KINEMATIC DIAGNOSTICS OF DISKS AROUND YOUNG STARS:
CO OVERTONE EMISSION FROM WL 16 AND 1548C27

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ABSTRACT

We report high spectral resolution observations of the CO vibrational overtone emission from the young stellar object 1548C27; our observations include both the $v = 2-0$ and $v = 5-3$ band head regions. These data and similar observations of the young stellar object WL 16, reported in a previous contribution to this journal, provide some of the most compelling evidence to date for the existence of inner disks around young stars. We describe the simple procedure that we use to synthesize band head emission from disks including the effect of thermal dissociation of CO and non-LTE excitation of the vibrational levels. Using this spectral synthesis procedure to extract the kinematics and physical properties of the emitting gas from the overtone data, we show how these high signal-to-noise ratio data are also powerful probes of the stellar and inner disk properties of these systems. Our modeling is consistent with the identification of WL 16 and 1548C27 as Herbig AeBe stars with stellar masses of approximately 2 and 4 $M_\odot$, respectively. Thus, the kinematic signature of rotating disks in the overtone spectra of these sources provides strong support for the role of accretion disks in the formation of intermediate-mass stars. For both WL 16 and 1548C27, we interpret our modeling results as indicating that the overtone emission arises from a temperature inversion region in the inner disk atmosphere. We also find evidence for suprathermally broadened lines and are able to place useful constraints on the radial temperature and column density distributions of the CO line formation region of the disk atmosphere. Given these deduced properties, we discuss the constraints that our observations place on the physical processes responsible for the overtone emission in these sources.

Subject headings: accretion, accretion disks — circumstellar matter — infrared: stars — stars: individual (1548C27, WL 16) — stars: pre-main-sequence

1. INTRODUCTION

Circumstellar disks play a fundamental role in the formation of stars and planets. Stars are believed to build up a substantial fraction of their main-sequence mass by accretion through circumstellar disks, and these disks constitute the reservoirs of mass for the formation of both planets and close stellar companions. Various pieces of observational evidence assembled over the past decade have been used to argue for the presence of accretion disks around young stars. These include the spectral energy distributions of young stars (e.g., Adams, Lada, & Shu 1987), the asymmetry of their optical forbidden line profiles (e.g., Edwards et al. 1987), their millimeter and submillimeter line and continuum emission (cf. Ohashi & Hayashi 1995; see also Koerner & Sargent 1995, and references therein; Keene & Masson 1990; Lay et al. 1994), and the cross-correlation line widths of optical and near-infrared lines of FU Ori objects (e.g., Hartmann & Kenyon 1987a, b). More recently, direct kinematic evidence for the existence of circumstellar disks has been obtained from high spectral resolution observations of the CO overtone emission of young stars (Carr et al. 1993; Chandler et al. 1993; Chandler, Carlstrom, & Scoville 1995). Since the CO overtone lines arise from the inner regions of these disks, detailed studies of this emission can be used to determine the properties of disks within several tens of stellar radii of young stars.

This inner disk region (< 1 AU) is of particular interest as it is believed to be dynamically active, giving rise to energetic inflows and outflows of gas which remove substantial amounts of mass and angular momentum from the system and regulate the rotation periods of young stars (Shu et al. 1994; cf. Edwards et al. 1993; and Bouvier et al. 1993). Although regions of this size cannot be imaged directly with present observational techniques, they can be studied with high resolution near-infrared spectroscopy of lines that selectively probe the physical conditions of the region. The exceptional stability of the CO molecule makes it a versatile probe of this region, and the excitation temperatures and
critical densities of the overtones near 2.35 μm are particularly well suited to the study of the high densities (up to \(10^5 \text{ cm}^{-3}\)) and temperatures (1000–6000 K) that characterize the inner regions of young stellar objects (YSOs), e.g., as first demonstrated by Scoville et al. (1983) for the BN object.

With the advent of a new generation of sensitive near-infrared spectrometers (e.g., CSHELL; Tokunaga et al. 1990), it is now possible to study low- and intermediate-mass YSOs in the near-infrared at high spectral resolution. Using these spectrometers, even faint, low-luminosity (\(\sim 10 L_\odot\)) YSOs can now be observed with a resolution up to 40,000 and at a signal-to-noise ratio that was previously only feasible for bright, high-luminosity sources (e.g., Mitchell et al. 1989, 1990) and FU Ori objects (Hartmann & Kenyon 1987a, b). We have previously reported spectroscopic observations of the \(v = 2–0\) bandhead region of the embedded sources SVS 13 (Carr & Tokunaga 1992) and WL 16 (Carr et al. 1993). Similar observations have been reported by Chandler et al. (1993) for several additional sources. Our recent observations of the \(v = 2–0\) and \(v = 5–3\) band head regions of the Herbig AeBe star 1548C27 are presented here. A striking aspect of these studies is that the band head often shows the characteristic shape of emission from a rotating disk. The WL 16 and 1548C27 spectra provide particularly good examples and clearly demonstrate the existence of inner disks in these systems.

Winds from low-luminosity YSOs may also produce substantial CO overtone emission (Carr 1989; Chandler et al. 1995; Safier, Martin, & Königl 1996). Calculations of the thermal and chemical structure of winds from these sources show that CO is a robust chemical product and that the inner wind region is warm enough to excite overtone emission (\(\sim 2000 \text{ K}\); Ruden, Glassgold, & Shu 1990). Since the densities of winds rapidly drop below the critical density for the overtone bands (see Appendix B), a significant wind contribution to the overtone emission is expected to be restricted to YSOs with large mass-loss rates. It is less clear whether a large CO abundance can be expected to survive possible photodissociation by the stellar radiation field of the earlier stellar spectral type sources in the present study. In any case, overtone emission from winds is expected to show a line asymmetry due to the occultation of the receding flow by optically thick inner disks. Since the wind is likely to be cooler than the star and to have a temperature similar to that of the inner disk region (Ruden et al. 1990; Chandler et al. 1995), the absorption of stellar and inner disk continuum photons by the wind will introduce an additional line asymmetry. The absence of these asymmetries in the overtone spectra in the present study suggests that the emission from these sources is not dominated by the contribution from a wind.

In this paper, we demonstrate that our overtone observations of WL16 and 1548C27 are well explained by a disk interpretation for the emission. We develop and implement a simple spectral synthesis technique in order to illustrate the constraints that these high-resolution, high signal-to-noise ratio data place on both basic stellar properties (e.g., masses and radii) as well as detailed disk properties (e.g., radial variations of temperature, column density, and intrinsic line width). In previous publications, we used a spectral synthesis procedure developed by Carr et al. (1993), which generalized previous work addressing isolated lines (e.g., Huang 1972, Smak 1981) to the band head. The synthesis problem was formulated in terms of an effective emissivity as a function of disk radius; inner and outer disk radii were introduced as fitting parameters. In addition, the line-emitting region was assumed to correspond to the entire vertical disk column density and to be optically thin in the continuum (see also Carr 1989; Chandler et al. 1995).

In the present application, we develop a more general synthesis method formulated in terms of the physical properties of the CO line formation region of the disk atmosphere. We use a two-layer approximation to this atmospheric region: the continuum forms in the lower layer, and the CO lines form in the upper layer. Thus, the two-layer model approximates the situation of an optically thick disk with a temperature inversion in the upper disk atmosphere (cf. Calvet et al. 1991). We adopt radial power-law distributions of temperature in both layers and of column density in the upper layer. Our analysis then yields the radial variation of physical conditions that must be explained by a more fundamental future theory. We avoid the introduction of arbitrary inner and outer disk radii by explicitly considering the abundance and excitation of CO. In addition to placing tight constraints on the temperature and column density distributions of the emitting layer, we find that the intrinsic (local) line width of the emitting gas is significantly broader than thermal Doppler values, probably symptomatic of supersonic magnetohydrodynamic (MHD) turbulence.

We begin, in § 2, by defining the two-layer disk atmosphere model and outlining the procedure used to synthesize the CO spectrum of a rotating disk from known level populations. In the next three sections, we discuss the applicability of LTE level populations and the circumstances under which departures from LTE become important. In § 3, we model the \(v = 2–0\) band head spectrum of WL 16 assuming LTE in order to illustrate the parameter dependence of the two-layer model and extract the basic physical properties of the emitting gas. In § 4, we develop a non-LTE theory for the level populations and draw attention to the diagnostic power of jointly fitting both the \(v = 2–0\) and \(v = 5–3\) band head regions. The utility of modeling these two spectral regions is demonstrated in § 5, where we present our observations of the Herbig AeBe star 1548C27 in these spectral regions and use our spectral synthesis procedure to deduce the physical properties of the CO line formation region in this source. Given these deduced properties, in § 6 we discuss the evolutionary status of WL 16 and 1548C27 and comment on the physical processes that are likely to be responsible for the CO emission in these sources.

### 2. A TWO-LAYER DISK ATMOSPHERE MODEL

We model the CO overtone emission as arising from the inner (\(\lesssim 1 \text{ AU}\)) region of a geometrically thin disk in Keplerian rotation about a star of mass \(M_*\). We treat the atmosphere as two adjacent layers, each of which is vertically homogeneous. The overtone lines form in the upper ("line") layer, which extends vertically over a thermal scale height and has negligible continuum opacity. The latter assumption is checked posteriori by computing the \(H^+\) continuous opacity. There is no scattering in the upper layer, and the only line opacities are those of the CO lines themselves. The lower ("continuum") layer is optically thick in both the overtone lines and in the 2.3 μm continuum.
We adopt power-law radial temperature distributions for the continuum and line layers: $T_c(r) = T_{c0}(r/r_o)^{-\alpha}$ and $T_l(r) = T_{l0}(r/r_o)^{-\beta}$, where $r_o$ is a fiducial disk radius. The line layer is additionally described by a vertical mass column density distribution that also varies as a power law with radius: $\sigma_l(r) = \sigma_{l0}(r/r_o)^{-\gamma}$. We assume solar abundances of gaseous carbon and oxygen in calculating the chemical equilibrium abundance of CO in this layer. The additional parameters required to synthesize the absolute strength and shape of the spectrum near the CO band heads are the local line profile function ($\phi_l$), the distance and $K$-band extinction to the system ($d$, $A_K$), the inclination of the rotation axis of the disk to the line of sight ($i$), and the stellar luminosity and radius ($L_*$, $R_*$). The last two quantities determine the stellar contribution to the 2.3 $\mu$m continuum which is assumed to be blackbody. Of these, the stellar and system parameters—$M_*$, $L_*$, $R_*$, $d$, $A_K$, and $i$—are relatively well constrained by the gross properties of the CO overtone emission and other observational data (see, e.g., Carr et al. 1993). The disk properties—$T_i$, $T_l$, $\sigma_l$, and $\phi_l$—which are the quantities of primary interest here, are well constrained by the type of detailed modeling that we discuss below.

