and neutrals is important, we neglect the collisional terms between electrons and the other species. This can be accounted for as follows. As will be demonstrated, the slippage is proportional to ω/ν_{ni} with $\omega = 2\pi f$. Since the collision frequencies between electrons and the other species are much higher than the ion-neutral collision frequency, neglecting the collisions with electrons is warranted. Finally, we assume that $m_i = m_n$, i.e., we study ions and neutrals of the same atomic species.

We study a uniform plasma, and enter the wave ansatz for a plane polarized wave propagating in the z-direction:

$$(E_1, B_1, v_1) = (E_{1x}, B_{1y}, v_{1x,y})\, e^{i(k\,z-\omega t)}, \quad (7)$$

with k the wave vector and ω the wave frequency.
Linearizing the equations of momentum conservation for ions, electrons and neutrals, we find

$$v_{1nx} = \frac{v_{1ix}}{1-\frac{i\omega}{\nu_{ni}}}, \quad (8)$$

$$v_{1ix} = \frac{a\Omega_i}{a^2+{\Omega_i}^2}\frac{E_{1x}}{B_0}, \quad (9)$$

$$v_{1iy} = \frac{-{\Omega_i}^2}{a^2+{\Omega_i}^2}\frac{E_{1x}}{B_0}, \quad (10)$$

in which

$$a = a_2 - i\omega a_1, \quad (11)$$

$$a_1 = \frac{1+\tau^2+\frac{n_n}{n_i}}{1+\tau^2}, \quad (12)$$

$$a_2 = \frac{\omega\tau\frac{n_n}{n_i}}{1+\tau^2}, \quad (13)$$

$$\tau = \frac{\omega}{\nu_{ni}}, \quad (14)$$

and Ω_i the ion-gyrofrequency. Since we consider ${\Omega_e}^2 \gg \omega^2$ (Ω_e the electron-gyrofrequency) the electrons are completely magnetized and $v_{1ex} \rightarrow 0$ and $v_{1ey} = -\frac{E_{1x}}{B_0}$. This corresponds to neglecting the Larmor gyrations of the electrons, so that the only electron motion left is an $\boldsymbol{E} \times \boldsymbol{B}$ drift. It is now easy to calculate the current j_{1x} carried by the wave :

$$j_{1x} = n_i e(v_{1ix} - v_{1ex}). \quad (15)$$

From Faraday's induction law and Ampère's law, we derive the wave equation (assuming the background magnetic field $\boldsymbol{B_0}$ is homogeneous and time-independent) :

$$-\nabla \times (\nabla \times \boldsymbol{E_1}) = \mu_0 \frac{\partial \boldsymbol{j_1}}{\partial t} + \frac{1}{c^2}\frac{\partial^2 \boldsymbol{E_1}}{\partial t^2}. \quad (16)$$

Given the wave ansatz (7), we can reduce (16) to :

$$(\omega^2 - c^2k^2)E_{1x} = -i\omega\mu_0 c^2 j_{1x}. \quad (17)$$

Substituting (9) into (15) gives

$$j_{1x} = n_i e\Omega_i \frac{(a_2 - i\omega a_1)}{a^2 + {\Omega_i}^2}\frac{E_{1x}}{B_0}. \quad (18)$$

We now limit our study to waves for which $\tau \ll 1$, i.e., $\omega \ll \nu_{ni}$. Fig. 2 shows ν_{ni} as a function of height above the photosphere. The neutral-ion collision frequency in Fig. 2 is calculated using (2) for neutral hydrogen with protons and ionized helium, in a solar atmosphere identical to the VAL IIIC semi-empirical model of Vernazza et al. (1981). It is clear from Fig. 2 that for waves with $\nu \leq 1$ Hz, the assumption $\tau \ll 1$ holds throughout the chromosphere. Limiting our study to $\tau \ll 1$ does thus not constitute much of a restriction to the practical applicability of our results, since most theories for the generation of Alfvén waves in the solar atmosphere predict typical frequencies below 1 Hz.

