ALFVÉN WAVES IN THE SOLAR ATMOSPHERE

III. Nonlinear Waves on Open Flux Tubes

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Abstract. The nonlinear propagation of Alfvén waves on open solar magnetic flux tubes is considered. The flux tubes are taken to be vertical and axisymmetric, and they are initially untwisted. The Alfvén waves are time-dependent axisymmetric twists. Their propagation into the chromosphere and corona is investigated by solving numerically a set of nonlinear time-dependent equations, which couple the Alfvén waves into motions parallel to the initial magnetic field (motion in the third coordinate direction is artificially suppressed). The principal conclusions are: (1) Alfvén waves can steepen into fast shocks in the chromosphere. These shocks can pass through the transition region into the corona, and heat the corona. (2) As the fast shocks pass through the transition region, they produce large-velocity pulses in the direction transverse to $\mathbf{B}_0$. The pulses typically have amplitudes of 60 km s$^{-1}$ or so and durations of a few tens of seconds. Such features may have been observed, suggesting that the corona is in fact heated by fast shocks. (3) Alfvén waves exhibit a strong tendency to drive upward flows, with many of the properties of spicules. Spicules, and the observed corrugated nature of the transition region, may therefore be by-products of magnetic heating of the corona. (4) It is qualitatively suggested that Alfvén waves may heat the upper chromosphere indirectly by exerting time-dependent forces on the plasma, rather than by directly depositing heat into the plasma.

1. Introduction

The propagation of small-amplitude Alfvén waves in the solar atmosphere has been considered in previous papers (Hollweg, 1978a, 1981a; hereafter designated Papers I and II, respectively). Paper II is particularly relevant to the Sun, since it considers the propagation of waves on solar magnetic flux tubes. That paper suggests the importance of Alfvén waves in heating and accelerating the chromosphere and corona. Paper II shows that Alfvén waves can supply the energy required to balance the radiative, heat conduction, and solar wind losses of the chromosphere and corona. It does not explain how the Alfvén waves dissipate their energy, but it suggests that the waves could become of sufficiently large amplitude to undergo nonlinear dissipation processes. Moreover, it is shown that the forces which the waves exert on the solar atmosphere can become large enough to drive significant mass motions; it was suggested in particular that spicules might be a consequence of these forces.

The purpose of this paper is to consider some nonlinear aspects of the propagation of Alfvén waves on solar magnetic flux tubes. First, the steepening of Alfvén waves to

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form shocks in the chromosphere will be investigated. Second, the interaction of a chromospheric switch-on shock (a fast shock propagating parallel to the ambient upstream magnetic field) with the transition region (treated in this paper as a contact surface) will be examined. Finally, we will investigate the extent to which the nonlinear forces exerted on the plasma by the Alfvén waves can drive mass motions in the chromosphere and corona. The results of these computations will be examined in terms of their applicability to chromospheric heating, coronal heating, and spicule formation.

This paper considers only the Alfvén mode (and its nonlinear coupling into the slow, or sound, mode), for several reasons. The principal reason is tractability. We wish to consider the propagation of waves in the horizontally stratified solar atmosphere in which are embedded magnetic flux tubes whose magnetic field strength decreases rapidly with height. Since the small-wavelength WKB approximation is not valid for frequencies of interest, full solutions of the wave equation are mandatory. Even in a linearized theory, the Alfvén mode is the only mode which is at present amenable to exact analysis under these conditions. The large-amplitude Alfvén wave presents new difficulties which preclude an exact analysis, but it can be subjected to a reasonably detailed numerical study if some simplifying restrictions are introduced. By contrast, the other MHD modes appear to require more sophisticated techniques.

However, there are some indications that the Alfvén mode is of more than academic interest. First, Alfvén waves which are presumably of solar origin have been observed in the solar wind, and their possible role in heating and accelerating the wind has been widely discussed (e.g., Barnes, 1979; Hollweg, 1978b). We are motivated, therefore, to investigate whether they can play a similar role in the lower solar atmosphere. Second, there are suggestions that spicules rotate about their axes (e.g., Beckers, 1968; Livshits, 1967; Pasachoff et al., 1968). If the rotation is a dynamic time-dependent twisting motion, then it is an Alfvén wave obeying the equations of this paper. Third, B. Lites and O. R. White (private communication, 1980) have found evidence for strong compressive motions in the chromosphere over a sunspot umbra, but there is an absence of compressive activity in the photosphere. A possible, but unconfirmed, explanation for this behavior is to assume that the photospheric motions are noncompressive shears, but that the motions nonlinearly become compressive as they propagate to higher levels. The noncompressive shearing motions can be regarded as Alfvén waves, while the conversion to compressive motions at higher levels may occur nonlinearly, as will be shown below. Finally, this paper emphasizes the Alfvén mode because, except for the Alfvénic surface mode, it appears to be the only mode which is capable of transporting the required energies to the upper chromosphere and corona (see the review by Hollweg, 1981b).

A number of simplifying assumptions are employed in this paper. Only flux tubes which are open into interplanetary space will be considered. The flux tubes are assumed to be axisymmetric with vertical axes of symmetry, and except for the waves the flux tubes are taken to be untwisted and in static equilibrium. The wave motions
are assumed to be axisymmetric, and to consist of only two vector components: in the azimuthal direction and parallel to the average magnetic field of the flux tube. The azimuthal motions ('twists') can in lowest order be regarded as Alfvén waves (Paper II), while the parallel motions will here be regarded as being nonlinearly driven by the Alfvénic twists. The third component of motion (perpendicular to the average field and to the azimuthal direction) is suppressed. This is a vital simplification, as it reduces the problem to one space dimension. Unfortunately, there is no convincing justification for this Ansatz. Alfvénic twists can nonlinearly drive motions in the suppressed direction as well as along the average field; the motions are driven by the wave-associated centrifugal force $-\rho \mathbf{v} \cdot \nabla \mathbf{v}$ and the wave-associated Lorentz force $\mathbf{j} \times \mathbf{B}/c$ ($\rho$ is the mass density, $\mathbf{v}$ is the velocity, $\mathbf{j}$ is the current density, $\mathbf{B}$ is the magnetic field, and $c$ is the speed of light). Roughly speaking, the magnitude of a given component of the nonlinearly driven motion is proportional to that component of the gradient of the square of the Alfvén wave amplitude. In view of the short vertical scale heights of the photosphere and chromosphere, it is at least possible that the vertical gradient of the square of the wave amplitude will be larger than the horizontal gradient. Now field lines which are not too distant from the flux tube axis will be mainly vertical, and it is therefore possible that the nonlinearly driven motions parallel to the $\mathbf{B}$ will be larger than those normal to $\mathbf{B}$. (Exceptions to this argument are in the photosphere, where the diameter of the flux tube may be comparable to the vertical scale height, and in the corona, where the vertical scale height is large; however, the Alfvén wave nonlinearity is weak in both of these regions.) A further simplification is that only an adiabatic equation of state is used. This means that the entropy of a given parcel of gas remains constant for all times. The entropy of the initial state of the atmosphere is allowed to vary spatially, however. Thus we can consider an atmosphere which by some unspecified process has developed a corona and chromosphere, but the subsequent thermal evolution of those regions, due to heat input and radiative and conductive losses for example, is ignored. This study must therefore be construed primarily as investigating the dynamics of nonlinear Alfvénic motions, with some weaknesses in thermo-dynamics. Finally, we assume that the waves propagate into a solar atmosphere which is initially in hydrostatic equilibrium with a single temperature characterizing the photosphere and chromosphere, and a single higher temperature characterizing the corona.