The synthesis procedure allows the introduction of inner and outer disk truncation radii ($R_{in}$, $R_{out}$) as model parameters and explicitly considers the variation of CO abundance and excitation with disk radius. Both disk truncation and these latter processes may limit the radial extent of the CO emission. For example, disks may be physically truncated from the inside at distances of several stellar radii by stellar magnetospheres, as suggested by a variety of recent observational and theoretical results (cf. Bertout, Basri, & Bouvier 1988; Königl 1991; Bouvier et al. 1993; Edwards et al. 1993; Hartmann, Hewett, & Calvet 1994; Shu et al. 1994). In addition, the dissociation of CO at temperatures $T_i \gtrsim 5000$ K will naturally introduce an effective inner radius to the CO emission; in contrast to the case of physical disk truncation, contributions to the 2.3 $\mu$m continuum may arise from within the CO dissociation radius. Disks may also be physically truncated from the outside by the formation of giant planets or the presence of nearby stellar companions (e.g., Lin & Papaloizou 1993; Mathieu 1994). Since the YSOs in the present study are not close binary systems and giant planets are believed to form outside the water-ice condensation radius (at $T \sim 150$ K), well beyond the region responsible for the CO overtone emission, we generally set the outer radius of the disk $R_{out}$ at infinity. In this case, effective outer radii for the CO emission are produced by either dust formation or departures from LTE populations (see §§ 4 and 6).

To synthesize the band head emission, the disk surface is first divided radially and azimuthally into finite elements. For each element, we calculate the emergent intensity 

$$I_v = B_v(T) e^{-tv/\mu} + S_v(1 - e^{-tv/\mu}),$$

(1)

and then add a frequency shift appropriate to the projected rotational velocity of the element. In the equation above, $B_v(T)$ is the Planck function at the temperature of the continuum layer, $\mu = \cos i$, and $S_v$ and $\tau_v$ are the source function and total vertical optical depth of the element. Both the source function and optical depth of the upper layer may include contributions from discrete lines that overlap in frequency. The line layer is assumed to be spatially thin so that the properties of the disk atmosphere are approximately constant along the inclined ray path. Having calculated the intensity from each element, the contributions of all elements not shadowed by the star are then summed to obtain the disk contribution to the spectrum. Finally, we add the stellar (blackbody continuum) contribution to the spectrum, accounting for the shadowing of the star by the disk. The resulting spectrum is extinguished by the adopted value of $A_K$ to produce the final synthesized spectrum.

3. LTE Model

In order to illustrate the dependence of the band head emission on the parameters of the two-layer model, in this section we apply the above spectral synthesis procedure to our $v = 2-0$ data for WL 16 assuming LTE populations. The rotational and vibrational constants and the energy levels are taken from Mantz et al. (1975). The use of LTE level populations allows us to extract the basic kinematics of the emitting gas from the band head spectrum in a simple way, i.e., without having to address the additional complication of non-LTE effects which generally bear weakly on the deduced kinematics. The LTE results also form a basis of comparison against which we can measure the magnitude of the non-LTE corrections that are computed in §§ 4 and 5.

3.1. Spectral Synthesis of the $v = 2-0$ Band Head Region

The spectral region covered by our observations of the $v = 2-0$ band head emission from WL 16 (see Carr et al. 1993) includes the R39–R62 lines which approach and retreat from the band head located at a rest wavelength of 2.29353 $\mu$m (Fig. 1). As described by Carr et al. (1993), the band head emission from WL 16 displays the characteristic shape of emission from a rotating disk: a blue wing, a shoulder, and an intensity peak redward of the rest wavelength of the band head. That this is the general shape expected for band head emission from a rotating disk can be demonstrated.

![Fig. 1.](image-url) The $v = 2-0$ band head emission from WL 16 which shows the characteristic shape of band head emission from a rotating disk: a blue wing, a shoulder, and an intensity peak redward of the band head. The spectral region covered by our observations includes the R39–R62 lines.
strated by imagining the convolution of the profile of an isolated line from a rotating disk with the distribution of the lines near the $v = 2-0$ band head (see also Carr 1995).

This convolution procedure is shown schematically in Figure 2, where the double-horned rotational broadening function (heavy line) is the profile of an isolated line from an inclined Keplerian disk that has a monotonically decreasing intensity distribution between an inner and outer radius. The rest spectrum of lines near the band head (light line) includes line blending due to local line broadening sources and assumes LTE level populations at a typical inner disk temperature. The blue wing is built up as the wing of the isolated line profile moves past the band head (Fig. 2a). A shoulder is formed when one of the horns of the isolated line profile reaches the band head (Fig. 2b). As the horn moves past the band head it encounters other lines longward of the band head; the emission consequently remains approximately constant, producing the shoulder. The red peak forms when the other horn of the isolated line profile reaches the band head (Fig. 2c). When both horns have moved past the band head, the overlap of the isolated line profile with the lines redward of the band head produces the slow descent redward of the red peak (Fig. 2d). In the observed band head spectrum, the velocity separation between the blue shoulder and the red peak is approximately twice the projected rotational velocity at the outermost radius of the CO emission, and the maximum observed velocity in the blue wing is the projected velocity of CO at the innermost radius of the emission.

The convolution described above is essentially the procedure used by Carr et al. (1993) to estimate the stellar mass and radius of WL 16. While suitable for this purpose, this procedure is approximate in that the relative strengths of the lines near the band head are assumed to be independent of disk radius. The dependence of the emergent spectrum on radial variations of temperature and column density enters only in the rotational broadening function which is effectively the isolated profile of an "average" line at the band head. The present modeling procedure removes both these approximations and allows us to place physical constraints on the properties of the inner disk region.

### 3.2. Application to WL 16

WL 16 is a low-luminosity YSO in the $\rho$ Oph cloud ($d = 160$ pc). Although originally classified as an embedded (Class I) object (Wilking, Lada, & Young 1989), the measurement of strong 3.3 \(\mu\)m aromatic hydrocarbon emission (Tokunaga et al. 1991) and a preliminary analysis of our observations of the $v = 2-0$ band head emission (Carr et al. 1993) imply an earlier spectral type and higher mass for WL 16 than is usually ascribed to Class I objects (e.g., Adams et al. 1987). The bolometric luminosity of WL 16 has been estimated as 18–22 $L_\odot$ (Wilking et al. 1989; Cabrit & André 1991). Carr et al. (1993) estimate the extinction to WL 16 to be $A_\lambda = 2$ based on the observed $(J-H)$ color of WL 16 and assuming that its intrinsic color is similar to that of very active pre-main-sequence stellar systems.

The analysis by Carr et al. (1993) of the $v = 2-0$ band head emission revealed that parameters of $M_\star \approx 2.5\, M_\odot$, $R_\star \approx R_\text{in} \approx 5\, R_\odot$, and $i \approx 60^\circ$ are required to fit both the observed velocity extent and flux of the band head emission. As discussed by these authors, an inclination of $i = 60^\circ$ represents a compromise between the projected emitting area required to fit the band head flux and the projected rotational velocities required to fit the velocity extent of the emission. Several general considerations constrain the stellar mass to be $\sim 2.5\, M_\odot$ given the assumed extinction to the source. On the one hand, the stellar mass cannot much exceed $\sim 2.5\, M_\odot$ for any pre-main-sequence age given the bolometric luminosity of the system. On the other hand, a smaller stellar mass would require that the emission arise closer to the star to maintain the velocity extent of the band head; an origin at smaller radii will tend to reduce the strength of the emission since, for a given band head shape, the emission scales as $R_\text{in}^2$. The impact of the reduced emitting area on the band head flux can be offset only to a limited extent by adopting a higher temperature for the emitting gas, since collisional dissociation of CO becomes important above temperatures of $\sim 5000$ K.

To illustrate the additional information that can be extracted from the WL 16 $v = 2-0$ data with the present spectral synthesis method, we reanalyze the data using the two-layer LTE model described above. We assume that nearly all the system luminosity arises from the star and adopt $d = 160$ pc, $i = 60^\circ$, $M_\star = 2.5\, M_\odot$, $A_\lambda = 2$, and $L_\star = 22\, L_\odot$; the remaining parameters which describe the properties of the disk—$T_\odot$, $P_\tau$, $T_\text{in}$, $P_\text{in}$, $\sigma_0$, $q_1$, and $\phi_\text{in}$—are varied to fit the observed band head spectrum. For the continuum layer, we adopt a radial temperature dependence in which $p_\tau = \frac{\alpha}{\tau}$, as expected for an active viscous or passive reprocessing disk. The adopted disk continuum temperature corresponds to a disk luminosity of

$$L_{\odot} = 2 \int_{R_{\text{in}}}^{\infty} \sigma T_{\odot}^4 2\pi r \, dr = 4\pi \sigma \left(\frac{T_0}{R_{\text{in}}}\right)^4 T_{\odot}^4$$

where $T_0$ is the temperature at the inner disk radius.
and implies a disk mass accretion rate which can be obtained from the relation

$$\frac{GM_\star M_D}{2R_{\text{in}}^3} = L_D - \mathcal{F}(1-\alpha)L_\star,$$

where the second term accounts for passive reprocessing of incident stellar radiation (Adams & Shu 1986). The quantity \(\mathcal{F}\) is the fraction of the stellar luminosity \(L_\star\) that is intercepted by the disk \((\mathcal{F} = \frac{1}{2}\) if \(R_{\text{in}} = R_\star\)) and \(\alpha\) is the effective scattering albedo. Values for the other disk parameters are chosen to reproduce both the strength of the 2.3 \(\mu\)m continuum and the contrast of the band head above the continuum. More detailed spectral features such as the shape of the blue wing, the sharpness of the shoulder, the location of the red peak, and the spectral shape redward of the peak constrain \(p_1\) and \(q_1\).

With these assumptions, we find that we can fit the strength of the 2.3 \(\mu\)m continuum if the \(M_\star = 2.5\ M_\odot\), \(L_\star = 22\ L_\odot\) star lies near the stellar birthline (\(R_\star = 5.0\ R_\odot\); Palla & Stahler 1993) and the continuum layer has a temperature of \(T_\odot = 3000\ K\) at \(r_\odot = 5\ R_\odot\). In order to fit the velocity extent and strength of the blue wing, \(T_i\) must be \(\lesssim 5000\ K\) at \(5\ R_\odot\). These high temperatures correspond to large line source functions without appreciable dissociation of CO and thereby maximize the band head flux from these radii. The location of the shoulder and red peak indicates that most of emission arises from within \(30\ R_\odot\); the sharpness of these features indicates that the emission must decrease relatively sharply beyond this radius. If the rapid decrease in the emission beyond \(30\ R_\odot\) is due, in part, to the reduced sensitivity of the \(r = 2\ R_\odot\) band head lines to gas below \(1500\ K\) (see below), the constraints on the layer temperature are roughly \(T_i \approx 1500\ K\) at \(\sim 30\ R_\odot\) and \(T_i \approx 5000\ K\) at \(\sim 5\ R_\odot\), which correspond to a temperature profile with \(T_i \approx 5000\ K\) and \(p_1 = \frac{1}{4}\). Flatter temperature profiles that roughly preserve the locations of the shoulder and red peak and velocity extent of the blue wing have lower temperatures within \(\sim 30\ R_\odot\) and require somewhat steeper column density distributions to fit the sharpness of the shoulder and red peak.

