SOLAR TRANSITION REGION RESPONSE TO VARIATIONS IN THE HEATING RATE

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ABSTRACT

We have examined the response of a numerical model for the upper chromosphere, transition region, and corona to variations in the energy input. The numerical model solves the set of one-dimensional two-fluid hydrodynamic equations in a simple vertical magnetic flux tube. The atmosphere responds to both the increase and decrease in energy deposition by smoothly readjusting the temperature gradient and the amount of material in the region of peak radiating efficiency to radiate away energy being deposited. At no time during this readjustment do we see a departure from a thin laminar transition region structure. In addition, a time-dependent description of the nonequilibrium ionization of all of the ionization stages of oxygen has been included. This calculation is coupled with the self-consistent calculations of the dynamical variables. We find that the nonequilibrium ionization balance calculations for both heating and cooling small loops in the quiet Sun predict relative ionic abundances which differ substantially from those which would be predicted by an equilibrium calculation.

Subject headings: hydromagnetics — Sun: chromosphere — Sun: corona

I. INTRODUCTION

The chromosphere-corona transition region plays a major role in determining the structure and energy balance of the outer layers of the solar atmosphere. The conventional picture of this region in the quiet Sun (e.g., Withbroe and Noyes 1977) is one of a thin layer of plasma whose structure is determined by a balance between local radiative losses and the energy provided by thermal conduction from the overlying corona, possibly with some additional energy deposition directly in the layer. Since this heat flux is the major energy loss from the quiet corona, the ability of the chromosphere and transition region to radiate that energy determines how much material can be supported in the corona. Small changes in the coronal heat source can therefore produce significant readjustments in density and temperature in both the transition region and the overlying corona. These readjustments should be reflected in EUV emission lines formed in this region and hence are in principle amenable to direct observation. It is our intent in this paper to study some of the dynamics of transition region readjustment to changes in an assumed heating rate. In addition, by including a dynamic calculation of ionization levels for signature ions of oxygen, we can predict directly and in a time-dependent manner some of the observational consequences of our assumed initial conditions and heating functions.

While the driving concern must be an understanding of the still mysterious coronal heating mechanism, this is a phenomenon which we have not conclusively identified. The fastest path to understanding coronal heating is, therefore, to understand the response of structures which we can see, such as transition regions, spicules, flux tubes, coronal arches, filaments, and condensations, to variations in energy input. Matching predicted and observed behavior quantitatively allows us to narrow appreciably the range of possible properties of heating mechanisms.

In this paper we use the NRL Dynamic Flux Tube Model (Boris et al. 1980) to examine the response of the transition region in the quiet Sun to variations in the energy input. We first construct a static equilibrium model for an idealized magnetic flux tube such as might be observed in quiet network regions. The flux tube consists of an idealized chromosphere, transition region, and corona. A constant volumetric heat source balances the radiative losses. Choosing this simplest heating function has several advantages. Because we are not certain of the coronal heating mechanism, a simple heating function is the best way to begin to establish the response of a quiet Sun transition region flux tube to changes in heating. Further, since the heating rate does not vary in space and generally not in time, we do not see instabilities arising as a result of the form of the heating function, a phenomenon observed in earlier calculations (Antiochos 1979; Craig and McClymont 1981). In the calculations reported here, most of the...
volume of the flux tube lies in the corona and hence most of the system heating occurs there even though most of the system mass lies below the transition region.

We first examine the response of this system to decreasing and then increasing the magnitude of the heating term, each time starting at the equilibrium conditions. Finally, we discuss the response of the atmosphere to consecutive cooling and then heating to simulate a fluctuating or intermittent coronal heating such as attends many of the proposed dynamic coronal heating mechanisms.

The motivation for this work is severalfold. First we are interested in the dynamic response of the physical variables such as temperature, density, pressure, and velocity to variations in the heating rate. Since the model we are using incorporates many of the physical processes that we believe are important on the Sun, it can be used to determine whether these processes alone are adequate to predict solar behavior. Previous models for the transition region and inner corona in quiet regions and coronal holes (e.g., Rosner and Vaiana 1977; Athay 1981) have been used to investigate either a static atmosphere or at best atmospheres with a steady flow of mass. They are thus solving a limited set of conservation equations. Since the model used in this work solves a set of coupled time-dependent conservation equations for the solar atmosphere, it is self-consistent and the dynamic response of the atmosphere can be studied.

Another motivation behind this work is the wealth of observations of EUV emission lines formed in the transition region and inner corona. These exhibit brightness changes, nonthermal line broadening, and apparent systematic mass flows (e.g., Withbroe and Noyes 1977). These dynamic phenomena may in part be due to the adjustment of the transition region to variations in coronal heating. It is clear, however, that these effects make equilibrium ionization balance calculations of questionable value for predicting EUV emission line fluxes. Previous work on this problem has involved calculations of ionization balance in steady flows through predetermined temperature-density models (e.g., Raymond and Dupree 1978; Dupree, Moore, and Shapiro 1979; Francis 1981; Roussel-Dupré and Beerinan 1981). Thus these calculations are somewhat restricted and handle only a limited class of nonequilibrium spectroscopic effects.