Thus taking $\tau \ll 1$, we know that all terms in a^2 are of the order ω^2 or smaller. Alfvén waves have frequencies well below the ion gyration frequency ($\omega^2 \ll {\Omega_i}^2$). Using this assumption in (18), and neglecting all terms of order τ^2, we find

$$j_{1x} = \frac{n_i e}{\frac{eB_0}{m_i}}\left[\omega\tau\frac{n_n}{n_i} - i\omega\left(1 + \frac{n_n}{n_i}\right)\right]\frac{E_{1x}}{B_0}. \quad (19)$$

Alternatively:

$$j_{1x} = \frac{\rho_{tot}}{{B_0}^2}\left[-i\omega + \frac{\rho_n}{\rho_{tot}}\omega\tau\right] E_{1x}, \quad (20)$$

where ρ_n is the mass density of neutrals and ρ_{tot} the total mass density of the plasma. The first part in the brackets ($-i\omega$) is the inertial part of the current, 90 degrees out of phase with the electric field. The second term, however, is due to the collisions between ions and neutrals. It is in phase with the electric field and will thus lead to Ohmic dissipation, unlike the inertial current. The dissipation of the electrical current associated with the Alfvén wave is proportional to the ionization ratio ρ_n/ρ_{tot} and the ratio of wave and collision frequency $\frac{\omega}{\nu_{ni}}$. Also note that in (20), the mass density involved in the inertial current is the total mass density ρ_{tot}, whereas one would have expected it to be the ion mass density since only ions can carry electrical current. This is because the ions and neutrals are collisionally coupled to such a degree ($\omega \ll \nu_{ni}$) that the motion of the ions, as caused by the varying electromagnetic field associated with the Alfvén wave, is forced upon the neutral particles *almost* immediately. The plasma, including neutrals, thus reacts to the Alfvén wave almost as a whole. However, the neutral-ion coupling time ν_{ni}^{-1}, though tiny, is not exactly zero. There is a small amount of slippage between the ion and neutral populations. This slippage leads to the dissipation of the Alfvénic current through the perturbed motion of the ions - as evident from (20) - and thus to heating of the plasma.

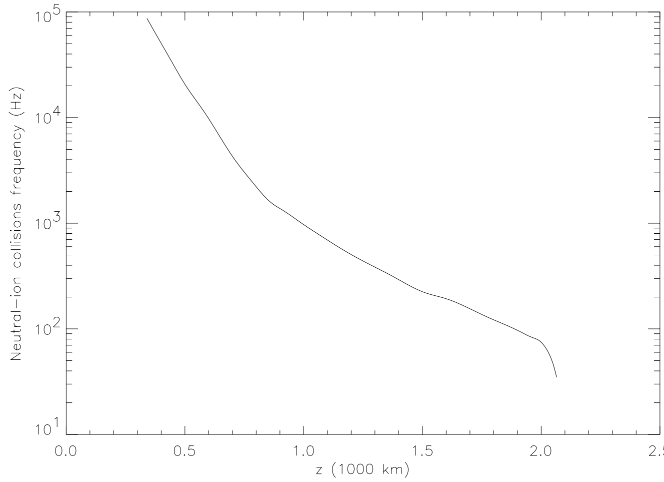

Fig. 2. Neutral-ion collision frequency (Hz) as a function of height above the photosphere (1000 km).

This is confirmed by solving (17) for the wave phase speed v_{ph} (using (20) and neglecting terms of the order τ^2) :

$$v_{ph} = \frac{\omega}{k} = v_A\left(1 - i\frac{\rho_n}{\rho_{tot}}\frac{\omega}{\nu_{ni}}\right)^{\frac{1}{2}}, \quad (21)$$

in which $v_A = B_0/\sqrt{\mu_0\rho_{tot}}$, again indicating that the plasma as a whole takes part in the Alfvénic motions. The imaginary term indicates that the wave is damped.

