I. INTRODUCTION

Since the early work of Alfvén (1947), Alfvén waves have been of interest as a possible means of heating the solar corona and chromosphere. Recent work has focused on the role of these waves in heating coronal loop structures, especially active region loops and X-ray bright points, which may act as resonant cavities for Alfvén waves (Hollweg 1979, 1981a, 1984a, b; Ison 1982, 1984; Zugzda and Locans 1982). The hypothesis that Alfvén waves heat coronal loops has recently been given observational support by Hollweg and Sterling (1984), who showed that the wave heating theory of Hollweg (1984a) is compatible with observations of fifteen loop structures as reported by Golub et al. (1980). The presence of resonances has also received observational support from a power spectral analysis of coronal motions in a single loop structure (Koutchmy, Zugzda, and Locans 1983).

The analyses cited above of Alfvén wave propagation in loops have utilized linearized theory. The Alfvén waves then do not couple into other modes, and thus they remain purely noncompressive and transversely polarized at all times. Moreover, the properties of the background atmosphere, in which the waves propagate, are unaffected by the waves. In particular, the locations of the transition regions at each end of the loop, which bound the coronal resonant cavity, are unaffected by the presence of the waves.

It is the purpose of this paper to investigate numerically some of the nonlinear aspects of Alfvénic pulses propagating in coronal loops and the underlying chromosphere. Heat conduction and radiation are included. The Alfvénic pulses are modeled as axisymmetric twists on a vertical cylindrical flux tube. They nonlinearly couple into acoustic-gravity waves propagating along the flux tube. A single Alfvénic pulse is found to leave two acoustic-gravity pulses in its wake. These pulses can result in significant motions of the transition region and underlying chromosphere. These motions do not resemble spicules, but they may correspond to a variety of observations indicating that the solar atmosphere is in a continual dynamic state. Indeed, we suggest that a dynamic chromosphere and transition region may be the inevitable consequence of the coronal heating process itself.

Subject headings: Sun: atmosphere — Sun: chromosphere — Sun: corona

ABSTRACT

We investigate numerically some nonlinear aspects of Alfvénic pulses propagating in coronal loops and the underlying chromosphere. Heat conduction and radiation are included. The Alfvénic pulses are modeled as axisymmetric twists on a vertical cylindrical flux tube. They nonlinearly couple into acoustic-gravity waves propagating along the flux tube. A single Alfvénic pulse is found to leave two acoustic-gravity pulses in its wake. These pulses can result in significant motions of the transition region and underlying chromosphere. These motions do not resemble spicules, but they may correspond to a variety of observations indicating that the solar atmosphere is in a continual dynamic state. Indeed, we suggest that a dynamic chromosphere and transition region may be the inevitable consequence of the coronal heating process itself.

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consider the possibility, raised by Hollweg (1982), that an Alfvénic pulse might serve as the initial impulsive deposition of energy which ultimately leads to spicule formation; one of the main conclusions of this paper will be that Alfvénic pulses are not effective in this regard.

This paper also includes a number of numerical improvements over the computations of HJG and Hollweg (1982). In those papers all sources of dissipation were ignored, and the chromosphere-corona transition region was represented by a contact discontinuity. In this paper we will extend the NRL dynamic flux tube model (Mariska et al. 1982; Mariska and Boris 1983) to include Alfvén waves. Radiation and heat conduction are included as dissipation mechanisms and the transition region is treated as a resolved structure whose properties are determined by energy balance.

We will find that the single Alfvénic pulses do not by themselves have a strong influence on the background atmosphere. Of more importance are two acoustic-gravity pulses which are left in the wake of the Alfvénic pulse. These pulses can result in significant motions of the transition region and underlying chromosphere. These motions do not resemble spicules, but they may correspond to a variety of observations indicating that the solar atmosphere is in a continual dynamic state. Indeed, we suggest that a dynamic chromosphere and transition region may be the inevitable consequence of the coronal heating process itself.

II. BASIC EQUATIONS

Except for the addition of dissipative terms, the equations considered here are very similar to those considered by HJG. Axisymmetric twists are imposed on a vertical magnetic flux tube which is initially untwisted and axisymmetric. We allow also for vertical flows, but flows in the third direction are suppressed.