In this paper we investigate departures from ionization equilibrium by using the flux tube model to calculate both equilibrium and nonequilibrium number densities of O i through O ix. This calculation is coupled with the self-consistent calculations of the dynamical variables. We then ask whether an equilibrium ionization balance model is adequate for predicting EUV emission line intensities. We show that in fact the equilibrium ionization balance assumption leads to significant departures from the actual nonequilibrium ionization balance.

High-resolution EUV observations also show that the transition region in the quiet Sun is highly structured. Understanding this structure and its implications for coronal heating is a third motivation for this work. The observed concentration of emission into the network regions (e.g., Reeves 1976) where the magnetic field is also concentrated has led to simple magnetic field models for the quiet transition region (e.g., Gabriel 1976). These models have magnetic field lines that diverge rapidly from the chromospheric network, become nearly radial in the low corona, and then close on a larger scale, resulting in the observed large-scale coronal structure.

On the other hand, the large amount of small-scale structure observed in the network at transition region temperatures could be due to small-scale closed structures. Lyman-alpha filtergrams indicate that the chromospheric network at the base of the transition region may be composed of many small loops (Bonnet et al. 1980). These loops would reach to heights of only a few thousand kilometers, the characteristic width of the network in transition region lines. Simple static models, such as those of Rosner, Tucker, and Vaiana (1978), would then indicate that for typical quiet Sun pressures, the maximum temperatures in these structures would be only a few hundred thousand kelvins.

Previous calculations of the formation of a transition region and its strong dynamical stability (Oran et al. 1980; Oran, Mariska, and Boris 1982) indicate that a single laminar transition region is considerably thinner than some observations would imply (e.g., Mariska, Feldman, and Doschek 1978). A conglomeration of ragged, patchy transition regions extending over many times the ideal laminar thickness is one possible way to reconcile the observations and calculations. Antiochos (1979) and An, McClymont, and Canfield (1980) explore the possibility of condensational instability in the transition region causing this raggedness. Such instabilities depend rather crucially on the form of the assumed heating function rather than on the intrinsic structure of the transition region. In fact, transition regions are physically extremely stiff structures because the heat flux and the radiation flux interact so strongly and the coronal sound speed is high. In our previous studies of the nonlinear condensational instability the structure of a rapidly varying transition region was always monotonic (Oran, Mariska, and Boris 1982). We did find that large condensations could be accelerated through the transition region in much the same manner as lasers can accelerate solid material by ablation (e.g., Bodner 1981; Max, McKee, and Mead 1980) without appreciably altering the transition region structure.

Previous indications, as well as calculations presented below, discourage the consideration of a ragged transi-
tion region in relatively quiescent situations. Thus the alternative explanation of a thin-skinned transition region hugging the contours of spicules, low-lying flux tubes, mass ejections, and small-scale closed structures seems more acceptable to explain the apparent thickness of the region of emission for transition region lines. The small loop studies presented below help us to examine whether such a model is consistent with the phenomena observed in the network at transition region temperatures.

II. THE MAGNETIC FLUX TUBE MODEL

a) Fluid Equations

We wish to isolate the effects of heating changes from other influences such as spatial energy deposition and geometry variations. So, for this study, we use the NRL Dynamic Flux Tube Model (Boris et al. 1980) and consider a vertical flux tube of constant cross sectional area. The plasma is composed of electrons and ions with total pressure $P$, mass density $\rho$, and fluid velocity along the flux tube $v$. The total pressure is the sum of the electron pressure $P_e$ and the ion pressure $P_i$, and is characterized by an electron temperature $T_e$ and an ion temperature $T_i$. We solve the equations for mass, momentum, and energy conservation,

$$\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x} (\rho v) = 0, \quad (1)$$

$$\rho \left( \frac{\partial v}{\partial t} + v \frac{\partial v}{\partial x} \right) + \frac{\partial P}{\partial x} = \rho g, \quad (2)$$

$$\frac{\partial}{\partial t} \left( \frac{1}{2} \rho v^2 + \frac{P}{\gamma - 1} \right) + \frac{\partial}{\partial x} \left[ \left( \frac{1}{2} \rho v^2 + \frac{P}{\gamma - 1} \right) v \right]$$

$$- \kappa_e \frac{\partial T_e}{\partial x} - \kappa_i \frac{\partial T_i}{\partial x} = \rho vg - L + S, \quad (3)$$

and

$$\frac{\partial}{\partial t} \left( \frac{P_e}{\gamma - 1} \right) + \frac{\partial}{\partial x} \left( \frac{P_e}{\gamma - 1} - \kappa_e \frac{\partial T_e}{\partial x} \right)$$

$$= - L + S + \gamma_{eq} (T_i - T_e) - P_e \frac{\partial v}{\partial x}. \quad (4)$$

