Origins of the H, He I and Ca II line emission in classical T Tauri stars

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ABSTRACT

We perform local excitation calculations to obtain line opacities and emissivity ratios, and compare them with observed properties of H, He I, O I, Ca II and Na I lines to determine the requisite conditions of density, temperature and photon ionization rate. We find that ultraviolet photoionization is the most probable excitation mechanism for generating the He λ10830 opacities that produce all the associated absorption features. We also calculate the specific line flux at an observed velocity of \( v_{\text{obs}} \) = ±150 km s\(^{-1}\) for both radial wind and infall models. All the model results, together with observed correlations between absorption and emission features and between narrow and broad emission components, are used to deduce the origins of the strong H, He I and Ca II broad line emission. We conclude that the first two arise primarily in a radial outflow that is highly clumpy. The bulk of the wind volume is filled by gas at a density \( \sim 10^9 \) cm\(^{-3}\) and optically thick to He λ10830 and Hα, but optically thin to He λ5876, Paγ and the Ca II infrared triplet. The optically thick He λ5876 emission occurs mostly in regions of density \( \geq 10^{11} \) cm\(^{-3}\) and temperature \( \geq 1.5 \times 10^4 \) K, while the optically thick Hα and Paγ emission occurs mostly in regions of density around \( 10^{11} \) cm\(^{-3}\) and temperature between 8750 and \( 1.25 \times 10^4 \) K. In producing the observed line fluxes at a given \( v_{\text{obs}} \), the covering factor of these emission clumps is sufficiently small to not incur significant absorption of the stellar and veiling continua in either He I or H lines. The strong Ca II broad line emission likely arises in both the magnetospheric accretion flow and the disc boundary layer where the gases dissipate part of their rotational energies before infalling along magnetic field lines. The needed density and temperature are \( \sim 10^{12} \) cm\(^{-3}\) and \( \leq 7500 \) K, respectively.

Key words: line: formation – stars: formation – stars: pre-main-sequence.

1 INTRODUCTION

Spectroscopic observations and analyses have been and will be an essential tool in discovering the intricate details in the formation of a classical T Tauri star, as the spectral lines carry information on the kinematics and physical conditions of the gases close to the star. Thus, red absorptions extending to velocities in excess of 100 km s\(^{-1}\) in Balmer lines and the Na I doublet (Appenzeller & Wolf 1977; Edwards et al. 1994) indicate active accretion. Blue absorptions are rare in optical lines, but are seen in Balmer lines, so there is also hint of outflows. In addition, the profiles of forbidden lines like [O I]\( \lambda \)6300 and [S II]\( \lambda \)6731 reveal the often presence of jet-like flows (Kwan & Tademaru 1988; Hirth, Mundt & Solf 1994). Even though the forbidden line emission originates at densities much lower than the expected values in the vicinity of the star, the high speed of the jet signals that the outflow likely starts from a deep potential well as either a stellar wind or an inner disc wind that is subsequently collimated. In that case, ejection of matter is as inherent a characteristic as accretion of matter in the star formation process.

Ultimately the unravelling of the contributions to the line intensities and profiles from different kinematical flows and the derivation of quantitative measures, like mass-flow rates, require modelling and analysis of the strong emission lines. Natta, Giovanardi & Palla (1988) first examined the excitation and ionization of hydrogen in a stellar wind, and found that the hydrogen line fluxes calculated for the same mass-loss rate can span a wide range, owing to differences in the gas temperature and in the stellar Balmer continuum. Hartmann et al. (1990) modelled both line fluxes and profiles in a stellar wind, which, driven by Alfvén waves, is characterized by large turbulent velocities in the accelerating portion of the flow. While the calculated fluxes cover the observed range, the Balmer line profiles are clearly unlike observed ones in being highly asymmetric with the red side much stronger than the blue side. The weakness of the blue side occurs because while the

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bulk of the hydrogen emission arises from the inner, denser and highly turbulent region, the outer, less strongly excited expanding envelope blocks the blue emission to an observer. Mitskevich, Natta & Grinin (1993) also modelled the Hα profile by postulating a flow that accelerates to a peak velocity and then decelerates towards zero. This double-valued velocity structure will produce an anomaly between the red and blue emission of an optically thick line at intermediate velocities, as the blue (red) emission will then arise from the outer (inner) part of the surface of constant observed velocity (v_{obs}), where the excitation temperature is lower (higher). Then, to account for the observed ranges in shape and depth of the apparent blue absorption, Mitskevich et al. (1993) advocated a clumpy flow for the flexibility in varying the degree of shielding between the outer and inner parts of the constant v_{obs} surface. They, however, adopted a parametric function for the dependence of the line excitation temperature on position, and it is not clear if the observed fluxes and profiles of several Balmer lines can be reproduced self-consistently.

The higher Balmer lines, as well as Paβ and Brγ, occasionally show red absorptions (Edwards et al. 1994; Folha & Emerson 2001). The clear indication of an infalling flow by the broad red absorption, and the realization that the optical/ultraviolet (optical/UV) continuum excess can arise from the impact footprints on the star (Calvet & Gullbring 1998), likely contributes to motivating the studies of hydrogen emission in an accretion flow from the disc along a dipolar trajectory (Muzerolle, Calvet & Hartmann 1998a, 2001; Kurosawa, Romanova & Harries 2008). The model hydrogen line fluxes generally agree with observed values and the model line profiles are centrally peaked with small blue centroids. The much better comparison of these calculated profiles with observed ones, in conjunction with findings of strong magnetic field strengths (Johns-Krull 2007) and theoretical investigations that probe the initiation of the accretion flow and the associated angular momentum exchange between the disc and star (Königl & Pudritz 2000; Romanova et al. 2007; Mohanty & Shu 2008), spurs the burst of activities on the present paradigm of magnetospheric accretion.

While magnetospheric accretion clearly occurs, the origin of the hydrogen emission in the accretion flow is not without question. On observational grounds, the main issue is the narrow width of the model line profile (Folha & Emerson 2001; Kurosawa et al. 2008). Other issues concern the observed high blue wing velocities, stronger blue emission at the line wings and sometimes large blue centroids. These characteristics are used as arguments against the He iλ5876 emission originating in the accretion flow (Beristain, Edwards & Kwan 2001, hereinafter BEK01), but are also present in observed hydrogen profiles (Folha & Emerson 2001). With the advance of infrared spectroscopy, it is natural to follow up the He iλ5876 study by observing the He iλ10830 line, which, being the transition immediately below He iλ5876, will have a higher opacity and be more effective in absorbing the stellar and veiling continua. The ensuing 1-μm spectroscopic survey of 38 Classical T Tauri Stars (CTTSs) (Edwards et al. 2006, hereinafter EFHK06) indeed produces additional information not conveyed by previously observed lines. It reveals that He iλ10830 has the strongest propensity of showing absorption features, including broad blue absorptions indicative of radial outflows, sharp, narrow, blue absorptions indicative of disc winds (Kwan, Edwards & Fischer 2007), red absorptions some of which are so broad and deep that challenge conventional assumptions of accretion flow structure (Fischer et al. 2008), and central absorptions.

The new 1-μm observational results bring forth several new insights into the problem of helium and hydrogen line formation in CTTSs. First, the significant He iλ10830 optical depth (∼1 or higher) in several different kinematic flows, together with the high excitation energy (∼20 eV) of its lower state, suggests that excitation via UV photoionization needs to be considered. Secondly, with the establishment of frequent presence of radial outflows, the issue regarding the comparatively rare occurrence of broad blue absorptions in Balmer lines now concerns the structures and physical conditions of those winds. Thirdly, He iλ10830 emission is common and comparable in strength to Paα emission among the CTTSs with the strongest hydrogen lines (EFHK06). Origin of the helium emission in an accretion flow faces high hurdles because of the arguments put forth earlier with respect to the He iλ5876 emission and because the physical conditions needed for strong He i emission (e.g., T > 10^4 K) may be taxing for a flow in which the gas is primarily in free fall. If the He i emission originates in a radial wind, hydrogen emission from the wind may also be significant, as it is more easily produced, and needs to be re-examined.

We will attempt to address the above issues in this paper. Unlike earlier investigations, which adopt a particular flow structure and evaluate how the model line fluxes and/or profiles fare with observations, we take a simpler but hopefully more general approach. We first calculate the atomic/ionic excitation at a local point and determine how the line opacities depend on the local physical conditions. We will use the observed relative opacities among the lines to shed light on the requisite physical conditions. The local excitation calculations also produce line emissivity ratios that can be compared with observed line flux ratios to further delimit the physical conditions. These local excitation results should be fairly independent of the flow structure. Then, to contrast between the outflow and infall velocity fields, we note that a major difference lies in the emission area contributing to the observed flux at a high |v_{obs}| and calculate the specific fluxes of the more important lines at |v_{obs}| = 150 km s^{-1} for both a radial wind and an accretion flow.

Crucial to our endeavour is an observational data set that covers simultaneously both optical and 1-μm spectral regions. He iλ10830 is a key line because of its propensity in showing absorption features, thereby indicating presence of particular kinematic structures. Deriving strong constraints on the physical conditions giving rise to He iλ10830 formation, however, needs another helium line, and He iλ5876 is ideal both for its being a fairly strong line and for its position as an antecedent of He iλ10830 in a radiative cascade. Such a data set of CTTSs, observed at optical and 1 μm wavelengths simultaneously or nearly simultaneously, has been procured by Edwards et al. (in preparation). Only a few of the objects are selected for use here to provide line ratios that are key to unravelling the origins of the line emission.

The outline of this paper is as follows. We describe the rationale and methodology of our model calculations in the next section and give details of the atomic/ionic models employed in Section 3. We present in Section 4 the relevant observational information. In Sections 5, 6 and 7, we present model results and compare them with observational data on line opacities, flux ratios and specific fluxes, respectively, to delimit the requisite physical conditions. We summarize the findings thus deduced in those three sections and utilize them, together with
other observational information, to decide on the locations of the H, He i and Ca II line emission in Section 8. We discuss the implications of our proposed origins of the line emission in Section 9 and review the major conclusions in Section 10.

2 LINE EXCITATION MODEL

The line emission from a gas depends on the local physical conditions and the radiative transfer. For three of the four kinematic structures revealed by the He ii 10830 absorptions, the high speed of each flow and the consequent large velocity gradients present isolate the radiative interaction to a region small in comparison with the overall size of the kinematic structure. Then, if the physical conditions within that region are taken to be uniform, both the excitation of the gas and the consequent line emission depend only on local physical quantities, namely density, kinetic temperature and velocity gradient. The spontaneous emission, stimulated absorption and emission of photons in a line together produce an effective emission rate given by $A\beta$, with $A$ being the spontaneous emission rate and $\beta$ the escape probability given by $(1 - e^{-\tau})/\tau$, where the line optical depth

$$\tau = \frac{g_u A\lambda^3}{g_l 8\pi} \left( N_l - N_u \frac{g_l}{g_u} \right) \frac{dl}{dv}$$

(1)

depends on the local population $N_l (N_u)$ of the lower (upper) level and the velocity gradient $dl/dv$ (Sobolev 1960). The local emissivity of the line ($\text{erg s}^{-1} \text{cm}^{-3}$) is then $N_u A\beta hv$.

We take advantage of the above reduction of a global radiative transfer problem to a local one in generating model results for comparison with observational data. We will be primarily interested in the stellar wind and the accretion flow in our attempt to ascertain the origin of the broad line emission. We consider a point roughly in the middle of the flow, specifically, at a distance $r$ from the star where the flow speed $|v|$ reaches 150 km s$^{-1}$. The calculated line opacity at this position is taken to be a representative and the ordering in magnitude of the opacities of different lines will be compared with the observational ordering of the lines in propensity of showing an absorption to delimit the requisite physical conditions. The local ratio of emissivities of two lines will also be compared with the observed line flux ratio. This appears to be, at first sight, a gross approximation, since each line flux is an integration of the emissivity over the entire kinematic structure and clearly the density at other parts of the flow, at least, is likely several times larger or smaller than the value at $|v| = 150 \text{ km s}^{-1}$. However, the same local excitation calculation applies to all positions, only for different physical parameters, so when the value of an emissivity ratio is presented as a function of the density and temperature, one can judge from the dependences how the results will be affected when averaged over a range of density and/or temperature. The local model calculation has the advantage of enabling us to explore a much broader parameter space, for example, over four orders of magnitude in density span, as well as to include many more lines, both from a given atom by incorporating more energy levels and from different atoms. It will be seen later that a comparison of different emissivity ratios with observed line flux ratios does indicate clearly enough the necessary physical conditions for us to infer the source of the line emission and that this deduction is not affected by the mentioned approximation.

For the stellar wind, we assume radial streamlines from the star. The velocity gradient transverse to the radial direction is $v/r$. We expect the velocity gradient in the radial direction to be larger, for an acceleration of the gas sufficiently strong to enable the gas to escape. When the velocity gradient is not isotropic, the escape probability is dependent on the direction. We will not be concerned with this nuance and simply assume an effective isotropic velocity gradient of $2v/r$. The error in this estimate is not significant, since the velocity gradient enters only as a factor in the line opacity and is always multiplied with the density, which varies over a much broader range. We take $r = 4R_*$ to illustrate the stellar wind model.

We also assume radial streamlines towards the star for the accretion flow. At $|v| \geq 150 \text{ km s}^{-1}$, the infall trajectory, even in the dipolar geometry, is approaching radial. The investigation of He ii 10830 red absorptions also finds that some red absorptions are so broad and strong that they are accounted for best by radial infall (Fischer et al. 2008). The same radial geometry also facilitates a comparison of the model results between the two flows. A gas particle infalling on to a star of 0.5 $M_\odot$ and 2 $R_\odot$ will attain a speed of 150 km s$^{-1}$ at 2.77 and 2.06 $R_*$ if it starts from 8 and 4 $R_*$, respectively, so we take $r$ to be 2.5 $R_*$ for the accretion flow, and assume an effective isotropic velocity gradient of $2v/r$. The latter assumption can be quite wrong, because, unlike the stellar wind, the accretion flow at high speeds fills only a very small solid angle and the above work on red absorptions also finds that the infall streamlines may not uniformly fill the solid angle. In this case, the velocity gradient is better given by $\delta v/\delta l$, where $\delta v$ is the thermal/turbulence linewidth and $\delta l$ is the transverse size of an infall bundle. We have no grasp on $\delta l$ and only note that $\delta v/\delta l$ may be closer to $2v/r$ even if $\delta l$ is very different from $r$. As mentioned before, the velocity gradient enters into the opacity with the density, which is varied over a wide range, and emissivity ratios will be presented also as a function of line opacity, so the effect of a very different velocity gradient can still be gleaned from the results.

The difference between the assumed velocity gradients for the two flows is not large. It will also be seen that the model results depend much more strongly on the density than the velocity gradient, so a change in the latter can be compensated by a much smaller change in the former. Thus, the calculated dependences of the local line opacities and emissivity ratios on physical conditions are applicable to all flow structures with a comparable velocity gradient to within an order of magnitude.

The specific flux of a line, on the other hand, depends on the volume of emission, so it is sensitive to $r$, the actual flow geometry, and the velocity at which it is calculated. In particular, because the correlation of the velocity with position is vastly different between the stellar wind and accretion flow, we expect their line fluxes to be quite different, even if line emissivity ratios are similar. It is with this consideration in mind that $|v| = 150 \text{ km s}^{-1}$ is specifically chosen. We think the line flux near this velocity or higher is a truer test of the accretion flow model than that near zero velocity, as it is much more constrained. Independent of the accretion flow geometry, the emission at high speeds must
arise from distances quite close to the star. Then, the focusing of the streamlines towards the star confines the solid angle of the accretion flow at high speeds, which is also constrained by the small area covering factor, typically \( \leq 0.03 \) (Calvet & Gullbring 1998), of the shocks marking the accretion footpoints. Thus, the emission volume of the gas at a high speed is well constrained, thereby making the corresponding line flux a more revealing diagnostic.

The observed line emission at \( v_{\text{obs}} = v \) arises from locations at distances other than \( r \), because emission from positions with higher speeds than \( v \) also contributes when their projected velocities along the line of sight equal \( v_{\text{obs}} \). To calculate it for the stellar wind, we assume that from \( r \) to 1.5\( r \) the flow speed increases linearly from \( v \) to 2\( v \). When the line is optically thick, the observed flux at \( v_{\text{obs}} \) depends on the excitation temperature and the projected area with an observed velocity of \( v_{\text{obs}} \). In Fig. 1, the right-hand dashed curve shows, in the \( x-z \) plane, the contour of \( v_{\text{obs}} = -150 \text{ km s}^{-1} \) for a spherical wind. It can be seen that the projected area equals \( \pi (1.5r \sin 60^\circ)^2 \) or \( 27(\pi/4)R^2 \). When the line is optically thick, the observed specific flux is then, in the simple case of a constant excitation temperature,

\[
F_{v_{\text{obs}}} = \frac{2h\nu^4}{c^4} \frac{1}{e^{\nu/kT_{\text{ex}}} - 1} \frac{27\pi R^2}{d^2},
\]

where \( T_{\text{ex}} \) is the excitation temperature between the upper and lower levels, and \( d \) is the distance to the star. When the line is optically thin, the observed flux is obtained from integrating the line emissivity over volume. If the latter varies with distance \( p \) from the star as \( p^{-a} \), the observed specific flux is

\[
F_{v_{\text{obs}}} = \frac{N_u A\beta h\nu}{4\pi d^2} \int_r^{1.5r} \left( \frac{r}{p} \right)^a \frac{4\pi p^2}{2v (\frac{2v}{c} - 1)} \, dp
\]

\[
= \frac{N_u A(1 - e^{-\tau})\nu^3}{2d^2\tau v} \int_1^{1.5} \frac{x^{2-a}}{2x - 1} \, dx.
\]

Both \( N_u \) and \( \tau \) are determined at \( r \). For \( a \) between 0 and 4, the integral ranges from 0.53 to 0.25. Substituting the expression for \( \tau \) given earlier in the denominator, the factor in front of the integral can be rewritten as

\[
(1 - e^{-\tau})\frac{2h\nu^4}{c^3} \frac{1}{e^{\nu/kT_{\text{ex}}} - 1} \frac{4\pi r^2}{d^2}.
\]

In order that, when a line transits from optically thin to optically thick, the observed specific flux increases smoothly to the earlier limit for an optically thick line, we choose the integral to be 0.422 or an \( a \) of 1. The error in deriving the flux of an optically thin line is then about \( \pm 40 \) per cent. It turns out that, with the exception of \( \text{O I} \lambda 8446 \), the other lines studied are all optically thick under the physical conditions responsible for the observed line emission. The error in the specific flux of an optically thick line arises mostly from the assumption of a constant excitation temperature, but one can gauge from the presented results how a distribution of density or temperature affects \( F_{v_{\text{obs}}} \).

For a spherical radial infall, the calculation of the specific flux when a line is optically thin is analogous to the spherical wind case. It involves integration from \( R_\star \) to \( r \) with an infall velocity distribution and is also made to ensure a smooth transition between the optically thin

Figure 1. Contour of \( v_{\text{obs}} = -150 \text{ km s}^{-1} \) for a spherical wind (right-hand side) reaching 150 km s\(^{-1}\) at 4\( R_\star \) and radial infall (left-hand side) reaching 150 km s\(^{-1}\) at 2.5\( R_\star \), for a viewer at \( x \to \infty \). The contour gives rise to a projected area of 27\( \pi R^2 \) and 1.286\( \pi R^2 \) for the wind and infall flow, respectively.
and Ca\textsuperscript{II} emissions. The contrast against the stellar wind value is even larger when the small solid angle of the accretion flow near the star is taken into account. When the wind/infall is not spherical, the projected area depends on the viewing angle, but we will not consider this nuance and simply obtain the observed line flux by multiplying the result for the spherical case by the filling factor of the flow in solid angle, $F_{\Omega}$. For the stellar wind and accretion flow, we adopt $F_{\Omega} = 0.5$ and 0.2, respectively, keeping in mind that the He\textsc{i}\,$\lambda$10830 line sometimes shows strong emission, but only a highly displaced shallow blue absorption (Kwan et al. 2007), and that $F_{\Omega} = 0.2$ is needed to model the few very strong red absorptions (Fischer et al. 2008). Occultation of the line emission by the star and the disc is not taken into account. If it were, it can be seen from Fig. 1 that in the infall case, the line flux at $v_{\text{obs}} \leq -150$ km s$^{-1}$ would be severely curtailed.

Three physical parameters are important for the local excitation calculations. They are density, temperature and ionization flux. The range of kinetic temperature, $T$, investigated is $0.5 \times 10^4$ to $3 \times 10^4$ K. We use the hydrogen nucleon number density, $N_H$, to indicate the number density, and assume a solar composition of the gas. With the number density of H : He : O : Ca : Na being in the ratio 1 : 0.0793 : 1.325 : 1 : 0.00017, the total nucleon number density is then $1.08 \times 10^8 n_{\text{H}}$, with $n_{\text{H}}$ being the mass of the hydrogen atom. The range of $N_H$ explored is $10^{12} \times 2 \times 10^{19}$ cm$^{-3}$. For a laminar flow, the mass flux is then given by $4\pi r^2 \rho v F_{\Omega}$ or $10^{-9} (N_H / 10^8 \text{ cm}^{-3}) (r / 4R_*)^2 (v / 150 \text{ km s}^{-1}) (F_{\Omega} / 0.5) M_\odot \text{ yr}^{-1}$.

We include photoionization as a means of excitation. As mentioned in Section 1, the high opacities of He\textsc{i}$\lambda$10830 in the accretion flow, stellar wind and disc wind suggest that photoionization of helium from its ground state is likely. As a rough estimate of this rate, we note that a luminosity of $10^{36} \text{ L}_\odot$ in photon energies above 24.6 eV, situated at the star, will produce at $4 R_*$ a He\textsc{i} photoionization rate of

$$\gamma_{\text{He}} = 2 \times 10^{-2} \left( \frac{\alpha - 1}{\alpha + 2} \right) \left( \frac{4R_*}{r} \right)^2 \frac{L}{10^{-4} \text{ L}_\odot} \text{ s}^{-1},$$

where $\alpha > 1$ is the power-law index of the luminosity energy distribution. The greatest uncertainty in this rate, however, is the attenuation between the UV luminosity source and the local point considered, since the mean free path for an optical depth of unity at the ionization threshold is small, ~1.3 $\times 10^4$ cm for a He\textsc{i} density of $10^4$ cm$^{-3}$. To circumvent this problem, we take $\gamma_{\text{He}}$ as a parameter. The UV source will also ionize hydrogen from its ground state. The ratio $\gamma_{\text{He}} / \gamma_{\text{He}}$ ranges from 1.6 to 4 for $\alpha$ between 1.5 and 3, if the luminosity source is not attenuated. The attenuations at the two thresholds can differ a lot and it is not clear which one is stronger. In most of the calculations, we simply adopt $\gamma_{\text{He}} = 2 \gamma_{\text{He}}$, but will comment on the effects of differential attenuation. The condition of $T_{\text{He}} / T_{\text{He}}(10830) \geq 1$ and $\tau_{\text{He}} \leq 1$, posed by the much more frequent occurrence of absorption features in He\textsc{i}$\lambda$10830 than Pa$\gamma$, will require a minimum $\gamma_{\text{He}}$ at temperatures low enough that collisional excitation of He\textsc{i} is ineffective. It turns out that this limit is ~$10^{-5}$ s$^{-1}$, so we will present results for $\gamma_{\text{He}} = 10^{-4}$ and $10^{-5}$ s$^{-1}$. The latter are much smaller than the unattenuated value given in the above equation for the hypothetical 10$^{37}$ L$\odot$ UV source.