The background magnetic field, \( B_0 \), is taken to be uniform. This contrasts with HJG, who allowed \( B_0 \) to decrease with height, from 1500 G in the photosphere to 10.5 G in the coronal hole. However, we feel that neglecting the height variation of \( B_0 \) is justifiable in the present context, for two reasons:

1. The stronger average magnetic field strengths in active regions implies that the height variation of \( B_0 \) is less than in coronal holes.

2. We are primarily concerned with examining the basic physics of Alfvénic pulses in an idealized chromospheric model and thus wish to simplify the problem as much as possible. Thus, the initial flux tube can be regarded as a cylinder; we therefore employ cylindrical coordinates \((r, \theta, z)\), where \( z \) is height in the atmosphere, and \( \theta \) is the azimuthal coordinate.

The basic equations express conservation of mass, momentum and energy, as well as Faraday’s law of induction in a perfectly electrically conducting fluid. Thus (see eqs. [2], [5], [6], [7], and [8] of HJG):

\[
\frac{\partial \rho}{\partial t} + \frac{\partial (\rho v_z)}{\partial z} = 0 , \tag{1}
\]

\[
\frac{\partial (\rho v_\theta)}{\partial t} + \frac{\partial (\rho v_\theta v_z)}{\partial z} = \frac{B_0}{4\pi} \frac{\partial B_\theta}{\partial z} , \tag{2}
\]

\[
\frac{\partial (\rho v_z)}{\partial t} + \frac{\partial (\rho v_z^2)}{\partial z} = -\frac{\partial p}{\partial z} + \rho g_z + \frac{1}{8\pi} \frac{\partial B_\theta^2}{\partial z} . \tag{3}
\]

\[
\frac{\partial}{\partial t} \left[ \frac{1}{2} \rho (v_\theta^2 + v_z^2) + \frac{B_\theta^2}{8\pi} + \frac{p}{\gamma - 1} + \rho \phi \right] + \frac{\partial}{\partial z} \left[ v_z \left( \frac{1}{2} \rho (v_\theta^2 + v_z^2) + \frac{B_\theta^2}{8\pi} + \frac{p}{\gamma - 1} + \rho \phi \right) \right] = -\frac{\partial}{\partial z} \left[ q_z - pv_z + \frac{1}{8\pi} (B_\theta^2 v_z - 2B_\theta v_\theta B_\theta) \right] - L , \tag{4}
\]

and

\[
\frac{\partial B_\theta}{\partial t} + \frac{\partial (B_\theta v_z)}{\partial z} = B_0 \frac{\partial v_\theta}{\partial z} . \tag{5}
\]

(all units are cgs). Here \( \rho \) is mass density, \( p \) is thermal pressure, \( v \) and \( B \) are velocity and magnetic field, \( g \) is the solar gravitational acceleration, \( \phi \) is the gravitational potential defined by

\[
g_z = -\frac{\partial \phi}{\partial z} , \tag{6}
\]

and \( \gamma = 5/3 \) is the ratio of specific heats. We will assume that our calculations refer to a semicircular loop which has been artificially straightened out. Thus

\[
g_z = -g_0 \cos (\pi z/L) ,
\]

where \( L \) is the full loop length, and \( g_0 = 2.72 \times 10^4 \) cm s\(^{-2}\).

The quantity \( q \) is the field-aligned heat conduction flux density, which we take to be due entirely to the electrons in the collision-dominated limit. Thus

\[
q_z = -K \frac{dT}{dz} , \tag{7}
\]

where the heat conductivity is

\[
K = 1.1 \times 10^{-6} T^{5/2} \tag{8}
\]

(Ulmschneider 1970), and \( T \) is temperature given by

\[
T = 6.93 \times 10^{-9} p/\rho . \tag{8}
\]

The molecular weight has been taken to be 0.5724, appropriate for fully-ionized hydrogen and helium, with the helium having 63.3\% of the hydrogen number density (Ross and Aller 1976). Equations (7)–(9) are inappropriate at low temperatures where the hydrogen and helium are not ionized. The resulting errors are inconsequential, however, since the heat conduction is totally negligible at those low temperatures.

