EXCITATION OF O I LINES IN THE SOLAR CHROMOSPHERE

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ABSTRACT

Observations of O I lines in the solar spectrum are examined to determine whether differences in behavior of lines of the quintet and triplet term systems are consistent with collisional excitation and/or photoexcitation of both quintets and triplets. Intensities, I_R, in near-infrared emission lines observed above the limb at total eclipse decrease exponentially with height h. The inverse scale heights (d ln I_R/ dh) for the triplet lines at 844.6 nm and quintet lines at 777.2 nm are found to be in the ratio of 1.45. Ultraviolet O I emission-line intensities I_UV observed on the solar disk show strong variations, and the distributions of triplet (130.4 nm) and quintet line intensities about the means are different. Variances in ln I_UV are found to have a triplet-to-quintet ratio of 1.50, in close agreement with the ratio of d ln I_R/ dh.

It is shown that the simple assumption of collisional excitation of quintets and triplets coupled with collisional de-excitation of the quintets leads to the correct ratios for both the UV variances and d ln I_R/ dh. Also, under this assumption d ln I_R/ dh for the quintet lines is predicted to have the same value as d ln I/ dh at the head of the hydrogen Balmer continuum, which, in fact, it does. On the other hand, Carlsson & Judge (1993) have shown that collision rates computed from the Vernazza, Avrett, & Loeser (1981, hereafter VAL) model chromosphere using current estimates of O I collision strengths are too low to produce the observed mean intensity in O I 130.4 nm. In a similar sense, we find that the predicted intensity of O I 130.4 nm is much too weak relative to O I 135.6 nm, and that the VAL mean models A–F cannot reproduce the observed behavior of these lines, even including photoexcitation by H Lyβ. These difficulties are removed by increasing specific electron-atom collision rates. Such increases could reflect large errors in atomic cross sections close to threshold and/or the inadequacy of the assumptions made by VAL for predicting line intensities. The nonlinear dependence of line intensities on temperature and density, especially for far-UV lines, makes the latter alternative a likely factor. We conclude that the O I UV lines are very sensitive to inhomogeneities, much more so than more traditional chromospheric lines (e.g., Mg II and k) which are formed over similar regions of the chromosphere. Such lines could therefore provide valuable diagnostics of departures of the chromospheric plasma from mean models and thereby place constraints upon heating mechanisms, once accurate atomic data become available.

Subject headings: atomic processes — eclipses — Sun: chromosphere

1. INTRODUCTION

Ultraviolet emission-line intensities in solar and stellar spectra provide one of the primary diagnostics of physical conditions in the chromospheres and transition regions in these objects. For such diagnostics to be useful, however, the line excitation mechanism must be properly identified and accurately calculated because of large departures from local thermodynamic equilibrium (LTE). Because these lines, collectively, contribute a net loss of radiative energy from the solar atmosphere, the ultimate channel through which that energy flows must involve collisions with gas particles, primarily with electrons. It seems clear, therefore, that the majority of the UV lines are collisionally excited. It does not follow that all of the lines are excited in this manner.

The chromospheric UV spectrum of O I is dominated by the resonance triplet \((2p^6 3P_{0,1,2} \rightarrow 2p^5 3s^2 \, ^3S_1)\) at 130.2, 130.4, and 130.6 nm, with weaker intersystem lines \((2p^6 3P_{1,2} \rightarrow 2p^5 3s^2 \, ^3S_2)\) at 135.5 and 135.6 nm. These lines are potentially important and almost unique, in the following sense. The lines represent some of the strongest high-energy transitions \((\Delta E/ k T_e \gg 1)\), where \(\Delta E\) is the excitation energy, \(k\) is Boltzmann's constant, and \(T_e\) is the electron temperature) formed in regions extending deep into the chromosphere. This makes the O I line intensities much more sensitive to high-energy electrons or photons, and thus inhomogeneous chromospheric structure, than more traditional chromospheric diagnostics (e.g., Ca II H and K, Mg II h and k), and continuum diagnostics used in semiempirical modeling efforts (e.g., Vernazza, Avrett, & Loeser 1981, hereafter VAL). Thus, these lines should provide important new clues to conditions in the solar chromosphere once the line excitation mechanisms have been identified.

It has long been argued that the strong triplet lines of O I near 130.4 nm may be excited through fluorescence with the hydrogen Lyβ line. Figure 1 shows a term diagram of O I including the levels and lines involved in the pumping process. Theoretical calculations using the best available atomic data and "mean" solar and stellar models have consistently shown that Lyβ fluorescence predominates over collisional excitation in solar and later type stars. The most thorough and most recent calculations are those of Carlsson & Judge (1993), who provide references to earlier work. On the other hand, those

1 The National Center for Atmospheric Research is sponsored by the National Science Foundation.
who approach the problem from a more empirical point of view frequently adopt the collisional excitation approach (cf. Athay & Dere 1991).

The fluorescent excitation mechanism for the O I lines has perhaps two arguments in its favor: (1) there is the supporting theoretical work (Carlsson & Judge 1993) and (2) these lines do exhibit characteristics that distinguish them from lines that are collisionally excited in low-gravity stars (Haisch et al. 1977; Jordan & Judge 1984; Carlsson & Judge 1993). However, the question of excitation mechanisms in the spectra of dwarf stars is by no means settled. Do mean chromospheric models provide a sufficient test? Are the needed atomic parameters sufficiently accurate? Specifically, Carlsson & Judge (1993) and Judge & Carlsson (1993) point to the need to reexamine electron-atom cross sections, since currently available data are based upon simpliﬁed atomic models.

Carlsson & Judge (1993) conclude that if Lyβ fluorescence indeed dominates, the emissivities (and hence emergent intensities) in the resonance triplet at 130.2 nm depend linearly on the n = 3 hydrogen density and roughly linearly on the local electron density. In their calculations, performed in several "mean" chromospheric models, the Lyβ radiation is nonlocally controlled. However, the Lyβ radiation depends on both electron density and temperature at the source or sources of the radiation. Furthermore, the diffusion scale of the Lyβ radiation and the mean separation of source locations in the real chromosphere may each be small enough that the Lyβ mean intensity itself remains closely correlated with local density and temperature. If such were the case, the O I lines would still show variations that correlate with local electron density and temperature even though the primary excitation is via Lyβ fluorescence. Nevertheless, one still expects that Lyβ fluorescence will not perfectly mimic collisional excitation. Since Lyβ selectively excites the triplet levels, and assuming the quintets and triplets are not tightly coupled by collisions, it seems most advantageous to examine differences in the behavior of the quintets and triplets for evidence of the Lyβ fluorescence.

