CORONAL CLOSED STRUCTURES. IV. HYDRODYNAMICAL STABILITY AND RESPONSE TO HEATING PERTURBATIONS

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ABSTRACT

We have studied the response of magnetically confined atmospheres (loops) to perturbations in (1) the temperature and density distribution, and (2) the local heating rate by means of a one-dimensional time-dependent hydrodynamical code which incorporates the full energy, momentum and mass conservation equations. These studies extend the linear instability analysis of Habbal and Rosner into the finite-amplitude regime, and generalize the confined atmosphere models of Serio et al. to the time-dependent domain.

The results show that (1) closed coronal atmospheres are stable against finite-amplitude perturbations if the chromospheric response is taken into account; and (2) observed correlated increases in coronal density and temperature can only be achieved under quiescent conditions by increasing the heat deposition rate relatively more in the chromosphere than in the corona. The relevance of these results to our current understanding of the physical mechanisms responsible for energy deposition in the solar chromosphere and corona is discussed.

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

1. INTRODUCTION

The general morphology of the solar corona is distinguished by two distinct classes of plasma structures, one confined by solar magnetic fields, and one from which steady mass loss to the interplanetary medium is allowed by the open topology of the associated magnetic fields. The confined structures from which the dominant radiative loss of the solar corona occurs are themselves characterized by continued activity, i.e., variations in plasma conditions, over a broad range of time scales (see Withbroe and Noyes 1977; Vaiana and Rosner 1978); and substantial theoretical efforts have been made to understand the physics which underlies this activity (cf. review by Priest 1980 and Withbroe 1980). In this paper, a hydrodynamic code is used to study the response of a model confined atmosphere to perturbations. Our aims are to investigate the stability of quasi-static loop structures on time scales comparable to, and somewhat longer than, the radiative cooling time scale, as well as the means by which structures which persist on such time scales may modulate the level of their radiative emission (see, for instance, Vaiana and Rosner 1978).

A number of recent theoretical studies have focused on the dynamics of the quiescent magnetically confined solar corona. Initial studies used analytical techniques to investigate quasi-static models of coronal “loop” structures (Rosner, Tucker, and Vaiana 1978; Craig et al. 1978), and more recent investigations have used numerical techniques (Vesecky, Antiochos, and Underwood 1979; Serio et al. 1981; see Priest 1980 for a review) to extend these studies, viz., to include gravitational stratification. In the course of this work, questions regarding the stability of quasi-static loop atmospheres to perturbations, and their response to changes in energy input, have arisen; and several recent studies have focused upon these issues, using both analytic and numerical techniques (for example, see Antiochos 1979; Habbal and Rosner 1979; Hood and Priest 1979). As has been emphasized and discussed by Habbal and Rosner (1979), the boundary layer separating the coronal portion of a loop structure from the photosphere is crucial to understanding loop dynamics; it is this boundary layer which determines the loop plasma pressure under quiescent conditions and (because it is the loop plasma pressure which largely fixes the coronal loop radiative loss; cf. Rosner, Tucker, and Vaiana 1978) establishes the quiescent level of radiative emission. For this reason, study of loop dynamics cannot be reasonably separated from the study of this boundary layer; and in the following analysis the boundary layer-loop atmosphere interaction will play a central role.

Our paper is organized as follows. In § II we establish the conditions for stability of solutions to the static equations of motion for loop atmospheres (which define our initial conditions). Section III is devoted to the

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response of the atmosphere to changes in heating rate; a summary and our conclusions are given in § IV. A brief discussion of the assumed chromospheric radiative losses and heating rate is given in an appendix.

II. STABILITY OF STEADY STATE MODELS

The starting point of any analysis of loop atmosphere evolution is an initial loop atmosphere. In the work described here, this initial atmosphere is always a solution to the hydrostatic equations of motion obtained by setting ∂/∂t = v = 0 in the single-fluid, one-dimensional equations of motion for a toroidal geometry with constant cross section.