To calculate the damping length and the heating and momentum transfer, we apply the WKB approximation. The above calculations are for a homogeneous medium. Since the wave is being damped as it propagates, it is certain that at least some quantities have a gradient in the direction of propagation. Moreover, we also know that the chromosphere is not homogeneous, because of the density and temperature gradients with height above the photosphere. If we now assume that the changes take place over scales much larger than one wavelength, we can use the WKB approximation (Swanson 1989). We again start from

(16), but this time we retain the spatial derivation ∇ explicitly, and use the wave ansatz (7) only partially (i.e. only the temporal part). We also use (20) and find that

$$\frac{\partial^2 E_{1x}}{\partial z^2} \;+\; ({k_1}^2 + i{k_2}^2)E_{1x} \;=\; 0, \tag{22}$$

in which ${k_1}^2 \;=\; \frac{\omega^2}{{v_A}^2}$ and ${k_2}^2 \;=\; {k_1}^2 \frac{\tau \rho_n}{\rho_t}$. We will show below that the typical damping lengths are much larger than one wavelength. The other condition for the validity of the WKB approximation is that the wavelength should not change (due to changes in the background parameters) much over one wavelength :

$$|\frac{1}{k}\frac{\partial k}{\partial z}| \ll \; k. \tag{23}$$

This reduces to $\frac{\partial v_A}{\partial z} \;\ll\; \omega$. We will for now assume that this is true in the chromosphere and for spicules, but we will come back to the validity of this assumption in the discussion.
We now use the ansatz

$$E_{1x}(z) \;=\; A(z)\; e^{iS(z)}, \tag{24}$$

where $A(z)$ is assumed to vary slowly and $S(z)$ may vary rapidly ($S(z) = kz$ in a homogeneous medium). If we insert (24) in (22), we get

$$\left[A'' + 2iA'S' + iAS'' - AS'^2 + {k_1}^2 A + i{k_2}^2 A\right] e^{iS} = 0. \tag{25}$$

where the derivation symbol $'$ stands for a spatial derivative. The dominant terms are the ones that do not contain derivatives of A. For the real terms, we thus get

$${k_1}^2 \;=\; S'^2. \tag{26}$$

Integrating (26), we get

$$S \;=\; \int^z k_1 dz. \tag{27}$$

Looking at the imaginary terms of (25), we get:

$$2i\,S'\,A' \;+\; iS''\,A \;+\; i{k_2}^2\,A \;=\; 0. \tag{28}$$

Using the result of (27) and integrating this equation, gives

$$A \;=\; A_0\sqrt{\frac{k_{1o}}{k_1}}\; e^{-\int^z \frac{{k_2}^2}{2k_1}dz}, \tag{29}$$

where the o subscript refers to conditions at the bottom of the atmosphere. We now see from (29) that the amplitude of the electric field associated with the Alfvén wave suffers an exponential attenuation with height, due to the ${k_2}^2/2k_1$ term. Assuming the background changes over one wavelength are small, we can calculate a local damping length L from (29):

$$L \;=\; \frac{2\,k_1}{{k_2}^2} \;=\; \frac{2\,v_A\,\nu_{ni}\,\rho_{tot}}{\omega^2\,\rho_n}. \tag{30}$$

Note that the damping length L_f for the wave energy density is equal to $\frac{L}{2}$.