The purpose of this paper is to examine the observational characteristics of the solar triplet and quintet multiplets of O I both in the UV and in the near-infrared. The latter spectral regions are well observed at total eclipse and give important information on both the rate of decrease of line intensities with height and the relative population of the excited quintet and triplet levels. The UV lines provide complementary data on the variations of quintet and triplet intensities with position on the solar surface. Also, they provide useful additional information on the relative populations of quintet and triplet levels.

2. ECLIPSE OBSERVATIONS

Fluorescent excitation of O I by Lyβ excites the 3d \(3\,D^\prime\) levels before making the resonance transition back to the ground state via the \(3\,p\ \ 3\,P\) level (Fig. 1). The transitions from the \(3\,d\ \ 3\,D^\prime\) to the \(3\,p\ \ 3\,P\) level and thence to \(3\,s\ \ 3\,S^\prime\) produce the 844.6 nm triplet. A similar transition in the quintet system from \(3\,p\ \ 3\,P\) to \(3\,s\ \ 3\,S^\prime\) produces the 777.2 nm multiplet. Only a quarter of an electron volt energy difference exists between the \(3\,p\ \ 3\,P\) and the \(3\,p\ \ 3\,P\) level.

Both the 844.6 nm triplet lines and the 777.2 nm quintet lines are moderately strong in the low chromosphere (e.g., Athay & House 1962). At total eclipse they reach maximum strength very near the limb of the Sun, then decrease exponentially with height. "Intensities" \(I\) observed at eclipse with slitless spectrographs are given by

\[
I' = \int_{h_0}^{\infty} I dh ,
\]

where \(I\) is the specific intensity in the viewing direction tangential to the limb, and \(h\) is the height above the limb. For an optically thin atmosphere in which the emissivity, \(\epsilon\), is assumed to vary exponentially with height, i.e.,

\[
\epsilon = \epsilon_0 \exp (-\beta h) ,
\]

the specific intensity \(I\) is simply the line-of-sight integral of expression (2), with \(h < R\) (the solar radius):

\[
I = \left(\frac{2\pi R}{\beta}\right)^{1/2} \epsilon_0 \exp (-\beta h) ,
\]

and with equation (1),

\[
I' = \left(\frac{2\pi R}{\beta}\right)^{1/2} \epsilon_0 \exp (-\beta h) .
\]

The observed reciprocal scale heights for the infrared quintets and triplets of O I, denoted here as \(\beta_3\) and \(\beta_4\), are \(\beta_3 = 2.0 \times 10^{-8}\) cm\(^{-1}\) and \(\beta_4 = 2.9 \times 10^{-8}\) cm\(^{-1}\) (Athay & House 1962). The difference between \(\beta_3\) and \(\beta_4\) is potentially an important clue to the excitation mechanism, and it demon-

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states that the population of the 3p \(^3P\) level decreases more rapidly with height than does the population of the 3p \(^1P\) level. In other words, the 3p \(^3P\) level increases in population relative to the 3p \(^1P\) level with increasing height.

At a height of 600 km above the limb, the observed intensity \(I'\) in the 777.2 nm line exceeds the intensity in the 844.6 nm line by a factor of 6. Standard chromospheric models predict that these lines are optically thin because of their high excitation energy, and they show none of the characteristics of lines that are self-absorbed, viz., the ratios of intensities within the quintet lines are close to the expected optically thin ratios, and the values of \(\beta_3\) and \(\beta_3\) do not decrease in the low chromosphere as they do in optically thick lines such as H\(\alpha\) and the Ca \(\text{II}\) infrared triplet. Under optically thin conditions, the ratio of the 777.2 nm quintet intensity to the 844.6 nm triplet intensity should be close to the ratio of \(A_3 n(3p \^3P)/A_3 n(3p \^3P)\), where \(A_3\) and \(A_3\) are the Ehrlich-Arnold A-coefficients and the \(n\)'s are population densities, which reduces in LTE to \((gf)_{3s}/(gf)_{3p}\), since the wavelengths of the transitions are similar (we use \(gf\) to denote the weighted average product of statistical weights and oscillator strengths for the multiplets). Using the \(gf\)-values tabulated by Wiese, Smith, & Glennon (1966), and observed values of \(\beta_3\) and \(\beta_3\) in equation (4), we obtain an expected LTE ratio of line intensities of approximately 1.4. (The more recent oscillator strengths adopted by Carlsson &Judge 1993 yield a value of 1.88). This differs by a factor of 4 from the observed ratio at 600 km above the limb. Clearly, therefore, the number density ratios \(n(3p \^3P)/n(3p \^3P)\) are out of LTE. The quintet levels are overpopulated relative to the 3p triplets, and the relative overpopulation increases with height.

Finally, we note that the model chromospheres of VAL, while clearly unsuccessful at describing many important features of the solar chromosphere (e.g., the very fine structure or dynamics), can successfully reproduce the emission gradients seen in eclipse spectra (Athay et al. 1955). We interpret this to mean that the density structure with height in these models is not, on average, unreasonable. We will therefore use the VAL models below to estimate mean densities in the solar chromosphere, recognizing that departures from the mean densities and temperatures are expected to have significant effects on the O I spectrum.

3. DISK OBSERVATIONS

For disk data of O I emission lines, we use the Naval Research Laboratory High Resolution Telescope and Spectrograph (HRTS) data from Spacelab 2 given by Athay & Dede (1991). From their line intensity data extending from center to limb, we select a rectangular area roughly 30° wide extending from Sun center to approximately 0.6 R\(_{\odot}\). This area includes mostly quiet solar regions, but also contains a small sunspot and several small plages. Values of \(\log I_0\), \(\sigma_i\) (mean and rms variance in \(\log I\)) and \(m_{ij}\) (= \(d \log I_0/d \log I\)) are given in Table 1. \(I_i\) is the frequency-integrated intensity of line \(i\) (minus the background intensity) in ergs cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\).

Values of \(m_{ij}\) are derived from scatter plots of \(\log I_i\) versus \(\log I_j\) using only log \(I_i\) \(\leq \log I_0\) and log \(I_j\) \(\leq \log I_0\). Reasons for these restrictions on log \(I_i\) and log \(I_j\) are given in the following paragraphs.