The stellar and veiling continua, which peak at optical wavelengths, are effective in ionizing the excited states of hydrogen and helium, as well as the ground state and excited states of Ca\textsc{II} and NaI. We assume a CTTS of temperature $T_*$ = 4000 K and radius $R_* = 2 R_\odot$, and a veiling continuum given by a blackbody of temperature 8000 K, covering 3 per cent of the stellar surface area. At 5000Å, the veiling continuum is then about as strong as the stellar continuum.

To summarize, the physical parameters in our local excitation model are primarily $T$, $N_H$, $\gamma_{\text{He}}$, and 2$v/r$. The first three parameters are explored over broad enough ranges to cover all expected possibilities, such as line optical depths from 10$^{-2}$ to 10$^{2}$, and emissivity ratios from values smaller to values larger than corresponding observed line flux ratios. From these local excitation calculations, line opacities and emissivity ratios are compared with observational data to delimit the requisite physical conditions. These results are fairly general and not strongly dependent on the kinematic structure. The line flux, on the other hand, is sensitive to the actual flow geometry. To illustrate the contrast between the stellar wind and the accretion flow, the line fluxes at $|v_{\text{obs}}| = 150$ km s$^{-1}$ are calculated for a stellar wind and an accretion flow reaching a speed of $v = |v_{\text{obs}}|$ at $r = 4$ and $2.5 R_*$, respectively. All these results aid in deciphering the observed line emission for their origin.

### 3 ATOMIC MODELS

CTTSs show many emission lines. The forbidden lines, such as [O\textsc{i}]$\lambda$6300 and [S\textsc{ii}]$\lambda$6731, are likely formed at more than 10 stellar radii away and not germane to probing the accretion flow and inner structures of disc and stellar winds. Among the permitted lines, their profiles can be narrow (FWHM $\sim 20$ km s$^{-1}$, where FWHM stands for full width at half-maximum), broad (FWHM $\sim 200$ km s$^{-1}$) or composite with a narrow component (NC) atop a broad component (BC) (BEK01). The narrow lines/components are most likely formed at the sites where accreting streamlines impact the star, so their fluxes convey information on the summed area of such regions and the cooling history of the shocked gas. Here we are concerned with the broad lines/components, as their widths indicate that they are likely formed as the gas accelerates either away from the star in a wind or towards the star in an accretion flow. In the optical and infrared domains, these broad lines/components are, ordered roughly in decreasing emission strength, Balmer lines, Ca\textsc{II} H\&K, and infrared triplet, He\textsc{i}$\lambda$10830, Paschen and Brackett lines, Fe\textsc{II} lines, and He\textsc{i}$\lambda$5876, NaI D, O\textsc{II}$\lambda$4446, He\textsc{i}$\lambda$6678, O\textsc{I}$\lambda$7773 and Fe\textsc{I} lines.

Several intrigues posed by the observed strengths of the above-mentioned lines indicate that an understanding of the CTTS spectra must involve examining the H\textsc{i}, He\textsc{i} and Ca\textsc{II} line excitations altogether. One puzzle is the relative strength between the He\textsc{i}$\lambda$10830 and Pa$\gamma$
emission. The observed He $\lambda$10830 spectra (EFHK06), with prominent absorption features, clearly demonstrate the high $\lambda$10830 opacity, a consequence of the metastability of its lower state, 2s $^3$S. With $\tau_{\text{He}\alpha 10830}$ at least as large as $\tau_{\text{He}\beta}$, as inferred from their relative propensity in showing an absorption, a He $\text{I}$ 2s $^3$S level population comparable to or larger than the H$\text{I}$ $n = 2$ level population is implied, and while Pa$\gamma$ competes against Pa$\delta$, Br$\gamma$ and H$\delta$ for the de-excitation of $n = 6$, He $\lambda$10830 is the sole permitted radiative decay channel for its upper level, 2p $^3$P. Then, the He $\text{I}$ 2s $^3$S $\rightarrow$ 2p $^3$P collisional excitation rate is larger than the hydrogen $n = 2 \rightarrow 6$ collisional excitation rate. Thus, collisional excitation should highly favour He $\lambda$10830 over Pa$\gamma$ emission, yet the two observed fluxes in emission are almost the same.

The second puzzle is the very strong Ca $\text{II}$ infrared triplet emission in those CTTSs with strong He $\lambda$10830 and $\lambda$5876 emission. The summed flux of the Ca $\text{II}$ triplet rivals that of H$\alpha$. Collisional excitation of the triplet has the advantage that their upper state is only 3.1 eV above the ground state, so the triplet can be strong, relative to H$\alpha$, at low temperatures, despite the low Ca abundance. However, Ca $\text{II}$ has an ionization potential of 11.9 eV, much less than the He $\text{I}$ ionization potential of 24.6 eV. Furthermore, the lower state of the Ca $\text{II}$ infrared triplet is metastable and only 1.7 eV above the ground state, so it is well populated. Ionization from this level takes 10.2 eV, very slightly less than the energy of the Ly$\alpha$ photon. A strong build-up of the Ly$\alpha$ intensity, through radiative trapping, is required to sustain population in levels $n \geq 2$ of hydrogen for strong Balmer, Paschen and Brackett line emission. It will at the same time ionize Ca $\text{II}$ from its metastable state and reduce the Ca $\text{II}$ fraction. Thus, the strong Ca $\text{II}$ triplet emission appears incongruous with not only the He $\lambda$10830 and $\lambda$5876 emission, but also the hydrogen Paschen and Brackett line emission. In addition, the infrared triplet are optically very thick, as their fluxes are nearly equal, despite a factor of 10 difference among their oscillator strengths. Yet they rarely show any absorption feature, even in CTTSs where their peak fluxes are only comparable to the continuum flux.

In addition to calculating the H$\text{I}$, He $\text{I}$ and Ca $\text{II}$ line excitations, we will also include those of O $\text{I}$ and Na $\text{I}$. The O $\text{I}$ λ10287 and λ8446 lines are produced primarily through the Ly$\beta$ fluorescence process (Bowen 1947), in which the O $\text{I}$ 3d $^3$D state is excited upon absorption of a Ly$\beta$ photon by the ground state, and decays via emission of a λλλλ11287, 8446 and 1303 photon in succession. Thus, the O $\text{I}$ λ8446 flux is expected to correlate with the Pa$\gamma$ flux. The O $\text{I}$ λ7773 triplet, on the other hand, are formed via recombinations and cascade, and collisional excitation, in a fashion similar to the Ca $\text{II}$ infrared triplet. So, inclusion of the above O $\text{I}$ lines will further check on the H$\text{I}$ and Ca $\text{II}$ excitation conditions. The Na $\text{I}$ atom, with an ionization potential of 5.14 eV, is easily photoionized by the stellar and veiling continua and requires a much lower temperature for collisional ionization than helium. The often appearance of absorption in the Na $\text{I}$ doublet, in conjunction with absorption in He $\lambda$10830, then sheds light on the pertinent physical conditions.

We will not examine the Fe $\text{II}$ and Fe $\text{I}$ lines here. There are many of them, from infrared to UV wavelengths. Their summed emission strength may even exceed that of the hydrogen lines, so they are an important heat sink, and must be counted in deriving the total energy generation rate. With ionization potentials of 7.9 and 16.2 eV for Fe $\text{II}$ and Fe $\text{I}$, respectively, the Fe $\text{II}$-to-Fe $\text{I}$ flux ratio will also provide a constraint on the ionization condition. A proper study of the Fe $\text{II}$ and Fe $\text{I}$ line excitations, however, needs good observational data on the fluxes of the many UV multiplets, which are not presently available. The large number of lines and multiplets also makes it more suitable to study them in a separate paper. We will also not consider the lines of highly ionized metals, such as C $\text{IV}$,1549 and O $\text{IV}$,1034, partly because of the lack of simultaneous UV spectra and partly because understanding the excitations of these lines will likely benefit from understanding the H$\text{I}$, He $\text{I}$ and Ca $\text{II}$ excitations first, rather than vice versa.

In the following sections, we model atom/ion for H$\text{I}$, He $\text{I}$, O $\text{I}$, Ca $\text{II}$ and Na $\text{I}$, but leave the references for the atomic parameters to an appendix.

### 3.1 H$\text{I}$

In modelling the hydrogen atom, we assume that the level population among degenerate energy states are thermally distributed and use 15 distinct energy levels, labelled $n = 1$–15. Limiting the hydrogen atom to a fixed number of levels can be questionable. This is because the $n \rightarrow n + 1$ collisional rate coefficient, $C_{n,n+1}$, increases while the spontaneous emission rate of level $n$, $A_n$, decreases as $n$ increases, so once the electron density $N_e$ exceeds the value $A_n/C_{n,n+1} \sim 5 \times 10^7$ cm$^{-3}$ for $n = 15$, or an even lesser value if some of the radiative transitions are optically thick, successive collisional excitations to higher levels, effectively leading to ionization, provide the quickest route of depopulating level $n$. Ignoring this population transfer to higher levels produces error in the population of not only level $n$, but also lower levels, since there is less population return to the lower levels from collisional and radiative de-excitations of upper levels. The proper number of levels to use is clearly dependent on $N_e$, which, in turn, is dependent on $N_{\text{H}}$, $T$ and $\tau_{\text{H}}$. To ensure that 15 levels are adequate over our explored density range, we have performed the following test. The net rate of population transfer from level 15 to 16 via collisions is given by $R_{15 \rightarrow 16} = (N_{15}N_eC_{15,16} - N_{16}N_eC_{16,15})$, where $N_{15}$ and $N_{16}$ are the populations in level 15 and level 16, respectively. The local excitation calculations with 15 levels show that the $n = 13 \rightarrow 14$ and $n = 14 \rightarrow 15$ excitation temperatures are higher than 6000 K for $T$ between 7500 and $1.5 \times 10^4$ K, so, if the $n = 15 \rightarrow 16$ excitation temperature is $\geq 6000$ K, $R_{15 \rightarrow 16}$ has a top value of $\sim 0.01N_{15}N_eC_{15,16}$, which is about one-half the direct collisional ionization rate, $N_eC_{15,\infty}$, from level 15. To estimate the effect of including this additional route of population transfer for level 15, we have artificially doubled $C_{15,\infty}$ and repeated the calculations. The differences in the hydrogen level population and line fluxes are less than a few per cent between the two sets of calculations and we are confident that the use of 15 levels is adequate to obtain reliable Paschen line fluxes up to the transition from $n = 13$ to 3.
3.2 He I

Our He I model atom consists of 19 lowest energy states. The spin of the two electrons in each state can add up to 1 or 0, so the 19 states can be separated, according to the degeneracies of the total spin, into a ladder of triplets and one of singlets, since radiative transitions across the ladders are forbidden by the electric dipole selection rules. For the same energy quantum number \( n \), the angular momentum states \( l = 1 - n \), unlike those in the case of H I, are separated in energy by many thermal Doppler widths, so we treat them as distinct energy levels. Thus, our 19-level atom consists of 10 singlets with energy quantum number \( n = 1-4 \) and nine triplets with energy quantum number \( n = 2-4 \). Fig. 2 shows a schematic diagram of the 11 \( n = 1-3 \) levels and indicates several important transitions.

The lower level of the \( \lambda 10830 \) transition is highly metastable. Its radiative decay rate is only \( 1.7 \times 10^{-4} \) s\(^{-1}\). For our explored density and temperature ranges, collisions with electrons provide the speediest way of returning population to the ground state. A more important route than direct collisional de-excitation is collisional excitation to the singlet \( 2s \, ^3S \), followed by another collisional excitation to \( 2p \, ^1P \) and then emission of a \( \lambda 584 \) photon, which can either escape or photoionize hydrogen. The latter means of depleting \( \lambda 584 \) photons is particularly important, once \( \tau_{\text{He},\lambda 584} \) becomes large. To determine its rate of depopulating \( 2p \), we note that the mean free path for hydrogen ionization is \( l_{\text{abs}} = 1/(N_{\text{HI}}, \sigma_{\text{584A}}) \), where \( \sigma_{\text{584A}} \) is the photoionization cross-section at 584 Å, while the absorption mean free path, averaged over a thermally broadened profile, is

\[
l_{\text{abs}} \approx \frac{8\pi \Delta V_D [\ln(t_{\text{He},\lambda 584} + 2.72)]^{0.5}}{B_{\text{He}} c N_{\text{He}}}.
\]

where \( \Delta V_D = (2kT/m_\text{He})^{0.5} \) is the helium thermal Doppler velocity width and \( B_{\text{He}} \) is the Einstein stimulated absorption coefficient. The rate \( (s^{-1}) \) of depopulating \( 2p \, ^1P \) due to \( \lambda 584 \) ionizing hydrogen is then given by the product of the spontaneous emission rate, \( A_{\text{He},\lambda 584} \), and the ratio \( l_{\text{abs}}/l_{\text{He}} \). As an illustration, in the case when hydrogen and helium are mainly neutral, this rate is about \( 2 \times 10^{-3} A_{\text{He},\lambda 584} \), but almost \( 2A_{\text{He},\lambda 20581} \). This process is important in determining not only the population in excited singlets, but also, through the reduced population flow from singlets to triplets via collisions, the triplet population.

The ionization of hydrogen by \( \lambda 584 \) photons also affects the hydrogen ionization structure, but only slightly. If helium excitation is produced primarily by UV continuum photoionization, the production rate \( (s^{-1} \text{ cm}^{-3}) \) of \( \lambda 584 \) photons must be less than the rate of continuum photoionization or \( N_{\text{He},\gamma H} \), the rate of hydrogen ionization by \( \lambda 584 \) photons is then less than \( N_{\text{He},\gamma H} = 0.079 N_{\text{HII},\gamma H}, \) which is less than the rate of continuum photoionization of hydrogen, \( N_{\text{HII},\gamma H} \). If thermal motion is the energy source for helium excitation, then collisional excitation and ionization are much more efficient for hydrogen than for helium. Therefore, we have not implemented this coupling between the two ionization structures, which would require an iterative procedure.

In analogy to \( \lambda 584 \) photons ionizing hydrogen, Ly \( \alpha \) photons can ionize helium from its two metastable states, \( 2s \, ^3S \) and \( 2s \, ^1S \), which will have the bulk of the excited-state population. This ionization rate \( (s^{-1}) \) is readily deduced from the earlier derivation. It equals \( N_{\text{H}} A_{\text{He},\lambda 20581} \), where \( N_{\text{H}} \) is the hydrogen population in \( n = 2 \), \( l_{\text{He}} \) is the analogous Ly \( \alpha \) absorption mean free path and \( \sigma_{\text{1216A}} \) is the \( 2s \, ^3S \) or \( 2s \, ^1S \) photoionization cross-section at 1216 Å. We have included this process in addition to the usual photoionizations by the stellar and veiling continua. The corollary effect of depopulating the hydrogen \( n = 2 \) level is insignificant, as we have verified from the results of the calculations.

![Energy level diagram of He I](https://example.com/energy-level-diagram-he-1.png)

**Figure 2.** Energy level diagram of He I, not drawn to scale.
3.3 OI

Fig. 3 shows a schematic diagram of the energy levels of OI pertinent to our calculations and the important radiative transitions. The ground state 2p$^3$P has three fine-structure levels, with energy separations such that the forbidden transitions $^3P_1 \rightarrow ^3P_2$ and $^3P_0 \rightarrow ^3P_1$ have wavelengths of 63 and 146 $\mu$m, respectively. The upper state 3d$^3$D likewise has three fine-structure levels, but with much smaller energy separations. The 2p$^3$P transitions to 3d$^3$D$_1$, $^3$D$_2$, and $^3$D$_3$ have wavelengths almost the same as Ly$\beta$ wavelength.

With charge exchanges maintaining O/\(\text{O}^+\) close to H/\(\text{H}^+\), absorption of Ly$\beta$ photons is a fortuitous enhancement of 3d$^3$D excitation, thereby leading to strong $\lambda$11287 and $\lambda$8446 emission.

The $\lambda$11287 photon emissivity (cm$^{-3}$s$^{-1}$) resulting from this Ly$\beta$ fluorescence process (Bowen 1947) depends on the Ly$\beta$ intensity and the probability that 3d$^3$D decays via emission of a $\lambda$11287 photon. Because Ly$\beta$ is expected to be very optically thick, we assume its specific intensity to be given by the blackbody value at the hydrogen n = 1 → 3 excitation temperature, that is, $I_{\text{Ly}\beta} = (2\pi h^3/\lambda^2 c^2)(9N_1/N_3 - 1)^{-1}$, where $N_1$ and $N_3$ are the hydrogen n = 3 and 1 level population, respectively. For the probability of $\lambda$11287 emission, we consider first the upper level 3d$^3$D$_1$. It can decay to all three fine-structure levels of the ground state as well as to 3p$^3$P, so the probability of $\lambda$11287 emission upon absorption of a Ly$\beta$ photon through the $2p^3P_2 \rightarrow 3d^3D_1$ transition is

$$P = \frac{A_{\lambda11287}I_{\text{Ly}\beta}B_{\lambda11287}}{A_{3d^3D_1}B_{3d^3D_1} + A_{3d^3D_2}B_{3d^3D_2} + A_{3d^3D_3}B_{3d^3D_3} + A_{3p^3P}B_{3p^3P}},$$

where $A$ signifies the spontaneous emission rate and $\beta$ the escape probability. The $^3D_1 \rightarrow ^3P_2$ escape probability equals 1, even though the transition is most likely very optically thick, because the Ly$\beta$ opacity is even larger, so the emitted $\lambda$1025.77 photon will not be absorbed by OI. The probabilities of $\lambda$11287 emission via Ly$\beta$ absorptions through the $^3P_0 \rightarrow ^3D_2$ and $^3P_0 \rightarrow ^3D_3$ transitions can be written out analogously. The most common situation is that $\lambda$11287 is optically thin and all the $2p^3P \rightarrow 3d^3D$ transitions are sufficiently optically thick that 3d$^3$D decays to the ground state via emission of $\lambda$1026 photons that are absorbed by hydrogen. In this case, the above expression reduces to $P = A_{\lambda11287}/(A_{3d^3D_1} + A_{3d^3D_2} + A_{3d^3D_3})$ and the $\lambda$11287 photon emissivity (s$^{-1}$ cm$^{-3}$) via Ly$\beta$ pumping is simply

$$R_{\text{Ly}\beta} = \frac{5}{9}N_1 N_3 / (9N_1 - N_3) - 2.077A_{\lambda11287},$$

with the approximation that 5/9 of the OI population is in the $2p^3P_2$ level. Even if $N_3/N_1$ is as low as $10^{-2}$, $R_{\text{Ly}\beta}$ is still larger than what can be brought about through UV continuum photionization or direct collisional excitation at $T \leq 10^4$ K.

The absorption of Ly$\beta$ photons by OI has the corollary effect of depopulating the $\text{H}^+ n = 3$ population. We have included this process in the hydrogen excitation calculations, but only with the simple expression enumerated above for the most common situation in order to avoid an iterative procedure involving $\text{H}^+$ and OI calculations. With $N_0/N_1 = 4.9 \times 10^{-4}$, it is readily seen from the above expression for $R_{\text{Ly}\beta}$ that the corresponding $n = 3$ de-excitation rate is 690 s$^{-1}$. While it can rival Ly$\beta$ escape, it is much weaker than $\text{H}\alpha$ escape at low densities and collisional de-excitation at high densities, so the actual effect on the hydrogen level population is insignificant.

The Ly$\beta$ fluorescence process enhances not only $\lambda$11287 emission but, through the subsequent radiative cascade back to the ground state, also $\lambda\lambda$8446 and 1303 emission. The resulting $\lambda$11287/$\lambda$8446 flux ratio is then just the ratio of the photon energies or 0.75. However, collisional excitation favours $\lambda$8446 over $\lambda$11287 emission, so a smaller observed flux ratio is a measure of the relative contribution between...
the two processes. Absorption in \(\lambda 8446\) is occasionally seen. It is facilitated by the population built up in \(3s^3S\) through Ly\(\beta\) fluorescence and sustained via radiative trapping of \(3s1303\) photons.

The \(\lambda\lambda\lambda 7772, 7774, 7775\) triplet emission is produced from recombination and cascade, and collisional excitation. The metastable nature of the lower level \(3s^3S\), whose spontaneous emission rate is only \(5 \times 10^7\) s\(^{-1}\), sustains a comparatively large population in that level and enhances the collisional pathway of triplet emission. It likewise helps to bring about absorption of the stellar and veiling continua through the larger line opacity.

In performing the \(O\) \(I\) excitation calculations, we assume that the fine-structure states are populated in proportion to their degeneracies, and use a single level to represent them. This simplification is appropriate when the transitions between those fine-structure states and another level are separated in energy by about one thermal Doppler width or less. When this is not the case, as for the \(3s^3S \rightarrow 2p^3P\) or \(3p^3P \rightarrow 3s^3S\) transitions, the simplified procedure works fine when the transitions are optically thin but, when they are optically thick, does not take into account the availability of several radiative channels for de-excitation. Also, when collisional de-excitation dominates over radiative de-excitation, it underestimates the total emitted flux by a factor equal to the number of distinct (Doppler-width separated) lines. To remedy this situation, we re-define the line optical depth as the usual definition divided by the above factor. It represents a sort of average of the optical depths of the separate lines. For example, in our procedure, this optical depth between \(2p^3P\) and \(3s^3S\), \(\tau_{s1303}\), equals 0.6, 1.0 and 3.0 of \(\tau_{s1302.2}, \tau_{s1304.9} \) and \(\tau_{s1306}\), respectively. This modification also produces the correct total emitted flux when the distinct lines are all optically thin and when they all have optical depths greater than \(\sim 2\). There may be a small error in the level population when some of the line optical depths lie between 0.5 and 2, but this regime occupies a very narrow strip of our explored density range, and we are not overly concerned.