\[ \frac{dn}{dt} = -n \frac{\partial}{\partial s}(\nu) \]

\[ n m_H \frac{dv}{dt} = -\frac{\partial}{\partial s}(\nu n m_H + \frac{\partial}{\partial s}(\mu \frac{\partial v}{\partial s})) \]

\[ \frac{d\xi}{dt} + w \frac{\partial v}{\partial s} = E_H - n^2 \beta P(T) + \mu \left( \frac{\partial v}{\partial s} \right)^2 - \frac{\partial}{\partial s}(\kappa \frac{\partial T}{\partial s}) \]

\[ p = (1 + \beta) nk_B T \]

\[ \xi = \frac{1}{2} (1 + \beta) nk_B T + n \beta \chi \]

\[ w = \frac{1}{2} (1 + \beta) nk_B T + n \beta \chi. \tag{2.1} \]

where \( n \) is the hydrogen number density, \( \nu \) the plasma bulk speed, \( m_H \) the hydrogen atom mass, \( g \) the local gravitational acceleration, \( s \) the field line coordinate \((s_0 = \) location of temperature minimum = \(4.5 \times 10^4 \) cm, \( s_{\text{max}} = \) location of temperature maximum = \(2 \times 10^4 \) cm), \( \mu \) the effective plasma viscosity, \( \beta \) the fractional ionization \( n_e/n_H \) (viz., Brown 1973), \( k_B \) the Boltzmann constant, \( P(T) \) the radiative loss function (identical to that of Rosner, Tucker, and Vaiana 1978 for \( T < 10^4 \), and for \( T < 10^4 \) derived from the Vernazza, Avrett, and Loeser 1980 atmosphere model as shown in the Appendix), \( T \) the plasma temperature, \( \kappa \) the thermal conduction coefficient (Spitzer 1962), \( \chi \) the hydrogen ionization potential, and \( E_H \) the local heating function. In light of our ignorance regarding the local heating rate, the heating function \( E_H \) is assumed to be uniform in the corona (without any attempt at modeling its functional dependence on plasma density and temperature), and is derived from the Vernazza, Avrett, and Loeser (1980) models for the chromosphere, as discussed in more detail in the Appendix.

Static solutions (obtained by setting ∂/∂t = v = 0) have been previously obtained in Serio et al. (1981) for loops whose base temperature is \(2 \times 10^4\) K; the initial model atmosphere considered here uses the Vernazza, Avrett, and Loeser (1980) models to extend the Serio et al. (1981) static calculations to chromospheric temperatures. Because we shall be studying perturbations which involve only moderate changes in the local heating rates (in contrast to flare studies), the appropriate boundary conditions at \(s_0\) are \(n(s_0, t)\) and \((T(s_0, t))\) constant. These boundary conditions apply because the atmosphere perturbations introduced in the present study are too weak to influence plasma conditions at the assumed lower boundary. We further assume that \(s = s_{\text{max}}\) (the loop apex) is a point of symmetry so that \(v(s_{\text{max}}, t) = 0\) and \(\nabla T = \nabla n = 0\) (e.g., only symmetric solutions are allowed); we therefore eliminate the possibility of steady, persistent flows.

A crucial point in the study of loop and chromosphere dynamics is the appropriate analysis of the transition region, whose steep gradients place severe restrictions on the grid spacing and the stability of the algorithm. For example, a uniform spatial grid whose spacing is fine enough to resolve the transition region well gives unnecessary spatial resolution in the corona and requires prohibitive amounts of computation time. These problems have been generally recognized. For example, Kral and Antiochos (1980) have bypassed such problems by not including the chromosphere in their model and instead simulating its possible contributions to the atmosphere dynamics by means of ad hoc boundary conditions. In contrast, Craig and McClymont (1981) have modeled the transition region and chromosphere, as well as the corona, but with low accuracy in the transition region. Nagai (1980) has modeled flares in loop atmospheres and included the transition region response, using coarse spacing in the corona and transition region and fine spacing in the chromosphere; the transition region of the initial static configuration thus seems to be treated with low accuracy. Kostyuk and Pikel’ner (1975) and Kostyuk (1976) do not report details of their computational grid.

Our solution to the steep gradient problem is to perform a change of variable \( \xi = f(s) \), where \( s \) is the field line coordinate; \( f(s) \) can be chosen so that the resulting uniform grid in \( \xi \) satisfies \( |T_0/(dT_0/d\xi)| \geq \Delta \xi \), where \( T_0(\xi) = T(\xi, t = 0) \) is the initial temperature profile. In particular, we have chosen \( \xi = \log(s/s_0) \).