Each of the lines listed in Table 1 has a line-center optical depth exceeding unity at some depth in the chromosphere. This does not necessarily mean, however, that the frequency-integrated line intensities are diminished by saturation effects. In the permitted multiplets at 130.6 (O I), 156.1, (C I), and 156.3 (Fe II) nm, the thermalization depth, \(\tau_{\text{TH}}\), can be large, and the only effect of \(1 \leq \tau \leq \tau_{\text{TH}}\) on the frequency-integrated intensity is to force photons into the line wings, where they readily escape. In the case of the intersystem line at 135.5 nm, however, collisional de-excitation from 3s \(^3S\) to the ground state, 2p \(^3P\), and intersystem collisional transitions to 3s \(^3S\) may approach, or even exceed, the spontaneous transition probability from 3s \(^3S\) to 2p \(^3P\) in the chromosphere. In this case, saturation occurs because photons are collisionally destroyed before escaping. Saturation will also occur when the optical depth in the background continuum exceeds unity. This was found to be the dominant process thermalizing the O I resonance lines in the calculations of Carlsson & Judge (1993).

Histograms of the number of pixels in the Spacelab 2 data as a function of log \(I_i\) (Athay & Dede 1991) suggest that intensities in the lines at 156.3, 130.6, and 135.5 nm are truncated on their high-intensity sides, whereas 156.1 exhibits a rather symmetric distribution. This trend is reflected in the values of \(\sigma_i\) in Table 1, with 156.1 exhibiting the largest values of \(\sigma_i\). The truncations on the high-intensity sides of the O I and Fe II line distributions are most likely due to saturation effects, as suggested by Athay & Dede (1991). For this reason, we use only the data for which log \(I_0\) and log \(I_i\) are below their mean value to determine \(m_{ij}\), since otherwise we cannot relate values of \(I_i\) and \(I_j\) linearly to properties of the emitting gas in approximate calculations (§ 4).

Two results in Table 1 are of interest in connection with excitation mechanisms for the emission lines of O I. These are (1) that both \(\sigma_i/\sigma_j\) and \(m_{ij}\) are near 1.5 for the O I lines at 130.6 and 135.5 nm, and (2) that both \(\sigma_i/\sigma_j\) and \(m_{ij}\) are near unity for any pair of the lines at 130.6, 156.1, and 156.3 nm. The equality of \(m_{ij}\) with \(\sigma_i/\sigma_j\) is expected, since each quantity measures the ratio of the variances in the intensities of the two lines on the disk. Even though both the 130.6 and 156.3 nm lines show evidence of saturation in the brighter solar features, the saturation is not strong enough to make the ratio \(\sigma_i/\sigma_j\) significantly different from \(m_{ij}\).

A third interesting feature of the O I emission lines in the UV is the intensity ratio \(R = (I_{135.5} + I_{135.8})/(I_{130.6} + I_{130.4} + I_{130.2})\). In the HRTS data used by Athay & Dede (1991) the three members of the triplet have very nearly the same intensity, whereas the 135.8 nm members of the quintet intersystem lines is only about 30% as strong as the one at 135.5 nm, consistent with the optically thin branching ratio of Zeippen, Seaton, & Morton (1977). The ratio \(R\) is observed to be near 0.07 under the average solar conditions (i.e., near the peak of the distributions of log \(I_i\)). This ratio is related to the conclusion in § 2 that the 3p \(^3P\) level is overpopulated with respect to the 3p \(^1P\) level by a factor of about 4 compared to LTE.

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calculations discussed by Carlsson & Judge (1993), the 3s and 3p levels within each system are strongly coupled through the transitions $3s \rightarrow 3p$ by photospheric radiation having a color temperature near 6000 K. This is nearly the same as the average chromospheric temperature at a height of 600 km above the limb (Athay et al. 1955), where the infrared lines were used to estimate the relative populations of the $3p^2P$ and $3p^2P$ levels, and where the empirically derived contribution functions for the UV O I lines are at their maxima (Athay & Dure 1991). Thus, we expect the relative populations of terms of the 3s and 3p configurations to be near their LTE ratio within each system, even though the two systems do not have LTE population ratios with respect to each other.

A fourth point of interest is the near equality of $\sigma_{13.0,6}/\sigma_{135.5}$ with the value of $\beta_3/\beta_5$ obtained from the infrared eclipse data. The first ratio is a measure of the relative variances between triplets and quintets across the solar disk, and the second is a measure of the relative variances with height in the chromosphere. At first glance, the equality of $\sigma_3/\sigma_5$ to $\beta_3/\beta_5$ is surprising. In mean models for the chromosphere, density and temperature are negatively correlated in height but most often positively correlated for different disk features. The chromospheric temperature increases outward, but the gas pressure and mean density decrease. In the case of different disk features, however, increased heating expands the atmosphere outward and results in increased densities and temperatures. Thus, for $\sigma_3/\sigma_5$ to closely approximate $\beta_3/\beta_5$, the sign of the correlation between density and temperature must be of minor importance. We investigate these interesting features of the O I variances in the following section.

4. ANALYSIS

4.1. Some Factors Influencing Line Ratios

Variances in line intensities can arise from a number of factors, including temperature, density, optical thickness, relative abundances, and unresolved plasma structure. The overall trend for direct proportionality between the intensities of the 130.6, 156.1, and 156.3 nm lines of O I, C I, and Fe II, respectively, is an indication that each of these lines has a similar dependence on temperature and density. Conversely, the reduced intensity variance in the 135.5 interstellar line of O I, together with the reduced height gradients in the infrared O I triplet intensities relative to the triplet intensities implies a reduced dependence of the triplet intensities on density and/or temperature. Among the UV lines we are considering, the 135.5 line is unique in two respects. It is optically thin in the middle chromosphere, and its upper level is metastable. There is no particular reason for low optical depth to reduce the variance in intensity, but there are straightforward reasons for reduced variance associated with lines from metastable levels. The situation is entirely analogous to that found when height gradients of coronal forbidden lines are compared to height gradients of permitted lines. The former always show smaller gradients than the latter.

The long accepted explanation for the different variances in the coronal lines is that permitted lines are (to first order) collisionally excited followed by radiative de-excitation, whereas the forbidden lines are both collisionally excited and collisionally de-excited. The line intensities are then represented by the functional forms

$$I_p \propto n_1 n_e e^{-h\nu/kT_e}$$

and

$$I_p \propto n_1 n_e e^{-h\nu/kT_e}$$

where the subscripts $P$ and $F$ refer to the permitted and forbidden cases, $n_1$ is the ground-state population density, $n_e$ is the electron density, $T_e$ is the electron temperature, and $\nu$ is the frequency of the line transition.