Given these restrictive requirements, it is very difficult to fit the strength and shape of the band head with thermal Doppler profiles. Within the context of the two-layer LTE model, we can obtain a good fit only within a very restricted range of parameters about \(\sigma_{10} = 5000\ g\ \text{cm}^{-2}\), \(q_1 = 4.5\), \(T_\odot = 6000\ K\), \(p_1 = \frac{1}{3}\), and \(R_{\text{in}} = R_\star\) (model L12; see Fig. 3a and Table 1). Since the wings of Gaussian profiles are weak, it is difficult to emit much flux between the line cores without these large column densities and the higher line layer temperature. These column densities are, in fact, so large as to violate our assumption of a negligible 2.3 \(\mu\)m continuum optical depth \(\tau\) in the line layer: the opacity due to H\(^-\) alone corresponds to \(\tau = 700\) at \(1.1\ R_\star\) and \(\tau > 1\) within \(1.8\ R_\star\).

Including this continuum optical depth in the model would significantly worsen the fit. In Figure 3a, we are able to fit the strength and extent of the blue wing and red peak because the high column density close to the star yields a sufficient abundance of CO to keep the optical depth of the overtone lines appreciable even at \(T_i = 6000\ K\). If we were to include the line layer continuum optical depth in the model, the high optical depth close to the star would significantly reduce the contrast of the line emission from this region and result in an excessively large 2.3 \(\mu\)m continuum as well as both an inadequate velocity extent and overall strength for the band head emission (see Fig. 3d). In general, somewhat smaller line layer column densities can produce the same band head flux if the line layer represents the entire disk column density within a few stellar radii, i.e., if the continuum layer is absent. In the present case, however, even the absence of a continuum layer cannot significantly reduce the required column density to values corresponding to reasonable H\(^-\) opacities within \(\sim 1.6\ R_\star\). We can conclude that, given the adopted extinction to the source, thermally broadened lines cannot account for the CO emission from WL 16.

Alternatively, line layers with column densities of \(\sim 100\ g\ \text{cm}^{-2}\) at \(10\ R_\odot\) and underlying optically thick continuum layers can produce good fits to the strength of both the band head and the 2.3 \(\mu\)m continuum if the intrinsic line profiles are broader than thermal. With Lorentzian lines, these fits require systematically lower column densities than in the case of thermal Doppler lines because more flux can be emitted in the stronger wings of Lorentzian lines. Figure 3b (model L14) shows the fit obtained with LTE populations, Lorentzian line profiles with a constant line width (half-width at half-maximum of \(\gamma = 0.5\ \text{km s}^{-1}\), \(\sigma_{10} = 370\ g\ \text{cm}^{-2}\), \(q_1 = 3.8\), \(p_1 = 0.78\), and all other parameters the same as in model L12 (see Table 1). At the lower line layer column densities of this model, the CO dissociation radius lies at a somewhat larger distance than in model L12: in model L14 the abundance and optical depth in the overtone lines become substantial beyond \(1.2\ R_\star\). Beyond this distance, the continuum optical depth due to H\(^-\) is negligible, consistent with our assumptions. We are able to fit the extent of the blue wing despite the larger dissociation radius because of the stronger wing of the Lorentzian line. With these parameters, the shadowing of the star by the disk reduces the stellar contribution to the 2.3 \(\mu\)m continuum to \(\sim 0.29\ Jy\). The star, in turn, shadows the part of the disk that contributes line emission at low projected velocities.

<table>
<thead>
<tr>
<th>Model*</th>
<th>(R_\star) ((R_\odot))</th>
<th>(T_\odot) (K)</th>
<th>(T_i) (K)</th>
<th>(p_1)</th>
<th>(q_1)</th>
<th>(\alpha)</th>
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<td>0.75</td>
<td>6000</td>
<td>0.75</td>
<td>5000</td>
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<tr>
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<td>3000</td>
<td>0.75</td>
<td>6000</td>
<td>0.78</td>
<td>370</td>
</tr>
<tr>
<td>L15</td>
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<td>3000</td>
<td>0.75</td>
<td>5810</td>
<td>0.78</td>
<td>310</td>
</tr>
<tr>
<td>U2</td>
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<td>3000</td>
<td>0.75</td>
<td>5500</td>
<td>0.75</td>
<td>200</td>
</tr>
<tr>
<td>V2</td>
<td>2.08</td>
<td>3000</td>
<td>0.75</td>
<td>5900</td>
<td>0.78</td>
<td>280</td>
</tr>
</tbody>
</table>

Note. — \(M_\star = 2.5\ M_\odot\), \(L_\star = 22\ L_\odot\), \(d = 160\ \text{pc}\), \(\theta = 20^\circ\), \(i = 60^\circ\), \(R_\star = 1\ R_\odot\), and \(r_\odot = 5\ R_\odot\).

* Models L12, L14, and L15 use LTE level populations. Models U2 and V2 use non-LTE level populations.
enhancing the sharpness of the shoulder. The (partially shadowed) disk continuum layer contributes an additional \( \sim 0.12 \) Jy to the 2.3 \( \mu \)m continuum and \( \sim 2 L_\odot \) to the system luminosity (eq. [2]).

Because the lines in this model are much less optically thick than in the previous case, there is a larger region of parameter space within which we can find acceptable model fits. For example, a similar fit is possible with a column density distribution as flat as \( q_1 = 2 \) (with \( \sigma_{10} = 90 \) g cm\(^{-2} \), \( T_{10} = 5500 \) K, and \( p_1 = \frac{1}{3} \)). Temperature distributions as flat as \( p_1 = 0.5 \) can also fit the data if the column density distribution is steeper (e.g., \( \sigma_{10} = 400 \) g cm\(^{-2} \), \( q_1 = 4.5 \), and \( T_{10} = 4700 \) K). Larger Lorentzian line widths further decrease the column density required to produce the strength of the band head and tend to flatten the slope of the spectrum redward of the peak. Using somewhat steeper temperature gradients to compensate for this flattening will produce a good fit. For example, for \( \gamma = 2.5 \) km s\(^{-1} \), adopting \( T_{10} = 5900 \) K, \( p_1 = 0.78 \), \( \sigma_{10} = 75 \) g cm\(^{-2} \), and \( q_1 = 3 \) will also fit the data.

WL 16 is unusual among low-luminosity YSOs in showing strong aromatic hydrocarbon emission features (e.g., Hanner, Tokunaga, & Geballe 1992), a property that has been interpreted as evidence for an early stellar spectral type. High resolution \( H \)-band spectroscopy of WL 16 is also consistent with an early spectral type (see § 3.4). The system luminosity of 22 \( L_\odot \) and deduced stellar mass restrict the evolutionary state of such a star. In order to investigate the possibility of an earlier stellar spectral type for WL 16, we consider whether WL 16 is an A-star with mass \( M_\ast = 2.5 M_\odot \), radius \( R_\ast = 2.08 R_\odot \), and luminosity \( L_\ast = 22 L_\odot \).

Figure 3c (model L15) shows the synthesized spectrum obtained with these stellar parameters, Lorentzian line profiles, and the following set of disk parameters: \( r_0 = 5 R_\odot \), \( T_{10} = 3000 \) K, \( T_{15} = 5810 \) K, \( p_1 = 0.78 \), \( \sigma_{10} = 310 \) g cm\(^{-2} \), \( q_1 = 3.7 \), and \( \gamma = 0.5 \) km s\(^{-1} \) (see Table I). The star contributes \( \sim 0.1 \) Jy of the continuum; the disk extends up to the surface of the star and contributes \( \sim 0.2 \) Jy to the 2.3 \( \mu \)m continuum and \( \sim 4 L_\odot \) to the system luminosity. The (inner) dissociation radius is effectively \( \sim 5 R_\odot \), as in model L14. It would be difficult to distinguish between this dissociation radius and a disk truncation radius \( R_{\text{in}} \leq 5 R_\odot \) because the additional (unshadowed) stellar contribution to the continuum in the latter case is comparable to the contribution to the continuum from the truncated disk region. Although the fit to the band head emission is comparable to
that obtained with model L14, the continuum is reduced by a third due to the earlier stellar spectral type.4

3.3. Discussion of WL 16 Results

The detailed modeling procedure presented above confirms the result obtained by Carr et al. (1993) that the parameters \( M_N = 2.5 \ M_\odot \), \( R_N \leq R_{\text{in}} = 5 \ R_\odot \), and \( i = 60^\circ \) provide a good fit to both the observed velocity extent and flux of the band head emission. We are able to explain the inner disk radius of 5 \( R_\odot \) as either a CO dissociation radius (model L15) or a physical truncation radius. The disk may either extend up to the surface of a 5 \( R_\odot \) star on the stellar birthline (model L14) or it may be truncated due to the presence of a stellar magnetosphere that extends to a similar distance (a variation of model L15).

By fitting the detailed shape of the band head emission, e.g., the location and sharpness of the shoulder, we find that the \( v = 2-0 \) emission must be truncated relatively rapidly beyond 30 \( R_\odot \). This effect can arise for a variety of reasons and probably does not indicate a physical outer edge to the disk. For example, at temperatures below 1500 K, the lines that make up the \( v = 2-0 \) band head are not well excited even if the level populations are in LTE. In the models shown, a steep column density gradient is used to emphasize this outer radius to the emission. Flatter column density distributions are possible if dust is present in the upper disk atmosphere (i.e., the line layer) below 1500 K. Because the CO lines and the 2.3 \( \mu \text{m} \) continuum would then form in the same layer this would eliminate the contrast of the CO emission above the continuum beyond this radius, creating an effective outer radius to the emission. The role of non-LTE effects in creating effective outer radii to the overtone emission is discussed in the next section.

We find that suprathermal line widths are needed to fit the strength and shape of the band head emission. The band head emission is well fit with Lorentzian lines characterized by \( \gamma = 0.5-2.5 \ \text{km s}^{-1} \). In §6 we offer a physical interpretation for these enhanced line widths. We also find that the required line layer column densities (\( \lesssim 100 \ \text{g cm}^{-2} \) at 10 \( R_\odot \)) are much lower than the values (\( \lesssim 10^5 \ \text{g cm}^{-2} \)) expected by extrapolating models of the "minimum solar nebula" to these radii. We interpret this result as indicating that the line layer represents a temperature inversion region in the upper atmosphere of a disk that has a much larger total column density, and note that this result also indicates the sensitivity of the CO overtone lines to relatively small amounts of circumstellar matter. Similar column densities may be present in the continuum "gaps" or "inner holes" of YSO disks (cf. Strom et al. 1989; Hillenbrand et al. 1992; Lada & Adams 1992; Marsh & Mahoney 1992; Mathieu 1994).