This method resolves the resolution problem as long as the spatial location of the steep gradient region does not depart from the high-resolution portion of the computational grid; in the present case, the applied perturbations are in fact sufficiently weak that this restriction is satisfied. In fact, because solutions with increasingly...
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a) Relaxation of Initial Atmosphere

As a first step, we have analyzed the relaxation of an initial configuration obtained from the hydrostatic equations. The initial atmosphere has semilength $2 \times 10^9$ cm, a coronal pressure of $6 \text{ dyn cm}^{-2}$, and it is uniformly heated in the corona at $1.3 \times 10^{-2}$ ergs s$^{-1}$ cm$^{-3}$. The solution to the initial static atmosphere is obtained via fourth-order Runge-Kutta integration of the static equilibrium equation (cf. Serio et al. 1981). The spatial grid spacing for the static calculations is fixed so that temperature variations between two consecutive space steps obey the constraint $\Delta T/T \leq 0.01$. Because the spatial grids for the static and time-dependent calculations differ substantially, it is necessary to interpolate the static solution in order to find the initial values of temperature and density on the (Eulerian) spatial grid of the time-dependent code. This initial interpolated atmosphere is not necessarily a static solution of the time-dependent finite-difference equations on the $\xi$ grid. It is therefore necessary to (1) relax the initial atmospheres on the $\xi$ grid and (2) demonstrate that the time-dependent code can successfully evolve the unperturbed static atmosphere.

As a first step, we therefore evolve the initial atmosphere given in Figure 1, whose characteristic plasma parameters are appropriate for a solar active region loop (cf. Rosner, Tucker, and Vaiana 1978). With the exception of the interpolation described above, the initial atmosphere has not been perturbed. We have followed the temporal evolution for several times $10^3$ s, and we find that the initial static solution indeed relaxes to a stationary solution of the time-dependent code on this time scale. Small amplitude disturbances, approximately periodic (period $\leq 250$ s), are observed in the corona; these appear to be generated near the base of the model transition region, and they damp as they propagate upward into the corona. We attribute the generation of these disturbances to the interpolation procedure, i.e., the initially interpolated atmosphere does not in fact satisfy the hydrostatic equations of motion on the transformed spatial grid, the errors introduced by the interpolation being most significant in the transition region. The observed disturbances thus reflect the relaxation of the atmosphere toward a static state and are generated by the readjustments in the transition region. The period of the observed motions is of the order of the sound travel time in the loop; and we interpret these disturbances as long wavelength coronal acoustic modes (see Craig and McClymont 1981 for a similar phenomenon). The relaxed atmosphere is used as the initial condition for all of our further studies.

b) Isobaric Perturbation

We next consider the effect of mild perturbations in temperature and density (but not in heating rate) of the relaxed atmosphere. The perturbations are chosen to be isobaric, to vanish at the base of the loop atmosphere, and to have initially constant fractional amplitude (i.e., $|\Delta T/T| = \text{constant}$) along the field line coordinate. The aim of this study is to see whether, in the absence of possible stabilization by the actual plasma heating process, finite amplitude fluctuations can destroy the static

Recall that in the present investigation the coronal heating function $E_H$ is specified a priori as a function of position only and hence has neither an explicit nor an implicit dependence on plasma conditions (cf. Habbal and Rosner 1979 for discussion of the case in which $E_H$ does respond to perturbations).
Fig. 2.—Evolution of the relaxed stationary solution after a short wavelength perturbation is applied (eq. [2.2] with \( m = 100, A = 0.1 \)). Temperature (a) and density (b) are plotted vs. field line coordinate (horizontal axis) and time (oblique axis). The results shown are sampled every 200 s. The perturbation does not appear sinusoidal in the corona because of the beating between the perturbation wavelength and the plotted logarithmic spatial grid.

initial state. This case, as well as the previous relaxation of the unperturbed atmosphere, can also be viewed as tests of code stability because the coronal portion (i.e., plasma at \( T \geq 10^6 \) K) of the loop atmosphere lies on the stable branch of the stability diagram given by Hood and Priest (1980).

Figures 2 and 3 show the effect of an isobaric perturbation on the relaxed static atmosphere. The analytical form of the perturbation chosen is

\[
\delta T(s) = A T_0(s) \sin(2\pi m s / s_{\text{max}})
\]

with

\[
n(s) = \frac{p(s)}{[(1 + \beta) k (T_0 + \delta T)]}.
\]
We have explored a wide range of wave numbers for the perturbations. We show two representative cases, in which $m = 100$ (Fig. 2) and $m = 1/4$ (Fig. 3), with $A = 10^{-1}$ in both cases. We see that the initial configuration is stable, i.e., in no case do we see a secular disruption of the perturbed configuration. Instead, there is a relaxation towards the initial conditions of temperature and density. As the atmosphere returns toward the steady configuration, it undergoes a series of fluctuations which manifest themselves as trains of sound waves propagating upward through the transition region and the low corona. These modes appear similar to those seen during the relaxation of the static loop (see § 1a above).

