THE ROLE OF NONCLASSICAL ELECTRON TRANSPORT IN THE LOWER SOLAR TRANSITION REGION

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ABSTRACT

Using a linear BGK method to study electron transport in one-dimensional, planar models of the solar transition region, we investigate the dependence of nonclassical electron transport effects on the assumed transition region temperature structure, with particular focus on how these nonclassical effects might influence line emission from the lower transition region. We find that nonclassical transport effects do not significantly alter the ionization or excitation balance of ions whose emission lines form predominantly in the lower transition region. However, the nonclassical tail heat flux that arises from the upper transition region can be of sufficient magnitude and can dissipate over a broad enough region to balance the observed energy loss in these lower transition region lines. The computed magnitude of the nonclassical tail heat flux in the lower transition region depends sensitively on the assumed steepness of the temperature gradient in the upper transition region, and so quantitative determination of the importance of such a nonclassical modification in the electron heat transport will require self-consistent modeling of the transition region temperature structure.

Subject headings: line formation — particle acceleration — Sun: atmosphere — Sun: chromosphere

I. INTRODUCTION

A problem of continuing interest in solar physics has been to understand the observed brightness of the quiet solar atmosphere in spectral lines that are formed in the lower solar transition region (i.e., at temperatures between \(10^4\) and \(10^5\) K). Unlike lines formed in the upper transition region (i.e., between \(10^5\) and \(10^6\) K), these lines are brighter than can easily be explained by models in which radiative energy loss is balanced by backheating from the corona, whether by classical conduction (Athay 1966; Moore and Fung 1972), or by the enthalpy flux associated with a downflow (Athay 1981, 1982). Observations indicate a combined quiet-Sun emission-line flux of several times \(10^2\) ergs cm\(^{-2}\) s\(^{-1}\) (e.g., Dupree 1972), of which about half is emitted in the Ly\(\alpha\) line of neutral hydrogen (Basri et al. 1979). Based on the calculated line emissivities (e.g., Gabriel 1976; Chambe 1978; Vernazza, Avrett, and Loeser 1981, hereafter VAL), these observed fluxes imply a temperature scale height in the lower transition region of \(~10\text{--}100\) km; with such a scale height, the classical conductive heat flux (which is proportional to the local temperature gradient) is several orders of magnitude too small to balance the observed line energy losses.

This difficulty has led to the widespread view that the radiative energy loss in the lower transition region must be balanced by an extension of the in situ heating mechanism for the underlying chromosphere, rather than by any form of back-heating from the overlying corona. For example, Withbroe (1981) has shown that, given the freedom to specify the in situ heating function arbitrarily, static one-dimensional transition region models can readily reproduce the observed line emission. However, as shown by VAL (see their Figs. 48 and 49), in the case of Ly\(\alpha\) this requires a heating mechanism with the seemingly unnatural property that it is about an order of magnitude larger in the Ly\(\alpha\)-forming region than in the contiguous upper chromosphere. Rabin and Moore (1984) have proposed Ohmic dissipation of filamentary currents as a specific heating mechanism which can balance the observed line emission, but this requires a rather detailed description of the inhomogeneous structure of the atmosphere. In light of the uncertainties and complexities in modeling such inhomogeneous structures, it is of interest to know if the observed line emission could be understood on more fundamental physical grounds with simpler one-dimensional models. In this paper we therefore examine the alternative possibility that the observed line emission results from nonclassical electron transport effects that are associated with the inherently steep temperature gradients in the solar transition region.