It should be noted that, even though we are formally considering the nonlinear evolution of axisymmetric twisting motions, it is possible that our solutions are qualitatively representative of the response of a flux tube to shaking, such as due to granular buffeting, for example (see Figure 1 of Paper II). Waves generated in this fashion might have time scales of the order of minutes. For this reason, this paper concentrates on waves with such time scales.

2. Equations of the Problem

Consider a single field line of the average magnetic field, $\mathbf{B}_0$. Let distance measured along this field line be denoted by $s$, and let $r(s)$ be the distance of any point on this
field line from the axis of symmetry of the flux tube. In the vicinity of this field line we consider a local orthogonal curvilinear coordinate system defined by \( s \), by the axialmuthal angle \( \theta \) measured about the axis of symmetry, and by a coordinate \( \xi \) measured in the \( \hat{e}_z \times \hat{e}_\theta \) direction (\( \hat{e} \) denotes a unit vector). The curvilinear scale factors along the given field line are \( h_s = 1 \), \( h_\theta = r \), and \( h_\xi \) can be arbitrary; note, however, that \( \partial h_s / \partial \xi \neq 0 \) and \( \partial h_\xi / \partial s \neq 0 \) in general.

As discussed in the introduction, we take \( \partial / \partial \theta = 0 \) and \( v_\xi = 0 \). The latter implies \( B_\xi = 0 \). The equation \( \nabla \cdot \mathbf{B} = 0 \) becomes then

\[
\rho h_\xi B_s = \text{constant along field lines}. \tag{1}
\]

Similarly, the equation of mass conservation, \( \partial \rho / \partial t + \nabla \cdot (\rho \mathbf{v}) = 0 \), becomes

\[
\frac{\partial}{\partial t} \left( \frac{\rho}{B_s} \right) + \frac{\partial}{\partial s} \left( \frac{\rho}{B_s} v_s \right) = 0, \tag{2}
\]

where (1) has been used and where \( B_s \) is independent of time. The adiabatic energy equation is \( \partial p / \partial t + \mathbf{v} \cdot \nabla p + \gamma p \nabla \cdot \mathbf{v} = 0 \), where \( p \) is thermal pressure and \( \gamma \) is the ratio of specific heats; if there are 3 degrees of freedom, then \( \gamma = \frac{5}{3} \). This equation becomes

\[
\frac{\partial}{\partial t} \left( \frac{p}{B_s} \right) + \frac{\partial}{\partial s} \left( \frac{p}{B_s} v_s \right) = -(\gamma - 1)p \frac{\partial}{\partial s} \left( \frac{v_s}{B_s} \right). \tag{3}
\]

The momentum equation is

\[
\rho \left( \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla \right) \mathbf{v} = -\nabla \left( p + \frac{B^2}{8\pi} \right) + \rho g + \frac{1}{4\pi} \mathbf{B} \cdot \nabla \mathbf{B}, \tag{4}
\]

where \( g \) is the gravitational acceleration, which will be taken to be constant over the height range we are considering. The \( \theta \)-component of (4) is

\[
\frac{\partial}{\partial t} \left( \frac{\rho v_\theta}{B_s} \right) + \frac{\partial}{\partial s} \left( \frac{\rho v_\theta}{B_s} v_s \right) = \frac{1}{4\pi} \frac{\partial}{\partial s} (rB_\theta). \tag{5}
\]

Equation (5) describes conservation of the angular momentum of the twisting motions, in the presence of a source term due to the magnetic field line tension. Similarly, the \( s \)-component of (4) is

\[
\frac{\partial}{\partial t} \left( \frac{\rho v_s}{B_s} \right) + \frac{\partial}{\partial s} \left( \frac{\rho v_s}{B_s} v_s \right) = -\frac{1}{B_s} \frac{\partial p}{\partial s} + \frac{\rho g_s}{B_s} + \frac{1}{B_s} \left[ \left( \frac{\rho v_\theta^2}{4\pi} - \frac{B_\theta^2}{8\pi} \right) \frac{\partial \ln r}{\partial s} - \frac{\partial}{\partial s} \left( \frac{B_\theta^2}{8\pi} \right) \right]. \tag{6}
\]

The three terms in square brackets are the terms whereby the Alfvén waves drive motions parallel to \( B_\theta \). The first term is the \( s \)-component of the centrifugal force associated with the twisting motions, the second term arises from the field line tension, while the third term is the gradient of the magnetic field pressure associated with the twisted field; the second and third terms together are the \( s \)-component of the Lorentz force.
Finally, the induction equation is $\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B})$. The $\theta$-component becomes

$$\frac{\partial}{\partial t} \left( \frac{B_\theta}{rB_s} \right) + \frac{\partial}{\partial s} \left( \frac{B_\theta}{rB_s} v_s \right) = \frac{\partial}{\partial s} \left( \frac{v_\theta}{r} \right).$$  \hspace{1cm} (7)

If $v_s = 0$, Equations (5) and (7) are identical to the equations of Paper II. Equations (5) and (7) may therefore be thought of as describing the Alfvén wave, while Equations (2), (3), and (6) describe the parallel motions which are nonlinearly driven by the Alfvén wave. The terms involving $v_s$ in Equations (5) and (7) represent higher-order modifications of the Alfvén waves, such as their steepening to form shocks. (See Barnes and Hollweg (1974), Cohen and Kulsrud (1974), Hollweg (1971), Montgomery (1959), Parker (1958), and Wray (1972) for discussions of nonlinear Alfvén waves when $g = 0$ and the WKB approximation can be used.)