Let us now calculate the wave magnetic field B_{1y}. From Faraday's induction law and using the temporal wave ansatz and (23), we find

$$B_{1y} = \frac{1}{i\omega}\frac{\partial E_{1x}}{\partial z} \;=\; \frac{E_{1x}}{v_A}(1 \,+\, i\frac{\omega}{2\nu_{ni}}\frac{\rho_n}{\rho_{tot}}). \tag{31}$$

We can now calculate the Lorentz force averaged over one wave period $\overline{F_z}$, from (31), (24) and (20):

$$\overline{F_z} = \overline{j_x\,B_y} = \rho_n\frac{\overline{{E_{1x}}^2}}{{B_0}^2}\frac{\omega^2}{2\nu_{ni}v_A}. \tag{32}$$

Since $\overline{F_z}$ depends on the square of E_{1x}, the damping length of this force is L_f, not L. From (10) we know that for magnetized ions $E_{1x} = -v_{1y}B_0$, so that (32) becomes:

$$\overline{F_z} = \rho_n\overline{{v_{1y}}^2}\frac{\omega^2}{2\nu_{ni}v_A}. \tag{33}$$

This result is a factor of two smaller than that found by Haerendel (1992a,1992b) in his calculations. The error in Haerendel's original calculations was caused by not taking the complex conjugate of B_y in (32).

In an undamped Alfvén wave, $\overline{F_z}$ is 0. The fact that it does not average out to 0 in (33) is due to damping by the ion-neutral collisions. That is why $\overline{F_z}$ depends on the neutral density ρ_n and is inversely proportional to ν_{ni}. The net force is always in the direction of the wave propagation, since a negative gradient of transverse magnetic pressure is set up in that direction. Through this mechanism momentum is transferred from the transverse direction to the direction of wave propagation.
It is now straightforward to calculate the energy dissipation rate averaged over one wave period :

$$\dot{\epsilon_w} = \overline{j_{1x}E_{1x}} = \rho_n\overline{v_{1y}^2}\frac{\omega^2}{\nu_{ni}} = 2v_A\overline{F_z}. \tag{34}$$

The factor of two in the last equation arises because both kinetic and magnetic energies are dissipated, but only the magnetic energy corresponds to a pressure gradient that can impart momentum to the gas.

3. Alfvén wave damping in the chromosphere and in spicules

We now study the effects of Alfvén wave damping on the chromosphere and spicules. For this we employ a thermodynamical model for the solar atmosphere based on the VAL IIIC semi-empirical model, as mentioned in the preceding section. We consider a magnetic field typical for the regions of flux concentration at the walls of the supergranulation cells. For this magnetic field model we take a flux tube with $B = 1500$ G at photospheric level, smoothly falling off with height to a coronal field value of 10 G (similar to (Haerendel 1992a), see (De Pontieu 1996) for details).

From (30) we calculate the damping length L_f as a function of height above the photosphere. L_f is plotted in Fig. 3 for a

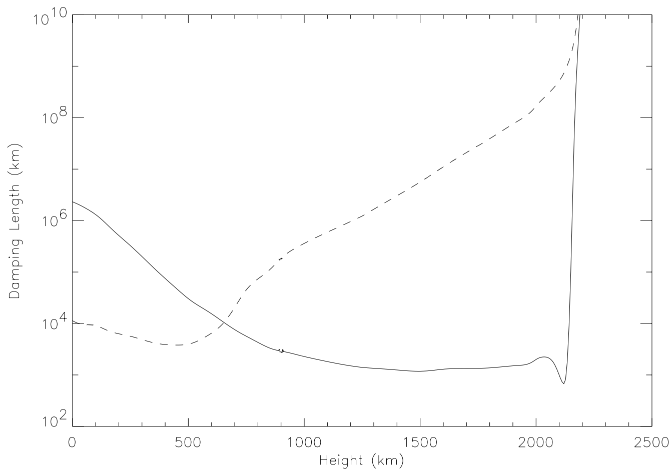

Fig. 3. WKB damping lengths due to Joule (dashed line) and ion-neutral damping (full line) as a function of height above the photosphere (1000 km); for an Alfvén wave with $f = 0.5$ Hz.

plane Alfvén wave with $f = 0.5$ Hz. Also plotted is the damping length L_j for normal Ohmic dissipation, i.e., due to the electrical conductivity depending on the collisional rate of electrons with ions and neutrals. The Joule damping length L_j is defined by the Ohmic dissipation :

$$L_j = v_A \, \lambda^2 \, \sigma \mu_0, \tag{35}$$

where λ is the wavelength of the wave and σ is the electrical conductivity. Osterbrock calculated damping lengths for both Joule and ion-neutral damping, but he used an unrealistic model for both the thermodynamical parameters of the solar atmosphere and for the magnetic field (Osterbrock 1961).