The quantity \( L \) in equation (4) represents the radiative loss rate and is given by the expression

\[
L = N_e N_\rho \Phi(T) , \tag{10}
\]

where \( N_e \) is the electron number density, \( N_\rho \) is the hydrogen number density, and \( \Phi(T) \) is a modified form of the radiative loss function due to Raymond (1979). Above \( 10^4 \) K, we use the Raymond rates. Between 9500 and \( 10^4 \) K, the radiative losses decrease linearly to zero. These rates are thus nearly identical to the rates used in Rosner, Tucker, and Vaiana (1978).

The last term on the right-hand side of equation (3) can be regarded as the term which couples the Alfvén wave \((v_\theta, B_\theta)\) into the acoustic-gravity wave \((v_z\). This term can be interpreted as due to the radiation pressure of the Alfvén wave, in close analogy with the well-known radiation pressure of light. There is evidence that the radiation pressure term is important to the solar wind, where it may account for the high-speed streams.
Equations (1)-(5) 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 flux-corrected transport scheme of Boris and Book (1973, 1976). An updated version of this scheme with the Eulerian grid option was used for all the computations described in this paper. The heat conduction flux $q_x$, and the radiation have been calculated implicitly. The calculations were performed on a finite-difference grid of 200 computational cells covering a height range of 6000 km. The grid is variably spaced to insure adequate resolution in the chromosphere, transition region and corona; see Mariska et al. (1982) for details.

Initially, the solar atmosphere is taken to be in hydrostatic equilibrium. The initial chromospheric temperature has been chosen to have the height-independent value of 9500 K. The initial temperature profile in the corona and transition region is computed by balancing electron heat conduction and radiation against a height-independent volumetric heating rate of $1.23 \times 10^{-2}$ ergs cm$^{-3}$ s$^{-1}$. Such a heating rate is postulated as being representative of the unknown coronal heating mechanism. A similar procedure has been used previously by Mariska et al. (1982) and Mariska and Boris (1983) in studies of other aspects of coronal dynamics. At a height of 4400 km above the base of the transition region, we obtain a coronal pressure of 1.91 dynes cm$^{-2}$ and a coronal density of $1.03 \times 10^{-14}$ g cm$^{-3}$. The corresponding temperature is $1.3 \times 10^6$ K. These coronal parameters are taken to be reasonably representative of coronal active region loops. In all the computations reported in this paper, the initial magnetic field strength has been taken to be 20 G. This value is somewhat less than the typical field strengths in active region loops reported by Golub et al. (1980) but is adequate for elucidating the physics of the evolution of Alfvénic pulses.

A wave launched at one end of a loop is partially reflected as it traverses the inhomogeneous chromosphere and transition region, and it is partially transmitted into the more uniform corona. However, the transmitted wave is strongly reflected by the transition region and chromosphere at the "far end" of the loop. For the purposes of this paper, we have verified that it is sufficient to replace the transition region and chromosphere at the far end by a reflecting boundary condition in the numerical code. In the calculations reported here, the reflecting boundary condition is applied at a height of 6000 km above the base of the model, i.e., at 4400 km above the base of the transition region. Thus the coronal part of the loop is only 4400 km long in these calculations. Both coronal active region loops and X-ray bright points tend to be much larger than this value (e.g., Golub et al. 1980; Rosner, Tucker, and Vaiana 1978). We have taken an artificially short loop primarily in order to reduce the number of grid points in the numerical computation. However, we have performed a few calculations for longer loops, and we have verified that the loop length has no important effect on our results. We also use a solid end wall at the base of the model. Reflections from this end wall have been verified to have a negligible effect on the solutions presented here. Thus mass is conserved. And apart from the source of $v_g$ at $z = 0$, radiation, and the assumed background heating rate, the net momentum and energy of the system are conserved.

In all of the calculations, an Alfvén wave is launched by imposing a time-dependent twist, $v_g$, as a boundary condition at $z = 0$. In all cases the time dependence of $v_g$ at $z = 0$ has been taken to be sinusoidal, truncated after one or more half-cycles.