Equations (2), (3), (5), (6), and (7) can be combined to yield an equation of energy conservation:

$$\frac{\partial}{\partial t} \left[ \frac{1}{B_s} \left( \frac{1}{2} \rho v^2 + \frac{B_\theta^2}{8\pi} + \frac{p}{\gamma - 1} + \rho \phi \right) \right] + \frac{\partial}{\partial s} \left[ \frac{v_s}{B_s} \left( \frac{1}{2} \rho v^2 + \frac{\gamma p}{\gamma - 1} + \rho \phi \right) \right] + \frac{1}{4\pi B_s} \left( -B_s v_\theta B_\theta + B_\theta^2 v_s \right) = 0,$$  \hspace{1cm} (8)

where $\phi$ is the gravitational potential, defined by

$$g_s = -\frac{\partial \phi}{\partial s}.$$  \hspace{1cm} (9)

The first square brackets in (8) contain the densities of kinetic, magnetic, thermal and gravitational energy. The first three terms in the second set of square brackets in (8) represent the convection of kinetic energy, enthalpy, and gravitational energy, while the last two terms in the second square brackets give the Poynting flux in the inertial frame.

Equations (2), (3), (5), (6), and (7) are the basic equations of this paper. On its left-hand side, each equation contains the time derivative of some quantity, plus the divergence of the flux of that quantity; source terms appear on the right-hand side. Equations of this form are appropriate for solution by the numerical flux-corrected transport scheme of Boris and Book (1973, 1976). The code published by Boris (1976) was adapted to the Cray-1 computer at NCAR and was used with the Eulerian grid option for all the calculations described in this paper.

For convenience we have studied only vertical field lines near the axis of the flux tube; for these field lines

$$B_s r^2 = \text{constant}$$  \hspace{1cm} (10)

and $g_s = -g$, where $g$ has the solar value $2.722 \times 10^4 \text{ cm s}^{-2}$. (At present, the detailed structure of solar magnetic flux tubes is unknown, and a more detailed treatment does not seem warranted.) The level $s = 0$ is chosen to be at zero height in
the HSRA/VAL model atmospheres (Gingerich et al., 1971; Vernazza et al., 1973),
where \( \rho = 3 \times 10^{-7} \text{ g cm}^{-3} \). The pressure at \( s = 0 \) is taken to be \( 1.188 \times 10^5 \text{ dyne cm}^{-2} \), and \( p/\rho = \text{constant} \) from \( s = 0 \) up to the transition region, which is
located at \( s = 2200 \text{ km} \), where \( \rho = 9.18 \times 10^{-14} \text{ g cm}^{-3} \). This corresponds to a pho-
tospheric and chromospheric temperature of 6230 K if the molecular weight of the
gas is 1.3. At the transition region the density drops by a factor \( 10^{-2} \), while the
pressure is held constant; the coronal density and pressure subsequently fall off
according to hydrostatic equilibrium. The top of the atmosphere is at \( s = 13200 \text{ km} \),
where an absorbing boundary condition is imposed. (It would have been desirable to
consider even smaller coronal densities, as in coronal holes. However, a smaller
coronal density results in a larger coronal Alfvén speed, which in turn imposes
unacceptably small time steps on the numerical integration, in virtue of a Courant
condition which is required by the flux-corrected transport scheme. A few runs with
smaller coronal densities have been made, and it has been found that the results are
not strongly different from those in the model described above.)

An Alfvén wave is generated near the base of the atmosphere by adding an
artificial body force, \( \rho f_\theta \), to the right-hand side of Equation (5). This body force acts
only in the near vicinity of \( s = 110 \text{ km} \) (it was inconvenient to locate the body force at
\( s = 0 \), but our choice does not affect the results). \( f_\theta \) is taken to be zero for \( t < 0 \), and
subsequently to vary as \( \sin \omega t \) for a finite number of half-cycles, after which \( f_\theta \) is again
zero. Generating the Alfvén wave by this means results in both a downward-
propagating wave, which is absorbed via a flow-through boundary condition in the
code, and an upward-propagating wave, which is the concern of this paper.

3. Chromospheric Shock Formation

In this section we consider the wave propagation for small times after \( t = 0 \), before
the initial wave front reaches the transition region.

The chosen variation of magnetic field strength with height is shown in Figure 1.
The behavior of the field strength is arbitrarily chosen to mimic that in Gabriel (1976)
and in Paper II. At \( s = 0 \) the field strength is 1500 G, falling to a value of 10.5 G in the
corona.

The variations with height of the Alfvén speed, \( v_A = B_i (4 \pi \rho)^{-1/2} \), and the sound
speed, \( v_{\text{sound}} = (\gamma p/\rho)^{1/2} \), are shown in Figure 2 at \( t = 0 \). (Both \( v_A \) and \( v_{\text{sound}} \) increase
by a factor 10 at \( s = 2200 \text{ km} \), and are nearly constant in the corona.) Below
\( s \approx 800 \text{ km} \), \( v_A/v_{\text{sound}} \) is nearly constant with height. This results from the balance
between the thermal plus magnetic pressures inside the flux tube and the thermal
pressure outside the flux tube. Above \( s = 800 \text{ km} \), \( v_A/v_{\text{sound}} \) increases, because the
magnetic field becomes increasingly more uniform while \( \rho \) continues to fall off
exponentially with height.

Consider first the case where \( f_\theta \) is turned on for a single half cycle, from \( t = 0 \) to
\( t = \pi/\omega \). Figures 3 to 5 show results for \( 2\pi/\omega = 90.26 \text{ s} \). Results are displayed for a
variety of values of \(|f_\theta|\), which result in a variety of values of \(|v_\theta|\) at \( s = 110 \text{ km} \).
Figure 3 displays snapshots of $v_\theta(s)$ at several values of $t$. The amplitude of the twisting motions increases with height. This is a consequence of the waves propagating into more rarefied regions of the solar atmosphere. The increase of $v_\theta$ with height is less rapid than would be predicted on the basis of the WKB approximation, in which case $|v_\theta| \sim \rho^{-1/4}$. For example, the WKB approximation predicts that $|v_\theta|$ should increase by a factor 11.4 between $s = 330 \text{ km}$ and $s = 1760 \text{ km}$, whereas the second panel from the bottom in Figure 3 shows an increase by the factor 8.4. The discrepancy occurs in part because of the inaccuracy of the WKB approximation at these wavelengths, and in part because the waves are converting some of their energy into motion parallel to $B_0$. The waves in the lowest two panels of Figure 3 move through this region without forming a shock. This is in part because their velocity amplitudes are too small. But it should be noted too that the increase of $v_A$ with height (Figure 2) hinders shock formation above $s \approx 1300 \text{ km}$ by allowing the foremost portion of the waveform to run ahead of the crest of the waveform, thus preventing the steepening of the wave and shock formation. If $|v_{\theta,\text{source}}|$ is large enough, however, a shock can form below $s \approx 1300 \text{ km}$, as shown by the upper two panels of Figure 3. In this case a switch-on shock appears at the leading edge of the waveform.