As shown by Roussel-Dupre (1980a, b) and Shoub (1982, 1983), the electron velocity distribution function in the solar chromosphere-corona transition region should form an enhanced "tail" from high-velocity electrons that propagate, nearly without collisions, from the corona toward the chromosphere. These authors note that this tail can have the dual effect of (1) modifying electron-ion collisional excitation/ionization rates and (2) modifying the nature of electron heat transport. Both these effects offer potential solutions to the problem of understanding the observed line brightness. First, the excitation and ionization of ions by collisions with tail electrons could increase the line emission, or extend it over a broader range in temperature; the observed brightness of the lines would then no longer necessarily require a small temperature gradient, and the classical conductive heat flux might then be large enough to balance the observed radiative energy loss. Second, the heat flux associated with the enhanced tail may itself be large enough to balance the line radiative energy loss. Unlike the classical conductive heat flux, this tail heat flux is not simply proportional to the temperature gradient; it could therefore balance the radiative energy loss from a region where the temperature scale height is large.

The most sophisticated of these calculations of the electron velocity distribution in the transition region is that by Shoub (1983), who computed detailed numerical solutions of the full Fokker-Planck equation for electron transport with explicit inclusion of a variety of physical effects (e.g., inelastic collisions,
polarization electric field, etc.). Nonetheless, despite its extensive nature, the results of this calculation cannot easily be used to estimate the importance of nonclassical transport effects on the line emission. The difficulty lies with the assumed models, which assume constant classical heat flux throughout and thus have an unrealistically steep temperature gradient in the lower transition region, where Lyα and other lines of interest are formed. Although Shoub claimed that the computed distribution functions in the lower part of his model are insensitive to assumptions about the temperature structure there, we find otherwise (see § III). In fact, assuming a temperature structure that is more consistent with that inferred empirically, we find that the effects of nonclassical transport in the lower transition region are substantially reduced, although not entirely eliminated.

In order to examine this question of how the assumed temperature structure affects the electron transport, we parameterize in § II the height variation of the electron temperature to enable correspondence with a variety of one-dimensional, constant pressure transition region models, including empirical models, theoretical models, as well as the constant classical heat flux model used by Shoub (1983). In § III we then calculate for these models the electron velocity distribution function using a linearized form of the so-called BGK (after Bhatnagar, Gross, and Krook 1954) equation for electron transport. As shown by Shoub (1982), this phenomenological model for electron transport yields distribution functions that are in good agreement with results of his full Fokker-Planck treatment (see Shoub 1983), and so it provides an appropriate basis for this assessment of the potential importance of nonclassical transport effects. As in these earlier calculations, we use a linearized approach of solving the electron transport in an assumed transition region model with fixed variation of temperature with height. Thus we do not address the much more difficult problem of solving for the atmospheric temperature structure when the nonclassical effects are self-consistently taken into account; our goal is rather to estimate from a linear calculation the nature and magnitude of these effects in order to establish a sound basis and justification for future work on this more difficult nonlinear problem. Accordingly, we use our linear calculations of the electron distribution function to estimate the importance on nonclassical modifications in, respectively, the collisional ionization and excitation balance (§ IV) and the electron heat flux (§ V) in the lower transition region. Finally (§ VI), we summarize our results and outline directions for future work.

II. TRANSITION REGION MODELS

To examine the importance of nonclassical electron transport effects in the solar transition region we must first specify an assumed transition region model. In this paper, we assume a fully ionized, isobaric, one-dimensional, planar model in which the variation of electron temperature $T$ with height $z$ is parameterized as described below. In the previous transport calculations by Shoub (1982, 1983), this height dependence of the temperature was very close to that which is obtained by assuming a constant value of the quantity $F_0 \equiv T^{5/2}dT/dz$, which (apart from a factor which varies logarithmically with temperature and density) is proportional to the classical electron conductive heat flux (Spitzer 1962). In the upper transition region, i.e., for $T > 200,000$ K, such a height variation of the temperature agrees quite well with that inferred empirically from observations of ultraviolet spectral lines, but in the lower transition region it is much steeper than is inferred empirically. It is also much steeper than is implied by more physically realizable theoretical models.