It is obvious from Fig. 3 that the Joule damping is negligible throughout both photosphere and chromosphere, as the damping length does not become smaller than 7000 km (at the temperature minimum). So, though Joule damping dominates ion-neutral damping below 600 km, it can effectively be neglected for the waves under consideration. Ion-neutral damping, on the other hand, becomes important just below the transition region, where the damping length is less than 1000 km, and plasma densities become low enough to be influenced by the wave force. In the corona, the electrical conductivity is so high that Joule damping is almost nonexistent. Additionally, coronal temperatures are so high that no neutrals are present in the corona, hence ion-neutral damping is nonexistent as well.

Fig. 3 thus shows that Alfvén waves with $f \approx 0.5$ Hz can reach the middle chromosphere without too much damping. In the middle to upper chromosphere damping by ion-neutral collisions becomes more important and the upward force $\overline{F_z}$ and accompanying heating rate $\dot{\epsilon_w}$ can influence the plasma there.

The latter is confirmed by calculating $\overline{F_z}$ and $\dot{\epsilon_w}$ for a typical spicular environment, which is very similar to upper chromospheric conditions. In doing so, we follow Haerendel (1992a), using the corrected expression (33) for the upward force $\overline{F_z}$. We use the same parameters for a typical spicule as Haerendel took: $B_0 = 10\,G$, $n_i = 3.5 \cdot 10^{16}$ m^{-3}, $n_n = 1.5 \cdot 10^{16}$ m^{-3} and $\nu_{ni} \approx 70$ Hz. If we equalize the gravity and the force exerted on the plasma by the wave, we get

$$\overline{v_{1y}^2} = \left(\frac{17.8 \text{ kms}^{-1}}{f} \right)^2 , \tag{36}$$

where $f = \frac{\omega}{2\pi}$ is the frequency of the wave. If we want the wave velocity amplitude to stay below v_A (taken to be 100 km/s in this model), the frequency f must exceed 0.18 Hz. On the other hand, the Alfvén wave should not be damped over too small a distance, since the length of spicules can be as high as 7000 km (measured above 2000 km).

$$L = 2 \, \frac{\rho_{tot}}{\rho_n} \frac{\nu_{ni} v_A}{\omega^2} \approx 3500 \text{ km}. \tag{37}$$

This implies $f \leq 0.6$ Hz. So, there seems to be a range of frequencies in which damped Alfvén waves can support spicular matter against gravity.

Of course, these values are all highly dependent on the chosen parameters. From (33) we can see that for a constant ratio $\overline{v_{1y}^2}/{v_A}^2$, the net upwards force, $\overline{F_z}$, is proportional to the magnetic field. If the coronal magnetic field in our model is taken as 30 G instead of 10 G, which does not seem unrealistic, the force could be three times as big! Another parameter $\overline{F_z}$ is very sensitive to, is the ionization ratio. In the spicular environment, which is in a state of non-LTE, calculating the ionization degree is fairly model dependent.