If we postulate that relations (5) and (6) apply to the O I lines, we find all of the observed properties for the variances. To demonstrate this, we need to determine the scaling between electron density $n_e$ and electron temperature $T_e$ in the chromosphere. In the LTE approximation, $n_e$ and $T_e$ are related through the ionization of hydrogen as given by the Saha equation (we assume the number of protons $\approx n_p$)

$$n_1(H) \propto n_e^{2/3} e^{3/2 h\nu/kT_e},$$

where $h\nu$ in equation (7) represents the ionization energy. However, in the low chromosphere hydrogen is out of LTE. The ionization is dominated by photoionization in the Balmer continuum by optically thin radiation whose intensity is controlled by the photosphere, following collisional or radiative excitation of the $n = 2$ level. The Lyman continuum plays essentially no role in the chromosphere, since the ionization and recombination rates almost cancel (VAL). As a result, the exponential term in the temperature in equation (7) is given by $h\nu_{21}/kT_e$, where $\nu_{21}$ is the energy between the $n = 2$ level and the ground level ($n = 1$) of hydrogen, and the other temperature term varies as $T_e^{-\alpha}$, where $\alpha \approx 3/2$. The exponential term in equation (7) therefore happens to be close to the exponential terms for the emissivities of the O I UV lines. Ignoring these small differences in energy, we can combine equations (5), (6), and (7), replacing $3/2$ by $\alpha$, to obtain the approximate scaling laws

$$I_p \propto n_e^{2/3} T_e^{-\alpha}$$

and

$$I_p \propto n_e^{2/3} T_e^{-\alpha}$$

In the low chromosphere where the O I lines are formed, hydrogen and oxygen are only weakly ionized and $n_1(O)/n_1(H)$ is constant (Carlsson & Judge 1993). Also, small changes in $T_e$ produce relatively large changes in $n_e$, (eq [7]), so that nearly all of the temperature and density dependence in equations (8) and (9) can be parametrically described by the terms in $n_e$. It follows that $d\ln I_p/d\ln T_e = 3/2$. Furthermore, the weak dependence of $d\ln I_p/d\ln T_e$ on temperature means that the sign of the correlation between temperature and density is unimportant. Equations (8) and (9) are equally valid for both height variances (eclipse data) and disk variances (UV disk data). The reader should note that this result is unique to O I. It arises because of the similarity in energy-level structure between O I and H I and is not expected in atoms or ions that do not share in this similarity.

Equation (9) for $I_p$, without the factor $n_1(O)/n_1(H)$ but with $\alpha = 3/2$, is of exactly the same form as the Balmer continuum intensity near the series limit, observed above the solar limb. Equation (9) predicts, therefore, that $\beta_3$ for the infrared quintets of O I should be close to $\beta_{bal}$ for the head of the free-bound Balmer continuum near 364.0 nm. The measured value for $\beta_{bal}$ at the same point on the limb where $\beta_3$ and $\beta_5$ were
measured is $2.1 \times 10^{-8} \text{ cm}^{-1}$ (Athay 1976), which agrees very well with $\beta_5 = 2.0 \times 10^{-8} \text{ cm}^{-1}$.

The striking success of equations (5) and (6) in predicting the correct absolute value for $\beta_5$ together with the correct ratios for both $\beta_3/\beta_5$ and $\sigma_3/\sigma_5$ provides an argument in favor of collisional excitation of both the quintets and triplets of O I together with collisional de-excitation of the quintets. As noted earlier, however, these conclusions are inconsistent with current atomic cross sections for O I and current chromospheric models. In the following, we consider physically reasonable modifications to O I cross sections and/or model chromospheres necessary to satisfy the observed behavior of the line intensities, and include treatments of the effects of radiative transfer.

4.2. Approximate Calculations of Line Intensities

In order to quantify the discussion of § 4.1, we adopt a three-level model atom. Level 1 represents the $2p^3P$ ground state. Level 2 represents the metastable $3s^3S^o$ level, and level 3 represents the $3s^3S$ level. For notational convenience we initially place levels 2 and 3 at the same excitation energy (they differ only by $\Delta W \approx 4000 \text{ K}$, and thus this does not influence our conclusions). Our goal is to examine the dependence of the 3–1 and 2–1 line intensities on collision rates to see whether the transitions can have the same total and relative intensities as those observed in the chromosphere without destroying the validity of relations (5) and (6). Observations require that, on average, $I_{31} + I_{21} \approx 800 \text{ ergs cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ and that $I_{31}/I_{21} \approx 4$, near disk center, for the quiet Sun.

We begin by considering energy conservation in an atmosphere in which background absorption of photons can be ignored (this is discussed in § 4.4). Thus, the ultimate fate of the energy in the O I UV line photons must either be collisional destruction (thermalization) or escape from the atmosphere. The energy represented by line intensities $I_{31}$ and $I_{21}$ comes from thermal energy in the chromosphere. It is converted into photon energy through collisional excitations to levels 3 and 2 in the model atom. Thus, energy conservation requires

$$I_{31} + I_{21} = \frac{hv}{4\pi} \int_{0}^{\infty} \left[(C_{13} + C_{12})n_1 - C_{21}n_2 - C_{31}n_3\right] dh,$$

or

$$I_{31} = \frac{hv}{4\pi} \int_{0}^{\infty} \left[(C_{13} + C_{12})n_1 - C_{21}n_2 - C_{31}n_3\right] dh,$$

and

$$I_{21} = \frac{hv}{4\pi} \int_{0}^{\infty} \left[(C_{12}n_1 + C_{32}n_2) - (C_{21} + C_{23})n_3\right] dh.$$

The terms $C_{23}$ and $C_{32}$ do not alter the total energy balance, but they are important in partitioning the energy balance between $I_{31}$ and $I_{21}$. It is important to note that no determinations of the cross sections needed to evaluate $C_{23}$ and $C_{32}$ have been made, and yet similar intersystem transitions in other neutrals (e.g., He I) can have quite large cross sections (e.g., Berrington & Kingston 1987) which are of great significance in the rate equations. In view of the lack of these data, we first note from equation (10) that

$$I_{31} + I_{21} \leq \frac{hv}{4\pi} \int_{0}^{\infty} (C_{13} + C_{12})n_1 dh.$$  