These results confirm the instability analysis results of Habbal and Rosner (1979) carried out in the linear growth (infinitesimal amplitude) regime and demonstrate that the reaction of the lower "boundary" (usually taken at a temperature $T \geq 2 \times 10^4$ K in previous instability analyses) is essential to determining actual stability or instability. This effect is crucial in reconciling our result with the analyses of Antiochos (1979) and Hood and Priest (1980): these authors chose to fix the loop...
temperature at a fixed coordinate located at the base of the transition region (Antiochos 1979) or at the top of the transition region (Hood and Priest 1980). The constraints

$$T_{\text{base}} = \text{constant}$$

$$s(T_{\text{base}}) = \text{constant}$$

(2.3)

are thus applied within portions of the loop atmosphere which our calculations show to be highly dynamic, i.e., the location of the temperatures assumed as base temperatures by these authors is highly variable, and furthermore it varies in a systematic way as loops evolve to higher or lower pressure states under the influence of heat deposition fluctuations (see § III below). We have avoided this problem by placing our loop base sufficiently deep in the atmosphere that fluctuations of the amplitude studied here have negligible effects on the boundary. Furthermore, if we instead apply constraints (2.3) at the base of the transition region (or higher), our code consistently shows instability for initial hydrostatic atmospheres with vanishing thermal conductive flux at the base, in accordance with the results of Antiochos (1979) and Hood and Priest (1980).

III. ATOMSPHERIC RESPONSE TO WEAK CHANGES IN LOOP HEATING RATE

We now turn to a study of the effects of changing the local heating rate within a loop atmosphere. We consider three distinct cases of temporally variable heat input: (i) variability restricted to the corona alone ($\Delta E_H/E_H = \text{constant}$ for $T > 10^6$ K, with $\Delta E_H = 0$ elsewhere); (ii) variability restricted to the upper chromosphere alone (i.e., $\Delta E_H/E_H = 0$ everywhere except for the temperature range between $7 \times 10^3$ K to $2 \times 10^4$ K); and (iii) variability throughout the loop (i.e., both in the chromospheric region and in the corona). In all cases, we begin our calculations with the relaxed static atmosphere discussed in § IIa above.

Case (i) ($\Delta E_H = E_H$ for $T > 10^6$ K and zero elsewhere): The results of the first type of heating variability are shown in Figure 4, where the evolution of coronal temperature, pressure, and density is plotted as a function of time. We assume an initial doubling of heat deposition in the coronal portion of the loop. It is evident that while the loop’s temperature increases on a time scale of a few hundred seconds, the density is not subject to any substantial change (although, on the basis of scaling laws for static loops (cf. Rosner, Tucker, and Vaiana 1978), one would expect a considerable increase in coronal density, i.e. “evaporation” from the base). Why is the $(T, p, L)$ scaling violated? The scaling laws define the possible loci of allowed values of $(T, p)$ or $(E_{\text{H, } \text{cromonal}}, p)$ pairs for a given fixed loop length under the assumption of hydrostatic conditions. In the present case, we have imposed a doubling of the coronal heating rate but kept the loop length and pressure constant (i.e., because the coronal heating perturbation is too small to alter the chromospheric energy balance via thermal conduction, the pressure—which is determined at the chromospheric level—will not change). We therefore expect that the perturbed loop does not correspond to a hydrostatic configuration which satisfies the scaling laws. In fact, our calculations show that the final configuration, although stable and stationary, is not static. In particular, we find persistent fluctuations in plasma conditions in the lower transition region; both $n(T_{\text{tr base}})$ and $s(T_{\text{tr base}})$ fluctuate, with $\delta n/n_0 \sim 0.9$ at the TR base (which lies above the loop base). Estimating the local time-averaged radiative losses $\langle E_R \rangle$ at the TR base leads to the result that $\langle E_R \rangle = \langle n^2 P(T) \rangle = (1 + [\delta n]^2/2n_0^2)n_0^2 P(T) \approx 1.4 E_R$ (unperturbed); thus the increased thermal conductive flux towards the chromosphere may be dissipated at the base of the TR by the increased effective radiative losses even in the absence of a pressure increase. As an aside, we note that the TR base fluctuations do propagate into the corona, but damp rapidly and consequently are of far lower amplitude in the corona; coronal conditions thus remain close to the hydrostatic regime.