A final consistency check comes from the energy balance. The force $\overline{F_z}$ is needed to overcome gravity, so we can easily estimate the energy input from the wave damping :

$$\dot{\epsilon_w} \approx \; 2 \, v_A \rho g_\odot = 4.4 \cdot 10^{-3} \text{ J m}^{-3} \text{ s}^{-1} \, . \tag{38}$$

with $g_\odot$ the solar gravity at photospheric level. This is the heat input to the plasma from the wave dissipation. The heat input must be able to balance the losses due to radiative cooling. Using an approximative formula for the radiative losses of the form

$$\dot{\epsilon_r} \; = \; n_e \, n_H \, Q(T), \tag{39}$$

where Q(T) is an empirical formula for the temperature dependence of optically thin losses (see (Rosner et al. 1978) for a detailed form). We can now balance $\dot{\epsilon_w}$ with $\dot{\epsilon_r}$ and find that the equilibrium temperature is of the order 15000 K. This is in good agreement with the observations, taking into account that our estimates are rather rough. Radiative losses in spicules are probably more chromospheric in nature and not necessarily as severe as the optically thin losses indicate (Anderson & Athay 1989). This does not invalidate the results found above. The heating depends on ion-neutral damping, and thus on the presence of neutrals. Excessive heating not balanced by radiative losses would lead to ionization of neutrals, less damping and thus less heating. This 'negative feedback' in the heating process ensures that the wave damping is ultimately adjusted to an equilibrium with whatever type of radiative losses is present. Additionally, the partially ionized chromospheric plasma acts in a very isothermal way, such that much of the energy input due to the wave damping is put into ionizing the neutral population instead of into increasing the temperature. We will study this dependence of our results on the type of radiative losses in more detail in a subsequent paper.

Our analytical approach thus indicates that ion-neutral damping of Alfvén waves can support spicular structures, in both a dynamic and thermodynamic sense. Of course, time-dependent calculations are needed to see whether this mechanism can also form spicular structures. We will explore the formation of spicular structures in a series of articles about the results of numerical simulations based on the ideas developed in this article.

4. Discussion

4.1. Alfvén waves in the solar atmosphere

Our proposed mechanism for spicular formation is based on the following sequence of events. A train of Alfvén waves is generated below the middle chromosphere. As the waves travel upward on a vertical flux tube through the partially ionized middle to upper chromosphere, they are damped because of collisions between ions and neutrals. This type of damping results in a net upward-directed force on the chromospheric plasma. It also produces enough heat liberation to sustain the observed radiation losses in spicules. Our driving mechanism is based on the generation and presence in the chromosphere of Alfvén waves with frequencies between 0.2 and 0.6 Hz and an energy flux of about 10^3 J m^{-2} s^{-1}.

Alfvén waves, or more precisely Alfvénic disturbances, have been observed in the solar wind by various spacecraft (of which the closest was at 0.3 AU from the sun). Most of the energy in these waves is contained in oscillations with long periods of several hours (Hollweg 1991, Belcher & Davis, 1971). Radio observations of Faraday rotation fluctuations closer to the sun indicate that there are also Alfvén waves just a few solar radii above the photosphere. It thus seems probable that these Alfvén waves are of solar origin. If that is the case, then it is not clear why the observed periods are so large. Perhaps the long periods are determined by the time scales of mesogranulation (Hollweg 1991). It seems most probable that the ultimate source of energy for waves in the solar atmosphere lies in the convective motions in and below the photosphere. However, it cannot be completely ruled out that the waves are due to reconnection events above the convective zone (Axford & McKenzie 1991, McKenzie et al. 1995). In several papers, Hollweg has also studied how high the Alfvén wave energy flux at the base of the solar atmosphere is, extrapolating from observed fluxes in the solar wind. This analysis is complicated because parts of the solar atmosphere act as resonant cavities for the waves, only transmitting waves with certain resonant frequencies (Hollweg 1978). Nevertheless, Hollweg has suggested that the solar wind fluxes are but a remnant of a much larger coronal flux. If so, the high-frequency portion might be detectable by some means (Zirker 1993).

There have been various reports of fluctuations in the intensity of coronal lines with periods as short as 0.5 s (Zirker 1993, Pasachoff 1991). Line intensity fluctuations indicate density variations, though, and Alfvén waves are incompressible, so these fluctuations, if real, are probably caused by magnetoacoustic waves. Nevertheless, these observations prove that there is some wave power generated in high frequencies.