The disk continuum temperature used in all fits corresponds to a disk mass accretion rate that is consistent with zero but is in all cases \( \lesssim 2 \times 10^{-7} \ M_\odot \ \text{yr}^{-1} \). The 2 \( L_\odot \) radiated by the disk in model L14 can be entirely attributed to reprocessed starlight if the effective albedo is 0.68 (eq. [3] with \( \mathcal{F} = \frac{1}{2} \)). Similarly, the 4 \( L_\odot \) disk luminosity in model L15 could be due to reprocessed starlight alone if the effective disk albedo is 0.2. Alternatively, the disk luminosity in both models can be attributed to energy release by disk accretion at \( \gtrsim 2 \times 10^{-7} \ M_\odot \ \text{yr}^{-1} \) if the effective albedo is unity. This accretion rate is low compared to typical values for both Class I objects (e.g., Adams et al. 1987) and Group II Herbig Ae stars (Hillenbrand et al. 1992), consistent with the lack of evidence for either an outflow from WL 16 or other phenomena commonly associated with active Class I sources such as SSV 13 and L1551 IR55.

The synthetic spectra shown in this section indicate the degree to which modeling the \( v = 2-0 \) band head spectral region assuming LTE level populations allows us to constrain the physical properties of the disk. We find that there is some trade-off between the line layer temperature and column density profiles required to obtain a good fit. This nonuniqueness is due to both the assumption of LTE and the choice of the \( v = 2-0 \) band head as the region to model. Since the assumption of LTE fixes the source function, this leaves the optical depth as the only variable in the radiative transfer that is affected by a given temperature and column density. Since the lines in the \( v = 2-0 \) band head region are excited at similar temperatures and moreover probe a restricted range in optical depth at a given temperature and column density, the trade-off between temperature and column density in the optical depth will work almost equally well for all the lines in the spectral region, producing a limited effect on the shape of the band head. In comparison, the consideration of non-LTE effects and spectral fitting of multiple band head regions provides powerful additional diagnostics of the detailed physical properties of the inner disk. Following a brief discussion of the evolutionary status of WL 16 (§ 3.4), we investigate these additional diagnostics. In §4, we construct a non-LTE model, and in §5 we apply the model to \( v = 2-0 \) and \( v = 5-3 \) observations of another YSO, 1548C27, in order to demonstrate the constraints that these additional considerations allow us to place on the physical properties of inner YSO disks.

3.4. WL 16: On the Stellar Birthline or Main-Sequence Star?

In fitting the \( v = 2-0 \) band head emission, we find that we can model the stellar component of WL 16 as either a G star on the birthline or an A star near the main sequence. The presence of aromatic hydrocarbon features in the spectrum of WL 16 and the fit to the 2.3 \( \mu \text{m} \) continuum provide other potential discriminants between these two possibilities. Of the various aromatic hydrocarbon features observed in the spectrum of WL 16, the 3.3 \( \mu \text{m} \) feature is the strongest indicator of an early spectral type because it is excited by photons shortward of 3000 \( \AA \) (Schutte, Tielens, & Allamandola 1993). We can compare the efficiency required to excite the observed 3.3 \( \mu \text{m} \) feature with UV radiation from either stellar spectral type with the efficiencies measured for other astronomical sources. For the models considered above, the fraction of the stellar (blackbody) luminosity that is emitted between 912 \( \AA \) and 3000 \( \AA \) is 1% for \( T_\ast = 5600 \ K \) (model L14) and 6% for \( T_\ast = 8700 \ K \) (model L15). For \( d = 160 \ pc \), the strength of the 3.3 \( \mu \text{m} \) feature in WL 16 is \( 2 \times 10^{-3} \ L_\odot \) (cf. Tokunaga et al. 1991) where we have assumed an extinction correction at 3.3 \( \mu \text{m} \) appropriate for \( A_V = 2 \). This luminosity is 0.9% of the available UV radiation for a 22 \( L_\odot \) star with \( T_\ast = 5600 \ K \) and 0.14% for \( T_\ast = 8700 \ K \). Although the latter value is more typical of those measured for Herbig AeBe stars, the former is not unreasonable given the large range in the measured values (Brooke, Tokunaga, & Strom 1993). Thus,
the strength of the 3.3 μm feature does not strongly imply a spectral type earlier than G.

While the continuum level is fit fairly well with a G star on the birthline and an optically thick inner disk, if WL 16 is instead an A star near the main sequence, we require an additional source of 2.3 μm continuum. The very small grains (VSGs) that are the larger counterparts of aromatic hydrocarbons are likely to accompany these macro-molecules in the near environment of WL 16 and may contribute to this continuum. Natta, Prusti, & Krügel (1993) have shown that VSGs surrounding Herbig AeBe stars can reprocess stellar radiation into the near- and mid-IR (≈2–20 μm), significantly enhancing the strength of the continuum at 2.3 μm. Compared to the aromatic hydrocarbons responsible for the 3.3 μm feature, VSGs are excited by a much broader spectrum of photons. Thus they may significantly enhance the 2.3 μm continuum of WL 16 over a range of stellar spectral types and may make up the deficit of 2.3 μm continuum noted above for model L15.

In an effort to better determine the spectral type of WL 16, we obtained a high-resolution CSHELL spectrum centered on the 1.711 μm Mg line. For stars of spectral type G and later, this line is among the strongest stellar absorption features in the H band. While WL 16 is substantially fainter at this wavelength (H = 10.5) than in the K band (K = 7.8), the detectability of stellar lines should be better since the infrared excess in pre-main-sequence stars generally decreases with decreasing wavelength. The total integration time was 70 minutes through a 2” slit, giving a spectral resolution of ≈23 km s⁻¹ and a signal-to-noise ratio of 70. No absorption features were present. The limiting factor on a line detection is the residual fringing in the spectrum, which has a peak-to-peak amplitude of 4% and a period of about 360 km s⁻¹.

The nondetection of the Mg line could be due to any of the following reasons: the infrared continuum excess significantly dilutes the stellar spectrum; the stellar rotational broadening of the line is large; the star is hot enough that the line is weak or absent; or some combination of these. In order to investigate the constraints that can be placed on the spectral type of WL 16 from the H-band data, we calculated synthetic stellar spectra for the 1.711 μm Mg line. We used an updated version of the spectrum synthesis program MOOG (Sneden 1973) and the Kurucz grid of model atmospheres (Kurucz 1992). The gf-value and damping constants for the Mg line were determined by fitting an observed solar intensity spectrum (Livingston & Wallace 1991).

In the G star model for WL 16 described in § 3.2 (model L14), ≈70% of the 2.3 μm continuum comes from the star; the star should contribute an even larger fraction of the continuum at 1.7 μm. For the assumed stellar parameters (T_eff = 5500 K and log g = 3.5), the Mg line is predicted to be 35% deep at our spectral resolution in the absence of rotational broadening and other sources of continuum emission. In order for rotational broadening alone to account for the nondetection of the Mg line, we require v sin i ≥ 200 km s⁻¹ which is a significant fraction of the breakup velocity for such a star (310 km s⁻¹). A larger 1.7 μm continuum excess reduces the required rotational velocity. For example, with a more moderate v sin i of 50 km s⁻¹ we require an excess flux of at least 3 times the stellar flux. We conclude that a late-type star cannot provide the major contribution to the near-infrared continuum unless the star is an extremely rapid rotator.

The Mg line strength decreases with increasing effective temperature. For the temperature of our A-star model (model L15, T_eff = 8500 K) the observed depth is predicted to be 17%. In this model, the A star contributes only 23% of the 2.3 μm continuum flux. Assuming that the star contributes ≈35% at 1.7 μm, a rather low v sin i of 25 km s⁻¹ would be sufficient to hide the line in our data. Thus an A star, with the majority of the continuum coming from excess emission, would be consistent with the nondetection of the Mg line.

4. NON-LTE MODEL AND MULTIWAVELENGTH FITS

4.1. Level Population Calculation

We calculate the population of the ro-vibrational levels using a 10 level CO molecule to represent the vibrational levels and assuming that the rotational levels within each vibrational level are thermally populated. This approximation is based on the relatively large rate coefficients for the collisional de-excitation of pure rotational transitions (Green & Thaddeus 1976; Flower & Launay 1985; Schinke et al. 1985). In the temperature range of interest, the collisional de-excitation rate for both H and H2, k(v, J → v, J − ΔJ) ≈ 5 × 10⁻¹⁰ cm³ s⁻¹ for ΔJ = 1, 2, ..., is much larger than the corresponding rate coefficients for collisional changes in the vibrational (as well as rotational) quantum numbers. This approximation is valid even in the important case in which collisions with atomic hydrogen dominate the collisional rates of the rovibrational transitions, where k(v, J → v − 1, J − ΔJ) ≈ 10⁻¹¹ cm³ s⁻¹ (see Appendix B). The assumption of rotational equilibrium greatly simplifies the determination of the level populations and allows us to include in a simple way the effects of radiative trapping on the level populations. In this case, the vibrational level populations can be determined by the solution of a coupled set of effective vibrational level population equations which we derive in the following way.

The large oscillator strengths and collision cross sections of the ΔJ = 1 transitions suggest that the population of a given vibrational level is primarily determined by radiative and collisional interactions with adjacent vibrational levels. Therefore, we begin by assuming that each vibrational level v is directly coupled only to the vibrational level above it, v = v + 1. As discussed by Scoville, Krotkov, & Wang (1980), in this approximation successively higher vibrationally excited states are populated by climbing the vibrational ladder “one rung at a time.” In steady state, the number of transitions into level u from level v is balanced by the number of transitions into level v' from level u:

$$\sum_{u,l} n_l (A_{ul} + B_{ul} \tilde{J}_v + C_{ul}) = \sum_{v',l} n_l (B_{lu} \tilde{J}_{v'} + C_{lu}),$$

where we have adopted the notation l = (v, J) and u = (v', J') and the sum is over all transitions permitted between the upper and lower vibrational levels (e.g., the R- and P-branches). In this equation, n_l = n(v, J) and q_l = n(v, J) are the number densities of the upper and lower levels; A_{ul},

5 There is, in principle, an important difference between the rotational equilibrium of ground and excited vibrational levels stemming from the Einstein A-values for the two cases. Since the A-value is proportional to the echo of the transition frequency, the critical density for an excited vibrational level can be very large, typically n_{crit}(v > 0, J) ≈ 10⁹ cm⁻³. Nevertheless, the large difference between vibrational and rotational collision rates means that the total vibrational populations are more readily disequilibrated, and we choose to focus on this aspect of a non-LTE calculation.


\[ B_{ul} \text{ and } B_{lu} \text{ are the Einstein } A-\text{ and } B\text{-values of the transitions connecting the upper and lower levels; } C_{lu} \text{ and } C_{ul} \text{ are the upward and downward collisional rates for these transitions assuming a sum over all collision partners (H, H_2, e, etc.). We have also denoted the mean intensity of the radiation field weighted by the line profile function } \phi_v \text{ as } \langle J_v \rangle \equiv \frac{1}{4\pi} \int J_v \phi_v dv.