In Figure 1a we compare the variation of $F_0$ versus $T$ for several transition region models. The dashed curve represents

![Graph showing temperature dependence of the quantity $T^{5/2}dT/dz$ in various theoretical (dashed curve) and empirical models (solid curves), and in models assumed in the previous transport calculations by Shoub (1982, 1983) (dot-dashed curves).](image)
results from a typical static theoretical model, in which the difference between radiative energy loss and local (in situ) heating in the transition region is balanced by the divergence of the classical electron conductive heat flux from the corona; in this model, both the classical conductive flux and its divergence are assumed to approach zero at a chromospheric boundary temperature $T_e = 8000$ K, the radiative cooling is according to the optically thin cooling law of McWhirter, Thonemam, and Wilson (1975), the local volume heating rate is proportional to the electron density $n$, and the pressure is such that $nT = 6 \times 10^{14}$ K cm$^{-3}$ is constant. The solid curves in Figure 1a represent empirical results derived from plots by Gabriel (1975) and Chambe (1982) of the mean quiet-Sun differential emission measure, $\Phi = n^2 / (d \ln T / dz) = (nT)^{2/3} / F_0$, also assuming a typical transition region pressure $nT = 6 \times 10^{14}$ K cm$^{-3}$. The difference between these empirical results, which were derived using the same set of EUV spectroscopic data but assuming different physical parameters (i.e., abundances), illustrates the uncertainty in our empirical knowledge of the transition region temperature structure. As summarized by Athay (1976; pp. 314–328), EUV data imply values of $F_0$ in the upper transition region ranging from $F_0 = 2.3 \times 10^{11}$ (Elwert and Raju 1972) to $F_0 = 8.8 \times 10^{11}$ (Dupree 1972), even when the same abundances are assumed. Athay (1976) thus adopts an intermediate value $F_0 = 5 \times 10^{11}$ which, for the upper transition region, corresponds closely to the values used by Shoub (1982, 1983) (see dot-dashed curves). In the lower transition region, however, both the empirical and theoretical results suggest a much less steep temperature structure than that implied by constant $F_0$.

In order to examine the effect of various transition region temperature structures on results for the electron transport, we assume models in which the temperature is given by the parameterization,

$$T^{5/2} dT / dz = \frac{F_0 [1 + (T_e / T)^b]}{[1 + (T_e / T)^b]} , \quad (1)$$

where $b > 0$ is a free parameter and $T_e$ is a “bridging” temperature that connects the different character of the temperature variation of the upper and lower parts of the model. For $T \gg T_e$, we then have $F_0 \approx F_{\text{ou}} = \text{constant}$, while for $T \ll T_e$ we have $F_0 \approx T^b$. The parameterization is normalized so that $F_0 = F_{\text{ou}}$ exactly at the upper boundary temperature $T_e$ of the model. In all calculations here, we choose the fixed values $T_e = 2 \times 10^6$ K and $T_0 = 2 \times 10^5$ K. For the upper transition region steepness and the transition region pressure, we choose the fiducial values $T_{\text{ou}} = 6 \times 10^{11}$ K$^{1/2}$ cm$^{-3}$ and $nT = 6 \times 10^{14}$ K cm$^{-3}$, about which we investigate the effects of variations by a factor of 2. The parameterization (1) is assumed to apply between chromospheric and coronal boundary temperatures of $8 \times 10^3$ and $2 \times 10^6$ K, beyond which the temperature is constant at these boundary values. Finally, comparison of Figures 1a and 1b shows that choices of $b = 0$, 2, and 3.5 enables us to simulate the temperature variations of, respectively, the constant classical heat flux model, the theoretical model, and the empirical model discussed above. For ease in later reference, we denote these models A, B, and C.