### 3.4 Ca \( II\), Na \( I\)

Ca \( II\), with an ionization potential of 11.87 eV, can be ionized quite readily by even the veiling continuum, if the latter has an energy distribution that extends towards the far-UV like a blackbody. For example, with our adopted veiling continuum of temperature 8000 K and area covering factor of 0.03, the photoionization rate of the \(4s^2S\) ground state at \(4R_e\) is \(3.3 \times 10^{-4}\) s\(^{-1}\), higher than our assumed UV continuum photoionization rate, which is \(\gamma_{\text{Ca}\ II} \sim (0.5/6.3)\gamma_{\text{H}}\). Moreover, the excited state \(3d^2D\), the lower state of the \(\lambda\lambda\lambda 8498, 8542, 8662\) triplet, is metastable, with an Einstein \(A\) rate of only \(1\) s\(^{-1}\), so it is most likely populated in thermal equilibrium with the ground state. Its lower ionization potential and larger photoionization cross-section produce an even stronger ionization rate, \(3.4 \times 10^{-2}\) s\(^{-1}\), with our adopted veiling continuum. As an example, if the \(4s^2S \rightarrow 3d^2D\) excitation temperature is 7000 K, ionization of Ca \( II\) from \(3d^2D\) is 30 times that from \(4s^2S\). Also, as noted in the beginning of Section 3, the \(3d^2D\) ionization potential is very slightly less than the Ly\(\alpha\) photon energy. This ionization rate by Ly\(\alpha\) photons (s\(^{-1}\)) is (cf. Section 3.2):

\[
\gamma_{\text{Ly}\alpha} = N_2 A_{\text{Ly}\alpha} \lambda_{\text{abs}} \sigma_{3d} \tau_d
\]

\[
= 5.54 \times 10^4 N_2 \left( \frac{T}{10^4 \text{K}} \right)^{0.5} \left[ \ln(\tau_{\text{Ly}\alpha} + 2.72) \right]^{0.5},
\]

where \(\sigma_{3d} \tau_d = 6.15 \times 10^{-18} \text{ cm}^2\) is the \(3d^2D\) photoionization cross-section at threshold. In the excitation calculations, the hydrogen \(n = 2\) to \(1\) population ratio, \(N_2/N_1\), has a value often larger than \(2 \times 10^{-17}\), that needed for \(\gamma_{\text{Ly}\alpha}\) to equal \(3.4 \times 10^{-2}\) s\(^{-1}\), the \(3d^2D\) photoionization rate by the veiling continuum. Thus, Ly\(\alpha\) photons constitute an even more potent source of Ca \( II\) ionization.

Our model Ca \( II\) atom consists of just the three levels \(4s^2S\), \(3d^2D\) and \(4p^3P\). Both \(3d^2D\) and \(4p^3P\) have two fine-structure states. The two \(4p^3P \rightarrow 4s^2S\) transitions are the well-known H and K lines at \(\lambda\lambda\lambda 3968\) and 3934, respectively, while the three \(4p^3P \rightarrow 3d^2D\) transitions constitute the infrared triplet. Like \(O\) \( I\), we assume the fine-structure states to be populated in proportion to their degeneracies and re-define the line optical depth to reflect the number of distinct lines. The observed line fluxes indicate that all five lines are very optically thick, so our simplified procedure of lumping into one single level the population in all associated fine-structure states is fine.

Our model Na \( I\) atom consists of just the ground level \(3s^2S\) and an excited level comprising the two fine-structure states of \(3p^2P\), whose radiative decays give rise to the well-known \(\lambda\lambda\lambda 5896, 5890\) doublet. The lower state of the doublet being the ground state helps to bring about strong absorption and/or emission in the doublet. On the other hand, Na \( I\), with an ionization potential of only 5.14 eV, is easily photoionized or collisionally ionized from the ground level, and even more so from the excited level, so the sodium ionization structure is key.

### 4 OBSERVATIONAL INFORMATION

The local excitation calculations can produce line opacities and emissivities over an extensive parameter space of physical conditions, but need information from observational data to demarcate the pertinent regions. Here we summarize the observational input on line ratios, line specific fluxes and line opacities that will be utilized to compare with model results.

For the information on hydrogen line ratios, we directly use Bary et al.’s (2008) collection of Pa \( n_s/\text{Pa}\beta\) and Br\(\gamma/\text{Pa}\alpha\) ratios, where \(n_s\), from 5 to 14, is the energy quantum number of the upper level of the Paschen transition. For line ratios involving the Ca \( II\) infrared triplet and \(\lambda\lambda\lambda 8446, 7773\), we utilize the data available in Muzerolle, Hartmann & Calvet (1999b). Three objects (DL, DG and BP Tau) are selected, because they are the only ones whose O \(\alpha\)7773 is not dominated by a strong red absorption. Because the lines considered have quite close wavelengths, we simply use the listed equivalent widths to obtain ratios among them. For the He \( I\) lines, we obtain the He \(\alpha\)10830 information from EFHK06 and the He \(\alpha\)5876 information from BEK01. To avoid the issue of estimating the true emission when an absorption feature is present, we assume the fine-structure states to be populated in proportion to their degeneracies and re-define

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present in He $\lambda$10830, only those objects among the reference sample (cf. fig. 4 of EFHK06) with emission much stronger than absorption are used. They (CW, DL, DG, HN, BP, GG and DG Tau, and RW Aur), totalling eight, are then looked up in the reference sample of BEK01 for the He $\lambda$5876 data. The He $\lambda$5876/$\lambda$10830 ratio is equal to $(1 + r_\gamma)EW_{5876}/(1 + r_\gamma)EW_{10830}$, where $r_\gamma$ denotes the veiling, $EW_\gamma$ denotes the equivalent width and $F_\gamma$ denotes the photospheric specific flux, assumed to be given by that of a 4000-K blackbody. For He $\lambda$5876, the EW of the BC is being used. In the same Echelle order of He $\lambda$10830, Pa$\gamma$ is observed, so the data set of EFHK06 conveniently provides the Pa$\gamma$/He $\lambda$10830 ratios for the same eight objects. Unfortunately, the He $\lambda$10830, 5876 data sets are procured at very different times, so the He $\lambda$5876/$\lambda$10830 ratios derived from them may have uncertainties associated with time-variation of the line emission. Partly to remedy this situation, we will also utilize the primary information on three objects (DL, DG and HN Tau) in the data set of Edwards et al. (in preparation). The He $\lambda$5876/$\lambda$10830 and Ca II $\lambda8498$/Pa$\gamma$ ratios of those three objects are particularly helpful, because the optical and 1-$\mu$m spectra are procured at the same time.

For the observed specific fluxes of lines, we make use of the He $\lambda$10830 and Pa$\gamma$ profiles in EFHK06, the He $\lambda$5876 profiles in BEK01, the Pa$\beta$ and Br$\gamma$ profiles in Folha & Emerson (2001), and the Ca II $\lambda 8542$ profiles in Muzerolle et al. (1998b). In the model calculations, the specific flux of a line at $|v_{\text{obs}}| = 150 \text{ km s}^{-1}$ is measured relative to the continuum level that is given by a 4000-K photosphere of radius 2 $R_\odot$ veiled by a 8000-K blackbody over 3 per cent of its surface. This is done to facilitate comparison with observed spectra, which are usually plotted with velocity as abscissa and strength relative to the continuum as ordinate. Normally, the veiling and photospheric temperature of the star are determined, so it is straightforward to take account of the differences in the veiling and photospheric temperature between the model and observed star. The remaining uncertainty lies in the unknown surface area of the stellar photosphere, which depends on both the stellar radius and, because of the presence of an opaque disc, the viewing angle.

The ratio of two emission lines conveys information on the line opacities in addition to the physical conditions of the density and temperature, but usually detailed excitation calculations are needed to disentangle the effects of opacity, density and temperature. An absorption feature, on the other hand, directly reveals that the line optical depth is $\tau$ or more, while an absence of a similar feature in another line indicates that its optical depth is much smaller, except when its emission is sufficiently strong to fill in the absorption. Thus, we can use the relative propensity among the observed lines in showing an absorption to constrain the physical conditions of the absorbing gas.

Clearly, He $\lambda$10830 is most proficient in showing absorption features. EFHK06 find that 47 per cent of 38 CTTSs observed show red absorptions at He $\lambda$10830 compared to 24 per cent at Pa$\gamma$ and 71 per cent show blue absorptions at He $\lambda$10830 compared to 0 per cent at Pa$\gamma$. If the blue absorptions are separated into broad and narrow ones, which are likely formed in a stellar wind and disc wind, respectively (Kwan et al. 2007), then the two kinds are present in 39 and 29 per cent, respectively, of the 38 objects. Among the 15 CTTSs observed by Edwards et al. (1994), red absorptions are seen in Na D and H$\delta$ in nine and eight objects, respectively. The rarity of H$\gamma$ red absorptions (2/15) is due to the strong H$\alpha$ emission. Blue absorptions are present in H$\alpha$ and Na D for about 50 per cent of the stars, but except for one star (As 353A), are narrow and sharp, indicative of a disc wind origin. Among the eight O I $\lambda7773$, 8446 profiles shown in Muzerolle et al. (1998b), two have red absorptions in both lines and another four have red absorptions in O $\lambda7773$ only, while none has a blue absorption. Among the 11 Ca II $\lambda8542$ profiles shown, one (DS Tau) has a red absorption and one (RW Aur) has a narrow blue absorption. BEK01 shows 31 He $\lambda$5876 profiles, of which three have red absorptions and none has a blue absorption.

From the above observational input on absorption features, we gather that a broad blue absorption indicative of a stellar wind is present quite often in He $\lambda$10830, only rarely in H$\alpha$ and Na D, and thus far, not seen in the Paschen lines, O $\lambda7773$, 8446, He $\lambda$5876 or the Ca II infrared triplet. In decreasing probability of showing a red absorption, the order of lines is roughly He $\lambda$10830, Balmer lines, Na D, O $\lambda7773$, Paschen lines, O $\lambda8446$, He $\lambda$5876 and Ca II infrared triplet. The rarity of red absorptions in the Ca II infrared triplet, however, does not indicate that the lines are optically thin. They are actually optically thick because of their nearly equal emission strengths, despite a factor of 5 between $\tau_{\text{CaII}}=6622$ and $\tau_{\text{CaII}}=8498$. Often their emission is very strong, as measured in relation to the continuum level, so it is likely the emission will fill in the red absorption produced from scattering the stellar and veiling continua. It also happens that Pa 13, Pa 15 and Pa 16 lie on the red side (at $\lambda 120 \text{ km s}^{-1}$) of the Ca II infrared triplet. In decreasing probability of showing a red absorption, the order of lines is roughly He $\lambda$10830, Balmer lines, Na D, O $\lambda7773$, Paschen lines, O $\lambda8446$, He $\lambda$5876 and Ca II infrared triplet. The rarity of red absorptions in the Ca II infrared triplet, however, does not indicate that the lines are optically thin. They are actually optically thick because of their nearly equal emission strengths, despite a factor of 5 between $\tau_{\text{CaII}}=6622$ and $\tau_{\text{CaII}}=8498$. Often their emission is very strong, as measured in relation to the continuum level, so it is likely the emission will fill in the red absorption produced from scattering the stellar and veiling continua. It also happens that Pa 13, Pa 15 and Pa 16 lie on the red side (at ~120 km s$^{-1}$) of Ca II $\lambda$8662, 8542 and 8498, respectively, and their contribution will further obliterate a red absorption.

## 5 LINE OPACITIES AND IMPORTANCE OF UV PHOTOIONIZATION

### 5.1 Model results

Fig. 4 shows the optical depth contours of four important lines in the ($N_\text{H}$, $T$) plane for $r = 4R_\odot$, and $\gamma_{\text{HeI}} = 10^{-4}$ (top panel) and $10^{-5}$ s$^{-1}$ (bottom panel). They are $\tau_{\text{HeI}} = 1$ and 3.16, $\tau_{\text{HeII}/\lambda10830} = 1$ and 3.16, $\tau_{\text{Pa}} = 0.1$ and 0.316, and $\tau_{\text{HeI}/\lambda5876} = 0.1$ and 0.316, with the right-hand contour of each pair having the higher value. The different contour levels between (H$\alpha$, He $\lambda$10830), on the one hand, and (Pa$\gamma$, He $\lambda$5876), on the other hand, reflect their relative tendencies to show an absorption. They help to locate the appropriate region in the ($N_\text{H}$, $T$) space causing absorption in one line, but not in another. For example, to produce a discernible blue absorption in He $\lambda$10830, but not in Pa$\gamma$, the bulk of the gas needs to have an $N_\text{H}$ between the $\tau_{\text{HeII}/\lambda10830} = 1$ and $\tau_{\text{Pa}} = 0.1$ contours. For $\gamma_{\text{HeI}} = 10^{-4}$ s$^{-1}$, the leeway in $N_\text{H}$ is a factor of between 10 and 50, depending on $T$.

The He I and H I optical depth contours demonstrate the role of photoionization in He I and H I excitation. Fig. 4 shows that at $T$ above 2.25 $\times$ 10$^4$ K (10$^5$ K), the He I (H I) contours are almost identical between the two cases of $\gamma_{\text{HeI}}$ ($\gamma_{\text{HeI}}$). This occurs because for $\gamma_{\text{HeI}} = 10^{-4}$ s$^{-1}$ ($\gamma_{\text{HeI}} = 2 \times 10^{-4}$ s$^{-1}$) and an $N_\text{e}$ of 10$^9$ cm$^{-3}$, the rate of collisional excitations to the He $\lambda 10830$(H$\alpha$) lower level will equal the rate of population via photoionization at $T = 2.25 \times 10^4$ K (10$^5$ K). Thus, above this temperature, He I (H I) photoionization at the assumed strength...
and Ca contours also depend on \( N_T \) line emission and \( \tau = 1 \) and 3.16 contours, and He \( \lambda 5876 \) and Pa\( \gamma \) have \( \tau = 0.1 \) and 0.316 contours.

The above rather surprising result is caused by the He \( \lambda 10830 \) lower level being extremely metastable, making spontaneous decay ineffective in comparison with depopulation via collisions with electrons for \( N_e > 10^4 \) cm\(^{-3} \). As mentioned in Section 3.2, collisional excitation to the singlet 2\( s \) is more important than direct collisional de-excitation to the ground state. The rate of this depopulation route is then \( \propto N_e \gamma_{9240}/T \), and the He \( \lambda 10830 \) optical depth, in the limit when photoionization dominates collisional excitation, can be expressed as

\[
\tau_{\text{He}\lambda 10830} \propto \frac{N_{\text{He}}}{N_e} \frac{\gamma_{9240}/T}{N_e} \gamma_{H \alpha} e^{-9240K/T}. \tag{10}
\]

As \( T \) decreases, both \( N_{\text{He}}/N_e \) and \( e^{9240K/T} \) increase, while \( N_e/N_{\text{He}} \) decreases, and they work in concert to make \( \tau_{\text{He}\lambda 10830} \) larger, thereby enabling a given optical depth to be obtained at a lower \( N_{\text{He}} \), hence the reversal of the \( \tau_{\text{He}\lambda 10830} = 1 \) contour direction seen distinctly in the bottom panel of Fig. 4.

An illustration of how the electron fraction \( N_e/N_{\text{He}} \) depends on the physical parameters will help understand the behaviours of the optical depth contours as well as results presented later. This fraction is contributed primarily by hydrogen ionization, but its dependences on \( T \) and \( N_{\text{He}} \) are not straightforward. In Fig. 5, we plot \( N_e/N_{\text{He}} \) versus log \( (N_{\text{He}}) \) for seven temperatures and \( \gamma_{\text{H \alpha}} = 2\gamma_{\text{He}\alpha} = 2 \times 10^{-8} \) s\(^{-1} \) (top panel) and 2 \( \times 10^{-5} \) s\(^{-1} \) (bottom panel). With the expectation that \( N_e/N_{\text{He}} \) is lower when \( \gamma_{\text{H \alpha}} \) is lower, the dependences of \( N_e/N_{\text{He}} \) on \( N_{\text{He}} \) and \( T \) in the two panels are qualitatively the same. At 8750 \( \leq T \leq 10^5 \) K, \( N_e/N_{\text{He}} \) initially falls with \( N_{\text{He}} \) increasing from \( 10^9 \) cm\(^{-3} \), as hydrogen ionization through UV photoionization dominates. It levels off and rises slowly when photoionization from \( n = 2 \), bolstered by \( N_{\text{He}}/N_{\text{H}} \) increasing with increasing \( N_{\text{He}} \) and Ly\( \alpha \) trapping, grows stronger and rises faster, once sufficient population is built up into the high \( n \) levels for collisional ionization to take over. At \( T \leq 7500 \) K, the contributions from excited state ionizations only slow down the rate at which \( N_e/N_{\text{He}} \) decreases with increasing \( N_{\text{He}} \).

\[\]
Figure 5. Dependence of electron fraction $N_e/N_H$ on $N_H$ for $r = 4R_*$, $\gamma_{H_I} = 2 \times 10^{-4}$ (top panel) and $2 \times 10^{-5}$ s$^{-1}$ (bottom panel), and seven temperatures (in units of $10^4$ K).

The quantitative values of $N_e/N_H$ make it easy to see the behaviour of the $\tau_{He\alpha,10830}$ contours. At $\gamma_{He_I} = 10^{-5}$ s$^{-1}$, the $\tau_{He\alpha,10830} = 1$ contour near $10^4$ K lies at $N_H \sim 10^{10}$ cm$^{-3}$. From Fig. 5 (bottom panel), it is seen that $N_e/N_H$ decreases with $T$ decreasing from $10^4$ to $7500$ K. This sharp decrease is largely responsible for the switchback of the $\tau_{He\alpha,10830} = 1$ contour towards a lower $N_H$. The same figure also helps to explain the large separation between the $\tau_{He\alpha,10830} = 1$ and 3.16 contours seen at $\gamma_{He_I} = 10^{-5}$ s$^{-1}$ and $8750 \leq T \leq 10^4$ K. This arises because in that temperature range, increasing $N_H$ above $10^{10}$ cm$^{-3}$ leads to $N_e/N_H$ increasing also, which counteracts the effect of $N_{He_I}/N_{He}$ in the expression for $\tau_{He\alpha,10830}$. As a result, $\tau_{He\alpha,10830}$ barely changes for $3 \times 10^6 \leq N_H \leq 10^{12}$ cm$^{-3}$. This large separation between the $\tau_{He\alpha,10830} = 1$ and 3.16 contours is not seen in the $\gamma_{He_I} = 10^{-4}$ s$^{-1}$ case, because in the same temperature range, the $\tau_{He\alpha,10830} = 3.16$ contour lies at $N_H < 10^6$ cm$^{-3}$.

At $T \leq 7500$ K, the $\tau_{He\alpha,10830}$ contours for the two values of $\gamma_{He_I}$ are not far apart. This insensitivity of $\tau_{He\alpha,10830}$ to decreasing $\gamma_{He_I}$ at low temperatures may explain why $He\alpha,10830$ shows such a penchant for absorption. Seen in $He\alpha,10830$ are not only the broad blue and red absorptions that arise from kinematic regions close to the star, but also the narrow, sharp blue absorption indicative of a disc wind and the central absorption that has been suggested as arising from a disc corona even farther away from the star (Kwan 1997). Presumably, the ionization sources are located close to the star and, even if intervening attenuation is ignored, $\gamma_{He_I}$ will decrease with distance. It turns out that the density required to produce $\tau_{He\alpha,10830}$ of unity is not prohibitive. To see this, we find that if $N_e$ is supplied through UV photoionization of hydrogen, as likely at $T < 7500$ K, then $N_e \propto (N_H \gamma_{H_I})^{0.5}$ and

$$\tau_{He\alpha,10830} \propto (N_H \gamma_{He_I})^{0.5} \left(\frac{\gamma_{He_I}}{\gamma_{H_I}}\right)^{0.5} e^{0.240K/T}.$$

Thus, $\tau_{He\alpha,10830}$ is more sensitive to $T$ for $T \leq 7500$ K than to $\gamma_{He_I}$. Also $\gamma_{He_I}/\gamma_{H_I}$ can be quite large if the UV continuum near the hydrogen ionization threshold is attenuated more severely by intervening hydrogen. In the limit that both helium and hydrogen are ionized by photons with energies $\geq 24.6$ eV, $\gamma_{He_I}/\gamma_{H_I}$ is $\sim 6$, which is a factor of 12 larger than the earlier assumed value of 0.5. This means that $\gamma_{He_I}$ can be smaller than $10^{-5}$ s$^{-1}$ by a factor of 12 and still produce at $T = 5000$ K $\tau_{He\alpha,10830}$ of 1 with $N_H = 5 \times 10^9$ cm$^{-3}$. Even lower values of $\gamma_{He_I}$ are possible, if $T$ is lower. The temperature at which direct collisional de-excitation of $3S$ occurs as rapidly as the postulated de-population route is $2700$ K, so $\gamma_{He_I}$ can be smaller by another factor of $\sim 20$ when direct collisional de-excitation dominates.

In Fig. 4, the pair of contours for each line purports to illustrate its sensitivity to varying $N_H$. It is seen that the $H\alpha$, $Pa\alpha$ and $He\alpha,5876$ optical depths increase faster than linearly with increasing $N_H$. This is because the population in each of their lower level is sustained via radiative trapping, so for $H\alpha$, for example, $\tau_{H\alpha} \propto N_2 \propto N_{H_I}(\gamma_{H\alpha} + N_{He} C_{12})$ $\tau_{Pa\alpha}/A_{He\alpha} \propto N_{H_I}(\gamma_{He_I} + N_{He} C_{12})$. The $He\alpha,10830$ optical depth, by comparison, is less responsive to increasing $N_H$, particularly at $T$ between 8750 and $10^4$ K when $\gamma_{He_I} = 10^{-5}$ s$^{-1}$, as discussed earlier.

The $H\alpha$ contours are dependent on $\gamma_{H_I}$, over its explored range only at $T < 10^4$ K, as mentioned before, and then only weakly so. The $Pa\gamma$ contours are more sensitive, because the $n = 3$ population relies on the $n = 2$ population being built up first. Above $10^4$ K, the hydrogen
contours move towards lower densities with increasing temperature, owing to the stronger collisional excitations from the ground level, but then reverse direction at $T \sim 1.5 \times 10^4$ K. This reversal is caused by the increase in collisional ionizations, which rapidly reduces the hydrogen neutral fraction.

Fig. 6 shows the contours $\tau_{\text{Na} \lambda 5892} = 1$, $\tau_{\text{Ca} \lambda 3945} = 1$, $\tau_{\text{Ca} \lambda 8498} = 1$, $\tau_{\text{O} \lambda 8446} = 0.1$ and $\tau_{\text{O} \lambda 7773} = 0.1$. The $\tau_{\text{Pa} \gamma} = 0.1$ contour is also re-plotted for comparison. For the line optical depth that represents a multiplet with well-separated components (in Doppler widths), it needs be mentioned that, with our single level stand-in, the plotted optical depth is a rough average of the individual component optical depths. To be specific,

$$
\begin{align*}
\left( \tau_{\text{Na} \lambda 5890}, \tau_{\text{Na} \lambda 5890} \right) / \tau_{\text{Na} \lambda 5892} & = (1.33, 0.67), \\
\left( \tau_{\text{Ca} \lambda 3934}, \tau_{\text{Ca} \lambda 3961} \right) / \tau_{\text{Ca} \lambda 3945} & = (1.33, 0.67), \\
\left( \tau_{\text{Ca} \lambda 8498}, \tau_{\text{Ca} \lambda 8542}, \tau_{\text{Ca} \lambda 8662} \right) / \tau_{\text{Ca} \lambda 8498} & = (0.2, 1.8, 1.0), \\
\left( \tau_{\text{O} \lambda 7772}, \tau_{\text{O} \lambda 7774}, \tau_{\text{O} \lambda 7775} \right) / \tau_{\text{O} \lambda 7773} & = (1.4, 1.0, 0.6).
\end{align*}
$$

The wavelength used to denote the Ca II infrared triplet as a group is the same as that of one member, because that particular member is used to produce observed line ratios for comparison with model results. A similar denotation is not adopted for the O I $\lambda 7773$ triplet, because its three members are separated by only $\sim 200$ km s$^{-1}$ and all are seen within one line profile, so their total flux is used.