Fig. 4.—Response of the relaxed atmosphere to a loop heating change of the form $\Delta E_H = E_H$ only for plasma with temperature above $10^6$ K. Density (a), pressure (b), and temperature (c), all at $10^9$ cm from loop base.

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A further interesting feature of the present calculation is that the physical parameters of the perturbed configuration are well within the scatter of observed values (e.g., within ~ a factor of two) about the scaling law prediction of Rosner, Tucker, and Vaiana (1978). The above analysis suggests that some of this scatter might well be due to perturbations of the kind described here. Put in its simplest terms, the effective TR pressure, as deduced from the radiative losses, is \( \langle p^2 \rangle^{1/2} \), which obeys the simple triangle inequality \( \langle p^2 \rangle^{1/2} \geq \langle p \rangle \); a loop whose TR pressure fluctuates about some mean value simulates—as far as its scaling behavior is concerned—a geometrically identical loop at higher mean pressure.

Case (ii) \( \Delta E_H = E_H \) for \( 7 \times 10^3 < T < 2 \times 10^4 \) K, \( \Delta E_H = 0 \) elsewhere: When doubling the heating rate only in the region between \( 7 \times 10^3 \) K and \( 2 \times 10^4 \) K (Fig. 5), we see that the coronal temperature and density remain essentially unchanged. This result suggests that the present heating perturbation, even if applied in the region in which the transition region base pressure is determined, is not large enough to significantly perturb the coronal portion of the loop. A substantially larger increase in chromospheric heating should increase the coronal pressure; this expectation is indeed verified by our code, which gives a twofold increase in coronal base pressure for the heating perturbation of Case (iii) below, but restricted to \( T < 2 \times 10^4 \) K.

Case (iii) \( \Delta E_H(T) = (6 \times 10^4 / T) E_H \) for \( 7 \times 10^3 < T < 2 \times 10^4 \) K, \( \Delta E_H = E_H \) for \( T > 2 \times 10^4 \) K): The case at hand (Fig. 6) shows that this form of the perturbation is effective in modifying both the temperature and the density in the coronal portion of the loop. In particular, the pressure at the transition region base is increased, leading to "evaporation" of plasma into the coronal part of the loop. The increased coronal heating in turn causes the coronal plasma temperature increase. The coronal part of the loop now reaches conditions close to those expected from the scaling laws of Rosner, Tucker, and Vaiana (1978).

IV. SUMMARY AND DISCUSSION

The aim of this paper has been to gain a better understanding of the dynamics underlying brightness fluctuations of the quiescent confined corona. A time-dependent, one-dimensional fluid code has been used to numerically model confined loop atmospheres, and has been tested against previous analytic and numerical time-independent calculations of stationary atmosphere solutions. Results show desired numerical stability, even
in the steep gradient transition region. A further series of model calculations show that (1) finite amplitude isobaric perturbations spanning a wide range of wavelengths do not drive loop instability; and (2) large-amplitude correlated changes in coronal density and temperature require change in both coronal and chromospheric heating rate. Our numerical instability calculations confirm the conclusions of Habbal and Rosner (1979), and reconcile their results with those of Antiochos (1979) and Hood and Priest (1980). We find that the constraints \( T_{\text{base}} = \text{constant} \), \( s(T_{\text{base}}) = \text{constant} \) do not lead to instability if applied deep in the atmosphere; if applied in the dynamic transition region or corona, instability results, as first shown by Antiochos (1979) and Hood and Priest (1980). The latter case, however, does not appear to be physically well motivated (although it is computationally convenient), because it is hard to see what physical process would ensure that \( s(T_{\text{base}}) \) remains fixed following a perturbation if \( T_{\text{base}} \) lies at transition region or higher temperatures.

The theoretical problem of accounting for correlated coronal density and temperature fluctuations has been generally addressed in the flare context (cf. Krall and Antiochos 1980 and references therein). In that case, Krall and Antiochos (1980) have shown that “chromospheric evaporation” (e.g., increase of coronal density due to injection of heated chromospheric matter) can result by imposing an ad hoc footpoint boundary condition which essentially ensures that excess thermal energy conducted to the footpoints is returned to the loop via upward-flowing matter, rather than being radiated away at the footpoints. Unfortunately, it is difficult to go any further; because of the extreme gradients encountered, it has not been possible as yet to unambiguously model this evaporation process in the chromosphere and transition region in any detail, and so to demonstrate its existence apart from the imposition of appropriate boundary conditions (see, however, Craig and McCartney 1981 for recent progress in this area). A very similar theoretical problem exists for density and temperature variations under the far less extreme conditions due to fluctuations in the solar chromosphere and transition region in any detail, and so to demonstrate its existence apart from the imposition of appropriate boundary conditions (see, however, Craig and McCartney 1981 for recent progress in this area). A very similar theoretical problem exists for density and temperature variations under the far less extreme conditions encountered in relatively quiescent active region loops (cf. Davis et al. 1975); our present calculations resolve the problem in this more restrictive regime as follows.