Figure 4 shows snapshots of $v_\varphi(s)$. In all cases, $v_\varphi$ is very small at low heights; this reflects the fact that the Alfvén waves only become large-amplitude (as measured by $B_\theta/B_\varphi$ or $v_\varphi/v_A$) as they propagate to greater heights. The waveform for $v_\varphi$ consists of upward-moving material in the front part of the waveform, and falling material in the rear. The upward motions are the result of the initial passage of the Alfvén wave,
which has a snowplow-like effect on the atmosphere, in virtue of the terms in square brackets in Equation (6). The downward motions arise in part from a change of sign of the driving terms in Equation (6), and in part from the atmosphere recovering after initially being lifted upward against gravity. The switch-on shocks are evident in the top two panels of Figure 4. These panels also show shocks at the trailing edges of the waveforms. These are upward-propagating slow shocks, which form in response to
Fig. 3. Snapshots of $v_\theta(s)$ for an Alfvén wave which consists of a single pulse lasting 45.13 s at the source. The four panels correspond to different values of the wave amplitude at the source. For each panel, the successive (reading from right to left) snapshots are at $t = 133$ s, 111 s, 89 s, 66 s, and 44 s, respectively; the right-most snapshot in the top panel is the only exception, being at $t = 127$ s. Switch-on shocks are present in the upper two panels.
the downflow of material at the rear of the waveform. (In addition, the top panel of Figure 4 contains a structure (near $s = 1485$ km) which appears to be a slow shock followed by a very rapid expansion (and plasma cooling). However, we cannot rule out the possibility that this structure is an artifact of the numerical scheme.)

Figure 5 displays the corresponding snapshots of $|B_\theta|/B_z$. This quantity at first increases with height because $B_z$ falls off more rapidly with height than $|B_\theta|$. At
greater heights $|B_\theta|/B_z$ decreases, because $B_z$ becomes increasingly more uniform while $|B_\theta|$ continues to decrease with height. (Recall that the WKB approximation predicts $|B_\theta| \sim \rho^{1/4}$.)

How do the parameters of the problem affect shock formation? Figures 3 to 5 show that the photospheric Alfvén wave amplitude plays an important role in determining...
whether shocks do or do not form. We have verified that $\omega$ is also an important parameter, as expected. For example, we have found that $|v_{\theta, \text{source}}| = 1.26 \text{ km s}^{-1}$ leads to shock formation when $2\pi/\omega = 45.13 \text{ s}$, where Figures 3 to 5 show that shock formation requires larger values of $|v_{\theta, \text{source}}|$ when $2\pi/\omega = 90.26 \text{ s}$. Similarly, we have found that a wave with $2\pi/\omega = 180.52 \text{ s}$ leads to a shock only if $|v_{\theta, \text{source}}| \geq 2.5 \text{ km s}^{-1}$, whereas Figures 3 to 5 show shock formation at smaller velocity amplitudes when $2\pi/\omega = 90.26 \text{ s}$. Another factor which influences shock formation is the length of the wave train. Consider Figure 4, where the Alfvén wave consists of a single pulse. The material at the front edge of the pulse is moving upward, while there is downflowing material at the rear. Now suppose that the Alfvén wave consists of two pulses. The upflowing material at the front of the second pulse then tends to run into the downflowing material at the rear of the first pulse, and this has the effect of enhancing the wave steepening and shock formation. For example, we have found that a continuous wave train with $2\pi/\omega = 90.26 \text{ s}$ can form shocks for $|v_{\theta, \text{source}}| \geq 0.9 \text{ km s}^{-1}$, in contrast to the situation in Figures 3 to 5 where larger values of $|v_{\theta, \text{source}}|$ are required for shock formation. Another important factor affecting shock formation is the behavior of $v_A(s)$. As suggested above, if $v_A$ increases strongly with $s$, then formation of fast shocks will be hindered. Conversely, if $v_A$ decreases with increasing $s$, then fast shocks can more easily form. The latter situation might pertain either in regions where the flux tubes are spaced widely apart and the field lines expand very rapidly with height, or near the edges of the flux tubes where the field lines may be nearly horizontal ($\rho \approx \text{constant}$) but expanding ($B_r$ decreasing with increasing $s$) resulting in $dv_A/ds < 0$. Unfortunately, not enough is known about the structure of magnetic flux tubes to assess these possibilities with any confidence. Finally, the length of the field line can affect shock formation. Clearly, a long field line (as at the edge of a flux tube, for example) is more favorable for shock formation than a short field line, simply because the long field line gives the wave more time to steepen into a shock. In this sense our treatment only of (short) field lines near the flux tube axis is a worst case for shock formation.

On the whole, then, it appears that Alfvén waves on solar flux tubes will be able to steepen into fast and slow shocks in the chromosphere, if their periods are less than a few minutes, and if their photospheric velocity amplitudes are in excess of some sizeable fraction of 1 km s$^{-1}$. It should be kept in mind, however, that these estimates are fairly model dependent.

If an Alfvén wave does not steepen into a shock in the chromosphere, it will probably not do so in the inner corona. The reason is simply that the nonlinearity of the wave, as measured by $B_0/B_z$ or $v_\theta/v_A$, is in general much smaller in the corona than in the chromosphere. (The solar wind beyond roughly 0.1 AU from the Sun may be an exception to this statement, but that region does not concern us here.) However, the fact that shocks can form in the chromosphere, even though they cannot form in the corona, may have important implications for coronal heating, as will be seen in the next two sections.
4. Shock-Transition Region Interaction

Shocks which form in the chromosphere will propagate upward, eventually interacting with the transition region. In this section we propose to isolate the shock-transition region interaction, and investigate it as a separate phenomenon. We shall confine our attention only to the switch-on shocks which appear in Figures 3 to 5. In addition, gravity, $dB_s/ds$, and the structure of the wave behind the shock front will be ignored; in effect, this means that we are studying only the initial instant when the switch-on shock first hits the transition region. The principal usefulness of this section will be as an aid to interpreting some features of the ‘full problem’, which will be studied in Section 5.