Here $g$ is the gravitational acceleration at the solar surface; $\gamma$ is the ratio of specific heats, taken to be 5/3 in our calculations; $\kappa_e$ and $\kappa_i$ are the electron and ion thermal conductivities and are given by

$$\kappa_e = 1.1 \times 10^{-6} T_e^{5/2} \quad (5)$$

and

$$\kappa_i = \kappa_e / 25; \quad (6)$$

$L$ is the energy loss rate due to radiation; $S$ is the heating rate; and $\gamma_{eq}(T_e)$, the rate of equilibration between electrons and ions, is given by

$$\gamma_{eq} = 1.4 \times 10^{-17} N_e^2 T_e^{-3/2}, \quad (7)$$

where $N_e$ is the electron number density. The constant in the electron thermal conductivity is from Ulmschneider (1970), while the relationship between the electron and ion thermal conductivities and the equilibration rate between the electrons and ions are from Braginskii (1965). The plasma consists primarily of fully ionized hydrogen and helium, with the helium having 6.3% of the hydrogen number density (Ross and Aller 1976). The electron number density $N_e$ and the ion number density $N_i$ are then related to the mass density by

$$N_e = \frac{\rho Z}{\mu m_p (1 + Z)} \quad (8)$$

and

$$N_i = N_e / Z, \quad (9)$$

where $Z$ is the mean ionic charge taken as 1.059 and $\mu$ is the mean mass per particle taken as 0.5724 in proton masses $m_p$. These values were calculated using only the hydrogen and helium abundances. Because of the large uncertainty in the solar helium abundance, they are, however, fully consistent with the net solar composition including trace elements. Assuming temperature-independent values of $\mu$ and $Z$ simplifies necessary equation-of-state calculations considerably. The electron and ion pressures satisfy the equations of state

$$P_e = N_e k_B T_e, \quad (10)$$

and

$$P_i = N_i k_B T_i, \quad (11)$$

where $k_B$ is Boltzmann's constant.

The radiation loss rate is given by the expression

$$L = N_p N_e \Phi(T_e), \quad (12)$$

where $N_p$ is the hydrogen number density and $\Phi(T_e)$ is the radiative loss function due to Raymond (1979). This loss function is nearly the same as that given by Rosner, Tucker, and Vaiana (1978) and assumes cosmic element abundances. We assume that the heating rate $S$ is uniform in space and that the heat goes into the electrons only, with thermal equilibration responsible for its reapportionment between the electrons and ions.

Equations (1)–(4) are solved using time-step splitting techniques (e.g., Oran and Boris 1981). Equations (1)
and (2) and the convective portions of equations (3) and (4) are solved using the Flux Corrected Transport (FCT) algorithm (see e.g., Boris and Book 1976, and references cited therein). The remaining terms in equations (3) and (4) are solved implicitly.

Although the FCT models in principle have an advantage because of their ability to follow energy-conserving shocks accurately, this advantage amounts to little in these flux tube heating and cooling calculations because no consequential shocks are seen to develop. The remaining advantage of FCT is an approximately twofold improvement in resolution for a given grid size because this method is monotonic and the underlying algorithm is fourth order accurate. This means that calculations performed with older nonmonotonic algorithms would require zones of ~2.5 km size to obtain answers comparable to those obtained here using 5 km resolution in the steepest part of the transition region.

A number of computational models of coronal loops, using both monotonic FCT techniques (e.g., Hollweg 1981; Smith and Auer 1980) and older nonmonotonic schemes (Craig and McClymont 1980; Krall and Antiochos 1980; Wu et al. 1978) have been developed and applied to coronal loop and flare problems. Representing the transition region accurately enough to say anything about its structure and response has always represented a serious computational problem, however, so connections of a coronal loop to the chromosphere have always been of questionable accuracy. Hollweg and Smith and Auer have focused on larger scale coronal problems and treated the transition region as a thin layer of indeterminant structure. Studying the effects of specific and localized large-scale energy additions reduces quiet Sun transition region dynamics to secondary importance during the short periods studied computationally. Krall and Antiochos attacked the resolution problem directly by developing an analytic transition-region boundary condition as the lower end condition on a coronal loop, but this approach did not permit treatment of several aspects of coronal response. In particular the chromospheric evaporation was postulated rather than calculated.

These calculations have generally had zones 10 km or larger (e.g., Nagai 1980; Craig and McClymont 1981) in the steepest portions of the transition region. This seems barely adequate at best, but represents a compromise between desired accuracy and computational expense. Craig and McClymont (1981), in presenting some similarly motivated heating calculations, discuss this aspect of error in their calculations as well as those of previous authors. In the calculations presented below we have been able to take advantage of both the increased accuracy of the FCT algorithms and the way these have been vectorized in a parallel processing computer. Thus we can consider nonequilibrium effects in the corona and increase the resolution appreciably in the transition region, both at no great increase in computational expense. A calculation including nonequilibrium ionization requires only a few minutes on the pipelined class VI Texas Instruments ASC at the Naval Research Laboratory.

b) Ionization Balance and Spectroscopy

In addition to solving the fluid dynamics equations, we also solve the ionization balance equations for all of the ionization stages of oxygen. The ion abundances are then used to calculate the radiation expected from a selection of EUV emission lines.

The ionization balance for oxygen ions is described by a set of equations of the form

$$\frac{1}{N_e} \left( \frac{\partial N_j}{\partial t} + v \frac{\partial N_j}{\partial x} \right) = - (Q_j N_j + \alpha_j N_j) + Q_{j-1} N_{j-1} + \alpha_{j+1} N_{j+1},$$

where $N_j$ is the number density of ion $j$, $Q_j$ is the ionization rate coefficient from ion $j$ to ion $j+1$, and $\alpha_j$ is the recombination rate coefficient to ion $j-1$ from ion $j$. For the ionization rate coefficients, we use an analytic fit to the data given by Lotz (1967). For the recombination rate coefficients, we use the results of Aldrovandi and Pécougniot (1973, 1976). Both radiative and dielectric recombination are included. For the recombination rate from O IX, we use the results of Burgess and Seaton (1964). Equilibrium ionization balance calculations made using these sources are in good agreement with the more recent results tabulated by Jacobs et al. (1978). The convective terms in the set of ionization balance equations are solved using the same FCT algorithm that is used for the fluid equations. The remaining portions of the equations are solved using the selected asymptotic integration method (Young 1980).