It is seen that the Ca II and Na I optical depth contours exhibit little difference between the two UV ionization fluxes, as the Ca II and Na I fractions depend only indirectly on the postulated UV continuum through the resulting $N_e$ that affects the recombination rate and, in the case of Ca II, also through the hydrogen $N_2/N_1$ ratio that affects the Ly$\alpha$ photoionization rate. The Ca II and Na I contours also respond similarly to $T$ increasing from 5000 K, owing to their common susceptibility to collisional ionization. The O I ionization structure is tied to hydrogen’s via charge exchange, so the O I optical depth contours at $T \lesssim 10^4$ K respond to varying $\gamma_{\text{H} \alpha}$ in much the same way as the hydrogen optical depth contours. Above $T = 10^4$ K, the $\tau_{\text{O} \lambda 7773}$ contour stays at nearly the same $N_{\text{H}}$, while the $\tau_{\text{O} \lambda 8446}$ contour moves towards much higher densities. This is due to the lower state of O I $\lambda 7773$ being metastable, while that of O I $\lambda 8446$ decaying rapidly with $\tau_{\text{O} \lambda 1303}$ decreasing with increasing temperature.

Both Figs 4 and 6 refer to local excitation calculations at $r = 4R_e$, a location appropriate for either a stellar wind or the farther portion of an accretion flow. Fig. 7 shows similar optical depth contours for $\gamma_{\text{He} \alpha} = 10^{-4}$ s$^{-1}$ and $r = 2.5R_e$ to evaluate the dependence on position. Comparing it with Figs 4 and 6 (top panels) indicates that the differences are quite small. For fixed values of $N_e$ and $T$, the H I, He I and O I optical depths are somewhat smaller at $r = 2.5R_e$, because of the factor $r/v$ in the opacity expression; hence, their contours are displaced towards higher densities, with the amounts dependent on the responses of the line optical depths to increasing $N_e$. For the Na I and Ca II optical depths, there is the additional effect of a stronger stellar and veiling continuum flux at $r = 2.5R_e$, which increases the photoionization.
rate and also leads to a displacement of their contours towards higher densities. From Figs 4, 6 and 7 we can infer that the line optical depths and, by analogy, line emissivity ratios, which will be discussed in the next section, are much more dependent on \( N_{\text{H}} \), \( T \) and \( \gamma_{\text{HeI}} \) than on \( r \).

5.2 Comparison with observations

Keeping in mind the observed propensities of the various lines in showing an absorption, we are in a position to draw the following conclusions from the model results presented in this section:

1. UV photoionization is necessary to produce the broad red and blue absorptions seen in He \( \lambda 10830 \), because both gases have a temperature \( \sim 10^4 \) K or less. The most direct argument comes from the observed red absorptions. They are seen in He \( \lambda 10830 \) and Na \( \lambda 5892 \) frequently, and Pa\( \gamma \) occasionally. For Na\( \lambda \) absorption to be more prevalent than Pa\( \gamma \) absorption, the temperature needs to be less than \( 10^4 \) K, as a comparison between the Pa\( \gamma \) and Na\( \lambda \) optical depth contours indicates. Then, for He \( \lambda 10830 \) absorption to be as prevalent as Na\( \lambda \) absorption at \( T < 10^4 \) K, UV continuum photoionization is paramount. In the next two sections, one result drawn from the analyses will be that the gas producing the broad blue absorption has a temperature not higher than \( \sim 10^4 \) K, which is also the condition needed for a broad blue absorption to be present in Na\( \lambda 5892 \) but not Pa\( \gamma \).

2. In the temperature range of \( T < 10^4 \) K, we can compare the observed order of lines in their propensity of showing a red absorption to the order gleaned from the model results by judging the available volume in the \((N_{\text{H}}, T, \gamma_{\text{HeI}})\) space, where the line optical depth exceeds 0.1, as reflecting the propensity. The latter decreases roughly in the order of He \( \lambda 10830 \), Balmer lines, Ca\( \lambda \) \( \lambda 3945 \), Na\( \lambda 5892 \), Paschen lines, Ca\( \lambda \) \( \lambda 8498 \), O\( \lambda \) \( \lambda 8446 \), O\( \lambda \) \( \lambda 7773 \) and He\( \lambda 5876 \). It needs to be noted that the range in \( \gamma_{\text{HeI}} \) is specifically chosen to place He \( \lambda 10830 \) close to the top of the order, so \( \gamma_{\text{HeI}} \sim 10^{-5} \) s\(^{-1} \) or higher is needed. Keeping in mind the explanation given in Section 4 regarding the lack of red absorptions in Ca\( \lambda \) \( \lambda 8498 \), this order is close to that observed except for the position of O\( \lambda 7773 \). This line, however, is a triplet and, with its members separated by only \( \sim 200 \) km s\(^{-1} \), is more effective than a single line in producing an absorption because of its greater velocity coverage, and this interesting property of O\( \lambda 7773 \) more than compensates for its smaller optical depth in relation to O\( \lambda 8446 \). The occasional presence of red absorptions in O\( \lambda 8446 \) and He\( \lambda 5876 \) also indicates a density \( N_{\text{H}} \sim 10^{11} \) cm\(^{-3} \) or higher for the red absorption gas.

3. The density of the gas producing the broad blue absorption is much lower than that of the accretion flow. This is obvious if the two flows have similar temperatures over most of their volumes, since Pa\( \beta \), Pa\( \gamma \) and O\( \lambda \) \( \lambda 8446 \), 7773 show red absorptions, but almost never blue absorptions. But even if that were not the case, the absence of Pa\( \beta \) and Pa\( \gamma \) blue absorptions would by themselves require, at \( T > 10^4 \) K, \( N_{\text{H}} < 3 \times 10^{8} \) cm\(^{-3} \). A rough estimate, based on Figs 4, 6 and 7, of the typical gas density occupying the bulk of the stellar wind and accretion flow is \( 10^9 \) and \( 10^{11} \) cm\(^{-3} \), respectively. The corresponding mass flux, however, can be quite similar, as the product of the factors \( r^2 \) and \( F_{\Omega} \) is likely 30–100 times larger for the stellar wind.
(4) When $\tau_{\text{He}I \lambda 10830} \sim 1$, the H$\alpha$ optical depth is close to or higher than 1 for $T \leq 2 \times 10^4$ K and $\gamma_{\text{He}I} \sim 10^{-4}$ s$^{-1}$ or less (cf. Fig. 4). Thus, we expect H$\alpha$ to be as effective as He$\lambda$10830 in absorbing the stellar and veiling continua. The rarer occurrence of broad blue or red absorption in H$\alpha$ is attributed to the strength of the H$\alpha$ emission, which is considerably greater than the continuum level.

6 LINE RATIOS AND DIFFERENT SETS OF PHYSICAL CONDITIONS

Line ratios are very good diagnostics of physical conditions, because while the observed flux of a line depends on the local line emissivity and the volume of emission, the flux ratio of two lines bypasses to a large extent the effect of emission volume and probes directly the local physical conditions. Here we examine several pairs of line ratios that are particularly illuminating in this regard.

The first pair is He$\lambda$5876/$\lambda 10830$ and Pay/He$\lambda / \lambda 10830$. The nature of He$\lambda 10830$ being the only allowed radiative transition following the $\lambda 5876$ transition means that emission of a $\lambda 5876$ photon is usually followed by emission of a $\lambda 10830$ photon, but not vice versa, while collisional excitation from $2s^3S$ strongly favours $\lambda 10830$ emission because of both the lower excitation energy and the larger collisional cross-section. The He$\lambda 5876$/$\lambda 10830$ ratio is then expected to be small at low densities. It will rise only when the $\lambda 10830$ transition becomes sufficiently optically thick that at the same time collisional de-excitation becomes competitive with $\lambda 10830$ escape, population is built up into $2p^3P$ from where collisional excitation of $\lambda 5876$ is more effective. It is therefore very sensitive to density and, with $\lambda \lambda \lambda 10830$, 5876 being the strongest observed helium lines, clearly the most important diagnostic of physical conditions giving rise to He emission. To find out how the hydrogen emission fares in the same conditions, we want to contrast He$\lambda 5876$/He$\lambda 10830$ against a ratio involving a hydrogen line and a helium line. The Pay/He$\lambda / \lambda 10830$ ratio is the most logical choice, since, as mentioned in Section 3, its observed value of close to unity appears incongruous with their perceived opacities. It also has the advantage of being quite easily and accurately measured, because both lines occur in the same echelle order.

It is possible and highly probable that hydrogen emission arises from more than one kinematic region, because the physical conditions conducive to strong hydrogen emission are likely to be more wide ranging than those conducive to either strong helium or Ca II emission. Ratios involving only hydrogen lines are therefore useful to ascertain the relevant region in the $(N_H, T, \gamma)$ parameter space. We will examine first Pay/Paβ and Pay/H$\alpha$ in tandem to see, for a given $\gamma$, how they depend on $N_H$ and T, and correlate with each other. The Pay/Paβ ratio is interesting, because there is a substantial amount of recent data pointing to a fairly uniform but somewhat surprising value (Bary et al. 2008), while Pay/H$\alpha$ involves both the strongest optical line observed and the line selected earlier to compare with He$\lambda 10830$. These two ratios involve hydrogen levels up to $n = 6$. We will also compare Pa$\lambda 5876$/Paβ with observed ratios for $n_u$ up to 12 so the observed hydrogen emission from high $n$ levels is also brought into play.

The next two pairs of ratios focus on the Ca II infrared triplet emission in relation to hydrogen and oxygen emission. We examine first Ca$\lambda 8498$/Pay and O$\lambda 8446$/Pay, and then Ca$\lambda 8498$/O$\lambda 8446$ and O$\lambda 7773$/Ca$\lambda 8446$. The latter pair has the added bonus that the three wavelengths involved are in close proximity, so the observed ratios have less uncertainties.

In the following three sections, we study in order the above-mentioned line ratios. In each section, we present first the results of the model calculations and then compare them with observations to demarcate the prevailing physical conditions.

6.1 He$\lambda 5876$/He$\lambda 10830$ and Pay/He$\lambda 10830$

6.1.1 Model results

Fig. 8 shows how He$\lambda 5876$/He$\lambda 10830$ and Pay/He$\lambda 10830$ depend on $N_H$ (top panel) or line optical depth (bottom panel) for $r = 4R_s$, $\gamma_{\text{He}I} = 10^{-4}$ s$^{-1}$ and various temperatures. In this and subsequent figures, we use different line types to denote different temperatures. Each of six line types signifies the same temperature in all the relevant figures, but a seventh line type, dot–long dashed, can signify a different temperature in a different figure in order to allow for more flexibility in the selection of appropriate temperatures to illustrate. Table 1 lists the designations between line types and temperatures, and serves as a reference. When a model track in a figure has space nearby, its temperature (in units of $10^4$ K) is marked for ready identification.

The most important physical process influencing the He$\lambda 5876$/He$\lambda 10830$ ratio is collisions with electrons. The summed rate of $2s^3S \rightarrow 3d^3D$ and $2s^3S \rightarrow 4f^3F$ collisions, which lead to He$\lambda 5876$ emission, is much weaker than the summed rate of collisions from $2s^3S$ that lead to He$\lambda 10830$ emission. In the regime where level de-excitation is dominated by radiative decay, the He$\lambda 5876$/He$\lambda 10830$ ratio resulting from collisional excitations ranges from $1.3 \times 10^{-3}$ at $T = 5000$ K to $4 \times 10^{-2}$ at $T = 2 \times 10^4$ K. The ratio resulting from recombination and cascade is higher. When the cascade ends at $2s^3S$, roughly two $\lambda 10830$ photons are produced for each $\lambda 5876$ photon. However, $2s^3S$ de-excites to the ground level primarily through collisions and the ratio of $2s^3S \rightarrow$ triplets to $2s^3S \rightarrow$ singlets collisional rate gives the number of additional $\lambda 10830$ photons emitted before each de-excitation of $2s^3S$. This number is 2.4, 7.1 and 12.2 at $T = 5000$, $10^4$ and $2 \times 10^4$ K, respectively. Thus, even if helium excitation were all caused by UV photoionization of the ground state, the He$\lambda 5876$/He$\lambda 10830$ ratio would only be 0.27, 0.1 and 0.07 at $T = 5000$, $10^4$ and $2 \times 10^4$ K, respectively. A larger ratio can only be obtained in the regime when the He$\lambda 10830$ optical depth is sufficiently high that collisional de-excitation begins to rival $\lambda 10830$ escape and effectively retards the growth of the He$\lambda 10830$ flux.

The detailed calculations bear out the above analysis. In Fig. 8, results for $T \leq 2 \times 10^4$ K are shown, because those for $T = 3 \times 10^4$ and $2 \times 10^4$ K are almost the same. At low densities, He$\lambda 5876$/He$\lambda 10830$ is very small because of the prominent contribution to the He$\lambda 10830$ emissivity from stimulated absorption of the stellar and veiling continua. With increasing $N_H$, He$\lambda 5876$/He$\lambda 10830$ roughly levels off at a value...
Figure 8. Dependences of $\text{He} \, \lambda$5876/\lambda10830 and $\text{Pa} \, \gamma$/He\,$\lambda$10830 on the density and line optical depth for $r = 4R_\odot$, $\gamma_{\text{He} \, I} = 10^{-4}$ s$^{-1}$ and various temperatures. See Table 1 for the designations between line types and temperatures.

determined by the relative contribution between collisional excitation, and recombination and cascade to the line emissivity. Then, depending on $T$, He$ \lambda$5876/\lambda10830 begins to rise sharply when $\tau_{\text{He} \, \lambda5876}$ exceeds between 1 and 10, while $\tau_{\text{He} \, \lambda10830}$ is between 30 and 350. The actual criterion for collisional quenching of He$ \lambda$10830 emission is $N_{\text{He}} C_{2p}^{1s_{-2s}} \sim A_{\text{He} \, \lambda10830} \beta_{\text{He} \, \lambda10830}$ and it is seen from Fig. 8 that the density $N_{\text{He}}$ needed to meet this criterion drops from $\sim 10^{12}$ cm$^{-3}$ at $T = 5000$ K to $\sim 3 \times 10^{10}$ cm$^{-3}$ at $T = 2 \times 10^4$ K.

Table 1. Line types and designated temperatures.

<table>
<thead>
<tr>
<th>Line type</th>
<th>$T$ ($10^4$ K)</th>
<th>Figure</th>
</tr>
</thead>
<tbody>
<tr>
<td>Short dashed–long dashed</td>
<td>0.5</td>
<td>5, 8, 9, 10, 11, 12, 15, 16, 19, 20, 21, 22, 23, 24</td>
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<tr>
<td>Solid</td>
<td>1.0</td>
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</tr>
<tr>
<td>Dotted</td>
<td>1.5</td>
<td>5, 8, 9, 10, 11, 12, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24</td>
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<tr>
<td>Dot–short dashed</td>
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<tr>
<td>Dot–long dashed</td>
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<tr>
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<td>3.0</td>
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</tr>
<tr>
<td>Dot–long dashed</td>
<td>0.625</td>
<td>5, 8, 9, 10, 11, 12, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24</td>
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Turning to Pay/He λ10830, it is seen that its dependences on \(N_H\) and \(T\) are drastically different. At a given temperature, Pay/He λ10830 peaks at a lower \(N_H\) than that marking the sharp rise of He λ5876/λ10830. Over the range of \(N_H\) explored, Pay/He λ10830 is strongest at \(T\) between 8750 and 10\(^{4}\) K and weakest at \(T \geq 2 \times 10^{4}\) K, while it is the exact opposite for He λ5876/λ10830. Thus, these two ratios anticorrelate. This relationship is readily understood when examined with the temperature range separated into two regimes.

At \(T \geq 10^{4}\) K, the He λ10830 emissivity is clearly very sensitive to temperature, since the 1s \(^3\)S \(\rightarrow\) 2s \(^3\)S collisional excitation rate, which supplements UV photoionization in populating 2s \(^3\)S, and collisional excitation of He λ10830 from 2s \(^3\)S are both strongly dependent on temperature. The He λ5876 emissivity is even more responsive to increasing temperature because of both a higher excitation energy from 2s \(^3\)S and its reliance on the build-up of the 2s \(^3\)S population. Consequently, the He λ5876/λ10830 ratio increases with \(T\) from 10\(^{4}\) to 2 \(\times\) 10\(^{4}\) K. It does not change much with \(T\) between 2 \(\times\) 10\(^{3}\) and 3 \(\times\) 10\(^{3}\) K, because the reduction in \(N_{He}/N_H\) due to ionization begins to take effect. The Pay emissivity also benefits from the higher collisional excitation rate when \(T\) increases from 10\(^{3}\) K, but the concurrent decreases in \(N_{He}/N_H\) (due to ionization) and Ly\(\alpha\) opacity retard the build-up of population into the excited levels. A confirmation of this very different behaviour of the excited-state population between \(\text{H} I\) and \(\text{He} I\) can be seen from Fig. 4, since the γ_{He,λ5876} and γ_{Pay} contours are essentially contours of the population in He \(2p^3\)P (lower state of the He λ5876 transition), \(N_{2p}\), and \(H1 n = 3\), respectively. As \(T\) increases above 10\(^{4}\) K, the \(N_{2p}\) contour shifts towards decreasing \(N_{He}\), indicating an increasing \(N_{2p}/N_{He}\) ratio at a fixed \(N_{He}\), and then hardly shifts between \(T = 2.25 \times 10^{4}\) and 3 \(\times\) 10\(^{4}\) K, while the \(N_{3}\) contour shifts very little between \(T = 10^{4}\) and 1.5 \(\times\) 10\(^{4}\) K, but then shifts towards increasing \(N_{He}\) with further increase in \(T\). This contrasting behaviour accounts for Pay/He λ10830 decreasing while He λ5876/λ10830 increases with \(T\) increasing above 10\(^{4}\) K.

As \(T\) decreases from 10\(^{4}\) K, the ionization of hydrogen begins to rely more on UV photoionization and the electron fraction \(N_e/N_H\) decreases with decreasing \(T\) (cf. Fig. 5). Both decreases in \(N_e/N_H\) and \(T\) reduce the He \(1s \^3S\) de-excitation rate, as alluded to earlier in Section 5.1, and the consequent higher 2s \(^3\)S population counteracts the decrease in the 2s \(^3\)S \(\rightarrow\) 2p \(^3\)P collisional excitation rate. As a result, the He λ10830 emissivity does not decrease as rapidly with decreasing temperature as the Pay emissivity and Pay/He λ10830 drops.

At a given temperature, Pay/He λ10830 rises rapidly as \(N_{He}\) increases from 10\(^{6}\) cm\(^{-3}\), peaks at an \(N_{He}\) between 10\(^{10}\) and 10\(^{11}\) cm\(^{-3}\), and then falls. The rapid rise is related to the high sensitivity to \(N_{He}\) of the hydrogen population build-up into upper levels, relying on radiative trapping of successive Δ\(n\) = 1 transitions. The rise to the peak corresponds to \(\tau_{Pay}\) increasing from less than 1 to a value between 2 and 20. Further increase in \(\tau_{Pay}\) reduces the photon escape probability and slows the growth of Pay emissivity, leading to a falling Pay/He λ10830 ratio. Curbing of Pay emissivity occurs before that of He λ10830 emissivity, because \(C_{6s-5s} \sim 35C_{He2p-1s-2s}\).

The earlier discussion in Section 5.1 on how \(\gamma_{He}\) affects the \(\text{H} I\) and \(\text{He} I\) optical depth contours also hints at the corresponding responses of the He λ5876/λ10830 and Pay/He λ10830 ratios, which are shown in Fig. 9 for \(\gamma_{He} = 10^{-5}\) s\(^{-1}\). As expected, He λ5876/λ10830 is affected little and then only at \(T < 10^{4}\) K, since both line emissivities, like the line optical depths, respond similarly to varying \(\gamma_{He}\), while Pay/He λ10830 clearly peaks higher at 8750 \(\leq\) \(T\) \(\leq\) 1.25 \(\times\) 10\(^{4}\) K, the temperature range over which the He 1s \(^3\)S population is most susceptible to \(\gamma_{He}\). The basic contrast between the two ratios in their dependences on \(N_H\) and \(T\), however, remains the same.

### 6.1.2 Comparison with observations

The anticorrelation between He λ5876/λ10830 and Pay/He λ10830 is more conspicuous when the two ratios are plotted against each other, as shown in Fig. 10. The logarithmic scale is used for greater clarity of the model tracks which will be bunched towards the x- or y-axis in a linear plot. He λ5876/λ10830 is clearly high when Pay/He λ10830 is low and vice versa. Many of the observed values that are marked lie at positions very much off the tracks that model results produce as \(N_{He}\) increases at a fixed \(T\). The track closest to them is generated at \(T = 5000\) K and it is more easily seen from Fig. 8 that He λ5876/λ10830 is slightly higher than 0.2 over the range 5 \(\times\) 10\(^{10}\) \(\leq\) \(N_{He}\) \(\leq\) 10\(^{12}\) cm\(^{-3}\) when Pay/He λ10830 drops from \(\sim\) 1.0 to 0.4. Several arguments, however, point to a temperature of the emitting gas much higher than 5000 K. They are based on ratios of hydrogen lines and on line fluxes, and will be propounded when the pertinent results are presented. The remaining way to explain the observed ratios is for He λ5876 and Pay emission to arise from two different locales, for example, a region of \(T = 2 \times 10^{6}\) K for the former and one of \(T = 10^{4}\) K for the latter. The He λ5876/λ10830 and Pay/He λ10830 ratios in their respective locales are then actually higher than the observed ratios, since both locales will produce He λ10830 emission efficiently. For example, Fig. 8 shows that at \(T = 2 \times 10^{6}\) K, Pay/He λ10830 is less than 0.125 for \(N_{He} > 3 \times 10^{10}\) cm\(^{-3}\), while at \(T = 10^{4}\) K, He λ5876/λ10830 is less than 0.06 for \(N_{He} < 3 \times 10^{13}\) cm\(^{-3}\). Then, if the observed He λ10830 emission is contributed equally by each locale, to produce observed He λ5876/λ10830 and Pay/λ10830 ratios of 0.23 and 0.6, respectively, the 2 \(\times\) 10\(^{4}\) K (10\(^{4}\) K) locale will need to produce a He λ5876/λ10830 (Pay/He λ10830) ratio of 0.4 (1.08).

In a way, the model results confirm the perceived enigma mentioned in Section 3 concerning He λ10830 and Pay. They demonstrate that a common set of physical conditions, a natural and reasonable premise implicit in formulating the enigma, cannot produce at the same time the He λ10830 and Pay emission strengths and their relative opacities as suggested by their propensities in showing an absorption.

We draw the following conclusions concerning the observed strong He λ5876, 10830 and Pay emission:

1. He λ5876 and Pay emission does not arise from the same locale. In particular, the gas emitting He λ5876 has a temperature close to 2 \(\times\) 10\(^{4}\) K, while that emitting Pay has a temperature close to 10\(^{4}\) K.
2. The gas producing the He λ5876, 10830 emission is very optically thick in those lines. Thus, \(\tau_{He,λ5876}\) needs to exceed 30 for He λ5876/λ10830 to exceed 0.3 (cf. Figs 8 and 9, left-hand bottom panel). The gas producing the He λ10830 blue absorption, which needs

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Figure 9. Same as Fig. 8, but for $\gamma_{\text{He}} = 10^{-5} \, \text{s}^{-1}$.

to be optically thin in He $\lambda$5876 to avoid a He $\lambda$5876 blue absorption, is then distinct from the gas producing the He $\text{I}$ emission. This does not exclude the emission to arise from a stellar wind. It only stipulates that the emitting regions fill a sufficiently small volume of the wind that their summed surface area with a particular projected $v_{\text{obs}}$ subtends a small solid angle.