\]

Consistent with our assumption of vertical homogeneity (e.g., of the level populations) in the line layer, we calculate \( \langle J_v \rangle \) at an "average" height in the slab, located where the vertical optical depth is half the total. In addition, we ignore the effect of overlapping lines in the determination of the mean intensity. While the issue of overlapping lines is important for the calculation of the intensity emergent from the slab, we show below that the line overlap does not significantly affect the level populations. Thus, in the calculation of the level populations we treat each line as isolated and assume that the mean intensity which excites the calculation is a function of the optical depth for that line alone (i.e., \( J_v = J_v^a \)).

With these assumptions, \( J_v^a \) for a given line is (see eq. [1])

\[ J_v^a = \frac{1}{4\pi} \int \phi_v dv \int I_v d\Omega = B_a(T_a) \int_0^\infty \phi_v x dx \int_0^1 e^{-\tau \phi_v x^2} d\mu + S_v \left( 1 - \frac{1}{2} \int_0^\infty \phi_v x dx \int_0^1 e^{-\tau \phi_v x^2} d\mu \right), \]

where all physical quantities such as \( B_a(T_a), \), \( S_v \), and \( \tau_v \) are assumed to be independent of \( \mu \) because the disk is geometrically thin. The integration limits reflect the difference between the background radiation field which originates from the \( 2\pi \) steradians subtended by the continuum layer and the diffuse field which originates from all \( 4\pi \) steradians. The optical depth

\[ \tau_v = \frac{1}{\Delta v} \int \tau_v dv \]

is an average over the line profile where \( \Delta v \) is the characteristic width of the line profile centered on the line frequency \( v_0 \).

We introduce a dimensionless line profile function \( \phi_v \equiv \phi_v, \Delta v \) such that \( \tau_v = \tau_v^a, \phi_v \), and \( x \equiv (v - v_0)/(\Delta v) \). Following Mihalas's treatment of the escape probability in a static plane-parallel atmosphere with vertically homogeneous properties (Mihalas 1978), we interpret quantities of the form

\[ \beta(\tau) = \int_0^\infty \phi_v x dx \int_0^1 e^{-\tau \phi_v x^2} d\mu = \int_0^\infty E_x(\tau \phi_v x) \phi_v x dx \]

as escape probabilities. The line profile function in the outer integral expresses the probability that at a vertical optical depth \( \tau \), a photon is emitted \( x = (v - v_0)/(\Delta v) \) line widths from line center. Such a photon has a probability \( E_x(\tau \phi_v x) \) of escaping the layer in any direction \( \mu \). Our expression for \( \beta(\tau) \) differs from Mihalas's expression for \( P_x(\tau) \) by a factor of 2 because we have assumed that the photon can escape out of both the top and bottom of the atmosphere and that the optical depth in either direction is \( \tau = \tau_v/2 \) on average. While the photon cannot really "escape" by propagating into the continuum layer below, we imagine that photons reaching the continuum layer are destroyed and reemitted as continuum radiation. With this identification, the mean intensity weighted by the line profile is

\[ \langle J_v \rangle = B_a(T_a) \beta(\tau_v/2)/2 + S_v (1 - \beta(\tau_v/2)). \]

Note that the limits of integration in dimensionless frequency \( x \) in equation (6) assume an isolated line. For the CO fundamental transitions, the \( P \)-transitions are separated by 4-8 cm\(^{-1}\), and the \( R \)-transitions are more closely spaced (0-6 cm\(^{-1}\)) since the transitions go through a band head. If we conservatively estimate the average spacing between neighboring lines to be \( \Delta \nu = 1 \text{ cm}^{-1} \), truncating the integral at \( x_{\nu} = \Delta \nu/2 \Delta \nu \approx 50 \) for a line width \( \Delta \nu \) of 2.5 km s\(^{-1}\) reduces the escape probability by less than 50% relative to the untruncated case. So overlapping lines do not have a dominant effect on the relevant optical depths and escape probabilities.

Using the definition of the line source function to rewrite equation (7) for a given line as

\[ J^{ul}_v = \frac{A_{ul}}{B_{ul}} \left[ \frac{(1 - \beta_{ul})}{n_l g_{ul} n_u g_l - 1} + \frac{\beta_{ul}/2}{e^{h \nu_k T_k} - 1} \right], \]

where \( \beta_{ul} = \beta(\tau_v/2) \) and \( n_l, n_u, g_l, \) and \( g_u \) are the number densities and statistical weights of CO in the lower and upper levels, equation (4) becomes

\[ \sum_{ul} (n_l C_{lu} - n_u C_{ul}) = \sum_{ul} n_u A_{ul} \beta_{ul} \left[ 1 - \frac{(n_l g_u/n_u g_l - 1)/2}{e^{h \nu_k T_k} - 1} \right]. \]

If the rotational levels are thermally populated at the gas kinetic temperature \( T \)

\[ \frac{n_{ej}}{n_v} = \frac{(2J + 1)}{Z} e^{-J(J + 1)/B/2T} \equiv f_J, \]

where \( n_{ej} \equiv n(v, J) \), \( n_v \equiv \sum_j n_{vJ} \), \( B \) is the rotational constant, and \( Z \approx T/B \) is the partition function for the rotational levels, the radiative terms in equation (8) can be further approximated as

\[ \left[ 1 - \frac{(n_{ul} n_{eJ} - 1)/2}{e^{h \nu_k T_k} - 1} \right] \sum_{ul} n_u A_{ul} \beta_{ul} \]

With this approximation, the assumption of thermally populated rotational levels, and a little algebra, equation (8) becomes

\[ n_v \left[ C_{vu} + \left( 1 + \frac{1/2}{e^{h \nu_k T_k} - 1} \right) A_{vu} \beta_{vu} \right] \]

\[ = n_v \left( C_{vu} + \frac{1/2}{e^{h \nu_k T_k} - 1} A_{vu} \beta_{vu} \right), \]

where

\[
A_{vu}^* \equiv \sum_{ul} f_{vu} A_{ul}, \quad C_{vu}^* \equiv \sum_{ul} f_{vu} C_{ul}, \\
\beta_{vu}^* \equiv \frac{1}{2} \sum_{ul} f_{vu} \beta_{ul}, \quad C_{vu}^* \equiv \sum_{ul} f_{vu} C_{ul}.
\]

The terms \( A_{vu}^*, \beta_{vu}^*, C_{vu}^* \), and \( C_{vu}^* \) can be regarded as effective Einstein \( A \)-values, escape probabilities, and collisional

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terms that relate the vibrational levels. Note that we have made the additional approximation of separately averaging the A-values and escape probabilities since the A-values usually do not vary substantially over the range of J in which the escape probabilities are substantial. The factor of \( \frac{1}{2} \) in the definition of \( \beta^*_v \) arises from the consideration of the R and P transitions out of the upper levels.

By detailed balance the upward and downward collision rates satisfy \( C^*_{v'v} = \frac{n_v}{n_{v'}} \frac{C^*_{v'v}}{e^{\left(v' - v\right)/T}} \). The assumption of thermal rotational level populations now implies that

\[
C^*_{v'v} = e^{\left(v' - v\right)/T} C^*_{v'v},
\]

where \( \theta \) is the vibrational constant. When collisions dominate, equations (10) and (12) yield the LTE result: \( n_v/n_{v'} = \exp \left[ -\left(v' - v\right)/T \right] \). When radiative transitions dominate, equation (10) implies \( n_v/n_{v'} = 1/[2 \exp (h\nu/kT_v) - 1] \), which differs from the Boltzmann relation for excitation temperature \( T_v \) because the underlying radiation field is, as noted earlier, only half a blackbody.

The system of coupled, statistical equilibrium equations for the vibrational level populations \( n_v \) is linked to the radiative transfer part of the problem through the determination of the escape probabilities \( \beta^*_v \equiv \beta(t^*_v/2) \), which requires the self-consistent determination of the vibrational level populations that enter in the optical depths \( t^*_v \). This suggests that the system of equations can be solved iteratively given the initial assumption of LTE level populations. The rates for vibrational excitation due to collisions with H and H\(_2\) are based on the rate coefficients discussed in Appendix B. The total collision rates, \( C^*_{v'v} \), and \( C^*_{v'v} \), are obtained by assuming chemical equilibrium abundances of H and H\(_2\). Although Einstein A-values for individual lines within a vibrational band can vary by factors of several, the average A-value between two vibrational levels \( A^*_{v} \) varies weakly with temperature (see also Appendix A). Therefore, we adopt a temperature-independent rate of \( A^*_{v} = A^*_{v}(T = 2000 \text{ K}) \). The set of equations (10) can easily be generalized to couple more than two neighboring vibrational levels together. It is usually not necessary to couple more than three levels together to solve the set of equations (10) to a fractional accuracy of \(< 10^{-7} \).

4.2. Spectral Synthesis of the \( v = 5-3 \) and \( v = 2-0 \)

Band Head Regions

The \( v = 5-3 \) band head region of the spectrum provides valuable complementary information to that contained in the \( v = 2-0 \) band head region. For example, the onset of non-LTE level populations is difficult to detect with measurements of the \( v = 2-0 \) band head alone. This is evident from a comparison of the fits to the WL 16 \( v = 2-0 \) band head region shown in Figures 3b–3c and 3e–3f. The latter set of model fits, obtained with the non-LTE model described in the previous section, shows that similar fits to the data can be obtained with nearly the same parameters as in the LTE case (see Table 1). This is because the \( v = 2-0 \) band head region consists of lines that are primarily sensitive to gas at temperatures larger than those at which departures from LTE become important. This is shown graphically in Figure 4 in terms of the relative optical depth of a line at the \( v = 2-0 \) band head in LTE as a function of temperature. Because lines in the \( v = 2-0 \) band head arise from relatively high-lying rotational levels, they are not well excited below 1500 K even in LTE. Thus, while the lines in the \( v = 2-0 \) band head provide useful information about the kinematics of the gas in the inner disk, they probe only a limited range in temperature.

In comparison, isolated lines from several other CO overtone transitions also fall in the \( v = 5-3 \) band head region (Fig. 5); these lines and the \( v = 5-3 \) band head together probe a much wider range of excitation temperatures (500

![Figure 4](image4.png)

**Fig. 4.**—Relative LTE optical depths (see text) for representative 5–3 (light solid line), 4–2 (dotted line), 3–1 (short dashed line), and 2–0 (long dashed line) lines in the 5–3 band head region compared with that for a line at the 2–0 band head (heavy solid line).

![Figure 5](image5.png)

**Fig. 5.**—The \( v = 2-0 \) and \( v = 5-3 \) band head regions of the spectrum of 1548C27.
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K < T < 5000 K) than those contributing to the v = 2–0 band head. In Figure 4, the relative LTE optical depths $f_n(2J' + 1) \exp (-E_n/kT)$ of representative lines from different overtone transitions in the v = 5–3 band head region are compared with that of a line at the v = 2–0 band head. The strongest lines in the v = 5–3 band head region are the 2–0 and 4–2 lines and the 5–3 band head itself. The 2–0 P transitions are the primary low-temperature constraint. Because the 4–2 and 2–0 lines do not strongly overlap, it is possible to establish their relative strengths. The relatively large separation in wavelength between these lines also makes it possible to measure the line profiles of individual lines for sources with projected rotational velocities $\lesssim 100$ km s$^{-1}$. In the next section, we present our high spectral resolution observations and modeling of the v = 2–0 and v = 5–3 band head emission from such a source, the probable Herbig AeBe star 1548C27.