While the ionization balance is solved in a fully time-dependent manner, it should be noted that this calculation is relatively decoupled from the fluid dynamics. The density and temperature from the fluid dynamics calculation are used to solve the set of equations given in equation (13). This information, however, is not used to calculate a time-dependent radiative loss rate for the loss terms in equations (3) and (4), nor is it used in an equation-of-state calculation for the plasma.

We calculate level populations and emission line intensities for the ions O III through O VII. The atomic parameters for these calculations have been obtained from a number of sources. For O III we use collision strengths and spontaneous radiative decay rates calculated by Bhatia, Doschek, and Feldman (1979) to obtain populations for a 20-level atomic model. For O IV we use collision strengths and spontaneous radiative decay
rates from Flower and Nussbaumer (1975) in a 15-level atomic model. For O v we use spontaneous radiative decay rates and collision rates from Duffon et al. (1978) in a 10-level atomic model. The collision rates have been converted to collision strengths for the level population calculations. For O vi we assume a simple two-level atomic model and use the absorption oscillator strengths and average Gaunt factors tabulated by Dupree (1972). For O vii we use rate coefficients from Pradhan, Norcross, and Hummer (1981a, b) and spontaneous radiative decay rates from Gabriel and Jordan (1972) to calculate resonance, intercombination, and forbidden line intensities. The O vii line intensities are calculated using these data in the manner outlined by Gabriel and Jordan (1969, 1972).

c) The Initial Conditions

Consider a vertical flux tube with a total length of $6 \times 10^8$ cm and a constant radius of $3 \times 10^7$ cm. The base of the tube consists of approximately $10^8$ cm of plasma in hydrostatic equilibrium at $10^4$ K. This region provides a dynamically accurate but simple approximation of the solar chromosphere. Above this there is a transition region in hydrostatic equilibrium whose properties are found by solving the static force and energy balance equations. The pressure at the base of the transition region is taken to be 0.2 dyn cm$^{-2}$, a number typical of the quiet Sun.

The finite difference grid laid down to represent this flux tube has variable resolution and consists of 100 cells. The idealized chromosphere, with a scale height of 500 km, is represented by 20 cells of characteristic depth 50 km. The narrow transition region is represented by 40 cells of characteristic thickness 5 km. The overlying corona has relatively small gradients, and therefore 40 cells varying from 5 km to 540 km in an exponential manner were used. Computational grid sizes were gradually transitioned from one size to another at the boundaries of these regions to maintain accuracy in the finite-difference formulations of equations (1)-(4). The system time step was generally determined in the vicinity of the highest 5 km wide cell and was typically chosen to be 0.026 s. The sound speed, radiation rate, and thermal conduction limits can all be important to this determination since the structure in this area of the transition region is the result of a balance of these effects. For all of the calculations described in this paper the computational grid was fixed in space.

The energy input, assumed to be uniform in space, was adjusted so that the temperature gradient was zero at the top. This results in a peak temperature of about $6.7 \times 10^5$ K and an energy input of $8.34 \times 10^{-4}$ ergs cm$^{-3}$ s$^{-1}$. The total coronal heating in this loop is about $1.2 \times 10^{29}$ ergs s$^{-1}$ and this appears primarily as a downward heat flux of about $4 \times 10^5$ ergs cm$^{-2}$ s$^{-1}$ into the transition region from the corona. At the lower boundary the temperature is fixed at $10^4$ K.

These initial conditions were used to begin a calculation which was run for several thousand time steps ($\sim 590$ s) with the equilibrium heating on to allow any small inconsistencies in the static model to settle before the more dynamic heating and cooling calculations were begun. The $10^4$ K region of the atmosphere was initially constructed without allowing for the energy deposition. The radiation rate at this temperature is greater than the energy deposition rate, so this region exhibits a slight cooling during this time period and reaches equilibrium with the energy input at a temperature slightly below $10^4$ K. Figure 1 shows the atmospheric structure after these initial calculations.

Even with 5 km resolution in the transition region, the resulting structure is only an approximation to the true structure implied by the equations and initial and boundary conditions. Figure 2 shows the temperature structure of the first 50 km of the transition region at the end of the initial calculations. Also plotted is the structure that results from the same calculation performed with 1 km resolution in this region. Only in the lowest 10 km of the transition region do differences between the two calculations become apparent, with the lower resolution calculation producing a slightly shallower temperature gradient.

The radiation losses at the base of the transition region are also dependent on the resolution adopted there. Figure 3 shows the radiative loss rate in the same atmospheric region shown in Figure 2. Again, only in
the lowest 10 km of the transition region are the differences between the two calculations apparent. Even with these differences, the integrated radiation rate out of the lowest 10 km of the transition region differs by less than 20%.