Maxwellian-averaged collision strengths $\Upsilon(T_e)$ computed from the cross sections of Rountree (1977) yield $\Upsilon_{31} = 0.55$ and $\Upsilon_{31} = 0.42$, at $T_e = 6000 \text{ K}$, with very weak dependences on temperature. Collisional de-excitation rates are given as $C_{ij} = 8.63 \times 10^{-2} (\Upsilon_{ji}/\Upsilon_{ij}) n_e T_e^{-1/2} \text{ s}^{-1}$ and $C_{ij} = (n_i/n_e) C_{ij}$, where $n_i$ is the LTE population of level $i$. Adopting conditions in the VAL mean model C at $T_e = 6000 \text{ K}$, where the contribution functions reach maxima for the O I lines, we find $C_{13} = 7 \times 10^6 \text{ s}^{-1}$ and $C_{12} = 1 \times 10^{6} \text{ s}^{-1}$. If the bulk of the emission comes from within a scale height of the $T_e = 6000 \text{ K}$ point of the VAL model, the integration in equation (13) can be approximated by

$$I_{31} + I_{21} \approx \frac{hv}{4\pi} (C_{13} + C_{12}) n_1 L_3,$$

where $L_3$ is a scale length of the order of $\beta_5^{-1} \text{ cm}$, and $n_1 \approx 2 \times 10^{10} \text{ cm}^{-3}$ is the population density of the $2p^3P$ ground term of O I at the height in the VAL mean model chromosphere where $T = 6000 \text{ K}$. Thus, we find $I_{31} + I_{21} \approx 10^{6} \text{ ergs cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$, which is a factor of 50 or so less than the observed value. It is clear, therefore, that the observed intensities cannot be matched without a very substantial increase in the collisional excitation rates.

In order to discuss the individual line intensities, we use the equilibrium rate equations for level populations to eliminate $n_3$ and $n_2$ from equations (11) and (12). We denote $R_i$ as the radiative rate in units of $\text{ s}^{-1}$ from level $i$ to level $j$. The equilibrium equations are then (with $R_{23} = R_{23} = 0$

$$n_1(R_{13} + C_{13} + R_{12} + C_{12}) = n_2(R_{23} + C_{23}) + n_2(R_{21} + C_{21}),$$

$$n_2(R_{23} + C_{23} + R_{21} + C_{21}) = n_3 C_{32} + n_1(R_{12} + C_{12}),$$

and

$$n_3(C_{32} + R_{31} + C_{31}) = n_2 C_{23} + n_1 R_{13} + C_{13} + C_{31}.$$

The radiative terms depend on radiation transport in the lines. Since 135.5 nm transitions are optically thin in the chromosphere, and the observed emission in the background continuum is weaker than the emission line, we can set $R_{21} = R_{13}, R_{13} \ll C_{12}$, and hence ignore $R_{13}$ in the rate equations. For the resonance lines, $R_{13} = J B_{13}$ and $R_{31} = J B_{31} + A_{31}$, where $J$ is the mean intensity in the line and $B$ and $A$ are Einstein coefficients (e.g., Mihalas 1978). To find a realistic approximation for these rates, we examine the net radiative rate between levels 2 and 1 by setting the “net radiative bracket” $\rho_{31} = (n_3 R_{31} - n_1 R_{13})/n_3 R_{31}$. Then the terms ($n_3 R_{31} - n_1 R_{13}$) can be replaced in the rate equations by $n_3 R_{31} \rho_{31}$, or by $n_3 R_{31} \rho_{31}$, since stimulated emission can be ignored ($A_{31} \gg B_{31}$). Physically, $\rho$ is the net radiative downward rate divided by the total radiative downward rate, i.e., it is the probability that a photon emitted at a certain depth will

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contribute to the local photon flux. In other words, \( \rho \) is the probability that this photon was created through processes other than scattering. Assuming statistical equilibrium in the rate equations, this creation probability is balanced by the destruction probability \( \epsilon \) (such as collisional de-excitation) and the probability \( P^e \) that a given photon will escape from the atmosphere. In general, \( \rho \) and \( P^e \) are functions of the source functions and optical depths in the entire atmosphere. However, deep in the atmosphere, where \( P^e \) is of the order of \( e, \rho \) is also of the order of \( P^e \) (cf. Athay 1972, 1984; Irons 1979a, b, c). In the case of the \( O I \) resonance lines, this condition prevails, and we may therefore use the approximation \( \rho_{31} \approx P^e_{31} \).

If \( P^e_{31} \) is controlled by single-flight escape (valid in the complete redistribution approximation), then analytic approximations can be used (e.g., Avrett & Hummer 1965) to determine \( P^e_{31} \), and its inverse \( N_{13} \) (the mean number of scatterings) in terms of line-center optical depth and the Voigt damping parameter \( a \). (Even with partial redistribution, approximate formulae can be derived for the mean number of scatterings, e.g., Adams 1972; Frisch 1984). With these simplifications, equations (15), (16), and (17) become

\[
n_{13}(C_{13} + C_{21}) = n_{23}(A_{31} P_{31}^e + C_{31}) + n_{2} (A_{21} + C_{21}),
\]

(18)

\[
n_{23}(A_{21} + C_{21} + C_{31}) = n_{3} C_{32} + n_{1} C_{12},
\]

(19)

and

\[
n_{3}(A_{31} P_{31}^e + C_{31} + C_{32}) = n_{2} C_{23} + n_{1} C_{13}.
\]

(20)

Denoting \( \eta_{i}^{j} \) as the LTE population density of level \( i \), and with \( \eta_{i}^{j} C_{ij} \), equations (18), (19), and (20) yield

\[
x_{31} = x_{31}^{\dagger} C_{31} \frac{\tau_{31}}{\tau_{3}} - C_{32} x_{32}^{\dagger} \frac{\tau_{31}}{\tau_{3}},
\]

(21)

and

\[
x_{21} = x_{21}^{\dagger} C_{21} \frac{\tau_{21}}{\tau_{2}} - C_{32} x_{31}^{\dagger} x_{32}^{\dagger},
\]

(22)

where we use the notation \( x_{ij} = n_{i}/n_{j}, x_{ij}^{\dagger} = \eta_{i}/\eta_{j}, \) and \( \tau \), as the summed rate out of level \( j \) (collisional plus net radiative rates). We also rewrite equations (11) and (12) as

\[
\frac{I_{31}}{4 \pi} = n_{1} [C_{13} + C_{21} x_{31} - (C_{31} + C_{32}) x_{31}],
\]

(23)

and

\[
\frac{I_{21}}{4 \pi} = n_{1} [C_{12} + C_{32} x_{31} - (C_{21} + C_{23}) x_{31}].
\]

(24)

We note that equations (23) and (24) reduce to

\[
\frac{I_{31}}{4 \pi} = \frac{P_{31}^e A_{31} n_{3} L_{3}}{4 \pi},
\]

(25)

and

\[
\frac{I_{21}}{4 \pi} = \frac{A_{21} n_{2} L_{5}}{4 \pi}.
\]

(26)

Equations (23)–(26) simply reflect conservation of energy for the photon/gas system: those photons which are not collisionally destroyed are emitted at a rate \( A_{31} \), of which a fraction \( P_{31}^e \) escape from the slab.