Fig. 6. The geometry appropriate for studying the interaction of a switch-on shock (the solid vertical line) with the contact surface representing the transition region (the dashed vertical line). Region ‘c’ is the corona, region ‘1’ is the unshocked chromosphere, and region ‘2’ is the shocked chromosphere.

Consider the situation illustrated in Figure 6. The contact discontinuity (the transition region) is indicated by the dashed line. The corona (region ‘c’) is to the right of the contact surface, and the chromosphere (region ‘1’) is to the left. The contact surface is initially taken to be at rest, and thus $v_1 = v_c = 0$. The switch-on shock is denoted by the solid line, and the post-shock flow is region ‘2’. The shock and contact surface are both normal to $\mathbf{B}$, so that $B_{\theta 1} = B_{\theta c} = 0$. Conditions in region ‘2’ are found by employing the standard jump conditions at a switch-on shock (e.g., Boyd and Sanderson, 1969); there results

\begin{align}
v_{s2} &= (S - 1)v_{A1}S^{-1/2}, \\
B_{\theta 2}^2 &= (S - 1)B_1^2\left\{[\gamma + 1 - S(\gamma - 1)] - 8\pi\gamma p_1/B_1^2\right\}, \\
v_{\theta 2} &= \frac{B_{\theta 2}}{(4\pi p_2)^{1/2}},
\end{align}

\((11)\) \hspace{1cm} \((12)\) \hspace{1cm} \((13)\)
\[ p_2 = p_1 - \frac{B_{\sigma_2}^2}{8\pi} + S\rho_1 v_{s2}^2/(S - 1), \]  

(14)

where

\[ S = \rho_2/\rho_1 \]

(15)

is a measure of the shock strength. In addition, there are the auxiliary relations

\[ p_1 = p_c \]

(16)

and

\[ B_c = B_1 = B_{s2} \]

(17)

while \( p_c/\rho_1 \) is arbitrary.

The numerical computations are initialized at the instant when the shock first hits the contact surface. At that instant region '1' becomes vanishingly small, and the plasma contains a single non-evolutionary discontinuity with coronal conditions on one side and post-shock conditions on the other. With these initial conditions, the numerical computation proceeds as described in Section 2.

Figure 7 shows results for \( B_1 = 7.5 \text{ G}, \rho_1 = 3 \times 10^{-14} \text{ g cm}^{-3}, \ p_1 = 0.025 \text{ dyne cm}^{-2}, \ S = 1.05, \) and \( p_c/\rho_1 = 0.02. \) These values are chosen for illustrative purposes only, but we have verified that the results are representative of a wide range of relevant parameters. Figure 7 shows that the shock-contact surface interaction leads to upward movement of the contact surface, a downward-propagating rarefaction wave (i.e., a pair of weak discontinuities), and both fast and slow shocks which propagate upward into the corona; the fast shock is a switch-on shock, which is not shown in Figure 7 because it propagates ahead of the slow shock and has already moved out of the right edge of the figure. The slow shock carries only a very small amount of energy into the corona, the fast shock being the main agent for carrying energy into the corona under these circumstances. A particularly interesting aspect of the interaction is shown in Figure 7a, where it can be seen that \( v_\theta \) and \( v_s \) in the vicinity of the transition region are considerably larger than \( v_{\theta 2} \) and \( v_{s2} \). In this example \( v_{s2} = 6 \text{ km s}^{-1} \) while near the transition region \( v_s = 11.5 \text{ km s}^{-1} \) and \( v_{\theta 2} = 37 \text{ km s}^{-1} \) while near the transition region \( v_\theta = 65 \text{ km s}^{-1} \). (For a wide range of relevant parameters, we have verified that it is a reasonably good rule-of-thumb to take \( v_\theta \) and \( v_s \) in the vicinity of the transition region to be roughly twice \( v_{\theta 2} \) and \( v_{s2} \), respectively.)

These results may have significant bearing on a number of solar phenomena. First, the upward motion of the transition region (and the chromospheric and coronal gas on either side of it) may play a role in spicule formation, and in producing a transition region which is not plane-stratified but instead strongly corrugated. Second, the fact that fast shocks which form in the chromosphere can enter into the corona offers an attractive mechanism for coronal heating. Finally, the large predicted values of \( v_\theta \) and \( v_s \) in the vicinity of the transition region may correspond to observations (Brueckner, 1980) of large wavelength shifts of coronal and transition region lines,
Fig. 7. The consequences of the interaction of a switch-on shock with the transition region. A fast shock (not shown) and a slow shock propagate into the corona, the transition region moves upward, and a rarefaction wave moves back down into the chromosphere. The fast shocks in the corona may be able to heat the corona.
and large increases of the line widths, lasting for some tens of seconds. These points will be discussed in more detail in Sections 5 and 6.

5. Spicules, and Coronal and Chromospheric Heating

In the previous two sections we have examined the formation of shocks in the chromosphere, and the interaction of switch-on shocks with the transition region. In this section we examine the 'full problem', beginning with generation of a sinusoidal wave at \( s = 110 \text{ km} \), which we allow to propagate through the chromosphere, transition region, and corona of the flux tube model described in Section 2 and by Figures 1 and 2.

The most interesting feature of the solutions is the upward displacement of the chromosphere and transition region in response to the forces exerted by the Alfvén wave. Figure 8 shows the height of the transition region as function of time for the case where \( 2\pi/\omega = 90.26 \text{ s} \), \( f_\theta \) is turned on from \( t = 0 \) to \( t = 4\pi/\omega \), and \( |v_\theta| = 2.52 \text{ km s}^{-1} \) at \( s = 110 \text{ km} \). The transition region is initially at a constant height of 2200 km, but it starts to move upward at \( t = 120 \text{ s} \), when the wavefront (in this case a

![Image](image_url)

Fig. 8. Motion of the transition region in response to an Alfvénic wavetrain which is generated in the photosphere beginning at \( t = 0 \). The wavetrain has \( 2\pi/\omega = 90.26 \text{ s} \), and is turned on for 180.52 s. The initial rise of the transition region is due to interaction with the Alfvén waves, while the rapid rise after \( t = 650 \text{ s} \) is due to interaction with a rebound slow shock. These phenomena may correspond to spicules and macrospicules.
switch-on shock) first reaches the transition region. The transition region moves upward at nearly constant velocity (approximately 18 km\,s\(^{-1}\) in this example), eventually reaching a maximum height (approximately 6500 km in this example), followed by a fall in response to gravity. It should be noted that the upward velocity, the maximum height, and the 'lifetime' of the feature, are not atypical of spicules.