Stronger evidence for the existence of Alfvén waves comes from non-thermal broadening of coronal and transition region lines. These observations are consistent with the presence of Alfvén waves with velocity amplitudes of 30 km/s at the coronal base (Zirker 1993). A coronal flux of 10^4 J m^{-2} s^{-1} for an active region and an *average* value of 10^3 J m^{-2} s^{-1} for the quiet sun can be derived from these observations. Giovanelli also claimed to have observed Alfvén waves in fibrils and supergranules, but this has not been confirmed (Giovanelli 1975).

In conclusion, it seems that from the observational side, the Alfvén waves that are needed for our driving mechanism cannot be ruled out. At the same time, it is generally difficult if not impossible to unambiguously identify Alfvén waves with such high frequencies.

From the theoretical side Alfvén waves cannot be ruled out either, even if, in general, not too much is known about their excitation. Osterbrock (1961) discussed the excitation of magnetosonic waves in the convection zone and their subsequent propagation upwards. They are refracted away from the vertical and develop into shocks as they travel upward, and couple with Alfvén waves, according to Osterbrock (1961). It is interesting to note that the refracted fast waves could cause (Alfvén) waves with resonant frequencies of the order $2R/v_A$ where R is the typical radius of a flux tube. Recent observations of flux tubes (Title et al. 1992, Berger et al. 1995) find that they could have diameters as small as 150 km. With $v_A = 40$ km/s (in the lower chromosphere) this could produce waves with periods as low as 5 seconds. Resonant wave excitation is likely to have substantial higher harmonics, which is also true for a breaking shock wave (Haerendel 1992b).

Stein & Nordlund (1981) performed simulations of the generation of magnetoacoustic waves by convection. They find that the convective flow generates acoustic waves which propagate upward until they reach the height (in the middle chromosphere) where the photospheric flux tubes spread out to fill the entire volume (the magnetic canopy). At this point, about 1/3 of the acoustic flux is converted to a fast magnetoacoustic wave flux and the rest is reflected. The transmitted fast magnetoacoustic waves are severely refracted and possibly totally internally reflected because of the increase in Alfvén speed with height. However, in the same article, the authors have found that waves ducted along the flux tubes can be generated as well, both directly by convection and by coupling to the acoustic wave flux incident on the magnetic canopy. Flux tubes can support three types of waves : kink (transverse), sausage (longitudinal) and Alfvén (torsional). So it seems that the convective motions, interacting with the flux tube, can directly generate torsional Alfvén waves. Also, at the canopy, acoustic waves from neighboring flux tubes can probably couple strongly to tube waves and convert a portion of their energy flux into tube wave flux. The tube waves do not suffer from refraction like acoustic waves do. They can, however, be reflected by the transition region, since the Alfvén speed changes rapidly with height there (Stein & Nordlund 1991, Edwin 1992). Stein roughly estimates the Alfvén wave flux caused by turbulent motions in the photosphere to be of the order 10^5 J m^{-2} s^{-1}. The Alfvén wave flux caused directly by granular and supergranular motions could be as high as 10^6 J m^{-2} s^{-1}. Both these estimates are based on simple theoretical considerations and are probably just indicative of the upper limit of the Alfvén wave flux generated in the lower solar atmosphere (Stein & Leibacher 1981).

Recent numerical simulations of the generation of longitudinal and transverse tube waves by the turbulent convection zone (Ulmschneider & Musielak 1998) (assuming a Kolmogorow type law for the turbulent energy spectrum) show that the energy flux carried by transverse tube waves in the frequency range 0.1 to 1 Hz can be of the order 10^6 J m^{-2} s^{-1}. This is several orders of magnitude higher than the energy flux (carried by Alfvén waves) needed for our spicular driving mechanism.

Concluding, both observations and theory indicate it is quite possible that the Alfvén waves needed for driving spicules are present in the chromosphere.