(3) The gas producing the red absorption is not producing the He $\lambda$5876 emission. Even though He $\lambda$5876 shows a red absorption in a few CTTSs, indicating in those objects an optical thickness that fulfills one requirement for strong emission, the deduced temperature of less than $10^4 \, \text{K}$ for that gas (cf. Section 5.2) does not meet the temperature requirement.

### 6.2 Pa$\gamma$/Pa$\beta$, Pa$\gamma$/H$\alpha$, Pa$\alpha$/Pa$\beta$ and Br$\gamma$/Pa$\alpha$

#### 6.2.1 Model results

The hydrogen line ratios are expected to correlate with one another, owing to the ladder-like energy level structure and very similar de-excitation pathways for each excited level. Fig. 11 shows how Pa$\gamma$/Pa$\beta$ and Pa$\gamma$/H$\alpha$ relate to $N_H$ and $T$ for $r = 4R_*$ and $\gamma_{\text{He}} = 2\gamma_{\text{He}} = 2 \times 10^{-4} \, \text{s}^{-1}$. We use the case of $T = 10^4 \, \text{K}$ to illustrate the Pa$\gamma$/H$\alpha$ dependence on $N_H$. As $N_H$ increases from $10^8 \, \text{cm}^{-3}$, Pa$\gamma$/H$\alpha$ falls because of the strong contribution to the H$\alpha$ emissivity from stimulated absorption of the stellar and veiling continua. This contribution diminishes when $\tau_{\text{H}\alpha}$ begins to exceed unity beyond $N_H \sim 6 \times 10^8 \, \text{cm}^{-3}$, whereupon Pa$\gamma$/H$\alpha$ increases with increasing $N_H$. The initial rise from $N_H = 6.3 \times 10^8$ to $3 \times 10^9 \, \text{cm}^{-3}$ is due to the continual contribution of Pa$\gamma$ stimulated absorption. At $N_H \sim 3 \times 10^9 \, \text{cm}^{-3}$, collisional...
excitation of Hα also becomes important and population is rapidly built up into $n > 2$ levels. The steeper response of the $n = 6$ population to $N_H$ increasing from $6.3 \times 10^9$ to $2.5 \times 10^{10}$ cm$^{-3}$ produces the rise of Paγ/Hα from $\sim 0.025$ to 0.06. The further rise of Paγ/Hα as $N_H$ exceeds $\sim 6.3 \times 10^{10}$ cm$^{-3}$ is caused by collisional de-excitation rivalling Hα escape, effectively stunting Hα emissivity growth. As seen from the accompanying plot with $\tau_{Pa\gamma}$ as the abscissa, the important rise of Paγ/Hα from 0.025 begins at $\tau_{Pa\gamma} \sim 1$. The Paγ/Pαβ ratio also increases from $\sim 0.45$ at about the same $\tau_{Pa\gamma}$. Once the lines involved are optically thick, their emissivity ratios are solely functions of the excitation temperatures. Thus, at $N_H = 1.6 \times 10^{10}$, $10^{11}$ and $6.3 \times 10^{11}$ cm$^{-3}$, when $\tau_{Pa\gamma} = 5.6, 33.4$ and 210, the (Hα, Pαβ, Paγ) excitation temperatures are (4710, 2570, 2350), (5930, 3240, 3020) and (6290, 3920, 3940), giving rise to Paγ/Hα and Paγ/Pαβ ratios of (0.05, 0.54), (0.065, 0.77) and (0.15, 1.14), respectively. Just like the case of HeIλ5876/λ10830, both Paγ/Hα and Paγ/Pαβ can, by means of collisional excitation, rise significantly above values produced from recombination and cascade only when the upper transition (Paγ here) is optically thick, and higher ratios require higher opacities and consequently higher excitation temperatures.

At a given $N_H$, Paγ/Pαβ is almost the same for $8750 \leq T \leq 3 \times 10^4$ K, but much smaller for $T \leq 7500$ K. As a function of $\tau_{Pa\gamma}$, the behaviour of Paγ/Pαβ at $T \leq 7500$ K is actually more similar to those at higher temperatures, only that the highest $\tau_{Pa\gamma}$ reached is much smaller. This stronger similarity is because, at the same $\tau_{Pa\gamma}$, the comparison between the rates of photon escape, $A_{Pa\gamma}/\tau_{Pa\gamma}$, and collisional de-excitation is almost independent of temperature.

Figure 10. Plot of log(HeIλ5876/λ10830) versus log(Paγ/HeIλ10830) for $r = 4R_\ast$, $\gamma_{He\ast} = 10^{-4}$ s$^{-1}$ and various temperatures (cf. Table 1). Data points marked by +s are determined from BEK01 and EFHK06. Those marked by ×s are determined from Edwards et al. (2010).
The dependences of $P_a^{\gamma}/P_a^{\beta}$ and $P_a^{\gamma}/H_\alpha$ on $N_H$ are only weakly sensitive to \( \gamma_{HI} \). At $T \geq 10^4$ K, there is no discernible difference between the cases of $\gamma_{HI} = 2 \times 10^{-5}$ and $2 \times 10^{-4}$ s$^{-1}$. At $T = (7500, 5000)$ K, $P_a^{\gamma}/P_a^{\beta}$ is lower by (0–10, 0–25) per cent and $P_a^{\gamma}/H_\alpha$ lower by (0–30, 20–40) per cent over the density range $3 \times 10^{10} \leq N_H \leq 10^{12}$ cm$^{-3}$ for $\gamma_{HI} = 2 \times 10^{-5}$ s$^{-1}$.

### 6.2.2 Comparison with observations

In Fig. 12, $P_a^{\gamma}/P_a^{\beta}$ is plotted against $P_a^{\gamma}/H_\alpha$ for $\gamma_{HI} = 2 \times 10^{-4}$ s$^{-1}$. As expected, $P_a^{\gamma}/P_a^{\beta}$ correlates with $P_a^{\gamma}/H_\alpha$. The figure shows that $P_a^{\gamma}/P_a^{\beta}$ and $P_a^{\gamma}/H_\alpha$ do not delineate temperature well, since the model tracks for different temperatures between 5000 and $3 \times 10^4$ K run similarly. They, however, are sensitive to density. From Fig. 11, it is seen that the dependence of $P_a^{\gamma}/P_a^{\beta}$ on $N_H$ is strong and quite similar for $8750 \leq T \leq 3 \times 10^4$ K, but changes rapidly as $T$ decreases from 8750 K. For example, to obtain $P_a^{\gamma}/P_a^{\beta} > 0.7$, the density needed at $T \leq 7500$ K is higher than that needed at $T \geq 8750$ K by a factor of $\geq 4$.

The data sample of Bary et al. (2008), including observations of the same object at multiple epochs, has 73 values of $P_a^{\gamma}/P_a^{\beta}$ that cluster closely about 0.86 with an estimated variance of $\pm 0.11$. Among case B models, the best fit to $P_a^{\gamma}/P_a^{\beta}$ as well as $P_a n_u/P_a^{\beta}$ for $n_u$ up to 14 has $T = 1000$ K and $N_e = 10^{10}$ cm$^{-3}$. While this best fit to eight data points has a reduced $\chi^2$ surpassing 99.9 per cent confidence, there is a glaring discrepancy at $P_a^{\gamma}/P_a^{\beta}$ where the fit value is only 0.73 and none of the other models explored has a higher value.
To see how the other Pa/Pa values from our calculations fare with the data, we show them in Fig. 13 for five densities at each of four temperatures. In general, as $N_H$ increases above $2.5 \times 10^{10}$ cm$^{-3}$, Pa/Pa (6 $\leq n_e \leq 12$) increases, with the higher $n_e$ increasing more. Although there is no single $N_H$ that produces an excellent match to data, the range in $N_H$ generating Pa/Pa ratios that bracket the data is quite small, of the order of less than a factor of 10. This range is almost the same for $T \geq 8750$ K, being $5 \times 10^{10} \leq N_H \leq 2 \times 10^{11}$ cm$^{-3}$, while that for $T = 7500$ K is $2 \times 10^{11} \leq N_H \leq 10^{12}$ cm$^{-3}$. Fig. 14 shows an analogous plot for Br/Pa, and the comparison between model results and data is similar.

Bary et al. (2008) has also compared their data with line ratios calculated with a constant excitation temperature (local thermal equilibrium), but find that in neither the optically thick nor thin case is there an acceptable fit. Our calculations are more similar to this set of models, only with the level excitations and line optical depths calculated self-consistently for a given $N_H$ and $T$. Thus, the Pa/PA excitation temperatures are not the same and the lines are not all completely optically thick. For example, at $T = 10^4$ K and $N_H = 10^{11}$ cm$^{-3}$, the excitation temperatures and line optical depths of (Pa, Pay, Pa 10, Pa 12) are (3240, 3020, 2910, 2940) K and (104, 33.4, 3.44, 1.78), respectively. The fair agreement between model results and observed values makes understanding the hydrogen emission as resulting from collisional excitation an attractive alternative to the proposal by Bary et al. (2008) of a recombining gas at $T = 1000$ K.

We draw the following conclusions from the comparison between model and observational results on the hydrogen line ratios:

1. Pay is optically thick in the hydrogen emission region. Thus, $\tau_{Pa}$ needs to exceed 20 for Pay/PA to exceed 0.7 at $T \leq 10^4$ K (Fig. 11, left-hand bottom panel). Even many of the higher order Paschen lines have optical depths exceeding unity.
Figure 13. Plot of $\frac{P_a}{P_\beta}$ versus $n_u$ for four temperatures, $r = 4R_*$, $y_{H_1} = 2 \times 10^{-4} \text{ s}^{-1}$ and various values of $N_H$ (in units of $10^{11} \text{ cm}^{-3}$). Data marked by open circles are from Bary et al. (2008).

(2) The temperature in the hydrogen emission region is higher than 5000 K. At 5000 K, $P_\gamma/P_\beta$ reaches a value of only 0.81 even at the highest $N_H$ explored, $2.5 \times 10^{12} \text{ cm}^{-3}$. The more limited and extreme range on $N_H$ needed to produce $P_\gamma/P_\beta \geq 0.86$ at $T \leq 5000$ K is one argument. Another is the difficulty in matching the observed $P_\gamma/\text{He I} \lambda 10830$ ratio. As seen from Fig. 8, $P_\gamma/\text{He I} \lambda 10830$ is only 0.38 at $1.6 \times 10^{12} \text{ cm}^{-3}$ and falls with increasing $N_H$. This match-up is better in the case of $\gamma_{\text{He I}} = 10^{-5} \text{ s}^{-1}$ (cf. Fig. 9), but then $P_\gamma/P_\beta$ is only 0.6 at $N_H = 1.6 \times 10^{12} \text{ cm}^{-3}$.

(3) The density in the hydrogen emission region is centred around $N_H \sim 10^{11} \text{ cm}^{-3}$, if $8750 \leq T \leq 3 \times 10^4 \text{ K}$ and around $N_H \geq 5 \times 10^{11} \text{ cm}^{-3}$ if $T \leq 7500 \text{ K}$, in order to produce the observed $P_a/n_u/P_\beta$, $n_u \geq 6$, ratios.

6.3 Ca $\lambda 8498$/Pay, O $\lambda 8446$/Pay, Ca $\lambda 8498$/O $\lambda 8446$ and O $\lambda 7773$/λ8446

6.3.1 Model results

We compare the Ca I infrared triplet to $P_a$ for its being a fairly strong hydrogen line not too far away in wavelength and to O $\lambda 8446$ for its tie to the hydrogen excitation through Ly$\beta$ fluorescence and its close proximity in wavelength. Fig. 15 shows the dependences of Ca $\lambda 8498$/Pay and O $\lambda 8446$/Pay on $N_H$ and $\tau_{P_a}$ for $r = 4R_*$, $y_{H_1} = 2 \times 10^{-4} \text{ s}^{-1}$ and seven temperatures. Fig. 16 is analogous, but with Ca $\lambda 8498$/O $\lambda 8446$ versus $N_H$ and $\tau_{\text{Ca I}}$, and O $\lambda 7773$/λ8446 versus $N_H$ and $\tau_{\text{O I}}$. Because of our single-level representation of multiple fine-structure levels, we divide the emissivity from our model Ca I ion by 3, thereby implicitly assuming that all three components...
of the triplet are optically thick, a reasonable assertion as all three components have nearly identical fluxes and profiles. Our model $\tau^*_{\text{Ca}^{\text{II}}\lambda 8498}$ is five times the optical depth of the component with the weakest oscillator strength (cf. Section 5.1), so the Ca II infrared emission requires $\tau^*_{\text{Ca}^{\text{II}}\lambda 8498}$ to exceed 10. Our O I 7773 line also represents a triplet, but the whole emissivity is employed here, since the observed flux also comprises contributions from all three components.

Fig. 15 shows that O I 8446/Pa $\gamma$ rises continually to beyond unity as $N_H$ increases. This is largely because the much lower O I 8446 optical depth (cf. Fig. 6) keeps $\lambda 8446$ escape more effective than collisional de-excitation. The dependence of O I 8446/Pa $\gamma$ on temperature depends slightly on $N_H$. At $2 \times 10^{10} \leq N_H \leq 2 \times 10^{11} \text{ cm}^{-3}$, O I 8446/Pa $\gamma$ increases as $T$ increases from 5000 K, peaks at $T \sim 8750$ K and then drops with further increase in $T$. This behaviour can be seen from the dependences of the O I 8446 emissivity, $\propto N_O/N_1$, and the Pa $\gamma$ emissivity, $\propto N_e A_{\text{Pa} \gamma} \beta_{\text{Pa} \gamma} \alpha N_1/N_3$, on the hydrogen level population. Thus, O I 8446/Pa $\gamma$ $\propto (N_O/N_1)(N_3^2/N_6)$. As $T$ increases from 5000 K at a fixed $N_H$, $N_3/N_6$ always decreases, while $N_1$ first increases but then decreases when hydrogen becomes more ionized. The counter-effects of $N_3/N_6$ and $N_3$ on O I 8446/Pa $\gamma$ in the beginning of the temperature rise and their concerted effects subsequently produce the mentioned behaviour. At $N_H > 2 \times 10^{11} \text{ cm}^{-3}$, the same behaviour holds except that O I 8446/Pa $\gamma$ increases again as $T$ increases beyond $1.5 \times 10^4$ K. This occurs because, with the rapid decrease in the hydrogen neutral fraction, radiative recombination of O II becomes important and $N_O/N_1$ increases from the constant value of 1.1$N_O/N_2$ maintained by charge-exchange reactions.

The two ratios Ca II 8498/Pa $\gamma$ and Ca II 8498/O I 8446 have very similar dependences on $N_H$ and $T$, as to be expected. The requirement of $\tau^*_{\text{Ca}^{\text{II}}\lambda 8498} > 10$ confines the pertinent densities to $N_H \geq 10^{11} \text{ cm}^{-3}$ for $T \leq 8750$ K and even higher limits for $T > 8750$ K. Thus, the large Ca II 8498/Pa $\gamma$ and Ca II 8498/O I 8446 model values at $N_H < 10^{10} \text{ cm}^{-3}$ are not relevant. They actually involve very small fluxes, but...
Figure 15. Dependences of \( \text{Ca}\ II\ \lambda 8498 / \text{Pa}\ \gamma \) and \( \text{O}\ I\ \lambda 8446 / \text{Pa}\ \gamma \) on the density and line optical depth for \( r = 4R_\odot, \gamma_{H I} = 2 \times 10^{-4} \) s\(^{-1}\), and various temperatures (cf. Table 1). The large \( \text{Ca}\ II\ \lambda 8498 / \text{Pa}\ \gamma \) values at \( N_{H} < 10^{10} \) cm\(^{-3}\) are not significant, because they occur at \( \tau_{\text{Pa}\ \gamma} \ll 1 \).

are shown for uniformity in the density range covered. As \( N_{H} \) increases from \( \sim 10^{11} \) cm\(^{-3}\) at a fixed temperature, both \( \text{Ca}\ II\ \lambda 8498 / \text{Pa}\ \gamma \) and \( \text{Ca}\ II\ \lambda 8498 / \text{O}\ I\ \lambda 8446 \) rise rapidly. Two factors contribute to this. First, while \( N_{H}/N_{H} \) is fairly constant in that density range (cf. Fig. 5), the fraction of calcium in \( \text{Ca}\ II \) increases with \( N_{H} \). Secondly, the onset of collisional de-excitation of the \( \text{Ca}\ II \) population occurs, by comparison with hydrogen, at a higher \( N_{H} \). When the latter is reached, the consequent collisional quenching of the \( \text{Ca}\ II \) triplet emissivity causes the decline of the ratio from its peak.

Both \( \text{Ca}\ II\ \lambda 8498 / \text{Pa}\ \gamma \) and \( \text{Ca}\ II\ \lambda 8498 / \text{O}\ I\ \lambda 8446 \) reach their highest value at \( T = 7500 \) K. This can be understood as follows. The hydrogen line emissivities are fairly insensitive to temperature for \( 8750 \leq T \leq 3 \times 10^{4} \) K (cf. Section 6.2). The \( \text{Ca}\ II\ \lambda 8498 \) emissivity, on the other hand, decreases with increasing \( T \) above 8750 K, owing to the rapidly decreasing \( \text{Ca}\ II \) fraction. Thus, the two ratios fall beyond \( T = 8750 \) K. At \( 5000 \leq T \leq 8750 \) K, collisional excitations of both \( \text{Ca}\ II \) and \( H I \) are sensitive functions of temperature because of the strong dependence of \( N_{e}/N_{H} \) on \( T \) (cf. Fig. 5) and the higher ratio of excitation potential to thermal energy. Above \( N_{H} = 10^{11} \) cm\(^{-3}\), however, collisional de-excitation comes into play for hydrogen and the \( \text{Pa}\ \gamma \) emissivity increases only modestly with increasing temperature in comparison with the \( \text{Ca}\ II \) triplet emissivity, which, until \( N_{H} \geq 6 \times 10^{11} \) cm\(^{-3}\), is still proportional to \( (N_{\text{Ca}\ II}/N_{\text{Ca}})(N_{e}/N_{H}) \), and the steeper rise of \( N_{e}/N_{H} \) with increasing \( T \) from 5000 K more than overcomes the counter-effect of a lower \( N_{\text{Ca}\ II}/N_{\text{Ca}} \) ratio. As a result, \( \text{Ca}\ II\ \lambda 8498 / \text{Pa}\ \gamma \) increases as \( T \) increases from 5000 K and, in conjunction with its fall at \( T > 8750 \) K, peaks at an intermediate temperature.

Regarding the \( \lambda 7773 / \lambda 8446 \) ratio, both recombination and cascade, and collisional excitation strongly favour \( \text{O}\ I\ \lambda 7773 \) over \( \text{O}\ I\ \lambda 8446 \) emission. This ratio is less than unity only because of the \( \text{Ly}\beta \) fluorescence process, which, as seen from Fig. 16, is most effective at \( N_{H} \).
between $10^{10}$ and $3 \times 10^{10}$ cm$^{-3}$. At $N_H < 10^{10}$ cm$^{-3}$, even though Ly$\beta$ fluorescence remains more effective than the other two processes, O$\lambda$7773/\lambda8446 increases as $N_H$ decreases. This is because stimulated absorption of the stellar and veiling continua contributes prominently to the O$\lambda$7773 emissivity. However, the line fluxes involved are very small. O$\lambda$7773/\lambda8446 also increases as $N_H$ increases above $\sim 3 \times 10^{10}$ cm$^{-3}$. This is because collisional de-excitation begins to affect the hydrogen level population, causing $N_3/N_1$, the important factor in the Ly$\beta$ fluorescence rate, to be nearly constant, while collisional excitation of O$\lambda$7773 continues to grow roughly as $N_H^2$. Thus, in this regime, O$\lambda$7773/\lambda8446 responds to increasing density in a fashion similar to Ca$\lambda$8498/O$\lambda$8446, but it has not reached a peak even at $N_H = 1.6 \times 10^{12}$ cm$^{-3}$, because collisional de-excitation of O$\lambda$7773 remains weaker than $\lambda7773$ escape.

The dependences of the four emissivity ratios on $\gamma_{HI}$, between $2 \times 10^{-4}$ and $2 \times 10^{-5}$ s$^{-1}$, are not strong, with Ca$\lambda$8498/Pa$\gamma$ and Ca$\lambda$8498/O$\lambda$8446 reaching somewhat higher values for $5000 \leq T \leq 7500$ K when $\gamma_{HI} = 2 \times 10^{-5}$ s$^{-1}$. There is also a weak dependence on $r$. Among the four cases of $\gamma_{HI}$ and $r$, the ($\gamma_{HI} = 2 \times 10^{-4}$ s$^{-1}$, $r = 4R_*$) and ($\gamma_{HI} = 2 \times 10^{-5}$ s$^{-1}$, $r = 2.5R_*$) cases show the largest difference.

### 6.3.2 Comparison with observations

Fig. 17 shows the relations between Ca$\lambda$8498/Pa$\gamma$ and O$\lambda$8446/Pa$\gamma$, and between Ca$\lambda$8498/O$\lambda$8446 and O$\lambda$7773/\lambda8446 for the case ($\gamma_{HI} = 2 \times 10^{-4}$ s$^{-1}$, $r = 4R_*$), and Fig. 18 shows those for the case ($\gamma_{HI} = 2 \times 10^{-5}$ s$^{-1}$, $r = 2.5R_*$). In addition to the data points marked, we
mention that Ca\textsc{ii}λ8498/O\textsc{i}λ8446 can be determined for the remaining five objects in the data set of Muzerolle et al. (1998b). It is 4.5, 12.7, 9.8, 0.81 and 1.31 for DR, CW, RW, DS and UY Tau, respectively. Thus, it appears that the observed values of Ca\textsc{ii}λ8498/O\textsc{i}λ8446 separate into two groups. The full data set of Edwards et al. (in preparation) also indicates this dichotomy, namely, one group with Ca\textsc{ii}λ8498/Paγ or Ca\textsc{ii}λ8498/O\textsc{i}λ8446 close to or less than 1 and the other with Ca\textsc{ii}λ8498/Paγ or Ca\textsc{ii}λ8498/O\textsc{i}λ8446 greater than ∼5. The physical conditions that will produce the first group of values are \(N_\text{H} \approx 10^{11}\) cm\(^{-3}\) and \(T \approx 5000\) and 1.25 \(\times\) \(10^4\) K (cf. Figs 15 and 16), but the temperature is narrowed to 8750 \(\leq T \leq 1.25 \times 10^4\) K when the constraints imposed by the observed Pa\textsc{n}/Paβ ratios are also implemented (cf. Section 6.2.2). To produce the second group of values, a considerably higher \(N_\text{H} \approx 10^{12}\) cm\(^{-3}\), is required, while \(T\) is restricted to \(\leq 7500\) K.