III. THE ELECTRON VELOCITY DISTRIBUTION FUNCTION

We now want to investigate the nature and significance of nonclassical transport phenomena that arise in the transition region models developed in the previous section. To do this we use a linearized BGK approach; i.e., we use the approximate collision operator of Bhatnagar, Gross, and Krook (1954), but we do not attempt the nonlinear problem of developing a self-consistent model transition region that conserves mass, momentum, and energy in the BGK approximation. Except for the important differences in assumed models (§ II), this approach follows closely that of Shoub (1982). We assume an isobaric, fully ionized, pure hydrogen plasma in which protons are idealized as infinitely massive and at rest. We also assume that any magnetic field that may be present is uniform and serves only as a passive guide for transport along a single longitudinal dimension (a magnetic "loop"). Ignoring further the effect of any external forces, such as gravity or a polarization electric field, the BGK equation is

$$\mu \nu [\partial f / \partial z] = v(n, z) f(v, n, z) - f(\mu, n, z) , \quad (2)$$

where $f$ is the electron velocity distribution function we wish to determine, $*^*$ is the Maxwellian source distribution specified by the electron density and temperature (as fixed in this linearized calculation by the parameterization described in § II), $z = \text{distance coordinate, } v = \text{electron velocity, and } \mu = v / v_\mu$.

For the electron collisional thermalization rate $\nu$ we use

$$\nu(n, z) = \frac{16 \pi e^2}{\Lambda \Lambda N (m^2 v^3)} , \quad (3)$$

where $e$ and $m$ are the electron charge and mass, and $\Lambda$ is the Coulomb logarithm. With $z = 1$, $v$ is the rate at which high-velocity test electrons are deflected by $90^\circ$ in a pure electron-proton plasma (see, e.g., Krall and Trivelpiece 1973). Strictly speaking, use of this form as a collisional thermalization rate is only appropriate in the high-velocity limit, $v > v_\mu \equiv \sqrt{(2kT/m)}$, but, since electrons with $v < v_\mu$ are in any case collisionally thermalized in the transition region, the high-velocity limit may formally be used for all velocities without loss of accuracy in the end results.

The need to include the free parameter, $\alpha$, in the definition of the collisional thermalization rate reflects the fact that the BGK equation provides only a phenomenological description of electron transport, based on our intuition that collisions should tend to thermalize electrons into a Maxwellian distribution. In the BGK picture, this thermalization occurs suddenly, so that, for example, a high-velocity test electron is transformed directly into a thermal one, without ever existing in any intermediate state. In the more realistic Fokker-Planck model (see, e.g., Rosenbluth, MacDonald, and Judd 1957), the thermalization is more accurately described as velocity diffusion, so that such a high-velocity electron is really thermalized gradually by the cumulative effect of many small-angle collisions, during which its mean speed steadily decreases (and its velocity dispersion increases). In spite of these differences, Shoub (1982) has shown that electron distribution functions derived from the BGK method can be qualitatively quite similar to those derived from the Fokker-Planck method, with the best quantitative agreement obtained by setting $\alpha = 2.2$. In the spirit of examining variations in uncertain model characteristics like the pressure, $nT$, or the upper transition region temperature steepness, $\nu_{\text{ou}}$ (see § II), we will consider $\alpha$ to be a free parameter and examine the effects of a factor of 2 variation about a fiducial value of $\alpha = 1$. Actually, it is not necessary to examine separately variations in each of these free parameters since, as can easily be shown, the BGK equation can be scaled in such a way that solutions depend only on the parameter combination $F_{\text{ou}} / 2nT$.

The motivation for using the BGK rather the Fokker-Planck equation is the relative ease with which solutions can
be obtained and understood. For example, for models with a constant classical heat flux (i.e., with \( b = 0 \)), Shoub (1982) has obtained approximate analytic solutions of the BGK equation (2) and used these to illustrate the physical nature of the electron transport. Similar analytic approximations can also be derived for models with the more general temperature parameterization of equation (1), but they are somewhat more complicated and hence less illuminating. In this study, we have chosen instead to obtain the electron velocity distribution function \( f(\mu, v, z) \) by straightforward numerical integration of equation (2). As a boundary condition we require that at the lower and upper boundaries, where the temperatures are \( 8 \times 10^3 \) and \( 2 \times 10^6 \) K, the incoming electrons with, respectively, \( \mu > 0 \) and \( \mu < 0 \) have \( \partial f(\mu, v, z)/\partial z = 0 \).