It should be noted that the temperatures in the idealized chromosphere and lower transition region are low enough that there may be some effects due to recombination and radiative transfer. Rather than perform a detailed radiative loss calculation for the chromosphere, we have allowed the radiation rate to smoothly go to zero at a temperature of 9500 K. This represents only a crude approximation to the actual radiation loss in the chromosphere, but produces a reasonable chromospheric region.

During this initial adjustment the chromospheric region of the atmosphere undergoes a series of small oscillations in response to perturbations applied to it. The upper boundary of the chromosphere, which we define to be the height at which the temperature is 20,000 K, oscillates over a height range of ~40 km with a period of ~225 s. To understand these initial oscillations, we have analyzed the linearized mass, momentum, and adiabatic energy equations for a vertical column of isothermal plasma in hydrostatic equilibrium. For boundary conditions we assume that the velocity is zero at the base and that the pressure is constant at the upper boundary. The roots of the resulting eigenequation for the oscillation frequency show that the lowest frequency mode of oscillation has a period of ~179 s. Considering the neglect of radiative losses in the linear analysis, we believe that the agreement between the linear analysis and the model calculations is good.

It is clear from the dependence of the oscillations on the conditions at the boundary with the transition region that the chromosphere is dynamically linked to the layers above it. The chromosphere behaves much like a loaded spring with overlying layers providing most of the loading. Thus rapid changes in pressure, such as might be caused by cooling or heating in the overlying layers, will cause the chromosphere to rise or fall in response. It is thus important to include the chromosphere in our calculations as a portion of the overall dynamical system and not, as in some earlier calculations, as just a source or sink for mass and energy.

III. RESULTS

As described above, the atmosphere is initially in a static equilibrium. Thus the structure shown in Figure 1 is determined by the equation of hydrostatic equilibrium and an energy equation which expresses the balance among the energy gain from the heating, the energy loss from radiation, and the divergence of the heat flux, which can be either a heat source or a heat sink depending on the atmospheric region. At the top of the atmosphere, the energy input is greater than the radiation rate and the excess energy is conducted away to the lower regions of the atmosphere. This is true down to a temperature of about 3.9×10⁵ K at a height of about 1500 km, where the radiation rate balances the energy input rate. Throughout the remainder of the transition region the radiation rate exceeds the energy input rate and energy balance is maintained by heat conducted in from the upper portion of the atmosphere. The optically thin radiation rate is largest in the lowest portion of the
transition region in the temperature range between $2 \times 10^4$ K and $10^5$ K. In this region all of the heat flux from the overlying regions is dissipated. In our simplified chromosphere classical conduction is relatively unimportant and the radiation rate of the much denser plasma is balanced by the energy input.

a) Cooling

The response of the initial model equilibrium to the interruption of the heating that maintains the atmospheric structure is shown in Figures 4a, 4b, 4c, and 4d. Here we have plotted the temperature, pressure, electron density, and velocity at each height in the atmosphere at selected times during the cooling. When the heating ceases, the energy losses are no longer balanced and the atmosphere begins to cool. In the equilibrium atmosphere both radiation and conduction cooled the region near the top of the atmosphere. With heating turned off, both mechanisms continue to cool the atmosphere, with 20%–30% of the temperature drop due to radiative losses and the remainder due to conduction. This apportionment of the cooling mechanisms continues through much of the initial cooling. Lower in the atmosphere, near the base of the transition region,
radiation is the dominant cooling mechanism and conduction from above was and remains the heating source. As the temperature drops, however, the heat flux decreases and eventually can no longer balance the radiative losses. Then the plasma rapidly cools and condenses onto the chromosphere. The drop in temperature stops at the point when the radiation loss function ceases to be important, at a temperature just below 10,000 K in our model.

The process of cooling lowers the pressure throughout the transition region, which causes substantial departures from gravitational equilibrium. As a result, mass motions develop which try to establish a new balance. In the transition region the newly created force imbalance produces downward motions of transition region material. The reduction of the pressure in the transition region also results in an unbalanced pressure gradient in the chromosphere which causes the chromosphere to expand upward. The result is a region of negative (downward) velocities in the upper transition region, followed at lower heights by a region in which the negative velocities decrease due to the expanding chromospheric material, and finally a region of small positive velocities in the upper chromosphere. The presence of finite velocities is reflected in the calculated density profiles in Figure 4c. The densities decrease with time at the top of the atmosphere, and we observe a small decrease in density with time at the base due to the expansion of the chromosphere. The analysis of the preceding section, in which we considered the vertical oscillations of a loaded chromosphere in conjunction with the establishment of an equilibrium configuration, also applies here. There are three characteristic times interacting in this transient cooling problem: (1) The lowest harmonic bounce time of ~56 s (one-fourth of the oscillation period), (2) the acoustic transit time of ~49 s (transition zone to loop apex), and (3) the energy drain time of the corona of ~350 s (coronal heat capacity over coronal heating rate). Inspection of Figure 4c shows that the adjustment of the chromospheric electron density is essentially complete by ~300 s, which is somewhat longer than the characteristic bounce time. Because the energy drain time is also longer than the bounce time, the pressure on the top of the chromosphere is released slowly enough that resonant oscillation of the height of the transition region is not expected or observed.