5. v = 5–3 AND v = 2–0 BAND HEAD SPECTROSCOPY OF 1548C27

5.1. What is 1548C27?

Discovered as a cometary nebula (Craine, Boeshaar, & Byard 1981), the YSO 1548C27 has an associated jet which is aligned with the nebular axis (Mundt et al. 1984). Energetic outflow activity is also indicated by the optical spectrum of the central object which is characterized by P Cygni H$\alpha$, H$\beta$ and broad blueshifted Na I D absorption (Mundt, Brugel, & Bührke 1987). In the near-infrared, 1548C27 has very prominent CO overtone and Br\gamma emission (Carr 1989, 1990). The U-band–100 $\mu$m spectral energy distribution indicates a system luminosity of $L \approx 128(kpc)^2 L_\odot$ (Vilchez et al. 1989) with the spectral shape indicating Herbig AeBe Group II membership (cf. Hillenbrand et al. 1992).

The distance to 1548C27 is controversial. It is located in front of the H II region S86 and in the same region of the sky as NGC 6823, both of which are located at a distance of $\sim 2.5–3$ kpc. Arguments based on the spectral energy distribution shortward of 3 $\mu$m suggest that 1548C27 is at a distance of $d = 0.91$ kpc, while CO (J = 2–1) and CO (J = 3–2) data have been used to derive a kinematic distance of 2.4 kpc (Dent & Aspin 1992).

5.2. CO Overtone Spectroscopy

Spectra of 1548C27 in the v = 2–0 and v = 5–3 band head regions were obtained on 1993 May 15 at the NASA Infrared Telescope Facility using the facility cryogenic echelle spectrograph (CSHELL; Tokunaga et al. 1990, Greene et al. 1993) with a NICMOS 3 256 $\times$ 256 HgCdTe array and a 1$''$ slit width that provided a resolution of 14 km s$^{-1}$. The data were acquired and reduced as described in Carr et al. (1993). The resulting spectra for the v = 2–0 and v = 5–3 band head regions are shown in Figure 5. As in the case of WL 16, the 2–0 band head of 1548C27 also shows the characteristic shape of band head emission from a rotating disk. Although the velocity extent of the band head emission (100–120 km s$^{-1}$) is less than in the case of WL 16, the characteristic blue wing, shoulder, and red peak are clearly present.

In the v = 5–3 band head region, the feature that peaks at 2.3837 $\mu$m is a complex blend of the 5–3 band head, located at 2.3829 $\mu$m, and neighboring 4–2 R and 2–0 P lines. Although the kinematics of the emitting gas are difficult to deduce from the shape of this feature, the double-peaked shape of the features at 2.3795 and 2.3810 $\mu$m strongly supports the rotating disk interpretation for the emission. The feature at 2.3795 $\mu$m is essentially a 4–2 R line, and the feature at 2.3810 $\mu$m is a blend of a 4–2 R line and a 2–0 P line. Based on the velocity extent of the v = 2–0 band head emission, these features overlap only slightly in their respective red and blue wings and are a reasonably good representation of an isolated line profile.

5.3. Modeling of the v = 2–0 and v = 5–3 Spectra

To determine the constraints that these observations place on the inner disk properties of the system, we assume that the distance to 1548C27 is 2.4 kpc (see Dent & Aspin 1992) and, therefore, the system luminosity is 740 $L_\odot$. Studies of the evolution of intermediate-mass pre–main-sequence stars restrict the stellar mass of such a system to $\lesssim 5 M_\odot$ (Palla & Stahler 1993). For this mass, the velocity extent of the 2–0 band head implies that the system must be viewed at relatively low inclination. For a K-band extinction of $A_K = 0.62$ mag (Vilchez et al. 1989) and $i = 30^\circ$, we can roughly match the velocity extent of the band head emission with $M_*, = 4 M_\odot$ and a CO dissociation radius at $\sim 20 R_\odot$. Assuming that the continuum layer has the usual radial temperature dependence $T_\odot = T_\odot (r/r_\odot)^{-3/4}$, we can fit the 2.3 $\mu$m continuum level and the system luminosity with $R_* = r_\odot = 5.9$ $R_\odot$ and $T_\odot = 9800$ K. These parameters correspond to a stellar luminosity of $L_\star = 200 L_\odot$, an age of $1.5 \times 10^7$ yr from the birthline (Palla & Stahler 1993), and a disk luminosity of $L_D = 4\pi R_D^2 \sigma T_D^4 = 280 L_\odot$. We assume that a boundary layer/hot spot contributes another $260 L_\odot$ so that the total system luminosity is 740 $L_\odot$.

The adopted stellar parameters place 1548C27 among the Herbig AeBe stars. In order to place a lower limit on the disk accretion rate, we conservatively assume that 25% of the nondisk luminosity ($460 L_\odot$) is reprocessed by the disk; the disk then has an accretion luminosity of $L_D^{\text{acc}} = (280 - 115) L_\odot = 165 L_\odot$. This value of $L_D$ corresponds to a disk accretion rate of $M_D \gtrsim 1 \times 10^{-5} M_\odot$ yr$^{-1}$ (eq. [3]), which is consistent with the evidence for energetic mass outflow from 1548C27 and with the large disk accretion rates characteristic of Herbig AeBe stars.

Applying the non-LTE model with these parameters, we obtain a good fit to the detailed properties of both spectral regions with Lorentzian line profiles and the parameters shown in Table 2. In each case, the synthesized spectrum for the 2–0 band head region is scaled by $\sim 95\%$ to match the continuum in both spectral regions (Fig. 6). In the 5–3 band head region, the true continuum is 0.246 Jy whereas the apparent continuum of 0.26 Jy (e.g., as measured at 2.382 $\mu$m) is formed by the overlap of rotationally broadened lines. The “wiggles” in the emission redward of the peak in the 2–0 band head region arise from the “beating” of the (double-peaked) rotational broadening function against the...
TABLE 2

1548C27 MODEL PARAMETERS

<table>
<thead>
<tr>
<th>Model</th>
<th>( R_{\ast} ) (R(_{\odot}))</th>
<th>( T_{\ast} ) (K)</th>
<th>( T_{\nu} ) (K)</th>
<th>( \sigma_{10} ) (g cm(^{-2}))</th>
<th>( q_{i} )</th>
<th>( \phi_{i} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>G1.10</td>
<td>5.9</td>
<td>9800</td>
<td>0.75</td>
<td>9700</td>
<td>0.6</td>
<td>7100</td>
</tr>
<tr>
<td>G5.9</td>
<td>5.9</td>
<td>9800</td>
<td>0.75</td>
<td>13000</td>
<td>0.7</td>
<td>800</td>
</tr>
</tbody>
</table>

Note.—\( M_{\ast} = 4.0 \, M_{\odot} \), \( L_{\ast} = 200 \, L_{\odot} \), \( d = 2400 \, \text{pc} \), \( A_{k} = 0.62 \), \( i = 30^\circ \), \( R_{\infty} = 1 \, R_{\ast} \), and \( r_{0} = 5.9 \, R_{\odot} \).

rest wavelength distribution of the overtone lines. For systems with reduced projected rotational motion (e.g., 1548C27 and SVS-13), this effect appears close to the 2–0 band head where the spacing of lines is comparable to the separation of the horns of the rotational broadening function. The high-resolution overtone spectroscopy by Carr & Tokunaga (1992) and Chandler et al. (1993, 1995) show the same effect.

As shown in Figure 6, we are able to fit the 2–0 shoulder, peak, blue wing, and slope of the emission redward of the peak reasonably well with \( \gamma \) in the range 1.0 km s\(^{-1}\) (model G1.10) to 5.0 km s\(^{-1}\) (model G5.9). The model “wiggles” would have reduced amplitudes if we convolved the synthetic spectra with the instrumental profile, but they would still appear too prominent when compared with the observations. This indicates perhaps a smoother inner cutoff to the CO abundance than given by the chemical and mechanical equilibria calculations of this paper. Apart from this problem, the relative heights of the emission at 2.3800 \( \mu \)m and 2.3815 \( \mu \)m are well fit by both models, as is the shape of the 5–3 band head at 2.3829 \( \mu \)m.

With \( \gamma = 1 \) km s\(^{-1}\), there is a relatively narrow range of \( T_{r} \) and \( \sigma_{i} \) within which acceptable fits can be obtained. If the outer radius to the emission is roughly held fixed, much flatter \( \sigma_{i} \) and steeper \( T_{r} \) distributions would steepen the slope redward of the 2–0 band head and weaken the 2–0 wing relative to the fit in model G1.10. This is because both a flatter \( \sigma_{i} \) and a steeper \( T_{r} \) reduce the emission from small radii (the latter by dissociation), weakening the blue wing and steepening the slope redward of the band head. Although equally good fits to the 2–0 band head are possible for somewhat flatter \( T_{r} \) and \( \sigma_{i} \) distributions, these parameters would also overemphasize the strength of the 2–0 \( P \) lines in the \( v = 5–3 \) band head region. Temperature distributions as flat as \( T_{r} \sim r^{-0.3} \) can be ruled out by the shape of the 2–0 band head alone. In this case, the flatter \( T_{r} \) distribution reduces the emission from small radii by reducing the line source function, steepening the slope redward of the band head and weakening the blue wing. Intrinsically broader lines that are less optically thick allow acceptable fits within a larger range of \( T_{r} \) and \( \sigma_{i} \) distributions. Model G5.9 with \( T_{r} = 9200(r/R_{\ast})^{-0.55} \) and \( \sigma_{i} = 620(r/R_{\ast})^{-2} \) also produces a good fit. Column density distributions as flat as \( r^{-1} \) can fit both band head regions reasonably well with \( T_{r} \sim r^{-0.6} \).

5.4. Discussion of the Non-LTE Effect

In these models, the strength of the feature at 2.381 \( \mu \)m and the shape of the band head at 2.384 \( \mu \)m are sensitive to the departure of the level populations from LTE beyond the radius at which \( T_{r} \) drops below 2000 K. This effective outer radius to the emission occurs somewhat beyond the radius at which H associates to form H\(_{2}\), since H\(_{2}\) has a cross section for collisionally exciting CO that is smaller than

![Fig. 6.—Non-LTE model fits to both the \( v = 2–0 \) and \( v = 5–3 \) band head regions of 1548C27 with (a, b) model G1.10 (\( \gamma = 1 \) km s\(^{-1}\)) and (c, d) model G5.9 (\( \gamma = 5 \) km s\(^{-1}\)). The \( v = 2–0 \) band head region has been scaled by 95% to match the continuum in both spectral regions.](https://example.com/fig6.png)
that of H by almost 2 orders of magnitude (see Appendix B). When non-LTE effects set in, the upper vibrational levels are depopulated first, followed by lower vibrational levels. This sequence produces effective outer radii that are larger for the lower overtone transitions. Radiative trapping plays an equally important role in determining the radii at which these effects occur because the line layer column densities required to fit the overtone emission correspond to very large optical depths in the fundamental transitions. The low escape probabilities associated with these optical depths result in significant radiative trapping of fundamental transition photons (cf. Scoville et al. 1980).