Figure 2 shows the variation of the angle-averaged distribution function,

\[
f^0(\mu, z) = \frac{1}{2} \int_{-1}^{1} d\mu f(\mu, v, z) ,
\]

as a function of the electron speed \( v \) (in units of thermal speed, \( \nu_{th} = 2kT/m \), at \( T = 10^6 \) K) for the parameter combination \( F_{\text{in}}/a = 6 \times 10^{11} \) K\(^{3/2}\) cm\(^{-1}\) and \( nT = 6 \times 10^{-14} \) K cm\(^{-3}\). To illustrate typical results for both the lower and upper portions of models A, B, and C, distributions are given for \( T = 20,000 \) and \( T = 200,000 \) K. For comparison, the dashed curves denote the corresponding Maxwellian source distribution \( f^* \) at the same temperatures.

The effects that will interest us in the remainder of this paper arise from the tendency for large departures of \( f^0(\mu) \) from \( f^*(\mu) \) to develop at large velocities. This tendency, which is most pronounced in the curves for the constant-flux model A, can readily be understood from the previous work by Roussel-Dupre (1980a, b) and Shoub (1982, 1983). For relatively low velocities near the local thermal speed, collisions keep \( f^0(\mu) \) close to the value given by the local Maxwellian source distribution \( f^*(\mu) \). However, for velocities above the thermal speed the collision cross section decreases as \( v^{-4} \), and so, at some critical velocity, the electron mean free path becomes comparable to the local temperature scale length. Near this velocity, \( f^0(\mu) \) becomes enhanced relative to a local Maxwellian because a substantial number of such electrons now originate from a thermal source distribution with a higher temperature. At very high velocities, nearly all the electrons originate in high-temperature source distribution in the upper portions of the atmosphere, from where they stream essentially without collisions downward into the cooler regions. The high-velocity tail of the electron distribution in these cooler regions therefore becomes virtually independent of the local temperature structure, and depends instead primarily on the temperature structure of the upper transition region. This explains why the outer tails of \( f^0(\mu) \) at \( T = 20,000 \) K and \( T = 200,000 \) K in models A, B, and C all merge for large \( v \), since these models all have very similar temperature structure in the upper transition region. In fact, because of this similarity, even the inner tails of \( f^0(\mu) \) are very alike at the relatively high temperature \( T = 200,000 \) K.

The phenomenon that most clearly distinguishes the results for the theoretical and empirical models B and C from those for the constant-flux model A of Shoub (1982, 1983) is the presence of the relative minimum in the inner tail \( [v \approx \nu_{th}(10^6 \text{ K})] \) of the distribution function at low temperatures.
electron velocity distribution function could significantly alter
where $v_0$ averaged electron distribution function through
rates $C$ for such collisional processes depend on the angle-
altering the collisional excitation/ionization balance there. The
the excitation/ionization threshold energy $\gamma$, and $a$ is the cross
over which an enhanced
for the maximum column depth $\Delta t$ results in
the relative strength of the high-velocity tail. The means that effects
which arise from the departures of the electron distribution
from a Maxwellian at large velocity, such as the nonclassical
modifications in ion excitation/ionization and electron heat
transport, may depend sensitively on the assumed temperature
structure. In the next sections, we shall examine these effects
more carefully.