### III. RESULTS: NO DISSIPATION

As discussed in the previous section, heat conduction and radiation are included in the specification of the initial equilibrium atmosphere. But since the initial atmosphere is in hydrostatic equilibrium once it is set up, it is possible to "turn off" the heat conduction, radiation, and coronal heat input and still maintain the initial atmosphere. This procedure has been followed for the calculations of this section. Thus, except for the possible formation of shocks, the waves propagate without dissipation. The case where the waves dissipate will be considered in the next section.

We consider first the effects of a short-period Alfvénic pulse. The twist at $z = 0$ is imposed for a single half-cycle, lasting for 36 s. The velocity amplitude at $z = 0$ is 13 km s$^{-1}$. For comparison, the Alfvén speed, $v_A$, at $z = 0$ is 10.2 km s$^{-1}$ for $B_0 = 20$ G. The initial Alfvénic twist is thus moderately nonlinear at $z = 0$; the nonlinearity, as measured by $v_g/v_A$, decreases slowly with height in the coronosphere. For an Alfvén wave which propagates according to the WKBJ approximation, it is easy to show that $v_g/v_A \propto \rho^{1/4}$ if $B_0$ is constant.

Figure 1 displays $v_g$ and $v_e$ in the chromosphere at 33 successive times, from $t = 0$ to $t = 128$ s in evenly spaced intervals of 4.0 s. These figures employ a sliding scale for the ordinate. At each successive time, the curve for $v_g$ is displaced downward by 6 km s$^{-1}$ and the curve for $v_e$ is displaced upward by 6 km s$^{-1}$. At early times, Figure 1a shows the upward propagation of the Alfvénic pulse. The pulse accelerates at greater heights, because $v_A$ is an increasing function of height due to the decreasing density. The front of the pulse shows a tendency to steepen. This is a consequence of the nonlinearity, but no shock front is obtained in this example. At $t \approx 68$ s, the Alfvénic pulse impinges on the transition region at $z \approx 1600$ km, which reflects most of the energy in the Alfvénic pulse. The downward propagation of the reflected pulse is evident at later times in Figure 1a. Figure 1b shows the evolution of the vertical motions which are nonlinearly driven by the Alfvénic pulse. The upward-propagating Alfvénic pulse directly drives a weak upward velocity pulse, which is not visible in Figure 1b. This pulse occurred also in the calculations of HJG. A similar downward-velocity pulse, driven by the reflected Alfvénic pulse, can, however, be seen following the dashed line in the upper left-hand corner of Figure 1b. Of more interest are the two upward-propagating pulses indicated by the arrows in Figure 1b. These pulses were not discussed by HJG. They can be thought of as moderately large amplitude acoustic-gravity waves which are left in the wake of the initial Alfvénic pulse. The leading pulse propagates faster than the following because it propagates into upward-moving gas (see below), while the following pulse propagates into the down-falling gas behind the leading pulse.

Figure 2 displays $\rho$ in the chromosphere at 33 successive times, from $t = 0$ to $t = 128$ s in evenly spaced intervals of 4.0 s. At each successive time the curve for $\rho$ is displaced upward by a multiplicative factor of 1.58. The rapid density
The behavior of the pressure in the chromosphere is displayed in Figure 3. The acoustic-gravity pulses are evident. The downward-propagating pulse which is directly driven by the reflected Alfvénic pulse is just discernible in the upper left portion of the figure.

We have already noted that the Alfvénic pulse pushes the transition region upward. This is similar to the effect noted by HJG, except that in their calculations the Alfvénic pulse had steepened into a switch-on shock. In the present calculations,
the Alfvénic pulse induces only a weak upward motion of the transition region.