The variability and fluctuations seen in the velocity profiles of Figure 4d exist because finite fluid accelerations arise from near but not exact cancellation of large terms in the momentum and energy equations. These variations do not, however, result in an unstable situation. Even dropping the heating instantly to zero is not considered to be a large perturbation by the transition region plasma. Continuous modest readjustments of the structure permit continued near balance of the heat fluxes, heat sources, and radiative losses throughout the cooling process. The system is never driven far enough from local equilibrium for nonmonotonic variations of density or temperature to occur. The temperature minimum of Figure 4a is a 20% decrease from the initial value arising from adiabatic expansion of the upper chromospheric layers as the coronal pressure from above is relieved.

While no departure from a simple laminar transition region or localized condensational instability is evident in this calculation to disrupt the structure of the chromosphere-corona transition, the base of the transition region has moved upward by 300 km under the combined action of chromospheric expansion and coronal condensation in the first 5 minutes after heating is interrupted. This corresponds to an apparent upward motion of about 1 km s\(^{-1}\). Roughly 30% of this apparent upward motion, however, is the accumulation of condensing coronal material. The physical variables shown in Figures 4a-4d do show that significant changes take place in the temperature-density structure and that these changes result in mass motions as the atmosphere cools. Observationally, these changes would be reflected in intensities and Doppler shifts of spectral lines formed in the transition region and low corona. Note that because of the upward chromospheric expansion, these shifts can be in opposite directions simultaneously, depending on where the various lines originate.

Doppler shifts indicating both downward and upward mass motions have been observed in the quiet Sun. Recent observations show an average downflow in bright network elements of about 4 km s\(^{-1}\) at the temperature of formation of the C \textsc{iv} lines (~10\(^{5}\) K) (Gebbie et al. 1980). Brueckner, Bartoe, and Dykton (1980) have reported velocities ranging from upflows of 20 km s\(^{-1}\) to downflows of 50 km s\(^{-1}\) in the same lines. Thus the calculated downflows fall within the range of the observations.

Figure 5a shows the way in which the cooling process affects the relative abundances of the ionization stages of oxygen, and Figure 5b shows the emission line luminosities for the EUV resonance lines of O \textsc{iv}, O \textsc{v}, and O \textsc{vi}. The relative abundances are calculated by integrating over the atmosphere in the model to determine the total number density of each ion. Since there is a large cool region at the bottom, O \textsc{i} and O \textsc{ii} are present in large quantities. This explains why most of the other ions shown in Figure 5a have relatively low integrated relative abundances. The line luminosities shown in Figure 5b are also integrated over the entire atmosphere in the model. Thus they indicate how an atmosphere with this structure would appear to a spectrograph that contains the entire structure, and only this structure, within its field of view.

Two sets of calculations are plotted in both Figures 5a and 5b. The solid lines represent the given quantity
Fig. 5.—Time evolution of (a) the integrated relative abundances of the ions O iv–O vii and (b) the integrated line luminosities of resonance lines of O iv (554 Å), O v (629 Å), and O vi (1032 Å) in response to the interruption of heating. The solid curves are for ionization equilibrium, and the dashed curves are for time-dependent ionization balance.

as a function of time for equilibrium ionization of oxygen. The dashed lines represent the results of the nonequilibrium oxygen ionization calculation. In the case of the nonequilibrium ionization balance calculations, because the recombination time is slower than the cooling time, the tendency of the species abundances to depart from equilibrium is apparent. For example, the recombination time for O vii is about 800 s at a temperature of $5 \times 10^5$ K and a density of $10^9$ cm$^{-3}$ (e.g., Kafatos 1973). Thus, as Figure 5a shows, the O vii abundance remains high throughout the cooling. The nonequilibrium behavior for the other ions shown in the figure is somewhat more complicated due to the variations in the dielectronic recombination rate with temperature for the intermediate ionization stages of oxygen. Interpretation of the relative abundances shown is further complicated by the integration over the entire atmosphere and, near the end of the calculation, by the development of mass motions. When we examine a particular location near the top of the atmosphere, we see that the ion abundances generally follow the cooling behavior discussed by Kafatos (1973) for a low density cosmic gas. That is, at high temperatures the abundance of an intermediate ionization stage will first build up due to recombination from the next highest stage. As the temperature drops, the abundance will decrease due to dielectronic recombination to the next lowest stage. Finally, dielectronic recombination ceases to be important, and the abundance again increases due to recombination from the next highest stage.

The integrated line luminosities shown in Figure 5b reflect the variations in abundance shown in Figure 5a as well as the temperature variations of the excitation rates. This temperature sensitivity is the reason why the ordering of the relative line intensities is opposite to that of the relative abundances at the end of the calculation. Specifically, the O vi 1032 Å line is easier to excite at low temperatures than the 629 Å line of O v or the 554 Å line of O iv.

b) Heating

The response of the initial conditions to an increase of a factor of 2 in the heating rate is shown in the four panels of Figure 6. Figures 6a, 6b, and 6c show the temperature, pressure, and electron density at each height in the atmosphere at the beginning and end of the calculation. Since the changes observed in these parameters are small, we have not plotted data at any intermediate times. Note that the chromosphere is compressed slightly in Figure 6c. Figure 6d shows the velocity at each height in the atmosphere at two intermediate times. Again substantial fluctuations are evident for the same reason given above: differences between large, nearly equal terms are involved in calculating the fluid acceleration. This calculation also demonstrates the strong dynamic stability of the transition region which was described earlier. The density and pressure profiles remain monotonic, and the transition region stays laminar.