IV. EFFECT ON COLLISIONAL EXCITATION AND
IONIZATION RATES

Let us now consider whether enhancements in the tail of the
electron velocity distribution function could significantly alter
the emission-line brightness of the lower transition region by
altering the collisional excitation/ionization balance there. The
rates $C$ for such collisional processes depend on the angle-
averaged electron distribution function through
$$ C = 4\pi n \int_{v_0}^{\infty} n_{\alpha}(v) f^\alpha(v) v^2 dv, \quad (4) $$
where $v_0$ is the speed of an electron with energy $m_0v_0^2/2$ equal to
the excitation/ionization threshold energy $\gamma$, and $\sigma$ is the cross
section. In order to examine the general effects of deviations of
the distribution function from a Maxwellian, let us assume a
generic cross section given by the simple analytic formula (see
Lotz 1967; Johnson 1972),
$$ \sigma(v) = A \ln \left(\frac{v}{v_0}\right)^2, \quad v > v_0. \quad (5) $$
We can then compute from equation (4) both the rate $C^*$ for a
Maxwellian distribution and the rate $C_{BGK}$ for the BGK
distributions calculated in the previous section.

Results for the ratio $C_{BGK}/C^*$ at various temperatures in
models A and B are compared in Table 1 for various values of $\gamma$
that are appropriate to lines formed in the lower transition
region. (Corresponding ratios for model C are all within 1% of
unity and so are not listed separately in Table 1.) For the
constant-flux model A, the ratio is near unity for high $T$ and
low $\gamma$. However, it can be quite large for relatively low $T$ and
high $\gamma$, for which ionization/excitation by electrons in the
enhanced, high-velocity tail of the BGK distribution dominates
that by electrons in the low-velocity, Maxwellian core. On
this basis, Shoub (1983) suggested that the ionization/
excitation balance, and hence the line emission, in the lower
transition region might be strongly affected by nonclassical
electron transport. However, for the models B and C with the
more realistic lower transition region structure, the BGK and
Maxwellian rates are must closer throughout the relevant
range in $\gamma$ and $T$. The reason (see § III) is that, in these models,
the temperature scale height in the lower transition region is
much larger than the mean free path for electrons with energies
somewhat above the ionization/excitation thresholds for lines
formed there. The excitation and ionization of ions that emit
such lines is therefore dominated by the local, collision-
dominated Maxwellian core of the distribution rather than by
the nonlocal, nearly collisionless tail. This implies that nonclas-
sical electron transport effects are not very important for the
direct excitation of lower transition region emission lines.
(Such effects may, however, be important for lines with higher
excitation/ionization threshold energies, which are typically
formed in the mid or upper transition region).

We may extend this reasoning to derive a general estimate
for the maximum column depth $N_{max}$ over which an enhanced
high-velocity tail in the distribution function can affect the
collisional excitation/ionization of an ion at a given threshold
energy $\gamma$. This is just the column depth over which the elec-
trons with kinetic energy $\frac{1}{2}mv^2$ somewhat above $\gamma$, say $\frac{1}{2}mv^2 \approx 2\gamma$, would be thermalized by Coulomb collisions. Using the
collision rate defined in equation (4), we have (for $\alpha = 1$)
$$ N_{max} \approx \left(\frac{\gamma^2}{\pi e^4 \ln \Lambda}\right). \quad (6) $$

For example, for the $Ly\alpha$ excitation energy $\gamma \approx 10$ eV, we
obtain $N_{max} \approx 10^{14}$ cm$^{-2}$, which is far less than the $3 \times 10^{17}$
cm$^{-2}$ column thickness of the $Ly\alpha$ plateau in VAL model C.