To investigate the effects of different collision terms, we have made various calculations of \( I \) (Table 2). The chromosphere is represented initially by \( T_{e} = 6000 \) K, \( \eta_{e} = 10^{11} \) cm \(^{-3} \), \( n_{s} = 2 \times 10^{10} \) cm \(^{-3} \), and length scales of \( L_{3} = 1/\beta_{3} = 3.4 \times 10^{7} \) cm and \( L_{5} = 1/\beta_{5} = 5 \times 10^{7} \) cm. Throughout these calculations the radiative data were taken as \( A_{31} = 6 \times 10^{8} \) s \(^{-1} \) and \( A_{21} = 5.7 \times 10^{3} \) s \(^{-1} \) (Zeippen et al. 1977). Avrett & Hummer (1965) give the following expression for the mean number of scatterings \( N \):

\[
N = \frac{3}{2 \pi^{1/4}} \left( \frac{\tau}{a} \right)^{1/2},
\]

(27)

for a line with the mean optical depth \( \tau \) and damping parameter \( a \). Using the VAL model, we estimate \( \tau_{13} = 2 \times 10^{4} \) and \( a = 3 \times 10^{-3} \), which give \( N_{13} = 3.7 \times 10^{12} \) and \( P_{31}^e = 2.7 \times 10^{-4} \).

We can now use equations (21), (22), (25), and (26) to examine the sensitivity of the number densities and computed

**TABLE 2**

<table>
<thead>
<tr>
<th>Model</th>
<th>( \Gamma_{31} )</th>
<th>( \Gamma_{21} )</th>
<th>( \Gamma_{23} )</th>
<th>( b_{32} )</th>
<th>( log I_{31}/I_{31} )</th>
<th>( I_{31}/I_{31} )</th>
<th>( b_{32} )</th>
<th>( log I_{31}/I_{31} )</th>
<th>( I_{31}/I_{31} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( a ).....</td>
<td>0.55</td>
<td>0.42</td>
<td>0</td>
<td>0.09</td>
<td>0.78</td>
<td>0.54</td>
<td>1.33</td>
<td>1.76</td>
<td>2.1</td>
</tr>
<tr>
<td>( b ).....</td>
<td>20</td>
<td>0.42</td>
<td>10</td>
<td>1.43</td>
<td>2.17</td>
<td>8.7</td>
<td>0.71</td>
<td>2.87</td>
<td>16</td>
</tr>
<tr>
<td>( c ).....</td>
<td>2.20</td>
<td>0.42</td>
<td>10</td>
<td>0.44</td>
<td>2.85</td>
<td>12</td>
<td>...</td>
<td>...</td>
<td>...</td>
</tr>
<tr>
<td>( d ).....</td>
<td>20</td>
<td>0.42</td>
<td>5</td>
<td>...</td>
<td>...</td>
<td>...</td>
<td>0.71</td>
<td>2.86</td>
<td>15</td>
</tr>
<tr>
<td>( e ) pumped</td>
<td>20</td>
<td>0.42</td>
<td>5</td>
<td>...</td>
<td>...</td>
<td>...</td>
<td>0.68</td>
<td>2.99</td>
<td>16</td>
</tr>
<tr>
<td>( f ) pumped</td>
<td>0.55</td>
<td>0.42</td>
<td>0</td>
<td>...</td>
<td>...</td>
<td>...</td>
<td>0.16</td>
<td>2.63</td>
<td>16</td>
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<td>...</td>
<td>...</td>
<td>...</td>
<td>0.3</td>
<td>2.90</td>
<td>14</td>
</tr>
</tbody>
</table>

**Notes**—Approximate calculations are from eqs. (21), (22), (25), and (26). Detailed calculations are for the VAL model C, and the table lists intensities computed at \( \mu = 0.88 \), where \( \mu \) is the cosine of the position angle of the solar disk; \( b_{32} \) is the computed ratio of \( n_{i}/n_{j} \) for comparison with the number derived from eclipse data, evaluated at 600 km above the solar limb (roughly 900 km above the photosphere observed at disk center). Model \( a \) is a three level + continuum model; model \( b \) is a three level + continuum model, with increased collision cross sections as described in the text; model \( c \) is a calculation with modified cross sections and mean temperatures and densities; model \( d \) has cross sections as in model \( b \), but with 12 levels plus continuum; model \( e \) is identical to model \( d \), but including photoexcitation by H Ly\( \beta \); model \( f \) has cross sections as for model \( a \), but with 12 levels and including photoexcitation by H Ly\( \beta \). Model \( f \) is the closest model to that used by Carlson & Judge 1993.
intensities to the atomic model. Our approach is to compare solutions to these equations with the observations to attempt to "determine" a set of collision strengths which can account for all of the observations.

First, we have already noticed from equation (14) that the total Maxwellian-averaged collision strength \( Y_{31} + Y_{21} \) would have to be on the order of 20, and not 1 as computed by Rountree (1977). Does this large increase contradict earlier work? Is it physically reasonable? Although Rountree provides arguments justifying his neglect of potentially important processes, he did use a restricted set of basis wave functions for his collision calculations, and collisional excitations from the singlet terms of the ground configuration were neglected. Thus Rountree may not have included important contributors to the collision cross sections. A total Maxwellian-averaged collision strength of 20 near 6000 K, although it seems very large, may not in principle be inconsistent with other theoretical work (see Carlsson & Judge 1993, Appendix A5), with observations of other plasmas—for example, the airglow spectrum or laboratory plasmas (e.g., Meier 1991). This is because the solar chromospheric spectrum is sensitive mostly to cross sections at energies less than \( kT_\gamma \sim 6000 \) K, which is equivalent to 0.5 eV of the collision threshold energy, whereas the airglow spectrum is controlled largely by populations of nonthermal electrons at higher energies (>20 eV). Similarly, most laboratory experiments and theoretical work have been conducted at relatively high impact energies. However, in practice such a large increase must require the presence of large resonance structure in the \( \text{O I} \) electron-atom collision cross sections near threshold. It remains to be seen whether such resonance structure is actually present.