In equilibrium the transition-region portion of the atmosphere is radiating energy at maximum efficiency. Thus the energy deposited by the increase in the energy deposition rate cannot be radiated away by the existing
Fig. 6.—The response of the (a) temperature, (b) pressure, (c) electron density, and (d) velocity to a factor of 2 increase in the energy input rate in the model atmosphere depicted in Fig. 1.

atmospheric structure. The additional energy input results in elevated local temperatures throughout the transition region and increases the heat flux conducted to the lower portions of the transition region. Since the lower transition region is already radiating at maximum efficiency, the additional energy is deposited at the top of the chromosphere, causing heating and convection of material upward into the transition region. This process has great similarities to evaporation because the transition region is so thin. The additional heating at the base of the transition region also increases the pressure and slightly compresses the chromospheric layers, producing small downflows. Eventually, enough material is convected into the transition region to allow the additional energy to be radiated away. A new equilibrium state then develops, with a steeper temperature gradient, higher pressure, and more mass in the transition region.

If the heating rate in the initial model atmosphere is increased by a factor larger than 2, even more material is convected into the transition region. The final configuration has a still higher peak temperature, and larger mass motions are produced. A factor of 4 increase in the
heating produces peak outflow velocities of about 5 km s\(^{-1}\), and a new relatively quiescent equilibrium is established after about 5 minutes.

Figures 7a and 7b show the manner in which the increased energy deposition affects the integrated relative abundances of the ionization stages of oxygen and the emission line luminosities for the EUV resonance lines of O \(\text{v} \) and O \(\text{vi} \). As with the cooling calculation, the solid lines represent equilibrium ionization balance calculations and the dashed lines represent nonequilibrium ionization balance calculations. As in the cooling calculation, most of the oxygen is in O \(\text{i} \) at the base of the atmosphere.

The differences between the equilibrium and nonequilibrium curves are the result of the inability of the oxygen ions to respond instantaneously to the changing temperature of the plasma. As the heating progresses, the temperature at each location in the atmosphere increases. However, the characteristic time for ionization of the most abundant ion present at any particular temperature to the next highest ion tends to be on the order of 100 s (e.g., Kafatos 1973). Thus as the temperature rises, the ion remains abundant for a considerable period of time past the peak temperature for the formation of the ion in equilibrium. In this same temperature region, however, the characteristic time for ionization to the most abundant ions is much shorter, on the order of 10 s. Thus as the region of the atmosphere with an initial temperature just below the peak temperature for an abundant ion is heated, that abundant ion is produced quite rapidly. The net result is that the integrated relative abundances of intermediate ions such as O \(\text{v} \) and O \(\text{vi} \) increase with time as shown in Figure 7a. The top portion of the atmosphere where most of the O \(\text{vii} \) is present does not change enough to significantly alter the integrated abundance of the ion. Almost all of the increase in line luminosity in the nonequilibrium calculation shown in Figure 7b is due to the increased relative ion abundance. A small part is due to changes in the Boltzmann factor in the collisional excitation rate due to the extra ion abundances at temperatures above the normal peak temperature of formation. After about 100 s the nonequilibrium curves begin to turn over and move toward the equilibrium ionization balance and line intensities that are characteristic of the new steady state toward which the model is moving.

It should be stressed that the changes in relative abundances and total line luminosities are due primarily to the rising temperatures at each location and not to the convection of material through a temperature gradient. Thus it is possible to have very small mass motions present and yet have a plasma that departs significantly from ionization balance. Larger mass motions will add even more complexity to the situation.

c) Rapid Fluctuations

In both the cooling and heating calculations discussed above, the changes in energy input do not change the basically thin nature of the transition region. Thus they do not reproduce the observed extended emission in transition region lines that is seen at the limb (Feldman, Doschek, and Mariska 1979). To see if higher frequency heating function fluctuations would produce larger deviations from a thin transition region structure, we have performed some smaller scale, shorter duration heating calculations.
Fig. 8.—Temperature structure used for higher frequency heating fluctuation calculations. The lower curve is the initial temperature structure and the upper curve is the result of increasing the upper boundary temperature by 20%.

Fig. 9.—Temperatures as a function of time at four heights in the atmosphere that results from varying the upper boundary temperature between the two extremes shown in Fig. 8 with a period of 3 s.

Figures 8 and 9 summarize the results of these calculations. At high frequency the energy must be deposited relatively close to the transition region or the bulk heat capacity of the overlying corona will effectively average out the effects of the fluctuations. We therefore begin with the idealized transition region shown as the lower curve in Figure 8. Note that the calculations extend only to a height of 100 km. The temperature distribution shown results from a balance between the heat flux conducted in from the top and radiative losses in the transition zone. If we instantaneously raise the upper boundary temperature by 20%, the atmosphere responds in the same manner as the larger scale heating calculations presented above, and the second temperature distribution shown in the figure is produced.

Figure 9 shows the result of varying the upper boundary temperature between the two extremes shown in Figure 8 with a period of 3 s. Here we plot the temperature variations as a function of time at four different heights in the model. All of the locations vary in temperature with the same period. Only the temperature structure at 5 km varies by a large amount, and this is due to the proximity of the chromospheric region at the base.