![Density profile graph](image)

**Fig. 9.** The density profile at \( t = 480 \) s for the model corresponding to Figure 8. The light line is the initial density profile. The density profile may correspond to that in spicules.

The heavy line in Figure 9 shows the density profile when the transition region is near its maximum height; the light line is the density profile at \( t = 0 \). The figure shows that the waves push dense chromospheric material up to large heights, and that a fairly flat density profile is formed above some 2200 km. The density there is about \( 1.7 \times 10^{-13} \) g cm\(^{-3}\), which would correspond to a proton concentration of about \( 10^{11} \) cm\(^{-3}\) if the atmosphere were pure hydrogen. It should be noted that both the value of the density above 2200 km and the flatness of the profile in Figure 9 are typical characteristics of spicules (Beckers, 1972).

Figures 8 and 9 show that Alfvén waves are able to lift chromospheric gas to great heights. Because the length of the lifted 'plug' of gas much exceeds the thickness of the chromosphere, it is clear that the gas has undergone appreciable expansion
parallel to $B_0$. This expansion cools the gas adiabatically, with the result that the lifted plug of gas is quite cold. The heavy line in Figure 10 shows the profile of $p/\rho$ at the same time as in Figure 9; the light line in Figure 10 is the $p/\rho$ profile at $t = 0$. The cooling of the gas above $s \approx 2860$ km is evident; the temperature of the gas has dropped more than an order-of-magnitude below the initial chromospheric temperature. At $s \approx 2860$ km there is a strong upward-propagating slow shock, behind which the gas is compressed to temperatures in excess of the original chromospheric temperature. The slow shock forms in the same way as the slow shocks appearing in Figures 3 to 5; to some extent it can be thought of as a ‘rebound shock’, such as appears in theories of supernovae. It should be noted that the adiabatic expansion has led to temperatures which are much less than those observed in spicules. The possible relationship of these solutions to spicules will be discussed further in Section 6.

We have investigated how the parameters of the problem affect the extent to which the Alfvén waves produce the initial rise of the transition region (e.g., for $t \approx 450$ s in Figure 8). Surprisingly, we have found that $|v_{\theta,\text{source}}|$ does not have a strong influence. The reason is as follows: Suppose that $|v_{\theta,\text{source}}|$ is increased. This results in a
larger-amplitude Alfvén wave which exerts a stronger force on the chromosphere and transition region. The first pulse of the wave train will therefore drive the transition region upward at a larger velocity than in the case of a weaker wave. Now consider the second pulse of the wave train. The transition region in the strong wave case is moving higher and faster than in the weak wave case, and this means that in the former case it will take longer for the second pulse to reach the transition region (this effect is analogous to the usual Doppler effect). This in turn means that a transition region moving in response to a strong wave will have more time to decelerate under gravity before receiving a ‘boost’ from the second pulse. The net result is that the average rate of rise of the transition region is fairly insensitive to the wave strength. The same reasoning leads us to predict that the wave frequency and the magnitude of $B_z$ will be important parameters in determining the initial rate of rise of the transition region. A higher wave frequency allows the transition region to receive closely spaced boosts from subsequent pulses, with smaller intervals of deceleration between the boosts; the net result is that higher frequency waves lead to a faster average rate of rise of the transition region. Similarly, a stronger magnetic field strength in the upper chromosphere and corona also leads to a faster rate of rise, because a larger value of $\nu_A$ reduces the Doppler effect referred to above. Overall, we have found that Alfvén waves can push the transition region up to spiculic velocities at spiculic velocities if the coronal magnetic field strength exceeds 10 G or so, if the wave period is less than about 2 min, if the wave train is generated for longer than about 3 min, and if $|\nu_{\theta, \text{source}}|$ is not much less than 1 km s$^{-1}$.

We have also examined a case which is identical to that shown in Figures 8 to 10, except that $\gamma$ was reduced to 1.1. The motion of the transition region in this case is very similar to that shown in Figure 8, except that the transition region tends to remain near its maximum height for a longer time (about 7 min in this case) before starting to fall under gravity. The different behavior is due to the fact that the smaller value of $\gamma$ results in a nearly isothermal expansion, and therefore hotter material below the transition region than in the adiabatic case. The pressure below the transition region is therefore greater, and this helps to hold the gas up. It is interesting to note that some spicules are observed to remain at nearly constant height for extended periods of time.

We turn our attention now to the hypothesis, presented in the preceding section, that the corona may be heated by fast shocks which form in the chromosphere. We consider waves propagating in the same model atmosphere as before. The waves again have $2\pi/\omega = 90.26$ s, but in this case $f_{\theta}$ is turned on at $t = 0$ and left on, and $|\nu_{\theta, \text{source}}| = 1.78$ km s$^{-1}$.

Figure 11a shows the total energy flux density (given by the second square brackets in Equation (8)) crossing the surface at $s = 8800$ km (this surface is always above the transition region in this example). The energy flux through the corona is almost always upward, and strongly pulsed. The large pulses correspond to the fast shocks which have formed in the chromosphere and entered the corona. The energy flux is pulsed as a consequence of the ‘spikes’ in $\nu_{\theta}$ associated with the fast shocks in the top
two panels of Figure 3. (Most of the energy flux appearing in Figure 11 is due to the fast shocks, although some of the flux is due to slow shocks or to more gradual motions of the corona which are not shock-associated.) The flux density reaches rather large values (several times $10^6$ erg cm$^{-2}$ s$^{-1}$) in the pulses, but the pulses are short-lived, lasting only several tens of seconds or less. Figure 11b shows the integrated area under the curve in Figure 11a, i.e., Figure 11b represents the energy, in ergs, which has crossed 1 cm$^2$ of area at $s = 8800$ km from $t = 0$ to $t$. The light line in Figure 11b represents an energy flux density of $5.38 \times 10^5$ erg cm$^{-2}$ s$^{-1}$, and this is roughly the average rate at which the fast shock pulses in Figure 11a supply energy to the corona. This is roughly the energy flux required by coronal holes or quiet coronal regions to balance radiative and heat conduction losses. We therefore conclude that our proposal, that the corona is heated by fast shocks generated in the chromosphere, is feasible energetically.