The two acoustic-gravity pulses shown in Figures 1–3 also interact with the transition region. They, however, have a much stronger effect on the transition region than does the Alfvénic pulse. To illustrate this, Figure 4 displays $\rho$ in the chromosphere and corona at 57 successive times, from $t = 36$ s to $t = 260$ s at evenly spaced intervals of 4 s. (At each successive time the curve for $\rho$ is displaced upward by a multiplicative factor of 1.58.) The lower part of the figure shows the upward motion of the transition region and underlying chromosphere, to a maximum height of approximately 2300 km. This upward motion is induced by the interaction of the leading acoustic-gravity pulse with the transition region. After reaching its maximum height, the transition region falls under the influence of gravity. This fall is then inhibited by the interaction of the second acoustic gravity pulse with the transition region, as can be seen in the middle part of Figure 4. The
downfalling transition region then rebounds and begins to move upward again in the top portion of the figure. In this example, the maximum upward velocities of the transition region are of the order of 21 km s\(^{-1}\).

The effectiveness of acoustic-gravity waves in producing upward motions of the transition region and underlying chromosphere agrees with the results of the numerical spicule model of Hollweg (1982). However, there are some important differences between the present results and those of Hollweg (1982). In the latter paper, the acoustic-gravity waves were generated by a spatially localized impulse in the photosphere or low chromosphere. In agreement with a linear calculation of Rae and Roberts (1982), this resulted in an extended train of acoustic-gravity waves, which nonlinearly became a train of hydrodynamic shocks. The repeated interaction of the wave train with the transition region was shown to be capable of elevating the transition region to heights characteristic of spicules, viz., 10^4 km or more. In the present calculation, the acoustic-gravity waves are not generated by a localized impulse. They are instead generated by the Alfvénic pulse, which is spatially extended over much of the chromosphere and rapidly changing its location in time. Our present results yield a pair of acoustic-gravity pulses, rather than an acoustic-gravity wave train (although some weak oscillations of the chromosphere do persist for very long times). Moreover, the two acoustic-gravity pulses are spaced so far apart in time that the transition region falls back down before encountering the second pulse. Thus the repeated interactions with a wave train, which were shown by Hollweg (1982) to be effective in producing spicule-like structures, do not occur in the problem being studied here. Even though Alfvénic pulses do lead to field-aligned motions of the transition region and underlying chromosphere, they do not seem to be capable of generating structures which resemble spicules.

This conclusion is fairly insensitive to the pulse length. For example, consider the effect of reducing the frequency of the initial Alfvénic twist, so that the twist at \(z = 0\) is imposed for a single half-cycle lasting 50.0 s. As in the previous example, the velocity amplitude at \(z = 0\) is 13 km s\(^{-1}\). Figure 5 shows the vertical velocity in the chromosphere at 26 successive times from \(t = 0\) s to \(t = 200\) s in equally spaced intervals of 8.0 s. (At each successive time the curve for \(v_z\) is displaced upward by 4.0 km s\(^{-1}\).) Figure 5 shows the same general features as were evident in Figure 1b; the two acoustic-gravity pulses are evident in the lower half of the figure. The corresponding response of the transition region is illustrated in Figure 6, which shows \(\rho(z)\) at 26 successive times, from \(t = 0\) s to \(t = 200\) s in equally spaced intervals of 8.0 s. (At each successive time the curve for \(\rho\) is displaced upward by a multiplicative factor of 1.58.) The results are quite similar to those displayed in Figure 4, except that the decreased frequency results in more energy in the wave available to push the transition region to a somewhat greater height.

The coupling of the Alfvénic twists to the field-aligned motions is quadratic in the amplitude of the Alfvénic twists. See equation (3), where the coupling explicitly depends on \(B^2\). We therefore anticipate that the amplitudes of the acoustic-gravity pulses, and their subsequent effects on the motion of the transition region and underlying chromosphere, will depend strongly on the amplitude of the Alfvénic pulse. For example, consider the same situation as was illustrated in Figures 5 and 6, but with the amplitude of the twist increased to 20 km s\(^{-1}\) at \(z = 0\). The motion of the transition region in this case is displayed in Figure 7, which shows \(\rho(z)\) at 39 successive times, from \(t = 0\) s to \(t = 304\) s in equally spaced intervals of 8.0 s. (At each successive time the curve for \(\rho\) is displaced upward by a multiplicative factor of 2.51.) In this case the transition region reaches a maximum height of
Fig. 5.—The evolution of $v_z$ at 26 times from $t = 0$ to $t = 200$ s for a pulse with an amplitude of 13 km s$^{-1}$ with a half-cycle lasting 50.0 s. At each time the curves are displaced upward by 4 km s$^{-1}$.