4.2. Limitations of the WKB model

Our analysis is only valid for waves that do not have too large amplitudes, since we linearized at some point. However, with Piddington, we feel that the results may be suggestive of the effects of stronger waves (Piddington 1956). Hollweg has shown that a large-amplitude Alfvén wave in an inhomogeneous medium (or one that is being damped) is governed by the same WKB approximation as that for the small-amplitude limit (Barnes & Hollweg 1974, Hollweg 1974). This seems to back up our assertion that our results may be valid for large-amplitude waves as well. In our analysis, the assumption of small-amplitude waves is not critical for the inclusion of damping, i.e., ion-neutral damping will also occur for large-amplitude Alfvén waves, and probably at similar rates as our analysis indicates. At the same time, it is clear that previous numerical simulations of Alfvén waves in the chromosphere indicate non-linear effects play an important role, and those should be taken into account into numerical simulations of our driving mechanism.

We have also assumed that the changes of the background model are not too large over one wavelength. For the frequencies of interest, this breaks down at the transition zone, where the rapid density decrease over a few tens of kilometers leads to a rapid rise of the Alfvén velocity v_A. But the Alfvén waves will already have been damped considerably when they reach the transition region, so that reflections may not be too important. Numerical simulations should take into account reflections off the transition region as well.

Additionally, even though in a dynamic model there may be problems with WKB initially, these are resolved as soon the spicule becomes taller, since then the waves will be damped over a larger range. And from observations it is known that there are no large gradients with height of any thermodynamic or magnetic parameters, so that the WKB assumption is most probably not violated in an already formed spicule. Of course a fully internally consistent model has to be studied to see where and when problems may arise.

Our analysis can easily be repeated for axisymmetric twists (in which δB_θ is the only non-zero wave magnetic field component and δv_θ is the only non-zero wave velocity component, with θ the azimuthal direction in a cylindrical flux tube). These axisymmetric twists are, similar to our linearly polarized Alfvén wave, Alfvénic only in the lowest order, since both agree with the condition that $\|\boldsymbol{B_0}+\boldsymbol{\delta B}\|$ remain constant only in the lowest order.

Our analysis, further, did not take into account any gradients of wave amplitude in the direction perpendicular to the background magnetic field. Gradients of transverse magnetic pressure in the horizontal direction might very well drive horizontal motions as well. Since Alfvén waves travel parallel to the magnetic field, the existence and size of such gradients is not defined by the wave mode, but rather by the boundary conditions of the solar atmosphere. Unfortunately, there is no observational knowledge about these gradients, so it would be difficult to model them. At least initially, the region where damping of the Alfvén wave occurs is smaller than the probable horizontal scales of spicules, so it is possible that damping (as opposed to horizontal gradients) will dominate the motion of the plasma.

One of the major assumptions was that $\tau = \frac{\omega}{\nu_{ni}} \ll 1$. This limits the applicability of our model to waves with frequencies lower than 1 Hz, since from Fig. 2 we know that ν_{ni} can become as low as a few dozen Hz. Note that the effects of ion-neutral collisions take on a qualitatively different form for $\tau \approx 1$.

Finally, throughout our calculations we studied the effects of only one neutral species, to facilitate tractability. The analysis can be extended to several neutral species if $\tau \ll 1$ by replacing the neutral-ion collision frequency by:

$$\nu_{ni} = \sum_j \frac{\rho_{n_j}}{\rho_n} \nu_{n_j i} \qquad (40)$$

in which the sum is taken over all neutral species and ρ_n is the total neutral density (Hartmann & MacGregor 1980).

5. Conclusion

The damping of upward traveling Alfvén waves with $0.2 \le f \le 0.6$ H, due to collisions between ions and neutral particles, can cause sufficient momentum and energy transfer from the Alfvénic motions to the plasma, that a spicular structure can probably be formed and sustained. The presence of such waves seems probable from both observational and theoretical evidence.

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