Figure 12 shows $v_{\theta}$ at $s = 8800$ km for the same case as illustrated in Figure 11. Pulses due to passage of the fast shocks are evident. The pulse amplitudes are of the order of 75 km s$^{-1}$ at $s = 8800$ km; the amplitudes at the transition region itself are somewhat smaller, of the order of 50 km s$^{-1}$ in this example. As in Figure 11a, the pulse duration is several tens of seconds or less. Now we would like to suggest that these large velocity pulses in the transition region and corona may correspond to the pulsed events reported by Brueckner (1980). Brueckner's events appear to be

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Fig. 11a. The pulses are the energy flux density at $s = 8800$ km in the corona due to the passage of fast shocks which have formed in the chromosphere. In this case a semi-infinite wave train has been generated in the photosphere with $2\pi/\omega = 90.26$ s and $|v_{\theta,\text{source}}| = 1.78$ km s$^{-1}$. 

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Fig. 11b. The area under the curve in Figure 11a. The pulses in Figure 11a carry an average energy flux density which is adequate to supply the coronal energy requirements.

associated with network regions on the Sun, which is consistent with their being magnetic in origin. The velocities implied by the spectroscopic data (line shifts but more often line broadenings) agree well with the predicted velocity amplitudes in Figure 12. And the fact that the observations imply the presence of pulse-like events with time scales of tens of seconds also agrees with the predictions of Figures 11 and 12. These various lines of evidence lead us to believe, therefore, that the events reported by Brueckner may in fact show that the corona is being heated by fast shocks which have formed in the chromosphere.

We have also investigated whether this model can shed any light on the mechanism by which the chromosphere (especially the upper chromosphere) is heated. We have only been able to find solutions in which the chromosphere and transition region are set into robust motion by the waves, while we have been unable to obtain solutions which result in quasi-stationary heating of a quasi-static chromosphere. It is possible that modification of the shock propagation by ion-neutral friction, radiation, ionization and recombination could alter this picture. However, we would like to suggest an alternative, namely that the Alfvén-wave-induced chromospheric motion is an integral part of the chromospheric heating process. Suppose that a single Alfvénic pulse propagates upward from the photosphere. This pulse will expend a significant
fraction of its energy in lifting the chromosphere and transition region against gravity. In this way the Alfvén wave energy can be temporarily 'stored' as gravitational potential energy. After the Alfvénic pulse is turned off, the lifted chromosphere will fall in response to gravity, and the stored potential energy will be converted into the kinetic energy of the falling gas and into thermal energy as the atmosphere is compressed as it falls. Radiation will act as an irreversible energy sink, and will tend to damp away the motion of the falling atmosphere. Overall, such a process will represent a conversion of Alfvén wave energy into radiation, with conversion of Alfvén wave energy into gravitational potential energy as an intermediate step. Of course, it will be necessary to properly include radiation in the model before the feasibility of this proposal can be assessed; this will be postponed to a future paper.

6. Discussion

In this paper we have extended our previous studies of Alfvén waves in the solar atmosphere to include nonlinear effects. Tractability has forced us to make a number
of simplifying assumptions, however. The most important of these are: the flux tubes are vertical in a horizontally stratified atmosphere; the flux tubes are initially untwisted, while any subsequently imposed twists propagate as Alfvén waves; the flux tubes are initially axisymmetric, as are the imposed twists; motions in the \( \hat{e}_z \times \hat{e}_\theta \) direction are artificially suppressed; only an adiabatic energy equation is used. It should be kept in mind that the Sun undoubtedly moves magnetic flux tubes in more complicated ways than the theoretical idealization of this paper.

A general conclusion of this paper is that it is vital to consider the evolution of the waves as they propagate through the entire structure of the solar atmosphere. For example, it has previously been concluded that Alfvén waves will not be effective in heating the corona, because they can only dissipate their energy into heat via nonlinear processes which are weak in the corona, in virtue of the smallness of \( B_\theta/B_z \) and \( v_\theta/v_A \) there (Wentzel, 1977, 1980). Our viewpoint is different in that we have shown that the essential nonlinearity, i.e., shock formation, can occur in the chromosphere, and that the Alfvén wave energy can enter the corona in a form which is already amenable to nonlinear dissipation. The entire structure of the solar atmosphere is important also in determining the rate at which the Alfvén waves steepen into shocks in the chromosphere. The role of the transition region represents another illustration of the importance of considering the entire atmospheric structure. The transition region plays an essential role in determining not only the wave properties, but also the response of the entire atmosphere to the waves. We believe that it may be impossible to understand coronal heating or the spicule phenomenon without properly modelling the role of the transition region. This point has been emphasized earlier by us and by Wentzel (1978, 1980), and it appears to be even more important when nonlinear phenomena are considered.

A second general conclusion of this paper is that the observed highly inhomogeneous nature of the solar atmosphere (as manifested by spicules and a strongly 'corrugated' transition region (Feldman et al., 1979; Huber et al., 1974; Mariska et al., 1978), for example) may in fact be an essential dynamic aspect of the chromospheric and coronal heating processes. This conclusion has been reached because we have found that if Alfvén waves are to carry a significant energy flux into the corona and chromosphere, then they must of necessity exert large forces on the chromosphere and transition region, and induce robust motions of the upper atmospheric layers. This implies that whatever deposits energy in the upper solar atmosphere must also deposit momentum, and that both aspects must be self-consistently considered if the properties of the solar atmosphere are to be correctly understood. This viewpoint differs from that of previous workers in several ways. First, it means that heating of the solar atmosphere does not necessarily have to occur via conversion of magnetic energy directly into heat, but it allows the possibility that magnetic energy is first deposited into kinetic (and potential) energy by doing work on the atmosphere, and that the heating results from subsequent dissipation of the kinetic energy. A case in point is our suggestion that the upper chromosphere may in fact be heated as a result of the tendency of Alfvén waves to lift that region, with the
heating occurring when the gas subsequently falls and becomes compressed. (A proper evaluation of this suggestion must await inclusion of radiation in the model, however.) Second, we depart from previous workers who view the inhomogeneities of the atmosphere as quasi-static consequences of boundary conditions at the photosphere. This point of view is of great theoretical convenience, because it still allows one to perturb about some well-defined ‘average atmosphere’, and calculate wave propagation in the presence of inhomogeneities. Our viewpoint disallows such a procedure, however, because it regards the inhomogeneities as essential aspects of the waves themselves. It is thus our belief that a full theoretical description of momentum and energy exchange in the solar atmosphere may be considerably more complex than heretofore realized, in essence because the processes are inherently nonlinear.