At these large optical depths, photons can only escape in the wings of the line, so the degree of radiative trapping depends significantly on the assumed behavior of the wings of the line profile function. For example, the escape probability for Gaussian lines depends asymptotically on the line optical depth as \( \sim \tau^{-1/2} \) whereas the escape probability for lines with Lorentzian or Voigt profiles depends asymptotically on the line optical depth as \( \sim \tau^{-1/2} \) (see Mihalas 1978, pp. 341–342). The shape of the line profile function also affects which transitions contribute most significantly to the escape of fundamental photons.

At the column densities required to fit the overtone emission, escapes from line profiles with weak line wings, such as Gaussian line profiles, are difficult because they must be made far out in the wings. As a result, only the high-\( J \) transitions have any substantial escape probability due to the reduced population of these levels and their consequently lower optical depths. At the same time, since the population of these levels is small, their contribution to the effective escape probability (eq. [10]) is also small because the weighting factor \( f_{J} \) is also reduced by the level population. This radiative trapping results in the extension of LTE level populations to radii significantly beyond the radius at which H associates to form H₂.

In comparison, the probability of escape from lines with Lorentzian profiles is larger enough that the transitions from relatively well populated \( J \) levels contribute significantly, resulting in substantially larger effective escape probabilities. In the fit shown in Figures 6a–6b, which has a Lorentzian line profile with a width of \( \gamma = 1.0 \text{ km s}^{-1} \), the vibrational level populations reach half their LTE value at \( \sim 1800 \text{ K} \) \( (v = 5) \) to \( \sim 1400 \text{ K} \) \( (v = 1) \). Correspondingly, the strong transitions in the \( v = 5–3 \) band head region have source functions that depart from LTE at \( \sim 1800 \text{ K} \) \( (v = 5–3) \), \( \sim 1750 \text{ K} \) \( (v = 4–2) \), and \( \sim 1550 \text{ K} \) \( (v = 2–0) \), introducing effective outer radii to the CO emission at these temperatures. In this particular case, this departure from LTE only affects the spectral fit through its impact on the strength of the \( 2–0 \) \( P \) lines in the \( v = 5–3 \) band head region. This is because only the \( 2–0 \) \( P \) lines still have a significant optical depth at the temperature at which the departure from LTE is significant: at \( \sim 1800 \text{ K} \) the optical depth of a representative \( 5–3 \) line is \( \tau(5–3) \approx 0.2 \), at \( \sim 1750 \text{ K} \) \( \tau(4–2) \approx 3 \), and at \( \sim 1550 \text{ K} \) \( \tau(2–0) \approx 11 \).

To illustrate the magnitude of the non-LTE effect on these lines, in Figures 7a–7b we also show the spectral fit obtained by the (artificial) extension of LTE level populations to 600 K \( (r = 100 \text{ R}_\odot) \). The synthesized \( v = 2–0 \) band head is somewhat stronger at the peak and the “wiggles” are filled in by emission from larger radii. The overall shape of the \( v = 2–0 \) band head is altered only slightly because the lines in the \( v = 2–0 \) band head have optical depth \( \tau \lesssim 1 \) at the radii at which the non-LTE effect would have set in. In contrast, the overall shape of the spectrum in the \( v = 5–3 \) band head region is altered significantly by the prominence of the \( 2–0 \) \( P \) lines which have

![Fig. 7](image_url) (a) The spectral fit obtained with the parameters of model G1_10 and the (artificial) extension of LTE level populations down to 600 K \( (r = 100 \text{ R}_\odot) \). (c) The spectral fit obtained with the same model parameters, non-LTE level populations, and an abrupt truncation of the emission beyond the radius at which \( T = 1500 \text{ K} \). As in Fig. 6, the \( v = 2–0 \) band head region has been scaled by 95% to match the continuum in both spectral regions. When compared with the fits shown in Fig. 6, these figures show that non-LTE level populations introduce a sharp outer radius to the emission, an effect that is easily diagnosed by fitting these band head regions.
We have successfully modeled the CO overtone band head regions of the observed spectra of the YSOs WL 16 and 1548C27 in terms of emission from Keplerian disks. Our observations provide some of the most compelling evidence to date for the existence of disks around young stars. Through the detailed study of this emission, we are able to place useful constraints on the detailed physical properties of the inner disk region. These results also indicate the possibility of dynamical mass estimates of young stars through the study of the emission from their surrounding disks. In the case of WL 16, we are able to place a useful constraint on the stellar mass due to the relatively high inclination of the system and the availability of additional observational constraints on the distance and system luminosity. In the case of 1548C27, observations of the \( v = 2-0 \) and \( v = 5-3 \) band head regions provide a particularly potent combination of diagnostics which reveal the onset of non-LTE level populations. In our non-LTE model for the emission from this source, the CO emission arises from gas at temperatures between ~5000–1500 K. The inner radius is due to the dissociation of CO; the outer radius is the result of the departure of the vibrational level population from LTE due to radiative escapes from suprathermally broadened lines. The need for this outer radius is clearly indicated by the strength of the 2–0 \( P \) lines in the \( v = 5-3 \) band head region.

Thermal Doppler line profiles appear unable to reproduce either the shape of the \( v = 2-0 \) and \( v = 5-3 \) band head regions of 1548C27 (§ 5.3) or the strength of the \( v = 2-0 \) band head emission from WL 16 given our adopted extinction to the source (§ 3.2). In contrast, Lorentzian line profiles with widths of \( \gamma = 0.5–5 \) km s\(^{-1}\) produce excellent fits to the overtone band head regions of both sources. These considerations argue against the presence of dust in the temperature inversion region of the disk atmosphere as the sole explanation for the effective outer radius to the over-

tone emission. In the absence of turbulence, dust is believed to settle very quickly from the surface layers of a circumstellar disk (Weidenschilling & Cuzzi 1993). In the presence of turbulence, which may mix dust into the surface layers, the non-LTE effect discussed in the text will come into play independent of any radiative transfer role for the dust.

The relatively large inner disk temperature inversion and the steep temperature and column density distributions required to fit the overtone emission restrict the possible physical explanations for the origin of the inversion. These include irradiation of the disk surface by the star and funnel flow accretion shock; turbulent heating in a wind-disk interface; and chromospheric heating that results from the vertical propagation of MHD waves into the upper disk atmosphere. Calvet et al. (1991) have used an approximate irradiation model to explore the ability of stars to heat their circumstellar disk surfaces and whether the CO overtone band heads consequently appear in emission or absorption as a function of stellar spectral type and disk accretion rate. However, in terms of the detailed properties of the inner disk region, the model predicts flatter temperature and column density distributions as well as smaller temperature inversions than are implied by our observations of WL 16 and 1548C27. The steeper temperature and column density distributions required to fit the observations could result, instead, from the interaction of a stellar wind with the circumstellar disk: as the wind blows over the disk, the disk atmosphere is stirred up and heated by turbulent dissipation, forming an aeolosphere (cf. Carr et al. 1993). Suprathermal widths may also characterize lines formed in disk chromospheres. If accretion in the inner disk is the result of the Balbus-Hawley instability (Balbus & Hawley 1991), the magnetic fields that participate in this instability may also propagate MHD waves into the upper disk atmosphere, producing a turbulent temperature inversion layer.

If the turbulent line broadening resulting from either of the latter two possibilities is similar in nature to that observed in spectral line profiles of turbulent interstellar gas (Falgarone & Phillips 1990), the Lorentzian line profiles adopted in our synthetic spectra may reasonably represent the expected non-Gaussian line wings. We have examined and rejected the alternative possibility that the enhanced line widths are due to pressure broadening. According to standard estimates for this broadening mechanism, the gas pressure in the model layers is too low by several orders of magnitude to be responsible for the line widths. We conclude that the line-emitting region in these sources is likely characterized by velocity widths much larger than those appropriate for a quiescent layer heated by the external radiation field of the star. Since external irradiation will tend to stabilize disk atmospheres against convection, we suggest that these turbulent motions arise from perturbing influences such as stellar winds or vertical MHD wave propagation.

The stellar properties and disk mass accretion rates implied by our observations also bear on the origin of the inner disk temperature inversion. In the case of WL 16, the implied disk accretion rate is low, \( \dot{M}_D \lesssim 2 \times 10^{-7} \, M_\odot \, yr^{-1} \), independent of whether we model the stellar component as a G star close to the stellar birthline, or as an A star close to the main sequence. Given the low disk accretion rate, heating due to irradiation by the relatively early type star may then provide a simple explanation for the temperature inversion in the upper atmosphere of the surrounding disk. The low disk accretion rate is also consistent with the lack of evidence for energetic outflow activity from this source; at such low wind mass loss rates, aeolospheric or chromospheric effects may be confined to a modest amount of line broadening which contributes to the strength of the \( v = 2-0 \) band head emission. However, the steep temperature and column density distributions derived for the line-emitting region may indicate that these effects are also responsible for additional nonnegligible heating.

In contrast, our models for 1548C27 indicate a large disk accretion rate, \( \dot{M}_D \sim 1 \times 10^{-5} \, M_\odot \, yr^{-1} \). Although this large mass accretion rate is consistent with both evidence for energetic mass outflow from 1548C27 and the disk accretion rates of other young stars with stellar parameters similar to those deduced for this object (Herbig AeBe stars), it also places 1548C27 in a parameter regime in which the models of Calvet et al. (1991) would nominally predict that
the CO lines should be in absorption rather than in emission. This discrepancy may be reduced by considering the impact of the funnel flows that are likely to accompany these large disk accretion rates if intermediate mass stars possess magnetospheres (cf. Najita, Carr, & Tokunaga 1996): the UV emission from the accretion hot spots at the base of the funnel flows is an energetic source of radiative heating. The possibility of a large contribution to the temperature inversion of the inner disk from funnel flow hot spots introduces the possibility of another source of energetic heating for the inner disk. As described by Shu et al. (1994), stellar magnetic fields strong enough to truncate the disk and to give rise to a funnel flow are also strong enough also to drive a magnetocentrifugal X-wind from the inner disk region. An aeolosphere may be created as the lowermost streamlines of this X-wind sweep over the inner disk region, producing the suprathernally broadened lines and steep temperature and column density distributions implied by our modeling of the overtone emission from this source. Similar line broadening and inner disk heating may arise from a disk chromosphere.

An improved understanding of the origin of the inner disk temperature inversion requires a detailed theoretical study of the temperatures and column densities that characterize aeolospheres and disk chromospheres. In addition to indicating the degree to which aeolospheres and chromospheres can heat the upper disk atmosphere, these studies may also indicate the observational signatures of aeolospheres and chromospheres. When combined with observations of the type presented in this paper, these distinguishing characteristics could be used to address other astrophysical issues. For example, the observational identification of aeolospheres as the origin of the disk temperature inversion would favor a stellar origin for the wind over a disk origin, whereas the identification of chromospheres as the origin of the temperature inversion would support the Balbus-Hawley instability as the physical process that gives rise to accretion in disks.