Correlated coronal density and temperature variations in a coronal loop imply a change in coronal loop pressure. As discussed in some detail by Habbal and Rosner (1979), the coronal pressure is essentially attained in the region within the loop corresponding to \( \sim 7500 \) K, e.g. at the location of one of several postulated temperature plateaus in the solar atmosphere (cf. Basri et al. 1979; Avrett 1980); any process which results in changing the pressure of the coronal portion of the loop must therefore affect the pressure at this low altitude. For example, varying the local chromospheric heating rate directly impacts the atmospheric structure at the coronal pressure-determining plateau; and sufficiently increased chromospheric heating yields increased chromospheric pressure, and consequently increased coronal pressure. In contrast, a moderate (\( \sim \) a factor of 2) increase in the coronal heating rate alone appears to drive large-amplitude density fluctuation at the transition region base, which effectively dissipates the additional heat flux conducted downwards from the corona. Such coronal heating rate changes thus cannot affect energy balance in the low chromosphere, and therefore do not result in coronal pressure changes. In this case the perturbed loop does not satisfy the assumptions underlying the Rosner, Tucker, and Vaiana (1978) scaling laws (viz., the loop is not hydrostatic); nevertheless, the mean loop parameters (i.e., pressure, temperature, and length) are such that if the scaling laws are applied, the result lies within the scatter of observed values. We conjecture that, in fact, some of the observed scatter may be a consequence of departures from hydrostatic conditions due to fluctuations in the coronal heating rate, as exemplified by our Case (ii) calculations.

Our results show that correlated increases in coronal density and temperature in which the density change predominates (demanded by X-ray observations; cf. Davis et al. 1975) occur only if both chromospheric and coronal heating rates are changed and suggest, as argued by Habbal and Rosner (1979), that the chromospheric and coronal heating processes must be strongly coupled in solar active regions. Whether this coupling exists because of a common underlying heating mechanism, or because changes in plasma conditions caused by chromospheric heating provide positive feedback to the coronal heating process, remains un resolved.

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APPENDIX

CHROMOSPHERIC RADIATION LOSSES AND STEADY-STATE ENERGY INPUT

Using the Vernazza, Avrett, and Loeser (1980) models for flare, plage, “very bright network element,” and “dark point within a cell” chromospheres, we calculate the quantities \( P_i(T) = E_i(T)/n_{H_{\alpha}} \), where \( E_i(T) \) is the radiative loss...
Fig. 7.—Plasma emissivity function $P(T)$ ergs cm$^{-3}$ s$^{-1}$ used in our model calculations.

(ergs s$^{-1}$ cm$^{-3}$) for the $i$th model, $n_H$ is the total hydrogen number density, and $n_e$ the electron number density. We have adopted an analytic expression which provides a reasonable fit to the $P_i(T)$, similar to that used by Hartmann and MacGregor (1980), of the form

$$P(T) = \begin{cases} (1.0606 \times 10^{-6} T)^{1.7} & 4.44 \times 10^3 < T < 8 \times 10^3 \\ (1.397 \times 10^{-8} T)^{6.15} & 8 \times 10^3 < T < 2 \times 10^4. \end{cases}$$

This approach allows a rather simple way of modeling the energy losses from the chromosphere on time scales which are long when compared with the radiative relaxation times, adequately accounting for the chromospheric energetics on such time scales. For higher values of temperatures we have used the piece-wise power law emissivity function given by Rosner, Tucker, and Vaiana (1978). The resulting composite radiative loss function is shown in Figure 7.

The steady state chromospheric heat input rate is calculated as follows. We choose an appropriate chromospheric model from Vemazza, Avrett, and Loeser (1980) which matches the pressure of the loop transition region at $\sim 2 \times 10^4$ K. The Vemazza, Avrett, and Loeser (1980) model allows us to calculate the local heat flux divergence and effective radiative loss and, by assuming local energy balance, the necessary local heating rate as a function of temperature $E_H(T)$.

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