Second, equations (21) and (22) naturally lead to "critical electron densities" \( n_{\text{crit}} \), densities at which the collisional de-excitation rates equal to the net radiative decay rates. For \( n_e \gg n_{\text{crit}} \), the number density ratios therefore approach LTE values. With the radiative rates given above, we find critical densities for the 130.4 and 135.5 nm transitions of

\[
n_{\text{crit}}(130.4) \sim \frac{43}{Y_{32} + Y_{31}} \frac{P_{31}}{2.7 \times 10^{-4} \times 10^{11}} \text{ cm}^{-3}
\]

and

\[
n_{\text{crit}}(135.5) \sim \frac{5.7}{Y_{32} + 2Y_{21}} \times 10^{11} \text{ cm}^{-3},
\]

respectively. In order for the solutions of equations (21)–(26) to be consistent with the observed value of 1.5 for \( m_\nu \), the values of the collision strengths must be such that, at \( n_e \sim 10^{12} \text{ cm}^{-3} \), the 135.5 nm transition is formed at densities above \( n_{\text{crit}}(135.5) \), and the 130.4 nm transition is formed below \( n_{\text{crit}}(130.4) \). These inequalities lead to \( (Y_{32} + Y_{31}) < 43(2.7 \times 10^{-4})/P_{31} \) and \( (Y_{32} + 2Y_{21}) > 5.7 \).

Third, using the above limits on the collision rates in equations (21)–(26), the requirement that \( I_{31}/I_{21} \sim 14 \) (when \( n_e \sim 10^{11} \text{ cm}^{-3} \)) leads to even more stringent constraints on acceptable collision strengths. Setting \( y = Y_{31}/Y_{32} \), equations (21), (22), (25), and (26) can be rearranged to obtain a quadratic equation for \( y \) as a function of \( Y_{21} \), given the observed value \( I_{31}/I_{21} = 14 \) and \( Y_{31} + Y_{21} = 20 \). The condition that \( Y_{32} \) and \( y \) must both be real and positive yields \( Y_{32} < 100 \). With \( (Y_{32} + Y_{31}) < 43 \) (i.e., adopting our estimated value of \( P_{31} \)), the acceptable parameter space is limited to \( 5 < Y_{32} < 20 \), \( Y_{31} \sim 20 \), and \( 0.4 < Y_{21} < 0.6 \). Physically, we would expect \( y < 1 \), since the 2–1 transition is spin-forbidden and the cross sections will be smaller than for the 3–1 transition, since an exchange of impacting and bound electrons is needed for the (2–1) transition to occur with a reasonable probability.

Last, we note that with these collision strengths, equations (21) and (22) yield

\[
b_{32} = \frac{(n_3/n_2)}{(n_3^2/n_2^2)} < 1
\]

instead of \( \sim 0.3 \) as inferred from the eclipse data. Thus, we cannot arrive at a set of collision strengths which are fully consistent with the observations in these approximate calculations. This can only be done by increasing \( P_{31}^e \) by an order of magnitude, thereby increasing \( \sigma_{3} \) and reducing the population of level 3 relative to level 2 (eqs. [21] and [22]).

Table 2 lists the results of these approximate calculations and the detailed calculations discussed below. Model \( a \) was computed with the cross sections of Rountree (1977) (\( Y_{31} = 0, Y_{31} = 0.55, Y_{21} = 0.42 \)), and model \( b \) was computed using \( Y_{32} = 5, Y_{31} = 20, Y_{21} = 0.42 \).

### 4.3. Effects of Changes in Thermal Structure of the Chromosphere

We now ask how the intensities respond to changes in the thermal structure of the chromosphere, since it is not clear that the VAL models, based largely on fits to continuum data, represent the variety of conditions which may be encountered in the real chromosphere and which can greatly influence the UV line emission.

An alternative to increasing collision strengths to account for the strength of the 130.4 nm line is to increase the mean temperature of line formation. A temperature increase enhances the electron density by further ionizing hydrogen as well as increasing the fraction of the electrons with energies above the collision threshold. As long as hydrogen is mainly neutral, the combined density and temperature effects increase the \( \text{O I} \) collisional excitation rates by approximately a factor \( n_e^2 \), where \( n_e \) is the factor by which \( n_e \) is increased. By comparison, collisional de-excitation rates are increased by approximately a factor \( n_e \) only.

There is a limit to how far the line intensities can logically be increased by enhanced temperature, since all UV emission is increased by such an effect, at least in the context of mean chromospheric models. Thus, further increases in the \( \text{O I} \) line intensities require selective increases in \( \text{O I} \) collision strengths. In order to increase \( I_{31} \) sufficiently relative to \( I_{21} \), the enhanced \( \text{O I} \) collision strengths are needed mainly in the 3–1 and 3–2 transitions.

By way of illustration only, we consider a hypothetical case (model \( c \)) in which temperature is increased sufficiently \((\sim 800 \) K) to raise \( n_e \) by a factor of 3. This increases \( C_{13} \) and \( C_{12} \) by factors of 27 and all other collision rates by a factor of 3. In addition, we selectively increase collision strengths by a factor of 4 in the 3–1 transition and set \( Y_{32} = 10 \), and we also increase \( P_{31}^e \) by a factor of 4. The added increase in \( C_{31} \) increases the total line intensities, and the added increase in \( C_{23} \) makes \( C_{23} \) the dominant rate out of level 3. The main effects of increasing \( P_{31}^e \) are too strengthen \( I_{31} \), relative to \( I_{21} \) and to keep \( P_{31}^e A_{31} \). The latter condition ensures a dominance of radiative de-excitation for level 2 and keeps \( n_e \) below its LTE value relative to \( n_e \) (eq. [21]), which is consistent with the \( \text{O I} \) infrared line intensities observed at eclipse.
The results in Table 2 using the modified atmospheric structure (model c) are fully consistent with observations. They are, however, achieved using crude representations of the chromosphere and ad hoc values for collision rates. In the following, we consider a more exact treatment of the chromosphere and of the radiative transfer problem, and then discuss likely implications of our work for the structure of the solar chromosphere.

4.4. Detailed Transfer Calculations

We have performed detailed transfer calculations in solar models of varying levels of activity using the program MULTI (Carlsson 1986). These calculations are identical to those described by Carlsson & Judge (1993), with the following exceptions: (1) We merged the three levels of the ground term (\(^3P_{2,1,0}\)) into one "level"; (2) we removed photoexcitation by H Ly\(\beta\); (3) we made calculations using a version of the code written to include the effects of partial redistribution (Uitenbroek 1989), modified by K. Gayley and P. Judge (1993, unpublished) to improve convergence; and (4) we adjusted collision rates to investigate the influence of inaccurate collision cross sections on the line fluxes, guided by result from §§ 4.2 and 4.3.