<table>
<thead>
<tr>
<th>Temperature (K)</th>
<th>$\gamma$ (eV) = 10</th>
<th>15</th>
<th>20</th>
<th>25</th>
<th>Model A</th>
<th>1</th>
<th>1.02</th>
<th>1.14</th>
<th>1.22(5)</th>
</tr>
</thead>
<tbody>
<tr>
<td>10,000</td>
<td>95.6</td>
<td>1.25(4)</td>
<td>2.17(6)</td>
<td>4.33(8)</td>
<td></td>
<td>1</td>
<td>1.02</td>
<td>1.14</td>
<td>1.22(5)</td>
</tr>
<tr>
<td>20,000</td>
<td>1.64</td>
<td>7.63</td>
<td>62.9</td>
<td>6.55(2)</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1.02</td>
<td>1.28</td>
</tr>
<tr>
<td>30,000</td>
<td>1.03</td>
<td>1.50</td>
<td>3.50</td>
<td>1.17(1)</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1.02</td>
</tr>
<tr>
<td>40,000</td>
<td>1</td>
<td>1.07</td>
<td>1.43</td>
<td>2.44</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1.01</td>
</tr>
<tr>
<td>50,000</td>
<td>1</td>
<td>1.01</td>
<td>1.09</td>
<td>1.36</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1</td>
</tr>
<tr>
<td>60,000</td>
<td>1</td>
<td>1</td>
<td>1.02</td>
<td>1.11</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1</td>
</tr>
<tr>
<td>70,000</td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1.03</td>
<td></td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1</td>
</tr>
</tbody>
</table>
We thus conclude that the observed brightness of Lyα cannot be explained by collisional production by an enhanced high-velocity tail. By similar arguments, this conclusion may be extended to other lines which, because of their relatively low ionization and excitation energies, are formed in the lower transition region.

V. EFFECT ON HEAT TRANSPORT

The electron heat flux $Q$ is obtained from the electron distribution function $f(\mu, v, z)$ from

$$Q(z) = 4\pi \int_0^{\infty} v f(\mu, v, z) v^2 dv,$$  \hspace{1cm} (7)

where $f^1$ is the first angular moment of $f$, defined by

$$f^1(v, z) \equiv \frac{1}{2} \int_{-1}^{1} f(\mu, v, z) d\mu.$$ \hspace{1cm} (8)

We can calculate a linear, BGK value of the heat flux in our models by applying to equations (7) and (8) the distribution function obtained from integrating the BGK equation (2). This heat flux contains both a classical, local component due to relatively low velocity electrons with collision mean free paths smaller than a local temperature scale height, as well as a nonclassical, nonlocal component due to the enhanced high-velocity tail. The classical component can be computed analytically by using a Chapman-Enskog perturbation expansion (see, e.g., Chapman and Cowling 1964) on the BGK equation (2) to obtain a local approximation to the electron distribution function,

$$f_c(\mu, v, z) \approx f^* \left[ 1 + \frac{v}{v^*} \right] v^2 \mu d\mu,$$ \hspace{1cm} (9)

Application of equations (8) and (9) to equation (7) then yields a linearized BGK form of the classical conductivity law (see Spitzer 1962).

We shall use the deviation of the nonclassical BGK heat flux from its classical value to gauge whether nonclassical heat transport is likely to be important in a given model. To isolate this nonclassical component, we define the BGK tail heat flux as

$$\Delta Q \equiv Q_{BGK} - Q_{cl}.$$ \hspace{1cm} (10)

Figure 3 shows the temperature variation of $\Delta Q$ for each of the transition region models A, B, and C. Curves are plotted for three values of the steepness parameter $F_{ou}/\alpha$, assuming the nominal value of the pressure $nT = 6 \times 10^{14}$ K cm$^{-3}$. The important features of Figure 3 can be summarized as follows: first, a factor of 2 change in $F_{ou}/\alpha$ results in about an order of magnitude change in $\Delta Q$; thus, the overall magnitude of the nonclassical component of the heat flux that one infers for the transition region is quite sensitive to the steepness of the assumed transition region model. Second, for a reasonable value of $F_{ou}/\alpha$, the nonclassical heat flux that arises in the upper transition region is of sufficient magnitude to balance the radiative energy loss observed in lower transition region emission lines. Third, for temperature decreasing below 200,000 K, $\Delta Q$ continues to increase in model A, but it decreases in models B and C; thus, the magnitude of the nonclassical heat flux at a temperature where such lower transition region lines form depends on the temperature structure of the lower transition region itself.