The CTTSs with Ca\textsc{ii}λ8498 emission much stronger than Paγ or O\textsc{i}λ8446 emission also tend to be the ones with the most prominent He\textsc{i}λ5876 emission. The physical conditions needed to produce the Ca\textsc{ii}λ8498/Paγ ratio cannot produce the He\textsc{i}λ5876/λ10830 ratio, so a separate region producing the bulk of the He\textsc{i}λ5876 emission is additionally required, as already inferred in Section 6.1. The Ca\textsc{ii} emission region itself appears to be capable of generating the observed ratios among the hydrogen lines (cf. Figs 13 and 14, \(T = 7500\) K panel).

The CTTSs with Ca\textsc{ii}λ8498 emission comparable to or less than Paγ emission have a broad spread in helium emission, from He\textsc{i}λ5876/λ10830 ratios close to those of the previous group to little He\textsc{i}λ5876 emission and He\textsc{i}λ10830 primarily in absorption (Edwards et al., in preparation). When He\textsc{i}λ5876 emission is weak, the physical conditions needed to produce the hydrogen line ratios

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure17.png}
\caption{Plot of Ca\textsc{ii}λ8498/Paγ versus O\textsc{i}λ8446/Paγ (top panel) and Ca\textsc{ii}λ8498/O\textsc{i}λ8446 versus O\textsc{i}λ7773/λ8446 (bottom panel) for \(r = 4R_\ast\), \(\gamma_\text{H}_1 = 2 \times 10^{-3}\) \(s^{-1}\) and various temperatures (cf. Table 1). Data points marked by \(\times\) and + are determined from Edwards et al. (2010) and Muzerolle et al. (1998b), respectively.}
\end{figure}
Figure 18. Same as Fig. 17, but for $r = 2.5 R_*$ and $\gamma_{\rm H_1} = 2 \times 10^{-5}$ s$^{-1}$.

(cf. Section 6.2) can also generate the requisite Ca II emission. When He I $\lambda 5876$ emission is strong, a separate region producing it is again called upon.

Based on the above understanding of the ratios involving He I, H and Ca II lines, we draw the following conclusion:

Three sets of physical conditions are identified. The first, characterized by a comparatively high temperature ($T$ close to $2 \times 10^4$ K), is necessary to produce strong He I $\lambda 5876$ emission. The second, characterized by a comparatively high density ($N_{\rm H} \sim 10^{12}$ cm$^{-3}$) and low temperature ($T \leq 7500$ K), is necessary to produce strong Ca II infrared triplet emission. The third, with less-stringent constraints ($8750 \leq T \leq 1.25 \times 10^4$ K, $N_{\rm H}$ around $10^{11}$ cm$^{-3}$), produce primarily hydrogen emission. When the second set is present, the first set is often present concurrently. The third set is present either by itself or concurrently with the first set. It is also likely to be present when both the first and second sets are present.

7 LINE FLUXES AND SIZES OF EMISSION REGIONS

As mentioned in Section 2, the specific flux $F_{\nu, \lambda}$ of an optically thick line depends on the line excitation temperature and the emission area with a projected velocity of $v_{\lambda, \text{obs}}$, $A_{\lambda, \text{obs}}$. The excitation temperature, in turn, depends on the local physical conditions ($N_{\text{H}}, T$, etc.) and the emission area depends on the specific kinematic model. In particular, for $|v_{\lambda, \text{obs}}| = 150$ km s$^{-1}$, a spherical wind model with $r = 4 R_*$ has $A_{\nu, \text{obs}} = 27(r/4 R_*)^2 \pi R_*^2$ and some leeway in its adjustment by varying $r$, while a spherical infall model with $r = 2.5 R_*$ has $A_{\nu, \text{obs}} = 1.286 \pi R_*^2$ and little room for maneuver. In a CTTS, either wind or accretion flow likely fills only a fraction, $f_{\text{fill}}$, of the $4\pi$ solid angle. This dilution
affects the actual observed area in a way that depends on the viewing angle, whose variation is not included here to keep the paper manageable. We simply assume an effective observed area given by $f_{\text{eff}} A_{\text{obs}}$. Thus, for the same excitation temperature, the line flux in the infall model will be $(0.2/0.5)(1.286/27)$ or $1/52.5$ of that in the wind model. We think this is a conservative estimate in that the true ratio is likely smaller.

For ease in comparison with an observed line, we scale the model $F_{\text{line}}$ to the continuum flux $F_\odot$, and present the ratio $y_{\text{line}} = F_{\text{line}}/F_\odot$ as a function of density or line optical depth. Our model continuum is represented by a $4000$-$K$ blackbody veiled by a $8000$-$K$ blackbody over $3$ per cent of its surface. Thus, our model continuum has a veiling contribution at He $\lambda\lambda 5876$ of $r_\gamma = 0.67$ and one at He $\lambda\lambda 10830$ or Pa$\gamma$ of $r_\gamma = 0.19$. When a CTTS has a different $F_\odot$, $r_\gamma$, or $r_\gamma$, the model $y_{\text{line}}$ can be easily rescaled to compare directly with the observed line profile.

We shall make use of the information deduced from line ratios to identify the appropriate physical conditions. Since we are calculating the specific flux at $|v_{\text{obs}}| = 150$ km s$^{-1}$, it is the ratio of the values of $F_{\text{line}}$ that is pertinent here. If the two lines have similar profiles, this ratio is close to the ratio of their integrated fluxes, which was used in Section 6 to identify the physical conditions responsible for the bulk of the emission. This appears to be the case among He $\lambda\lambda 10830$, 5876 and Pa$\gamma$. However, Ca $\lambda\lambda 8498$ is distinctly narrower than Pa$\gamma$. The contrast is obvious in DR Tau, notable in DG Tau, and measurable in DL Tau (Edwards et al., in preparation). When we come to discussing the Ca $\lambda\lambda 8498$ specific flux, we will take note of the distinction.

7.1 H$\alpha$, Pa$\gamma$, Pa$\beta$, Br$\gamma$

In this section, we select among the hydrogen lines for presentation H$\alpha$, the strongest optical member, and Pa$\gamma$, Pa$\beta$ and Br$\gamma$ for the availability of large data sets (EFHK06; Folha & Emerson 2001).

Figs 19 and 20 show, for the wind and infall model, respectively, how the H$\alpha$ and Pa$\gamma$ specific fluxes depend on $N_\text{H}$ and line optical depth. The $y$-axis limits in Fig. 20 are $1/52.48 (-1.72$ in log) of those in Fig. 19, so chosen that, if the specific flux in the infall model is $1/52.48$ of that in the wind model, the two figures will be visually identical. They do look very much alike. When one figure is overlaid on the other, the largest difference in the density-dependence plots occurs for $y_{\text{line}}$ (H$\alpha$) at $N_\text{H} \sim 10^5$ cm$^{-3}$, where the line emissivity is contributed primarily by stimulated absorption of the incident continuum flux, which is a factor of $2.56$ larger at $r = 2.5R_\ast$. The optical depth dependence plots, when overlaid together, show larger differences. This is because at the same $T (\propto N_{\text{H}}/v)$, the density in the $r = 2.5R_\ast$ case is higher and the line specific flux is more sensitive to density than to the velocity gradient, as will be propounded next.

Figs 19 and 20 show not only that the primary distinction in the specific flux produced between the wind and infall models arises from the emission area, but also that variation in the velocity gradient (a factor of $1.6$ larger in the $r = 2.5R_\ast$ model) generates only a small change. As a specific example, we consider the case of $N_\text{H} = 10^{14}$ cm$^{-3}$ and $T = 10^3$ K, conditions suitable to bringing about the observed Pa$\gamma$/Pa$\beta$ ratio. The value of $y_{\text{line}}$ (Pa$\gamma$) in the wind model is $3.81$, while that in the infall model is $0.7242 = 3.802/52.5$. Their nearly identical emissivity, $y_{\text{line}}/(f_{\text{eff}} A_{\text{obs}})$, is not surprising, since Pa$\gamma$ is very optically thick ($r_{\text{Pa}\gamma} = 33.44$ and $20.46$ in the wind and infall model, respectively, cf. Figs. 19 and 20, right-hand bottom panel). Even when Pa$\gamma$ is optically thin, the difference in emissivity is small. For example, at $N_\text{H} = 10^{10}$ cm$^{-3}$, $T = 10^4$ K, $y_{\text{line}} = 0.356$ and $r_{\text{Pa}\gamma} = 1.63$ in the wind model, while $y_{\text{line}} = 4.78 \times 10^{-3} = 0.251/52.5$ and $r_{\text{Pa}\gamma} = 0.355$ in the infall model. The Pa$\gamma$ emissivity is a factor of $1.41$ smaller with the larger velocity gradient. However, at the next density step of $1.585 \times 10^{10}$ cm$^{-3}$, $y_{\text{line}} = 0.0154$ in the infall model or a factor of $3.22$ larger than the earlier value. From the slope of the log $y_{\text{line}}$ dependence on log $N_{\text{H}}$ at $10^{10}$ cm$^{-3}$, it can be estimated that the larger velocity gradient by a factor of $1.6$ can be compensated by an increase in $N_{\text{H}}$ by a factor of only $1.147$ to produce the same line emissivity. For this reason, even though we emphasize the location where $|v| = 150$ km s$^{-1}$, the local excitation calculations of line emissivities and ratios are applicable elsewhere, because a large variation in the local velocity gradient is tantamount to only a modest variation in the density.

The two figures also show that over the density range $2 \times 10^{10} < N_{\text{H}} < 2 \times 10^{12}$ cm$^{-3}$, the dependences of $y_{\text{line}}$ (H$\alpha$) and $y_{\text{line}}$ (Pa$\gamma$) on temperature are strong for $5000 \leq T \leq 8750$ K, but weak for $8750 \leq T \leq 3 \times 10^4$ K. This is because, despite the stronger collisional excitation with $T$ increasing above $8750$ K, the concomitant more rapid build-up of population into high $n$ levels produces a stronger collisional ionization and reduces the neutral fraction.

The $38$ Pa$\gamma$ profiles in the reference sample of EFHK06 provide a convenient data set of the values of $y_{\text{line}}$ for comparison with our model results. To avoid the difficulty of extracting the underlying red emission when a red absorption is obviously or probably present, we concentrate on the blue wing emission and determine $y_{\text{line}}$ at $v_{\text{obs}} = -150$ km s$^{-1}$. Of the $38$ objects, nine (DR, CW, DL, HL, DG, DK and HN Tau, AS 353 A and RW Aur A) have high $1$-$\mu$m veiling ($r_\gamma > 0.5$), of whom eight have $y_{\text{line}} > 0.25$ and one (DK Tau) has $y_{\text{line}} \sim 0.15$; $11$ (DS, HK, BP, DF, DO, GG and GK Tau, YY and GW Ori, UY Aur and UZ Tau E) have intermediate $1$-$\mu$m veiling ($0.3 \leq r_\gamma \leq 0.4$), of whom $10$ (excluding GK Tau) have $0.1 \leq y_{\text{line}} \leq 0.25$. Of the remaining $18$ objects with low $1$-$\mu$m veiling ($r_\gamma \leq 0.2$), two (SU Aur and TW Hya) have $y_{\text{line}}$ between $0.1$ and $0.25$, and the rest have $y_{\text{line}} < 0.1$.

For members in the first veiling group, our model values of $y_{\text{line}}$, when scaled to the CTTS stellar plus veiling continuum, need to be reduced by a factor of $\sim 2$ or more. Then, at the physical conditions needed to produce the Pa$\beta$/Pa$\gamma$ ratios ($N_{\text{H}}$ centred around $10^{11}$ cm$^{-3}$, $T \geq 8750$ K, or $N_{\text{H}}$ centred around $5 \times 10^{14}$ cm$^{-3}$, $T = 7500$ K, cf. Section 6.2.2 and Fig. 13) and assuming that these ratios of integrated fluxes are the same as the ratios of their specific fluxes at $v_{\text{obs}} = -150$ km s$^{-1}$, the $y_{\text{line}}$ values in the infall model will be less than $0.05$. They fail to produce the values of those nine objects by factors ranging from $\sim 3$ to $\sim 15$. We therefore conclude that a kinematic flow other than accretion inflow is responsible for those hydrogen emission. We also draw the same conclusion for the majority of the members in the second $1$-$\mu$m veiling group. The infall model has difficulty in accounting not only for their specific fluxes at $v_{\text{obs}} \leq -150$ km s$^{-1}$, but also for their observed line wing asymmetry (say, at $|v_{\text{obs}}| \geq 150$ km s$^{-1}$). Thus, whereas in the infall model, the blue wing emission is occulted by the
The wind model poses the opposite conundrum in that the model values of $y_{vobs}$ are much larger than those observed in the high 1-μm veiling group, by factors of about 4–10. It can be responsible for the observed hydrogen emission, provided the optically-thick emitting gas ($\tau_{Pa\gamma} > 10$, cf. Fig. 19, right-hand bottom panel) fills only a portion of the postulated flow and covers a fraction of the contour surface shown in Fig. 1. This also happens to be the necessary proviso to avoid the occurrence of a blue absorption. Thus, the wind flow needs to be highly inhomogeneous, with the density of the emitting gas considerably higher than that pervading the bulk of the flow. The temperature of the emitting gas is probably $\lesssim 10^4$ K, since the line emissivity at a higher temperature is not higher and the filling factor likely decreases rapidly with increasing temperature above $10^4$ K, owing to the stronger cooling at a higher temperature. It is also probably $>7500$ K; otherwise, the density needed to produce the Pa $n_u$/Paβ ratios would be around $6 \times 10^{11}$ cm$^{-3}$, an order of magnitude greater than that needed at $T \geq 8750$ K, yet the Paγ emissivity is about the same, so the filling factor required of the gas at $N_H \sim 6 \times 10^{11}$ cm$^{-3}$ is as high and therefore more restrictive.

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Figure 19. Dependences of Hα and Paγ specific flux (measured relative to the local continuum), $y_{vobs}$, on the density and line optical depth for $y_{H1} = 2 \times 10^{-4}$ s$^{-1}$ and various temperatures (cf. Table 1) in the wind model.
Figure 20. Same as Fig. 19, but for the infall model.

Among the optical hydrogen lines, \( \text{H} \alpha \) has the highest optical depth, \( \sim 3 \times 10^3 \). This can be seen from Fig. 19, which shows that at \( N_H = 10^{11} \text{ cm}^{-2} \) and \( T = 10^4 \text{ K} \), \( \tau_{\text{H} \alpha} \) exceeds \( 10^3 \), the upper bound on the abscissa. The high \( \tau_{\text{H} \alpha} \) leads to a fair amount of emission in the damping wings and likely accounts for the conspicuous extension of the observed \( \text{H} \alpha \) profile at the base. Thus, assuming just natural broadening of the emission profile, the rate of emitting \( \text{H} \alpha \) photons that are displaced from line centre by \( \geq \sigma \) thermal widths is

\[
\frac{n_u A_{\text{H} \alpha} \alpha}{(\pi \sigma)},
\]

where \( a = \Gamma_r/(4\pi \Delta \nu_D) \), \( \Gamma_r \) is the summed radiative decay rate from the upper and lower states of \( \text{H} \alpha \) and \( \Delta \nu_D \) is the thermal Doppler frequency width. It is assumed that \( x \) is sufficiently large that the emission is not re-absorbed. With the total \( \text{H} \alpha \) emission rate being \( n_u A_{\text{H} \alpha} \beta_{\text{H} \alpha} \), the fraction of photons emitted beyond \( x \) thermal widths is then \( 2\alpha/(\pi x \beta_{\text{H} \alpha}) \). At \( 10^4 \text{ K} \), the thermal Doppler velocity width is \( \sim 13 \text{ km s}^{-1} \), so for \( x = 27 \), the fraction of \( \text{H} \alpha \) photons emitted at \( |v_{\text{obs}}| \geq 350 \text{ km s}^{-1} \) is \( \sim 0.16 \).

Fig. 21 shows the analogous plots of \( \text{Pa} \beta \) and \( \text{Br} \gamma \) for the infall model. The stellar plus veiling continuum at 1.28 \( \mu \text{m} \) of the CTTSs in the sample of Folha & Emerson (2001) is not listed but most likely significantly stronger than our model value because of the typically high 1-\( \mu \text{m} \) veiling observed and probable contribution from dust emission (Muzerolle et al. 2003). Even assuming that our model continuum is representative, the model \( v_{\text{obs}}(\text{Pa} \beta) \), at the appropriate conditions indicated by the \( \text{Pa} \beta /\text{Pa} \beta \) ratios, is \( \sim 0.15 \) and falls short of quite a few observed values. Thus, among the 49 objects with \( \text{Pa} \beta \) profiles, eight (DG, DL, DR, HL and CW Tau, GM and RW Aur, YY Ori) have, at \( v_{\text{obs}} = -150 \text{ km s}^{-1} \), \( y_{\text{obs}} \geq 0.5 \), eight (GG, RY, DO, DS and FS Tau, UY and SU Aur, V1331 Cyg) have \( 0.2 \leq y_{\text{obs}} < 0.5 \), and 10 (BM And, DE, DF, DK, HK, BP, GI, HP and T Tau, GW Ori) have \( 0.1 \leq y_{\text{obs}} < 0.2 \). Again then, a comparison between model results and \( \text{Pa} \beta \) data leads to the same conclusions deduced earlier. It is interesting that the objects with strong (moderate) \( \text{Pa} \beta \) flux at \( |v_{\text{obs}}| \geq 150 \text{ km s}^{-1} \)
also have strong (moderate) Paγ flux. This may simply be a confirmation of both the finding of Bary et al. (2008) that the Paγ/Pαβ ratio does not vary strongly among CTTSs and the assumption of similar Paγ and Paβ profiles.

In summary, based on the hydrogen-line specific fluxes at \(|v_{\text{obs}}| \geq 150\,\text{km s}^{-1}\), we arrive at the following conclusions regarding the infall accretion and wind flows:

1. The infall accretion flow is not responsible for the strong or moderate hydrogen emission at \(|v_{\text{obs}}| \geq 150\,\text{km s}^{-1}\). We can extend this inference to the bulk of the observed emission for the following reasons. First, the emission generated by gas with \(|v| \geq 150\,\text{km s}^{-1}\) is a substantial fraction of the total, since the same particles also contribute to observed emission at \(|v_{\text{obs}}| < 150\,\text{km s}^{-1}\) when their trajectories do not parallel the line of sight. A rough estimate of the line profile produced by them in an azimuthally symmetric distribution is a flat-top shape in the regime \(-150 \leq v_{\text{obs}} \leq 150\,\text{km s}^{-1}\). Using this guideline to determine the fraction of the blueward emission (to avoid the red absorption), we find from fig. 3 of EFHK06 that 19 (excepting GK Tau) out of 20 CTTSs in the high and intermediate veiling groups \((r_Y \geq 0.3)\) have this fraction \(\geq 0.5\). Secondly, with the velocity distribution expected to be a smooth function, the infall model will also have difficulty accounting for the observed specific fluxes at \(|v_{\text{obs}}|\) somewhat less than 150 km s\(^{-1}\). Thirdly, whatever kinematic structure called for to produce the observed emission at \(|v_{\text{obs}}| \geq 150\,\text{km s}^{-1}\) will likely have gas particles moving at \(<150\,\text{km s}^{-1}\). All these considerations led us to conclude that the bulk of the Paγ emission in these CTTSs arises from a kinematic structure other than the accretion flow.
(2) For the wind flow to be responsible for the moderate or strong hydrogen emission, the emitting gas, whose density is considerably higher than that of the gas occupying the bulk of the flow, can only fill a small fraction of the volume.

These conclusions remain true, of course, when the two flows are compared at the lower ionization rate of $\gamma_{HI} = 2 \times 10^{-5} \text{ s}^{-1}$. For either flow, the specific fluxes are almost identical to those in the $\gamma_{HI} = 2 \times 10^{-4} \text{ s}^{-1}$ case at $T \geq 8750 \text{ K}$ and are weaker at $T \leq 7500 \text{ K}$, as can be deduced from the contrast in the $\tau_{HI}$ contours between the two ionization cases (Fig. 4). Thus, for ($H\alpha$, $Pa\gamma$), the values of $Y_{v,obs}$ at $N_H = 10^{11}$ and $10^{12} \text{ cm}^{-2}$ with $\gamma_{HI} = 2 \times 10^{-5} \text{ s}^{-1}$ are smaller by a factor of (2.55, 4.14) and (1.06, 1.16), respectively, at $T = 7500 \text{ K}$, and (3.67, 6.64) and (1.92, 2.69), respectively, at $T = 5000 \text{ K}$. This dependence of $Y_{v,obs}$ on $\gamma_{HI}$, when $T \leq 7500 \text{ K}$, as well as rapid rise of $Y_{v,obs}$ with $T$ increasing above 7500 K at a given $\gamma_{HI}$, led us to favour the temperature in the hydrogen emission region being $\geq 8750 \text{ K}$.

### 7.2 $He\lambda 10830, 5876$

Figs 22 and 23 show the $He\lambda 10830, 5876$ values of $Y_{v,obs}$ as a function of density and line optical depth for the wind and infall model, respectively. Like the $He\lambda 5876/\lambda 10830$ ratio, the model results for the $He\alpha$ specific fluxes at $T = 3 \times 10^4 \text{ K}$ are very close to those at $T = 2 \times 10^4 \text{ K}$ and are not shown.

The 22 $He\alpha 5876$ BC profiles in fig. 7 of BEK01 present a convenient data set for comparison with our model results. We again look at the specific flux at $v_{obs} = -150 \text{ km s}^{-1}$ to avoid the red absorption, which appears to affect several profiles. The 22 optical veiling range

Figure 22. Dependences of $He\lambda 10830$ and $He\lambda 5876$ specific flux (measured relative to the local continuum), $Y_{v,obs}$, on the density and line optical depth for $\gamma_{He\alpha} = 10^{-4} \text{ s}^{-1}$ and various temperatures (cf. Table 1) in the wind model.
from 0.1 to 20 and the observed values of $y_{\text{vobs}}$ range from 0.01 to 0.5, but the correlation between $y_{\text{vobs}}$ and $r_V$ does not appear to be strong. Eight objects (CW, HN, DG, DL and DR Tau, AS 353 A, RW and UY Aur) have $0.2 \leq y_{\text{vobs}} \leq 0.5$ and $r_V$ between 0.8 and 20, four (GG, DQ and DF Tau, GM Aur) have $0.1 \leq y_{\text{vobs}} < 0.2$ and $r_V$ between 0.2 and 0.7, and the rest have $y_{\text{vobs}} < 0.1$ and $r_V$ between 0.1 and 4.7. With our model $r_V$ of 0.67, the model values of $y_{\text{vobs}}$ likely need to be lowered (raised) for comparison with the first (second) veiling group. Even for observed values of $y_{\text{vobs}}$ as high as 0.5, it appears that the infall model can accommodate them. As a specific example, we assume that the He I $\lambda 5876/\lambda 10830$ ratio is 0.4 and determine from Fig. 8 that the required physical conditions can be ($T = 2 \times 10^4$ K, $N_H = 1.1 \times 10^{11}$ cm$^{-3}$), ($T = 1.5 \times 10^4$ K, $N_H = 3.2 \times 10^{11}$ cm$^{-3}$) or ($T = 1.25 \times 10^4$ K, $N_H = 1.5 \times 10^{12}$ cm$^{-3}$), which produce, as seen from Fig. 23, a $y_{\text{vobs}} \sim 0.6$ in each case. The He I $\lambda 5876$ optical depth is greater than 50 (cf. Fig. 23, right-hand bottom panel), so a red absorption will be seen for certain lines of sight. For smaller observed values of $y_{\text{vobs}}$, however, a small filling factor of the emitting gas can be invoked to reduce the incidence of red absorption.