These calculations were made with three bound level atoms plus the O \(\Pi\) continuum and 12 bound level plus continuum atoms equivalent to the model atom discussed by Carlsson & Judge (1993). Our results concerning frequency-integrated line intensities are shown in Figure 2. For each atom we computed the emergent intensities of the UV lines in models A, C, F, of Vernazza et al. (1981), model P of Basri et al. (1979, plage model), and model F2 of Machado, Avrett, & Loeser (1980) (bright flare model). Once again, our aim is to investigate whether we can account for the relative intensities of the lines of the triplet and quintet series in the quiet Sun, and also to account for the relative behavior of these lines as a function of activity.

![Figure 2](image)

Figure 2—Computed behavior of the frequency-integrated multiplet intensities near disk center (\(\mu = 0.88\)) of the 135.5 and 130.4 nm emission multiplets (with background continuum subtracted). The figure shows calculations for a given atomic model (marked a–f) as a function of the chromospheric model (A, C, F, P, and FL) (off-scale). The area marked "observed" is taken from Fig. 7 of Athay & Dere (1991): it roughly delineates the regions where the bulk of the emission-line intensities lie. The observed multiplet intensities were taken to be 3 and 1.3 times the intensity of the 130.6 and 135.5 nm lines shown in Figure 7 of Athay & Dere (1991).

Figures 2 and 3 contain data from the five sets of atomic parameters discussed above: models a–c were performed using the same three level atoms (plus O \(\Pi\) continuum) as discussed in §§ 4.2 and 4.3. This includes only the UV lines and collisional rates between the three levels and continuum. Model d, a 12 level atom, serves as a check on the three-level "collisional" models. It includes the same terms as used by Carlsson & Judge (1993), plus all collisional and radiative rates with the higher levels, with collisions between the \(^5S_2, ^3P_1\), and \(^2S-^3P\) terms as in model b. Model e is identical to model d, except that the pumping by H Ly\(\beta\) was included. Model f is essentially identical to the model considered by Carlsson & Judge (1993), the only differences being the merging of the three levels of the ground term and the treatment of partial redistribution. The adopted collision strengths for these various calculations are listed in Table 2.

From studying Tables 1 and 2, and Figure 3, it is clear that, of the atomic models considered, models b, d, and e are the only ones which can satisfy most of the observational constraints. Models a, c, and f have too little emission in the O \(\Pi\) UV lines in the mean Sun. None of the models is capable of explaining the observed behavior as a function of different levels of activity seen on the Sun: If the VAL models really are capable of representing different regions on the solar disk, then the calculations A–F should fall within the shaded area which contains the data from Figure 7 of Athay & Dere (1991). Apparently, the VAL model A produces intensities which are very rarely seen, at least in the HRTS Spacelab data. Also, the more active models F and P produce weaker intensities of the 130.4 nm lines than observed. This is because (1) collisional destruction dominates over photon escape where \(Y_{\text{Si}} \approx 20\), or (2) absorption in the background continuum (pre-dominantly Si I photoionization) dominates when \(Y_{\text{Si}} \approx 0.5\). Clearly, modifications to the atomic collision cross sections, radiative rates, and/or assumed chromospheric structure are required to account for the observed properties of the chromospheric O \(\Pi\) spectrum.

Profiles of the "mean" resonance transition at 130.4 nm from the calculations using atomic model b for the VAL model
C are shown in Figure 3. These profiles differ substantially from the observed profiles in that the observed profiles are self-reversed: our computations are not. A likely explanation is that the filling in of the line core due to emission at the base of the transition region has been overestimated, owing to inaccuracies in modeling this part of the solar atmosphere.

These calculations show that (1) it is possible to construct atomic models, based solely upon electron collisions, and also with the pumping by H Lyβ, which can account for some of the observed behavior of the O I lines in the near-infrared and ultraviolet, and (2) serious problems exist in more detailed comparisons between computed and observed behavior of UV lines. This indicates that the use of "mean" models for the chromosphere based primarily upon continuum data lacks some essential elements influencing the formation of chromospheric UV lines like those of O I.

5. Conclusions

We have argued that collisional excitation is consistent with the solar data only if (1) the critical collision cross sections at thermal energies are too small by an order of magnitude and/or (2) if the solar chromosphere has thermal properties which are substantially different from the mean thermal models of Avrett and colleagues. The resolution of the discrepancy between observed and computed behavior will certainly provide interesting results. If the first is the cause, then we have identified lines which are valuable (and almost unique) diagnostics of the chromospheric electron density; if the second, then models of the chromosphere lack physical processes probably related to the elusive heating mechanism(s).

If the chromospheric models are to blame for the disagreement, as indicated by our detailed calculations, then we have suggested that increases in the "mean" temperature of line formation are needed. Such apparently arbitrary increases in temperature, seemingly not permitted by "mean" modeling procedures, may have a more solid physical basis than initially expected. Given that the chromospheric heating mechanism is not known, and in the light of recent work on chromospheric dynamics (Carlsson & Stein 1992), it is possible that the mean models constructed by VAL are of limited use in studying excitation mechanisms. Specifically, if (for example) acoustic shocks are important as heating agents, the bulk of the emission from UV lines such as O I will be physically limited to hot postshock regions, in which the time-averaged emission-line spectrum differs very substantially from a model constructed in the manner described by Vernazza et al. (1981). As we have already noted, UV lines of high excitation energy relative to the thermal energy, like those of O I, are expected to be much more sensitive to the presence of high-temperature plasma than more traditional diagnostics of the chromosphere (e.g., Ca ii H and K, Mg ii h and k) formed in similar atmospheric regions.

If the atomic models are also at fault, then the choice of atomic parameters needed to fit the observed behavior is not arbitrary. In order to describe the observed behavior of the multiplets of O I as a function of the level of solar activity, the calculations fail even qualitatively if current estimates of atomic cross sections are used. It seems that something is missing from the calculations of the intensities of the O I lines. Whether the physical processes that are missing are related to the atomic physics of O I or to the assumption of a mean thermal structure for the chromosphere remains to be seen. In either event, it seems that further investigations would be extremely worthwhile, since, on the one hand, we will finally have a potentially good diagnostic of the chromospheric electron density (a central component in understanding the thermal structure of the chromosphere) and, on the other, we must discover new and valuable information on the departures of the solar chromosphere from the assumptions imposed by the mean modeling procedures.

At a bare minimum, calculations of the low-energy electron-atom cross sections should be done including the 2p4 3P, 2p4 1D, 2p4 1S, 2p3 3S, 2p3 3S, 2p3 3S terms, with higher terms needed to account for resonance structure in the cross sections. Cross sections are especially needed for collisions between the 2p3 3S term and 2p3 3S term, as well as between these and the terms of the 2p4 configuration. The close-coupling method would seem to be the best approach to use (e.g., Berrington & Kingston 1987).

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