Figure 4 shows the variation of $\Delta Q$ with column depth $N = \int_0^z n(z)dz$, measured from a point where $T = 10^6$ K. For
Fig. 4.—Column depth dependence of the tail heat flux for the indicated parameter values. Curves also apply to other values of $nT$ and $a$ with the same $F_0$ if the ordinate is scaled in proportion to $nT$ and the abscissa in proportion to $1/a$.

This calculation the bottom boundary of each model has been extended to include a region with a constant temperature of 8000 K, to mimic the chromosphere. The results show that differences in nonclassical heat flux for the three models are confined to a narrow range in column depth near the lower transition region. These differences reflect the number of electrons in the inner tail of the distribution function. Beyond the relatively short range of column depth over which such low-velocity tail electrons are thermalized, the heat fluxes become equal because the outer tail of the distribution function is quite similar for all three models (see Fig. 2). Figure 4 also shows that, upon entering a layer in which the temperature is not steeply decreasing, the energy carried by these high-velocity tail electrons is typically dissipated over a column depth of $\sim 10^{17}$ cm$^{-2}$. Since this is also approximately the column thickness inferred empirically ($\sim 3 \times 10^{17}$ cm$^{-2}$; VAL) for the region emitting Ly$\alpha$ and other lower transition region lines, we see that, unlike the classical heat flux, the heat flux carried by these high-velocity electrons can be dissipated over a broad enough region to balance the radiative energy loss in these lines.

VI. SUMMARY

In this paper we have examined the dependence of nonclassical electron transport effects in the solar transition region on the assumed transition region temperature structure, with the focus on how these nonclassical effects might influence line emission from the lower transition region. Our calculations of electron transport have been based on the relatively simple, linearized BGK method in which, to simplify further, we have also neglected a number of possible effects (e.g., polarization electric field, nonvertical and/or nonuniform magnetic field, inelastic collisions, collective plasma effects, etc.).

Our primary results may be summarized as follows:

1. Upon entering the lower transition region of models with plausibly low temperature gradients, the enhanced high-velocity tail in the electron velocity distribution function, which arises in the upper transition region, tends to be thermalized first in its lower velocity (inner) end. This leads to an electron distribution function that has a relative minimum at intermediate velocities; since such a distribution function is subject to the two-stream instability, this implies that collective plasma effects may play an important role in the thermalization of electrons in the lower transition region.

2. Nonclassical transport effects are not very important for the excitation and ionization balance of ions whose lines form in the lower transition region, which characteristically have excitation and ionization potentials around 10 eV. This is because electrons with such relatively low energies are thermalized by Coulomb collisions over a column depth that is much smaller than the empirically inferred column depth of line-emitting material.

3. For models that reproduce the observed upper transition region line emission, the computed nonclassical heat flux arising in the upper transition region is of sufficient magnitude and dissipates over a broad enough region to balance the observed energy flux emitted in all lower transition region lines.

4. However, the magnitude of the computed nonclassical heat flux in the lower transition region depends sensitively on...
the assumed steepness of the temperature variation in the upper transition region.

The primary conclusion we draw from these results is that it is quite possible that nonclassical electron heat transport plays an important role in the line emission from the lower solar transition region; in order to be certain of its overall importance, it will be necessary to incorporate several other effects into transition region modeling. A logical first step would be to derive an improved transition region model on the basis of a nonlinear transport calculation in which energy is conserved. Second, collective plasma effects may be important in the lower transition region, where the inner tail of the electron distribution function tends to form a relative minimum. (However, recent work by Shoub [private communication], using the more rigorous Fokker-Planck description of electron collisions, indicates that energy diffusion effects, which have been neglected in the present BGK analysis, will prevent the formation of such a relative minimum.) Finally, this calculation should necessarily include the various factors that can influence the temperature structure, including, e.g., geometry, radiative losses, in situ heating, as well as electron heat transport.

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