The 1-$\mu$m spectroscopic survey of CTTSs by EFHK06 also procure a set of He I $\lambda 10830$ profiles. Most of them have strong absorption features, so we consider solely the high 1-$\mu$m veiling group ($r_V \geq 0.5$) whose members are among the objects with the strongest He I $\lambda 10830$ emission. Excluding DK Tau, which has only absorption features, we determine the values of $y_{\text{vobs}}$ at $v_{\text{obs}} = 150$ km s$^{-1}$ to avoid the blue absorptions. They range from 0.3 to 1.3. Our model $r_V$ is 0.19, but even lowering the model values of $y_{\text{vobs}}$ considerably for comparison with the CTTSs with particularly high values of $r_V$, it appears that the infall model can produce the highest observed He I $\lambda 10830$ specific fluxes.
For both He\,\lambda\,10830 and \lambda\,5876, it is also clear that their emission is far too strong in a laminar wind flow and that the gas producing the helium emission can only occupy a fraction \( \leq 0.01 \) of the contour area shown in Fig. 1 in order to match the observed specific fluxes.

While the above analysis does not directly reject either model, we favour the wind region as the production site of strong helium emission for the following reasons:

1. Except near the accretion shock where UV photons are produced, the required temperature of \( T \geq 1.25 \times 10^4 \) K is difficult to generate and maintain in an accretion flow, which is primarily in free fall. On the other hand, it is reasonable to expect that acceleration of the gas in a wind would produce heating. Helium emission in the wind is also compatible with the earlier deduced condition for hydrogen emission in the wind in that the decreasing filling factor of the gas with increasing temperature from \( 10^4 \) to \( \sim 2 \times 10^4 \) K is consistent with the cooling and expansion of gas at a higher temperature.

2. An infalling flow is hard pressed to explain the strong blue asymmetry in the He\,\lambda\,5876 line wings. Judging the emission at \( |v_{\text{obs}}| \geq 150 \text{ km s}^{-1} \) (cf. fig. 7 of BEK01), we find that 18 out of the 22 He\,\lambda\,5876 BCs are stronger on the blue side, although three of them (AS 353 A, RW and GM Tau) have red absorptions that clearly accentuate the asymmetry, three (DD, DF and DO Tau) are stronger on the red side, and one (DE Tau) is too weak to discern. While lack of azimuthal symmetry or strong local inhomogeneities can generate asymmetries in the wing emission, their numbers of blue and red asymmetries should, for random lines of sight, statistically balance out. On the other hand, stellar occultation of the approaching accretion flow and disc occultation of the receding wind flow always produce a red and blue asymmetry, respectively. As seen from Fig. 1, the projected area of the \( v_{\text{obs}} = -150 \text{ km s}^{-1} \) contour in the infall model is only slightly larger than the projected stellar surface area. There is little leeway in adjusting the accretion flow geometry to avoid diminishing the blue wing emission, let alone enhancing it above the red wing emission.

3. The accretion flow model has greater difficulty producing the He\,\lambda\,10830 profile morphologies of those objects with strong emission (cf. fig. 4 of EFHK06). First, while most of the 38 objects have blue and/or red absorptions that make it difficult to decipher the wing emission at \( |v_{\text{obs}}| \geq 150 \text{ km s}^{-1} \), four have mostly emission, of whom BP Tau’s wing emission is too weak to compare its two sides, and the remaining three (CW and HW Tau, RW Aur A) all have stronger blue wing than red wing emission and will pose a challenge to the infall model. Secondly, when emission above continuum is observed along with absorption, the objects with strong emission tend to associate with blue absorptions indicative of a stellar wind (cf. fig. 10 of Kwan et al. 2007), whereas the objects with strong absorptions on the red side only do not have comparably strong emission (e.g. AA Tau, BM And, RW Aur B, LkCa 8 in fig. 4 of EFHK06 or fig. 2 of Fischer et al. 2008). The latter appear to have both a lower \( r_v \) and a smaller equivalent width in the emission above the continuum. Bearing in mind that the observed emission is contributed in part already by scattering of continuum photons, it is surprising, if the accretion flow is the site of helium emission, that the \( \text{in situ} \) emission is relatively weak in those accretion flows with particularly large widths and sizes that are needed to produce the broad and deep red absorptions (cf. sections 4.2 and 4.3 of Fischer et al. 2008). On the other hand, the above-mentioned correlation naturally follows if the wind flow is the site of helium emission. Also, in producing the needed \( \text{in situ} \) emission, there is more leeway afforded the wind model in the choices of \( r \) and filling factor of the high-temperature regions. The Monte Carlo simulations by Kwan et al. (2007) already indicate that the He\,\lambda\,10830 profile morphologies can be accounted for by a stellar wind, depending on the opening angle of the wind and the viewing angle. However, in light of the present finding that the emission regions do not occupy the bulk of the volume, there is the constraint that these regions share the velocity distribution of the gas occupying the bulk of the volume and responsible for the He\,\lambda\,10830 absorption.

In Section 6.1, where the He\,\lambda\,5876/\lambda\,10830 and \( \text{Pa}\gamma/\text{He}\,\lambda\,10830 \) ratios are discussed, it was mentioned that, even though their model values generated at \( T = 5000 \) K, \( 5 \times 10^{10} \leq N_H \leq 10^{12} \text{ cm}^{-3} \), are closest to the observed values (cf. Fig. 10), several reasons argue for a higher temperature. One, given in Section 6.2.2, is based on the difficulty of simultaneously matching \( \text{Pa}\gamma/\text{Pa}\beta \) at \( T = 5000 \) K. We add a second reason here, based on the weakness of the He\,\lambda\,5876 flux at \( T = 5000 \) K. As seen from Fig. 22, its model \( y_{\text{obs}} \) in the wind model increases from 0.1 to 1.3 as \( N_H \) increases from \( 5 \times 10^{10} \) to \( 10^{12} \text{ cm}^{-3} \). Reducing it by a factor of 2 upon rescaling it to the continua of the eight objects with the \( r_v \) between 0.8 and 20 and observed values of \( y_{\text{obs}} \) between 0.2 and 0.5, it is clear that the emitting gas needs to fill a substantial portion of the contour surface shown in Fig. 1 in order to match the observed values. This would, however, produce a blue He\,\lambda\,5876 absorption, which is never seen.

As \( T \) increases from \( 5000 \) K, the He\,\lambda\,5876/\lambda\,10830 ratio first decreases (cf. Fig. 8) and then rises again when \( T \) exceeds \( 10^5 \) K, so we expect the helium emission to arise from regions of temperature above \( 10^4 \) K. With \( T \) increasing from \( 10^4 \) to \( 2 \times 10^4 \) K, the specific fluxes of both He\,\lambda\,5876 and \( \lambda\,10830 \) increase rapidly (cf. Fig. 22), but then stay nearly constant for \( T \) between \( 2 \times 10^4 \) and \( 3 \times 10^4 \) K. If regions in the wind flow are raised to temperatures above \( 2 \times 10^4 \) K before cooling and expansion come into play, we anticipate them to be most effective in contributing to the observed helium emission when they are at temperatures near \( 2 \times 10^4 \) K.

When \( y_{\text{He}} \beta \) is reduced from \( 10^{-4} \) to \( 10^{-5} \text{ s}^{-1} \), the specific fluxes of both He\,\lambda\,10830 and \( \lambda\,5876 \) remain almost the same for \( T \geq 1.5 \times 10^4 \) K, but fall rapidly for \( T < 1.5 \times 10^4 \) K. This result can be anticipated from the large difference in the \( T_{\text{He,10830}} \) or \( T_{\text{He,5876}} \) contour between the two cases of \( y_{\text{He}} \beta \) in the temperature range \( 10^4 < T < 1.5 \times 10^4 \) K (cf. Fig. 4). This strong dependence of He\,\lambda\,10830 on \( y_{\text{He}} \beta \) in that temperature regime and the need for a higher density to obtain the same He\,\lambda\,5876/\lambda\,10830 ratio as \( T \) decreases from \( 2 \times 10^4 \) K led us to favour a temperature range from \( 1.5 \times 10^4 \) to \( 2 \times 10^4 \) K for the He\,\lambda\,5876 emission.

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It was pointed out in Section 6.3.2 that the observed Ca\textsc{ii} λ8498/Pa\textsc{γ} ratios of integrated fluxes appear to fall into two groups and, earlier in Section 7, that the strong Ca\textsc{ii} lines are narrower than Pa\textsc{γ}. For those objects with Ca\textsc{ii}/Pa\textsc{γ} ∼ 1 or less, the Ca\textsc{ii} emission can be produced by the set of physical conditions (8750 ≤ T ≤ 1.25 × 10^4 K, N\textsubscript{H} around 10^{11} cm\textsuperscript{-3}) producing primarily hydrogen emission. The larger issue, of course, concerns the strong emitters with Ca\textsc{ii}/Pa\textsc{γ} greater than ∼5. Their Ca\textsc{ii} profiles appear narrower than the Pa\textsc{γ} profiles in the data set of Edwards et al. (in preparation). There may already be a hint of this property from a comparison of non-simultaneous Ca\textsc{ii} and Pa\textsc{γ} profiles. Thus, the ratio of specific flux at v\textsubscript{obs} = −150 km s\textsuperscript{-1} to that at v\textsubscript{obs} = 0 is smaller for Ca\textsc{ii}λ8542, being 0.27, 0.1 and 0.4 for DG, DR and DL Tau, respectively (Muzerolle et al. 1998b), than for Pa\textsc{γ}, being 0.46, 0.33 and 0.54 (EFHK06).

In view of the possibly different Ca\textsc{ii}λ8498 and Pa\textsc{γ} profiles, we will, in comparing between model and observed values of the Ca\textsc{ii}λ8498 specific flux at v\textsubscript{obs} = −150 km s\textsuperscript{-1}, not utilize the constraint imposed on the physical conditions by the Ca\textsc{ii}/Pa\textsc{γ} ratio of integrated fluxes and consider all densities in the regime 5000 ≤ T ≤ 1.25 × 10^4 K as possible. Fig. 24 shows the dependences of y\textsubscript{v\textsubscript{obs}} on N\textsubscript{H}, T and line optical depth for the infall model. The observed values of DG, DR, DL Tau and RW Aur are 1.4, 0.6, 2.5 and 3.7, respectively (Muzerolle et al. 1998b). Anticipating that the model values of y\textsubscript{v\textsubscript{obs}} have to be reduced by a factor of ≥2, owing to the strong veiling at λ8498 (between 0.7 and 5.1) of those objects, we find that, with the exception of RW Aur, the infall model can produce the wing emission. We do not consider this strong enough evidence to argue against the infall model.

**Figure 24.** Dependences of Ca\textsc{ii}λ8498 and O\textsc{i}λ8446 specific flux (measured relative to the local continuum), y\textsubscript{v\textsubscript{obs}}, on the density and line optical depth for μ\textsubscript{H} = 2 × 10\textsuperscript{−4} s\textsuperscript{−1} and various temperatures (cf. Table 1) in the infall model.
The issue concerning the strong Ca II emission is, of course, its emission near line centre. While our specific flux calculation provides no aid in discriminating between the wind and infall models, the earlier derivation of requisite physical conditions for strong Ca II emission (i.e. \( T \leq 7500 \text{ K}, N_\text{H} \text{ close to } 10^{12} \text{ cm}^{-3} \)) remains true because of the much stronger dependence of the specific fluxes and flux ratios on \( N_\text{H} \) or \( T \) than on the velocity gradient. With this in mind, we do not think the wind region is the site of strong Ca II emission for the following reasons:

1. The physical condition of high density and low temperature is incongruous with that inferred from the helium and hydrogen emission and with the constraint imposed by the lack of blue absorption in Pa\( \gamma \). With the bulk of the wind region occupied by gas at a density of \(<10^{10} \text{ cm}^{-3}\) (cf. Fig. 4), the hydrogen-emitting regions at \( T \leq 1.25 \times 10^{4} \text{ K} \) and \( N_\text{H} \sim 10^{11} \text{ cm}^{-3} \), and the indication that lower temperature is associated with lower density and larger filling factor, it does not appear that the wind region has the physical conditions needed for strong Ca II emission.

2. The strong Ca II emission has a linewidth narrower than the Pa\( \gamma \) width or the typical velocity extent seen in a broad, blue He\( \lambda 10830 \) absorption. If Ca II emission originates from the wind region, then it is surprising, particularly in the case of DR Tau, that, while the Pa\( \gamma \) and He\( \lambda 5876 \) line shapes are quite similar, the Ca II line shape is so different.

3. The strong Ca II line appears symmetrical about line centre. It does not have the characteristic feature of a stronger blue emission as exhibited often by Pa\( \gamma \) and He\( \lambda 5876 \), and attributable to disc occultation of the receding wind.

The strong Ca II emission near line centre can be produced in an accretion flow. Its narrower width and different line shape from Pa\( \gamma \) will not be issues. The condition of \( T \leq 7500 \text{ K} \) is not unreasonable in light of Martin’s (1996) calculation of the thermal structure of the infalling gas and the density criterion of close to \( 10^{10} \text{ cm}^{-3} \) may not be too restrictive a constraint on the mass accretion rate, if the flow is diluted. There is a second possible site, the disc boundary layer where the gas dissipates part of its rotational energy before infalling along a stellar magnetic field line. We will elaborate on our preference in the next section when all the deductions and arguments put forth in Sections 5–7 are synthesized.

We will not comment much on the O\( \lambda 8446 \) specific flux. Our calculations indicate that the specific flux of O\( \lambda 8446 \) rises more rapidly than Pa\( \gamma \) with \( N_\text{H} \) increasing from \( 10^9 \) to \( 10^{12} \text{ cm}^{-3} \), causing O\( \lambda 8446/Pa\gamma \) to be higher at the conditions suitable for strong Ca II emission. Thus, there is a larger contribution to the O\( \lambda 8446 \) line from the Ca II emission region than is the case for Pa\( \gamma \).

8 ORIGINS OF THE STRONG H, HE I AND CA II LINE EMISSION

In Sections 5 and 6, we have presented calculations that shed light on the physical conditions giving rise to the absorption and emission features of most of the prominent lines observed in CTTSs. These are excitation calculations that include all the important physical processes affecting the atomic/ionic level population. They are facilitated by the presence of large velocity gradients expected in a wind or infalling flow in that the level population depends only on the local density \( (N_\text{H}) \), temperature \( (T) \), photon ionization rate \( (\gamma_{\text{He}}) \) and velocity gradient \( (\Delta v/\delta l) \). Thus, the resulting line emissivity \( (\text{erg s}^{-1} \text{ cm}^{-3}) \) can be calculated in the \( (N_\text{H}, T, \gamma_{\text{He}}, \Delta v/\delta l) \) parameter space and the ratio of two line emissivities will demarcate the requisite physical conditions. One of our findings is that the H and He I emission regions can occupy only a very small fraction of the wind flow. Our use of \( \Delta v/\delta l \) for the velocity gradient is then inaccurate and a more suitable choice would be \( \delta v/\delta l \) where \( \delta l \) is the linear dimension of an emission region and \( \delta v \) is the thermal/turbulence velocity width. However, as demonstrated in Section 7.1, the line emissivity is much more sensitive to \( N_\text{H} \) than to \( \Delta v/\delta l \) or \( \delta v/\delta l \), particularly when the line is optically thick, as is true for the H, He I and Ca II lines studied, so a large change in \( \Delta v/\delta l \) or \( \delta v/\delta l \) occasions only a small change in \( N_\text{H} \). Because the requisite physical conditions deduced for the gases responsible for the He\( \lambda 10830 \) absorption, H emission, He I emission and Ca II emission are so disparate, we are confident our conclusions are not fundamentally altered by the uncertainty in this parameter. For the same reason, even though we emphasize the two locations of \( r = 4 \) and \( 2.5R_* \), and the velocity \( v = 150 \text{ km s}^{-1} \), the results on the ordering of line opacities and the responses of line-emissivity ratios to \( N_\text{H} \), \( T \) and \( \gamma_{\text{He}} \), are applicable to other locations and not sensitive to the kinematic structure. In Section 7, we then distinguish between the wind and accretion flows by evaluating the specific fluxes of the more important lines at \( |v_{\text{obs}}| = 150 \text{ km s}^{-1} \). Here we recapitulate and synthesize the deductions reached separately in Sections 5–7 and, in conjunction with other arguments, decide on the locations of the H, He I and Ca II line emission.

Our conclusions drawn with regard to the H, He I and Ca II broad line emission in the following sections apply only to the strong emitters, specific examples of which have been mentioned in Sections 6 and 7. This qualification arises for the following reasons. First, the determination of observed line ratios, particularly He\( \lambda 5876/\lambda 10830 \), is not as reliable for weak emitters because of the presence of absorption features and the contribution to emission from scattering of continuum photons. Secondly, a couple of our arguments concern the limitations of an infalling flow on the production of line photons and therefore hinge upon the observational bar placed by the strong emitters.

8.1 Physical conditions of the gases producing emission and absorption lines

Below we summarize and discuss the findings obtained from a comparison between model results and observational data on line opacities and line ratios. They shed light on the properties of the gases responsible for the emission and absorption lines.

1. Optically thick H and He I emission lines. The gases producing the He\( \lambda \lambda 10830, 5876 \) and Pa\( \beta \), Pa\( \beta \) emission are very optically thick in those lines, in order to produce the observed He\( \lambda 5876/\lambda 10830 \), Pa\( \beta/Pa\beta \) and Pa\( \gamma/He \lambda 10830 \) ratios (cf. Sections 6.1.2 and 6.2.2).
(2) Separate physical conditions for H, He I and Ca II emission. The need for distinct physical conditions conducive to H and He I emission is demonstrated clearly in Fig. 10, where it is seen that many of the marked CTTSs, which are among the strong He λ5876 emitters, have He λ5876/λ10830 and Paβ/He λ10830 ratios that cannot be simultaneously produced with a common temperature range, but must require that the bulk of the He λ5876 emission be produced at temperatures higher than those producing the bulk of the Paβ emission (cf. Section 6.1.2).

For the strong He λ5876 emission, we favour the temperature range $1.5 \times 10^4 \leq T \leq 2 \times 10^4$ K. The upper bound is adopted, because the line emissivity at a higher temperature is not higher (Section 7.2), while presumably the emission area is smaller. The lower bound is adopted, because at a lower temperature, the line emissivity is much more dependent on $\gamma_{\text{He}}$ and a higher density is needed to produce the same He λ5876/λ10830 ratio, as seen in Figs 8 and 9 (cf. Section 7.2). The same figures also show that within the favoured temperature range, He λ5876/λ10830 rises steeply with increasing density and the density needed to produce He λ5876/λ10830 $\geq 0.3$ is $\geq 10^{11}$ cm$^{-3}$.

The physical conditions responsible for H emission, as deciphered by examining the ratios among the hydrogen lines themselves, namely Pa$n_u$/Paβ and Bry/Pa$n_u$, indicate that $N_H$ is centred around $10^{11}$ cm$^{-3}$ for $T \geq 8750$ K and $5 \times 10^{10}$ cm$^{-3}$ for $T \leq 7500$ K (cf. Section 6.2.2). In conjunction with the observed Paβ/He λ10830 ratio, however, the temperature can be narrowed to the range $5000 < T < 1.25 \times 10^4$ K. The upper bound comes from the small Paβ/He λ10830 ratio ($\leq 0.2$) that can only be generated at $N_H$ around $10^{11}$ cm$^{-3}$ at $T \geq 1.5 \times 10^4$ K (cf. Figs 8 and 9, right-hand top panel). The lower bound is due to the need of very high densities ($> 2 \times 10^{12}$ cm$^{-3}$) to produce Paβ/He λ10830 $> 0.8$ (cf. Fig. 11, left-hand top panel) and the concomitant result of Paβ/He λ10830 lower than observed (cf. Figs 8 and 9, right-hand top panel, Section 6.2.2). We further favour the temperature range conducive to H emission being $8750 < T < 1.25 \times 10^4$ K, because the line emissivities, particularly Paβ and Bry, decrease rapidly with $T$ decreasing below 8750 K (cf. Figs 19–21, right-hand top panel, Section 7.1). In this temperature range and with $N_H$ around $10^{11}$ cm$^{-3}$, the Ca II λ8498/Paβ ratio is also in line with the observed values among the weak Ca II broad emitters (cf. Section 6.3.2).

The observed He λ10830/Paβ ratio, typically $\sim 1.5$ among the strong emitters, also indicates that the He I emission area is smaller than the H emission area. Thus, if the emission area contributing to H emission, primarily with $T \sim 2 \times 10^4$ K, were the same as that with $T \sim 10^4$ K, He λ10830/Paβ would be $\geq 10$, because while the Paβ emission from each area is about the same (cf. Fig. 19, right-hand top panel), the He λ10830 emission at $2 \times 10^4$ K and $N_H$ around $10^{11}$ cm$^{-3}$ is $\geq 20$ times the Paβ emission (cf. Figs 8 and 9, right-hand top panel).

The physical conditions responsible for strong broad Ca II emission are most evident in Figs 17 and 18 where the observed Ca II λ8498/λ8546 ratios greater than $\sim 7$ point to a temperature $T \leq 7500$ K and then from Fig. 16 (left-hand top panel) a density $N_H$ close to $10^{12}$ cm$^{-3}$ (cf. Section 6.3.2). They can also be inferred from the large Ca II λ8498/Paβ values. There are, however, also CTTSs observed to have Ca II λ8498/λ8546 and Ca II λ8498/Paβ of $\sim 1$ or less (e.g. BP Tau in Figs. 17 and 18), and Paβ emission strength not much weaker than those among the strong Ca II emitters. Their H and Ca II emission indicates $8750 < T < 1.25 \times 10^4$ K and $N_H$ around $10^{11}$ cm$^{-3}$ (cf. Section 6.3.2), as mentioned just earlier. Such physical conditions also naturally arise through cooling of the regions responsible for He I emission, which is often observed to be strong when Ca II emission is strong. Thus, we conclude that this set of physical conditions producing primarily H emission is also present. It may indeed be the most prevalent, since it is needed when both He λ5876 and Ca II emission is weak.

When the latter emission is strong, the other two sets, namely $(1.5 \times 10^4 \leq T \leq 2 \times 10^4$ K, $N_H \geq 10^{11}$ cm$^{-3}$) for the bulk of the He I emission and $(T \geq 7500$ K, $N_H \sim 10^{12}$ cm$^{-3}$) for the bulk of the Ca II emission, are also present.

The fundamental cause of the distinct sets of physical conditions conducive to H, He I and Ca II emission is the different H, He I and Ca II atomic/ionic structures, making their efficacies in photon production sensitive to different temperature ranges. Thus, while He λ10830 and λ5876 emission rises rapidly with increasing $T$ up to $2 \times 10^4$ K, hydrogen emission in Balmer, Paschen and higher order lines is not more efficient at $T > 10^4$ K, because the stronger collisional ionization reduces the neutral hydrogen fraction, and a temperature decreasing below $10^4$ K clearly favours Ca II over H emission. The separate physical conditions we identify for the H, He I and Ca II emission regions are therefore understandable and not particular to our excitation model, despite its simplifications and approximations.