3400 km, compared to approximately 2300 km in Figure 4 and 2340 km in Figure 6. Thus single Alfvénic pulses can result in significant upward motion of the transition region and underlying chromosphere, but the twisting velocities have to take on rather large values.

In Figure 7 it is interesting to note that the density profiles are quite flat for more than 1000 km below the elevated transition region. This is a feature of our results which agrees qualitatively with a well-known property of spicules. However, the densities obtained in Figure 7 are considerably larger than those commonly associated with spicules. For example, at the point indicated by the filled circle in Figure 7, the density is $1.83 \times 10^{-12}$ g cm$^{-3}$, or $10^{12}$ protons cm$^{-3}$. This is one or two orders of magnitude denser than the "typical spicules" reviewed by Beckers (1972). On the other hand, the computed density is comparable to the density reported by Landman.
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IV. RESULTS: DISSIPATION

When the effects of radiation and conduction are included in the calculations, the results obtained in the previous section change somewhat. In the chromosphere conduction is of little importance, but radiation is an important energy loss mechanism. In the corona both radiation and conduction are important.

Figure 8 displays the evolution of $v_e$ and $v_z$ for the same pulse that is presented in Figure 1. The plots are formatted in the same way as Figure 1. Comparison of Figure 8a with Figure 1a shows that the evolution of $v_e$ is not significantly altered by the presence of radiation and conduction in the atmosphere. This is the expected result since $v_e$ does not directly couple to the radiation or conduction terms in the energy equation. The acoustic-gravity waves that follow the Alfvénic pulse are significantly altered, however. Comparison of Figure 8b and Figure 1b shows the effect. The acoustic-gravity waves are now propagating as radiatively and conductively damped waves. As the wave compresses the chromospheric and coronal plasma the plasma radiates more efficiently and removes energy, reducing the amplitude of the wave. The conduction also acts to smooth the temperature gradient associated with the wave. Note in particular the reduced amplitude of the acoustic-gravity waves in the corona.

These effects are clear in both the density and pressure plots in Figures 9 and 10, which should be compared with Figures 2 and 3. The energy removed from the waves is of course not available to push the transition region upward. Comparison of Figure 9 and Figure 2 shows the result of these losses. The evolution of the pulse shown in Figure 9 results in a maximum height for the transition region of only about 1960 km, while the same pulse with no dissipation reached a height of about 2300 km. Similarly, the maximum velocity is also reduced. In this case, the maximum upward velocities of the transition region are on the order of 10 km s$^{-1}$, while for the same case with no dissipation they were on the order of 21 km s$^{-1}$.

Figures 11 and 12 show the evolution of the density profile in the chromosphere for the same cases presented in Figures 6 and 7, respectively. For the case shown in Figure 11 in which the frequency has been reduced from the calculation shown in Figure 10, the transition region reaches a maximum height of about 2050 km and a maximum upward velocity of about 12 km s$^{-1}$, compared to 2340 km and 20 km s$^{-1}$ in Figure 6. When the pulse amplitude is increased to 20 km s$^{-1}$, the transition region evolution shown in Figure 12 leads to a maximum height of about 2700 km and a maximum upward velocity of about 24 km s$^{-1}$, compared to 3400 km and 40 km s$^{-1}$ in Figure 7.

The density profile shown in Figure 12 shows the same flat region noted in Figure 7. The density in the flat region is about $2.1 \times 10^{-12}$ g cm$^{-3}$, roughly the same as the density in the flat region shown in Figure 7.

V. SUMMARY AND DISCUSSION

We have numerically studied several aspects of nonlinear Alfvén wave propagation in the solar atmosphere, which were not considered in the earlier study by HJG. In particular, we have emphasized the behavior of Alfvénic pulses and the acoustic-gravity waves which are nonlinearly produced by them. We have also included the effects of radiation and heat conduction, which turn out to have a strong damping effect on the acoustic-gravity waves, and we have considered conditions which are representative of solar active regions.