Perhaps the most important conclusion of this paper is our demonstration that Alfvén waves can heat the corona by being converted into fast shocks in the atmosphere, which then propagate through the transition region into the corona. There may be observational support for this suggestion. Our model predicts that fast shocks moving through the transition region from the chromosphere into the corona should induce upward motions of the transition region (at 10 km s\(^{-1}\) or more) accompanied by large motions (50 km s\(^{-1}\) or more) transverse to \(\mathbf{B}\). These motions are predicted to be pulsed, with time scales of some tens of seconds. Brueckner (1980) has reported observations of wavelength shifts and broadenings of transition region lines. Velocities in excess of 60 km s\(^{-1}\) are implied. The line profiles show strong fluctuations within 20 s, and appear to be pulsed. These observed features are in accord with our predictions for fast shocks passing through the transition region, if the large observed velocities are associated with the predicted values of \(v_B\). Another line of evidence in support of our suggestion is the presence of broad wings of coronal lines at the limb (Kjeldseth, Moe and Nicolas, 1977). These wings may indicate the presence of short-lived high velocity pulses such as those predicted in Figure 12 (see also Shine et al., 1976).

The model of this paper yields ambiguous results when it is applied to the spicule problem. We first take the viewpoint that spicules are the chromospheric gas below a transition region which has moved upward along some selected set of field lines. (This viewpoint has already been taken by Osterbrock (1961) but it differs from that of Parker (1964) and Wentzel and Solinger (1967) who regarded spicules as the material behind ordinary hydrodynamic shocks.) We have succeeded in showing that an Alfvénic wavetrain can lift chromospheric material to spicular heights at spicular velocities. The lifted material has the densities characteristic of spicules as well as the fairly flat density profiles which are indicated by spicule observations. Since Pikel'ner (1971) has argued that thermal pressure alone is insufficient to lift spicules to the heights observed, we believe that our demonstration of the adequacy of dynamic magnetic fields may represent a significant step toward an understanding of spicules. Observations indicating that spicules rotate, and that signals propagate through spicules at large velocities which are consistent with presumed values of \(v_A\) in the
spicule (Pasachoff et al., 1968), are also in accord with our suggestion that the material is lifted by Alfvénic disturbances.

However, we have been unable to achieve a satisfactory understanding of the spicule temperature. Our model predicts spicules which are initially cold, in virtue of the adiabatic cooling which results when a parcel of chromospheric gas expands in the course of being lifted. (This tendency for cooling might be common to a wide variety of proposed spicule mechanisms. It might also be a feature of any phenomenon in which chromospheric gas is lifted into the corona, such as sprays, surges, and coronal transients.) It therefore appears that this model does not allow us to identify the entire ‘plug’ of gas below the transition region in Figures 9 and 10 as being a spicule. This poses a severe dilemma for the Alfvén-wave-driven spicule model, and it may mean that our entire concept is wrong.

However, there are some other possibilities. First, absorption of radiation, and ionization and recombination, which are not included in our computations, may keep the spicules warmer than shown in Figure 10, but these effects will not be able to produce temperatures in excess of $10^4$ K, as suggested by Beckers (1972). (It should be noted that some workers, e.g., Nikol’skii and Sazanov (1967) have suggested spicule temperatures only of the order of 6000 K, but Beckers’ work appears the most definitive in this regard.) Second, it is possible that the Alfvén waves are responsible for the spicule dynamics, much as described by our model, but that some other unknown heating process is occurring and maintaining the spicules at high temperatures. This viewpoint, of course, adds but little to our understanding of spicules. Third, it is possible that the spicules are heated by some additional damping mechanism for the Alfvén waves, such as ion-neutral friction. Fourth, the spicules may be the gas behind the slow shock in Figure 10. In that case the spicules would be hot, and their temperatures would increase with height, as suggested by Beckers (1972). In this scenario the generation of spicules would be a two-stage process. In the first stage, the Alfvén waves lift the chromospheric gas to great heights, but this gas is cold and invisible at the limb. In the second stage, the gas is heated and made visible by the slow shock. The cold gas and hot gas seen against the disk might correspond to the dark and bright mottles, respectively. This possibility is at least consistent with the observations (see Beckers', 1972, review) that the dark mottles are higher than the bright mottles, but it cannot be properly evaluated until the radiation field is properly included in the model. This possibility would suggest, however, that dark mottles should evolve into bright mottles, which to our knowledge has never been observed. A final possibility, which is similar to the one just described, is that a spicule results when the slow shock in Figure 10 encounters the transition region. This occurs at $t \approx 650$ s in Figure 8, where it is seen that the transition region is ejected rapidly upward (at approximately 50 km s$^{-1}$ in this case) by the slow shock. Because of the increasing noisiness of the solutions with time, we have been unable to provide a fully convincing description of this aspect of the problem. However, the solutions obtained do show temperatures in excess of $10^4$ K below the transition region. The density profiles are not as flat as in Figure 9,
however, showing a gradual decline to values of \( n_H \) which are only somewhat in excess of \( 10^{10} \text{ cm}^{-3} \) at \( s \approx 10^4 \text{ km} \). Such a situation could also be compatible with Beckers' (1972) spicule model.

The motions induced by interaction of the transition region with a slow shock, such as for \( t \geq 650 \text{ s} \) in Figure 8, have an alternate interpretation, however. The upward velocities in these cases tend to be greater than spicule velocities, and the densities below the transition region tend to be somewhat less than typical spicule densities. The possibility exists that these rapidly moving low density structures correspond to the macrospicules (e.g., Withbroe et al., 1976).

The situation regarding spicules is clearly not yet resolved, but the partial successes of the model do offer encouragement, as do recent observational arguments that spicules are in fact driven from below, as in this work, rather than from above (Rabin and Moore, 1980). But it should be noted that our approach does not seem capable of explaining certain observational features of spicules, such as why they are so thin; this may be a fundamental consequence of the physics, or it may simply reflect the boundary conditions.

It should be noted that our results may indicate a natural explanation for the large 'downflows' observed in the chromospheric network. We have found that Alfvén waves induce cooling of the chromospheric material as it is lifted. Conversely, it follows that heating occurs as the material falls and compresses. It can therefore be concluded that downfalling material will tend to be hotter, brighter, and more visible than rising material, and that a spatial or temporal average will be biased in favor of the redshifted material. A proper evaluation of this effect requires that the thermodynamics be correctly treated, however.

For the future, an attempt must be made to properly include thermal effects, though the dominant transport mechanisms in such an inhomogeneous geometry are by no means clear. We also hope to treat closed loops in active regions (where, interestingly, there is an apparent absence of spicules); in this case arguments similar to those of Meyer (1976, p. 124) suggest the addition of a zero-order twist to the tube may make coronal heating easier. Finally, the 'suppression' approximation should ultimately be removed, though the resulting numerical problem seems currently quite formidable.

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