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APPENDIX A

RADIATIVE RATES

A1. APPROXIMATE A-VALUES

The Einstein $A$-values for individual fundamental transitions from an upper level ($v', J'$) to a lower level ($v, J$) can be obtained from

$$A_{v', v, J, J'} = \frac{64\pi^4}{3h} \bar{v}^3 \langle v' J' | \mu(x) | v J \rangle^2 \frac{|m|}{2J' + 1},$$

(A1)

given the transition frequency $\bar{v}$ and the electric dipole moment function $\mu(x)$ as a function of the reduced internuclear distance $x$. The quantity $m = [J'(J' + 1) - J(J + 1)]/2$. For the purpose of calculating the vibrational level populations, we approximate the expectation value of $\mu(x)$, $M_v^0 = \langle v' J | \mu(x) | v J \rangle$, with the rotationless transition dipole moment $M_v^0(0)$,

$$A_{v', v, J, J'} = \frac{64\pi^4}{3h} \bar{v}^3 [M_v^0(0)]^2 \frac{|m|}{2J' + 1},$$

(A2)

and use transition frequencies $\bar{v}$ calculated using energy levels from Mantz et al. (1975) and values of $M_v^0(0)$ from Goorrivitch & Chackerian (1994).

Values of $A_{v'}^R$ (cf. eq. [11]) can be calculated by adding the rates for the $R$- and $P$-transitions associated with each upper level, weighting by the population of the upper level, and summing over all upper levels:

$$A_{v'}^R = \sum_{J'} f_{v, J} (A_{v', v, J, J' - 1} + A_{v', v, J, J' + 1}).$$

For the purpose of calculating $A_{v'}^R$, we have approximated the rotational level populations as

$$f_{v, J} = \frac{(2J + 1)}{(T/B_v)} \exp \left[ -\frac{J(J + 1)B_v}{T} \right],$$

where the rotational constants $B_v$ are from Mantz et al. (1975). Since $A_{v'}^R$ varies weakly with temperature over the range $T = 200$–$8000$ K, we further approximate $A_{v'}^R$ with its value at 2000 K.

A2. APPROXIMATE ESCAPE PROBABILITIES

The escape probabilities of fundamental photons averaged over rotational levels (eq. [11]) are approximated in the following way. The Einstein $B$-values for the fundamental transitions are related to the $A$-values by

$$B_{v', v, J, J'} = \frac{c^2}{2h\bar{v}^3} A_{v', v, J, J'} \frac{2J' + 1}{2J + 1}.$$
With the approximate form for \( A_{v'J',vJ} \) given by equation (A2),

\[
B_{vJ,v'J'} = b_{vJ} \frac{|m|}{2J + 1},
\]

where

\[
b_{vJ} = (32\pi^4/3h^2c)[M_0(0)]^2.
\]

The sum of the \( B \)-values for the \( R \)- and \( P \)-transitions from the same lower level is independent of \( J \),

\[
B_{vJ,vJ+1} + B_{vJ,vJ-1} = b_{vJ}.
\]

Therefore, \( b_{vJ} \) is the effective \( B \)-value relating the vibrational levels \( v \) and \( v' \),

\[
B^*_{vJ} = b_{vJ}.
\]

The corresponding optical depths are obtained from these upward \( B \)-values (eq. [A3]) using

\[
\frac{\pi e^2}{m_e} f_a = \frac{\hbar \nu}{4\pi} B^*_{vJ} \frac{|m|}{2J + 1}
\]

for the absorption oscillator strength \( f_a \) in the usual expression for the line optical depth. We also use the following simple approximation for the vibration-rotational energy levels in order to calculate transition frequencies and stimulated emission terms:

\[
E_{vJ} = v\theta_0 + J(J + 1)B_0,
\]

where \( \theta_0 = 3084 \text{ K} \) and \( B_0 = 2.768 \text{ K} \). Accordingly, for the rotational level populations, we adopt

\[
f_J = \frac{(2J + 1)}{(T/B_0) \exp \left[ -J(J + 1)B_0/T \right]}.
\]

Given these optical depths, the effective escape probability \( \beta^*_{vJ} \) (cf. eq. [11]) is determined by averaging the escape probabilities for the \( R \)- and \( P \)-transitions associated with each upper level, weighting by the population of the upper level, and summing over all upper levels:

\[
\beta^*_{vJ} = \frac{1}{2} \sum_J f_J (\beta_{vJ,J-1} + \beta_{vJ,J+1}),
\]

where \( \beta_{vJ,J-1} = \beta vJ,J-1, \beta_{vJ,J+1} = \beta vJ,J+1 \), and the rotational level populations are as approximated in equation (A4).

**APPENDIX B**

**COLLISIONAL EXCITATION OF CO VIBRATIONAL LEVELS**

The excitation of the rotation-vibration bands of CO has been discussed earlier in a variety of astrophysical contexts, e.g., Thompson (1973; late-type stars); Scoville et al. (1980; young stellar objects); and Ayres & Wiedemann (1986; the Sun). Our discussion focuses mainly on collisional excitation by atomic hydrogen, where there are real differences between authors, in contrast to collisions with electrons and hydrogen molecules. We emphasize new experimental data on the role of \( \Delta v \geq 1 \) transitions and the results of Elitzur (1983) based on surprisal theory.

Ayres & Wiedemann (1986) gave a critical discussion of existing experiments and theory on \( H + CO \) vibrational excitation. They used the Landau-Teller parameters determined in the shock-tube experiment of Glass & Kironde (1983, hereafter GK) and standard theory (Herzfeld & Litovitz 1959) to obtain the \( \Delta v = 1-0 \) de-excitation rate coefficient \( k(1-0, H) \). GK found that the product of the postshock pressure \( P \) and the measured relaxation time \( \tau \) is essentially constant in the temperature range of the experiment \( T = 1000-3000 \text{ K} \),

\[
P\tau = 1.8 \pm 0.4 \times 10^{-8} \text{ atm s}.
\]

Substituting this into the standard theory for the relaxation time

\[
k(1-0) = k_n(T/\tau(1 - e^{-\theta_0/T})^{-1},
\]

we obtain a useful, direct representation of the GK data,

\[
k(1-0, H) = 7.57 \times 10^{-15} \text{ cm}^3 \text{ s}^{-1} T(1 - e^{-\theta_0/T})^{-1}.
\]

It must be emphasized that GK measured only one relaxation time and that they used the traditional method of analysis, based on the Landau-Teller theory, in which the collision rate coefficients obey the rules for dipole matrix elements of a simple harmonic oscillator (the same as for radiation), i.e., there are no \( \Delta v > 1 \) transitions. The GK result for \( H + CO \), i.e., equation (B3), is 100-20 times larger than the rate coefficient they measured for \( H_2 \) in the temperature range 1000-3000 K.

At the time of the Ayres & Wiedemann (1986) paper, there was little experimental information on the validity of the harmonic oscillator selection rules, which forbid \( \Delta v \geq 2 \) transitions. Hooker & Millikan (1963) did study the problem for rare
gas and molecule collisions with CO with shock-tube experiments at 2000 K. From the characteristics of the fundamental and first overtone time signals, they concluded that the CO population is largely built up by a sequence of $\Delta v = 1$ collisions, i.e., by "climbing the vibrational ladder." They estimated that the $\Delta v = 2$ collisions cannot occur as frequently as 1/10 as often as $\Delta v = 1$ collisions. In a high-pressure experiment on vibrational relaxation of electronically excited CO in the $A^1\Pi$ level, Fink & Gomes (1974) found that $\Delta v = 2$, 3 collisions were almost as important as $\Delta v = 1$ collisions. However, it is likely that physical processes other than those operative for the ground electronic level are at work in this experiment.

A new generation of molecular beam experiments is beginning to provide information on the dependence of the H + CO collisional excitation rate coefficients with selected final vibrational and rotational quantum numbers (e.g., Wight & Leone 1983; McBane et al. 1991). Although the experiments have only been done so far at energies higher than are relevant for astrophysical applications involving CO (1–3 eV), they do indicate that $\Delta v \geq 2$ transitions are significantly smaller than $\Delta v = 1$ transitions and that they decrease in importance as the energy is decreased. This is in conflict with Elitzur (1983), who correctly criticized the complete ignoring of $\Delta v \geq 2$ transitions but proposed rate coefficients which we believe are too large, especially for transitions involving large changes in v. Elitzur's work is based on a statistical theory known as "surprising theory" (e.g., Levine & Bernstein 1976). We suspect that the data underlying his application of this method are inappropriate to the H + CO system, i.e., they refer to different systems and situations, e.g., the one piece of evidence cited for CO involves rare gas atoms de-exciting an excited electronic level of CO (Fink & Gomes 1974).

In this paper, we adopt the following rate coefficients:

1. Electron collisions.---The rate coefficient for $T \leq 5000$ K that we derive from laboratory experiments (e.g., Morrison 1988) is

$$k(1-0, e) \approx 1.24 \times 10^{-11} \text{ cm}^3 \text{ s}^{-1} T^{1/2}.$$  \hspace{1cm} (B4)

On this basis, we can ignore electronic collisions compared with neutrals collisions for $x_e \leq 10^{-3}$, a restriction satisfied by the disk atmospheres and winds considered here.

2. Molecular hydrogen and helium.---We follow previous studies and use the results of shock-tube experiments in Landau-Teller form,

$$k(1-0, X) = n_X T_e^{-\frac{1}{2}} \left(1 - e^{-\frac{n_X}{T}}\right)^{-1} \text{ atm s}^{-1}.$$  \hspace{1cm} (B5)

The parameters by Milliken & White (1963) are $X = \text{He}$, $A = 99 \text{ atm s}$, and $B = 20.4 \text{ atm s}$; $X = \text{H}_2$, $A = 68 \text{ atm s}$, and $B = 19.1 \text{ atm s}$.

3. Atomic hydrogen.---In accord with the above discussion, we use equation (B3).

The available information suggests that $\Delta v \geq 2$ transitions are relatively unimportant although quantitative data are not yet available for the temperature range of interest. We ignore these collisions and, again consistent with available experimental information, we assume that all $\Delta v = 1$ transitions have the same rates. The rate coefficients summarized above indicate that for the disk layer analyzed in this paper, only atomic and molecular hydrogen collisions are important. Accordingly, in the level population calculation described in § 4, the total vibrational collisional excitation rate is

$$n C_v^* = n_H k(1-0, \text{H}) + n_{\text{H}_2} k(1-0, \text{H}_2).$$  \hspace{1cm} (B6)

This rate depends on the ratio $n_H/n_{\text{H}_2}$, which is assumed to be the chemical equilibrium value. Given the values for the vibrational levels are $n_H \sim 10^{12} - 10^{13}$ cm$^{-3}$, the over the temperature range $T = 2000$–$5000$ K. The critical densities increase toward both lower temperatures (due to reduced rate coefficients) and higher densities (which favor $\text{H}_2$ over $\text{H}$).

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