We have found that single (i.e., half-cycle) Alfvénic pulses do not in themselves have a strong influence on the chromosphere and transition region. Of much more significance are the...
Fig. 8.—The same as Fig. 1, but for a calculation that includes dissipation. In Fig. 8a, the minor velocity fluctuations at the transition region in the model are due to small imbalances between radiative losses and the divergence of the conductive flux, which is the dominant heat source in the lower transition region. This imbalance results in small velocity fluctuations which propagate upward into the corona, reflect off the upper boundary, and return. Their presence does not affect the interaction of the pulse with the transition region significantly.
acoustic-gravity waves which are left in the wake of the Alfvénic pulse. The single Alfvénic pulse produced two upward-propagating acoustic-gravity pulses, which basically constitute a full cycle. We have seen no indication of an extended oscillatory wake (or wave train) such as discussed by Rae and Roberts (1982). This is not unreasonable, since the wave trains result from a source which is impulsive in time and localized in space, in contrast to the situation here where the source of the acoustic-gravity waves is the $v_e$ pulse which is extended in space and time. This result effectively negates the suggestion by Hollweg (1982) that Alfvénic pulses might drive an acoustic-gravity wave train which could in turn produce spicules. And in fact the transition region and chromospheric motions obtained in this paper do not correspond to the motions characteristically associated with spicules.

Even though single Alfvénic pulses do not seem to lead to spicules, they can result in marked motion of the transition region and underlying chromosphere. More properly stated,
the Alfvénic pulse nonlinearly drives acoustic-gravity waves, and they in turn result in significant motions of the chromosphere and transition region. The extent of the upward motion of the transition region is governed primarily by the energy content of the original Alfvénic pulse. As a benchmark, an Alfvénic pulse lasting for 36 s with a maximum velocity amplitude of 13 km s\(^{-1}\) at \(z = 0\) results in an upward excursion of the transition region of some 360 km in the dissipative case. Longer period pulses result in greater excursions. And, by inference, Alfvénic wave trains will similarly lead to greater excursions; this inference is consistent with the results of HJG, but those authors were primarily concerned with trains of fast shocks.

Thus a dynamic transition region and chromosphere appears to be an inevitable consequence of the presence of Alfvén waves, and, indeed, of the process of coronal heating.

Fig. 11.—The evolution of \(\rho\) at 32 successive times from \(t = 0\) to \(t = 248\) s for the same pulse parameters presented in Fig. 6, but with dissipation. At each successive time the curve is displaced upward by a factor of 1.58.

Fig. 12.—The evolution of \(\rho\) at 39 successive times from \(t = 0\) to \(t = 304\) s for the same parameters presented in Fig. 7, but with dissipation. At each successive time the curve is displaced upward by a factor of 2.51.
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This viewpoint is not without its difficulties. We have found that radiative losses in the chromosphere and transition region can significantly dissipate the acoustic-gravity waves which are driven by Alfvén waves. Nonetheless, Figures 9, 11, and 12 show that appreciable motions of the transition region and chromosphere can still occur in this case but to a lesser extent than in the dissipationless case. Another possible difficulty is the sensitivity of the results to the amplitude of \( v_e \). We found that rather large values of \( v_e \) (in excess of 10 km s\(^{-1}\)) had to be imposed at \( z = 0 \), in order to produce significant transition region displacements (more than a few hundred km). On the other hand, inspection of Figures 1a and 8a shows that the magnitude of \( v_p \) does not increase dramatically in the chromosphere, and that the model velocities in the upper chromosphere are of the same order as the horizontal velocities reported by Canfield and Beckers (1975).

Other difficulties are associated with the limitations of the numerical model. Motions in the third spatial dimension have been artificially suppressed, we have only considered a vertical flux tube of constant cross section, and we have not been able to explore strong magnetic fields (\( > 20 \) G). And we have made no provision for direct dissipation of the \( \theta \)-motions (by viscosity or electrical resistivity, for example). These restrictions will have to be relaxed in future work.

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