We have performed similar calculations with heating fluctuations with a 9 s period and find similar behavior. Thus we find that for higher frequency heating function fluctuations there is no appreciable thickening of the transition region. Rather, the thin transition region rides smoothly up and down on top of the underlying cool chromospheric material. The extent to which this result is sensitive to the simplifications we have made in modeling the chromosphere is unclear. We expect to perform more sophisticated calculations on the atmospheric response to rapid heating fluctuations in the near future.

IV. DISCUSSION AND CONCLUSIONS

First we note that the nonequilibrium ionization balance calculations for both heating and cooling of small loops in the quiet Sun predict relative abundances which differ substantially from those which would be predicted by an equilibrium calculation. In the case of cooling, the nonequilibrium abundances for the oxygen ions are smaller early in the calculation. The abundances decay, however, much more slowly than the equilibrium calculation suggests, due to relative lags in recombination rates from higher to lower ionization states. At late times in the calculation, the nonequilibrium abundances dominate relative to the equilibrium abundances. For loop heating, the nonequilibrium ionization model always shows higher relative abundances for intermediate stages of ionization. This is due to differing time constants in the various ionization rates. Of course these abundance changes are reflected in the calculated line intensities.

The validity of using a radiative loss function based on ionization equilibrium has been questioned previously for cooling plasmas (e.g., Kafatos 1973; Dere et al. 1981). Our calculations show that there are also likely to be some departures in heating plasmas. The time scales
lower temperature, higher density regions, where the temperature is changing rapidly. In the transition region tend to maximize where the density is more than a factor of 2 to 4. Further, these deviations from equilibrium are likely to be no larger than that from the finite resolution (5 km) of the steepest portions of the transition region. Ideally, we would like to perform a detailed, time-dependent radiative loss calculation as part of the solution of the basic fluid-dynamic conservation equations. However, given the uncertainties in other input physics in the model, it does not appear to be necessary at this time.

Our calculations show that in both a heating and a cooling atmosphere, the changes in relative ion abundances and total line luminosities are due primarily to the changing temperatures at each location, with mass motions providing a smaller contribution. Larger mass motions will produce a larger contribution to the departures from ionization equilibrium. They will also probably significantly alter the structure of the atmosphere. Thus we feel that calculations of the effects of ionization balance of mass flows though a predetermined temperature gradient should be used with some caution. We wish to emphasize that it is possible to have situations in which little or no mass motion would be observed and yet there would be significant departures from ionization equilibrium.

At high spatial resolution (1"), EUV emission lines formed in the transition region show multicomponent line profiles that suggest emission from unresolved structures in the network (Brueckner 1979). Generally, the emission from the redshifted component is greater than from the blueshifted component. Our calculations support the idea that the redshifted emission is due to material that is cooling and falling toward the chromosphere and that the blueshifted emission is due to material that is being heated and is convecting upward into the loop. It is clear, however, from the calculated line luminosities that if the initial state for the heating and cooling is the same, then the blueshifted emission will be brighter than the redshifted emission.

The actual situation may be much more complicated. If the structures that are being heated start at a slightly lower pressure, then the line emission would be greatly reduced. Thus one can envision a set of loops that are constantly cycling through a process in which the heating first decreases and then increases. The heating decrease would lead to a redshifted emission component and a lower pressure in the loop as material drained out. The heating increase would lead to a blueshifted emission component of reduced intensity due to the lower initial pressure.

Variations of the cross sectional area of the loop with height might also play an important role. Doschek, Feldman, and Bohlin (1976) have suggested that plasma rising into a loop of increasing cross sectional area would expand and thus have a lower density than plasma that is falling and being compressed by the changing area of the flux tube. The falling plasma would then appear brighter than the rising plasma. In both of these explanations, it is not at all clear that there must be a net downward flow of mass in the transition region as is often assumed when discussing the redshift observations.

The outflow velocities produced by increases in the heating rate for the initial model by factors of 2 and 4 are quite small, generally 5 km s\(^{-1}\) or less. Thus it appears that a rather large increase in the heating rate would be required to produce the much larger upflows observed in some cases. An alternative method for producing large upflows is to begin with a much cooler atmosphere at a reduced pressure and increase the heating rate to that necessary to produce the initial atmosphere used in our calculations. For example, if we begin with a flux tube of the same length as that used in our calculation and containing the same amount of material, but at 10,000 K, then the increase in the heating rate to that used to produce our initial model will result in outflows of up to 60 km s\(^{-1}\) at transition region temperatures. The real upflow velocities, and probably also the coronal structure, lie somewhere between these two extreme cases.

In closing, we wish to stress that nowhere in these calculations, or indeed those of the previous paper (Oran, Mariska, and Boris 1982), do we find that nonmonotonic local condensations are a dynamically fluctuating transition region such as might be responsible for the spatially extended distribution of transition region material observed in the quiet Sun. We must argue, instead, that a combination of convoluted structures on the underlying chromosphere and/or small low-lying cool loops full of transition region plasma are required to explain both the observations and calculations. The questions arising from the possibility of multidimensional instabilities and localized mass ejections into the corona at high velocity we must leave to future papers.

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