(3) Higher densities for the gases producing the H and He I emission lines than that of the gas producing the broad, blue absorption. Broad, blue absorptions indicative of a stellar wind are often seen in He λ10830 but almost never seen in He λ5876, Paβ or Paγ. This observational constraint, in conjunction with finding 1 above, means that a laminar wind will not be able to produce simultaneously the observed absorption and emission features seen in either He I or H lines. This dilemma is the same as that realized in earlier wind models of hydrogen line emission (Hartmann et al. 1990). Anticipating that the temperature of the gas occupying the bulk of the wind volume is $10^4$ K, the constraint placed by the observed absorption features limits the density of this gas to be no more than $\sim 10^{10}$ cm$^{-3}$ (cf. Fig. 4), which is considerably lower than the density of the H or He I emission gas. Thus, a highly clumpy flow is called for if the H and He I emission regions also reside in the wind (cf. Sections 7.1 and 7.2).

(4) In order to produce a $\Delta$He λ10830 $\geq 1$ in the radial wind, the minimum density needed at the location where the wind reaches a speed of 150 km s$^{-1}$ is $\sim 5 \times 10^9$ cm$^{-3}$ (cf. Fig. 4) and the corresponding mass-loss rate is $\sim 0.5 \times 10^{-6}$ M⊙yr$^{-1}$.

(5) The red absorption gas has a temperature $T < 10^4$ K and a density $N_H$ greater than or about $10^{11}$ cm$^{-3}$. The former condition is needed to ensure a more prevalent occurrence of red absorptions in the Na I doublet than in Paβ (cf. Fig. 6, Section 5.2). The latter condition is inferred from the observed ordering of the lines in propensity of showing a red absorption.

(6) UV photoionization is necessary to produce the broad red and blue absorptions seen in He λ10830. This is demonstrated forcefully by the more prevalent occurrence of red absorptions in He λ10830 than in Paβ, given that the low temperature of the absorption gas ($T < 10^4$ K as noted above) renders collisional excitation futile (cf. Section 5.2). UV photoionization of He I in the gas occupying the bulk of the...
wind flow is also needed to produce He I λ10830 blue absorptions, since it is most likely that the temperature of that gas is also \(<10^4\, \text{K}\), if the hydrogen emission regions occupy only a small fraction of the volume and have a temperature \(8750 \leq T \leq 1.25 \times 10^4\, \text{K}\). With the main depopulation path of the highly metastable He I λ10830 lower level being collisional excitation to the 2s^1S state, whose rate decreases with increasing \(T\) and \(N_e\), a \(r_{\text{He I},10830} \sim 1\) is quite realizable at \(T \leq 6250\, \text{K}\) even for low-ionization fluxes (cf. Section 5.1). We therefore think that UV photoionization is also the excitation mechanism responsible for the sharp blue absorptions and central absorptions seen in He I λ10830.

UV photoionization of H I in the bulk of either the wind or the accretion flow is not as crucial, but it likely contributes, particularly if the gas temperature is \(<7500\, \text{K}\), given that UV photons are already present for He I ionization. In the regions producing either H or He I optical and infrared line emission, collisional excitation is the primary mechanism and UV photoionization plays a less-significant role.

We will make use of these findings and additional ones deduced in Section 7, as well as arguments based on correlations among absorption and emission features, to support our decisions on the emission sites of the strong He I, H and Ca II lines in the following two sections.

8.2 Wind region as the site of the strong He I and H emission

We first summarize the arguments against the accretion flow as the appropriate site:

1. Difficulty for an infalling flow to account for the stronger blue wing emission. In Sections 7.1 and 7.2, we have listed CTTSs to show that both Paβ and He I λ5876 have distinctly a stronger blue wing. In an infalling flow, the emission area at \(v_{\text{obs}} \leq -150\, \text{km s}^{-1}\) is not much larger than \(\pi R_e^2\) (cf. Fig. 1), so it is very difficult to avoid stellar occultation of the infalling flow approaching an observer. This preferential attenuation of the blue wing emission is intrinsic to the infall flow geometry, but it is counter-effective to the observed trend.

2. Difficulty for an infalling flow to produce the observed fluxes in the line wings. The finding that the H and He I emission lines are optically thick means that their specific fluxes depend on the emission area and the line excitation temperatures. For an infalling flow, its geometry severely confines the projected areas with large values of \(|v_{\text{obs}}|\). The hydrogen lines face, in addition, a strong limitation on their excitation temperatures through the \(n \to n + 1\) collisional excitations that lead to rapid ionization at \(T > 10^4\, \text{K}\). As a result, we find that the observed Paβ and Paγ fluxes at \(v_{\text{obs}} \geq 150\, \text{km s}^{-1}\) cannot be produced by our infall model over the broad ranges of physical parameters explored. This result is not surprising, as it is simply an alternative, albeit more explicit, way of phrasing previous findings that the hydrogen profiles calculated from accretion flow models are narrower than observed (Folha \& Emerson 2001). Interestingly, the steep rise of the helium line emissivity with increasing \(T\) from \(10^4\) to \(2 \times 10^4\, \text{K}\) (cf. Fig. 23) enables the infall model to reproduce the observed He I line fluxes. However, a temperature \(\geq 1.25 \times 10^4\, \text{K}\) is required and we think it is difficult to heat the gas, which is primarily in free fall, to those temperatures, except possibly near the impact sites where photoionization heating by photons generated in the shocks can be important. It has also been deduced earlier that the gas occupying the bulk of the infalling flow has a temperature \(<10^4\, \text{K}\), so we do not think the accretion flow generally has the physical conditions conducive to strong He I emission.

3. Weak or no correlation between strong He I λ10830 red absorptions and strong Paβ or He I λ5876 emission. If Paβ or Paγ or He I λ5876 is produced in the accretion flow, it is reasonable to expect that presence of a strong He I λ10830 red absorption should be accompanied by a strong Paγ or Paβ or He I λ5876 line. From figs 3 and 4 of EFHK06, it is seen that among the CTTSs with the strongest He I λ10830 red absorptions, two (DK Tau, YY Ori) have moderate Paβ emission, but seven (AA, GI, DN and V830 Tau, BM And, RW Aur B, LkCa 8) have only weak Paβ emission. In the simultaneous 1-μm and optical data sample of Edwards et al. (in preparation), of the six objects with strong He I λ10830 red absorptions, only one (DR Tau) has strong Paγ and He I λ5876 emission, while the rest (DK, GK, AA and GI Tau, BM And) have weak Paγ and He I λ5876 (BC) emission.

4. It is noted in Section 5.2 that when \(r_{\text{He I},10830} \sim 1\), the He opacity is close to or higher than \(r_{\text{He I},5876}\). Then, if the He emission originates in the accretion flow, which is probably within the radial wind, it will be scattered by the wind and the resulting observed profile will be highly asymmetric with the red side much stronger than the blue side. The collection of 31 Hz profiles in fig. 15 of BEK01 shows that such profiles are rare, with AS 353A, DR Tau and DO Tau being the lone examples.

5. Absence of correlation between He I λ5876 BC and NC emission. A narrow line, either on top of a broad one or by itself, is seen in Paγ (EFHK06), the Ca II infrared triplet (Muzerolle et al. 1998b), and most frequently and prominently in He I λ5876 (BEK01). This observational result is understandable if the narrow line emission is formed in the post-shock regions at the impact sites of the infalling flow on the star, since the He I lines form at higher temperatures and will likely reach higher excitation temperatures than the H and Ca II lines. The He I λ5876 NC strength, being sensitive to the surface area of impact and the pre-shock density, is then a good indicator of the accretion flow magnitude. Then, if the He I λ5876 BC arises from the accretion flow, one would expect a correlation between the BC and NC strengths. Even though, with the available information of BC and NC equivalent widths (\(W_\lambda\)) and red veiling (\(r_{\text{BC}}\)), a direct plot of the BC flux versus NC flux for the sample of 31 CTTSs cannot be made, we can obtain a sense of this correlation from fig. 1 of BEK01. There it is seen that among those CTTSs with the highest values of \((1 + r_{\text{BC}})W_\lambda\) in the BC four (CW, HN and DG Tau, AS 353A) have no or comparably much weaker NC, three (DL Tau, RW and UY Aur) have a comparably weaker NC (but typical, in terms of \((1 + r_{\text{BC}})W_\lambda\) in the NC, relative to the whole sample), and one (DR Tau) has a comparatively weaker NC (but strong relative to the whole sample). Then, among those CTTSs with the highest values of \((1 + r_{\text{BC}})W_\lambda\) in the NC, two (DD and DO Tau) have a comparable BC and two (FM and HK Tau) have no BC. We judge from this comparison between the BC and NC values of \((1 + r_{\text{BC}})W_\lambda\) that there is no correlation between their strengths and conclude that the broad helium line emission is generally not related to the accretion flow. We can extend this conclusion to the hydrogen line emission. Even
though the 38 Paγ profiles in EFHK06 and the 31 HeⅡλ5876 profiles in BEK01 are not simultaneous, we see that the group of CTTSs with strong (weak) HeⅡλ5876 BC is almost the same as the group of CTTSs with strong (weak) Paγ emission.

While the above-mentioned characteristics of the H and HeⅠ line profiles pose severe difficulties for an infalling flow, they favour a wind flow for the line origin. These include the stronger blue wing emission and the usually blue centroid (both owing to disc occultation of the wind flow receding from an observer), high blue wing velocities (as there is no limitation on the terminal wind speed) and the association of strong H and HeⅠ emission with blue absorptions in HeⅡλ10830 that are indicative of a stellar wind. The H and HeⅠ specific fluxes, either at the line wings or at the line centre, are not fundamental issues in the wind model, even though their observed values, as well as the absence of blue Paγ and HeⅡλ5876 absorptions, stipulate that the emission arises from only a small fraction of the wind volume. The physical conditions and filling factors deduced for the H and HeⅠ emission regions may indeed be brought about through the expansion and cooling of high-temperature, high-pressure clumps. It does remain to be demonstrated, however, that there is a viable acceleration process that produces a highly clumpy flow. Despite this uncertainty, the overall positive comparison between model and observed line profiles and strengths, together with earlier arguments against the accretion flow, led us to decide squarely on the wind region as the site of the strong H and HeⅠ emission.

The possibility of strong photoionization heating of the accretion flow near the impact shock, however, means that HeⅠ line emission, possibly strong, may occasionally arise from the accretion flow. In Section 7.2, it is pointed out that, statistically, the HeⅡλ5876 profiles show a definite preference for a stronger blue wing, but three (DD, DF and DO Tau) clearly have a red centroid (cf. fig. 7 of BEK01). Such a red asymmetry can be brought about if the emission arises primarily from the part of the accretion flow with a z-velocity component towards the disc, specifically the part at polar angle θ < 54.7° in a dipolar trajectory, which is the part close to the impact shock (BEK01). The three profiles are much weaker at v_20 ≤ −150 km s⁻¹ than at v_20 ≥ 150 km s⁻¹ and also have strong NCs. Thus, based on the HeⅡλ5876 profile morphologies, an origin of the helium emission in an accretion flow is favoured for those three objects. A possible caveat is the uncertain extent to which azimuthal asymmetries in flow geometry and physical conditions affect the profiles. This can only be answered by time-monitoring in both optical and 1-μm spectral regions, so that the relative contributions from the radial wind and accretion flow can be assessed through analysis of both HeⅡλ10830 and λ5876 profiles. The great majority of observed HeⅠ profiles, however, show characteristics that favour a wind origin.

8.3 Disc boundary layer and accretion flow as sites of CaⅡ line emission

In Section 7.3, we argue against the wind region as the site of the strong CaⅡ infrared triplet emission and mention that, in addition to the accretion flow, the disc boundary layer is a possible site. In essence, this boundary layer, where the accreting particles dissipate part of their rotational energies before falling along the stellar field lines, is the base of the accretion flow. Here we list the reasons why the CaⅡ emission from this base may be significant.

(1) The energy dissipated in the disc boundary layer is a significant fraction of the total gravitational potential energy released by the infalling gas. Most of this energy emerges in a photon continuum, which we identify as part of the observed continuum excess around 1 μm. Like the situation in a stellar atmosphere, the chromospheric region above the continuum formation zone in this boundary layer will produce CaⅡ infrared triplet emission, but the line profile, instead of being narrow, will be Doppler broadened by the strong rotational motion. There appears to be a correlation between the CaⅡλ498 line strength and the excess 1-μm continuum flux (Fischer et al., in preparation)

(2) The profiles of the strong CaⅡλ8498 emitters (Muzerolle et al. 1998b; Edwards et al., in preparation) appear fairly symmetrical. To be sure, the Pa 16 line in the red wing of CaⅡλ8498 needs to be subtracted and the profile contains uncertain amounts of contribution from the accretion flow and the hydrogen emission regions in the wind flow. However, if this profile characteristic, as well as the narrower width in comparison with the H and HeⅠ lines, holds up in more strong CaⅡ emitters, it is consistent with rotational broadening. The same broadening also produces a depression at line centre, but the centrally peaked CaⅡ emission from the accretion flow and from the hydrogen emission regions in the wind flow may fill it up.

(3) The continuum produced in the disc boundary layer, like that produced at the accretion footpoints, increases with mass accretion rate, but unlike the latter, also increases with decreasing distance of its location from the star. Thus, there may only be a weak correlation between the optical/UV and 1-μm continuum excesses. Then, if the CaⅡ narrow and broad emission arises from the accretion shocks and disc boundary layer, respectively, they will reflect a similar relation between their strengths. The 11 CaⅡ line profiles shown in fig. 1(d) of Muzerolle et al. (1998b) show that the weak broad CaⅡ emitters have small values of r_8600 and distinct NCs, while the strong broad emitters have larger values of r_8600 and at best weak NCs, although their appearances are rendered less conspicuous by the presence of the strong broad emission and the stronger underlying continuum. It does appear that the correlation between the CaⅡ broad and narrow emission strengths is weak, but more objects are needed for better statistics.

We are confident that the strong CaⅡ emission does not originate from the wind region and think that a good portion of it arises from the disc boundary layer.
9 DISCUSSION

In this section, we discuss the implications of the findings described earlier in Section 8. They include the need of a clumpy wind, identifying the source of ionizing photons and determining the underlying cause of the correlation between strong Ca II and He I line emission. Most of this discussion is speculative in nature.

A clumpy wind has also been advanced by Mitskevich et al. (1993). Their motivation is to explain the blue absorption seen in Hβ as being an intrinsic part of line formation in a stellar wind that accelerates to a peak velocity and then decelerates towards zero. Our arguments for a clumpy wind are not dependent on a specific velocity structure. One is based on the lack of blue absorptions in Paγ and He I λ5876, whose observed intensities are not strong enough to cover up underlying absorptions of the continuum. So, whatever be the velocity structure, the optically thick emission regions can only screen a small fraction of the stellar surface at each v\(\text{obs}\). The other is based on the need to reduce the model intensities of all the H and He I lines, calculated assuming a smooth flow with physical conditions suitable to produce the line ratios, in order to match the observed values.

With regard to the accelerating and then decelerating wind advocated by Mitskevich et al. (1993), we have several reservations for its general applicability to CTTSs. First, the Paβ and Paγ lines are also quite optically thick (\(τ_{\text{Paγ}} > 20\), cf. Fig. 11) but, unlike Hβ, they rarely show a blue absorption. Secondly, the maximum blue velocity seen in He I λ10830 or Paγ is often \(≥350\) km s\(^{-1}\) (EFHK06). If this is reached at \(3R_\ast\), the locale adopted by Mitskevich et al. (1993), the gravitational pull of a CTTS of \(1M_\odot\) and \(2R_\odot\) will only decelerate it to \(240\) km s\(^{-1}\), not low enough to explain most of the Hβ blue absorptions, whose maximum depths usually occur at \(v_{\text{obs}} ≥ 100\) km s\(^{-1}\) (cf. fig. 15 of BEK01). We think that the Hβ blue absorptions, other than the broad ones that originate in a radial wind, are caused by disc winds lying beyond the hydrogen emission zone. They are fairly sharp and narrow, similar to those produced by a disc wind scattering the stellar continuum (Kwan et al. 2007). However, because the hydrogen emission zone is more extensive than the stellar surface, a revised modelling is needed to ascertain the absorption profile.

One possible mechanism for accelerating the gas and initiating a clumpy flow is the occurrence of multiple coronal mass ejections. These ejections will have to be much more energetic than those in the solar corona and occur at a much higher frequency, since the wind density needs to be \(≥5 \times 10^9\) cm\(^{-3}\) to produce a \(τ_{\text{HeI}10830} \geq 1\), thereby implicating a minimum mass-loss rate of \(≥0.5 \times 10^{-9} M_\odot\) yr\(^{-1}\). The clumpy gas distribution can arise from the initial multiple ejection centres and the shocked regions produced when the different ejecta intersect.

There can be several sources for the photons needed for helium ionization. One, known already, is the shocked region at the footpoint of the accretion stream. For an impact velocity of \(300\) km s\(^{-1}\), the temperature in the post-shock region reaches \(10^6\) K. Cooling by free–free emission will produce many photons more energetic than 24.6 eV. A fraction of this radiative luminosity will propagate towards the star and will be re-processed into optical and UV continuum seen as veiling, but a fraction will propagate away from the star and may even escape absorption by the infalling flow if the gases fill only partially the overall accretion envelope (Fischer et al. 2008). For an accretion rate of \(3 \times 10^{-6} M_\odot\) yr\(^{-1}\), if 1 per cent of the radiative luminosity generated escapes, the luminosity in photons with energies \(≥24.6\) eV will be \(≈10^{-4} L_\odot\). Ionizing photons can also originate from the wind region, for example, at the sites of coronal mass ejections and the filaments formed by interacting ejecta, likely with comparable luminosities. As mentioned before, however, the issue with ionizing photons is not so much the availability of sources, but more the propagation distance from a given source. Both a clumpy accretion flow and a clumpy wind help to increase the propagation length, but a definitive understanding will probably need Monte Carlo simulation of the propagation of photons with energies both near and far away from ionization thresholds and of photons produced from subsequent recombinations to the ground state, as well as consideration of multiple ionizations when the ejected electron has an energy exceeding the H I ionization threshold.

We have identified the disc–magnetosphere interface where accreting gases dissipate part of their rotational energies before falling along stellar field lines as a source of energy for both the veiling continuum near 1-\(\mu\)m and the Ca II line emission and there appears to be a correlation between \(R_\gamma\) and the strength of the Ca II infrared triplet (Fischer et al., in preparation). The correlation of strong He I line emission and strong Ca II line emission is probably a consequence of the correlation between He I emission and \(R_\gamma\). BEK01 noted that when He I λ5876 has a signature indicative of a wind origin, the NC is comparatively weak and may be even absent, and suggested a scenario in which high accretion rates or weak stellar fields may cause the magnetosphere to be crushed sufficiently that the disc extends almost to the star, so impact velocities of the accreting matter will be smaller and less energy will be available for the NC emission. At the same time, a larger fraction of the rotational energy of the accreting matter needs to be dissipated, possibly leading to a larger \(R_\gamma\). Also, the equatorial region of the star will be preferentially torqued and the resulting differential rotation with latitude may induce stronger magnetic activities for field lines anchored at the polar region and produce a greater mass ejection, leading to stronger He I line emission.

The minimum coronal density we derive is \(≈5 \times 10^9\) cm\(^{-3}\) and the corresponding mass-loss rate is \(≈0.5 \times 10^{-9} M_\odot\) yr\(^{-1}\) in order to produce a \(τ_{\text{HeI}10830} \geq 1\). The density and mass-loss rate probably span a factor of \(≥10\) in range. They are germane to the investigation of the role of a stellar wind on the angular momentum evolution of a CTTS (Matt & Pudritz 2005) and the numerical simulation of magnetospheric accretion of matter in various field geometries (Romanova et al. 2009).

We have applied our model calculations to CTTSs because of their extensive information on line fluxes and profiles. These results on line optical depths, line emissivity ratios and specific line emissivities from local excitation calculations are usable for the broad lines of other pre-main-sequence stars, because the individual star affects the calculations only through its stellar and veiling continua. For H I, He I and O I excitations, the effects of these continua via photoionization of and stimulated absorption by excited states are not significant. The effects on photoionizing Ca II and Na I are stronger. They can be gauged from comparing Figs 6 and 7, which show results for the locations of 4 and
2.5R*, respectively. The Ca II and Na I line optical depths in the r = 2.5R*, calculation are affected by a smaller 2r/v factor in the τ expression and a larger photoionization rate by a factor of 2.56 (cf. Section 5.1). Together the two factors affect little the τ_{Ca II,8498} contour, but produce a displacement of the τ_{Na I,5892} contour by a factor of ~2 or less.

10 CONCLUSION

We first summarize the main results of this work and then comment on future studies that may shed further light on the formation and origins of the line emission in CTTSs. Our primary conclusions are as follows:

(1) UV photoionization is needed to produce the He II λ10830 opacities in both the accretion flow and radial wind that generate the observed red and broad blue absorptions, respectively. It is also the most probable excitation mechanism responsible for the narrow, sharp blue absorptions and the central absorptions.

(2) The strong He I and H line emission originates primarily in a radial outflow that is highly clumpy. The bulk of the wind volume is filled by gas at a density ~10^3 cm^{-3} and optically thick to He II λ10830 and Hα, but optically thin to He II λ5876, Paγ and the Ca II infrared triplet. The optically thick He II λ5876 emission occurs mostly in regions of density ≥10^{11} cm^{-3} and temperature ≥1.5 × 10^5 K, while the optically thick Hα and Paγ emission occurs mostly in regions of density around 10^{11} cm^{-3} and temperature between 8750 and 1.25 × 10^5 K. In producing the observed line flux at a given v_{obs}, the area-covering factor of these emission clumps is sufficiently small to not incur significant absorption of the stellar and veiling continua in either He I and H lines. He I emission, possibly strong, may occasionally arise from the accretion flow close to the impact shock as a result of photoionization heating by the UV photons.

(3) The strong Ca II line emission likely arises in both the magnetospheric accretion flow and the disc boundary layer where the gases dissipate part of their rotational energies before falling along magnetic field lines. The needed density and temperature are ~10^{12} cm^{-3} and ≤7500 K, respectively. Weak Ca II line emission, on the other hand, can come from the clumps producing the H emission in the wind.

We plan to follow up this work by presenting all the optical and 1-μm spectral data procured simultaneously or near-simultaneously (Edwards et al., in preparation), with an aim to provide more comparisons with model results on the ratios and specific fluxes of not only lines alluded to here, but also others, such as Hβ, Hγ and He II λ6678, that are useful as consistency checks. The larger sample will also convey information on the ranges of variation in the line properties and related physical conditions. It will also apprise of possible correlations among continuum veiling, line ratios and line fluxes that may elucidate the relationship between accretion and stellar wind activities.

A future project that will be enlightening is time-monitoring at optical and 1-μm spectral regions simultaneously more than a rotational period for several CTTSs. The simultaneous coverage of both spectral regions is needed to include He II λ10830, the key indicator of intervening kinematic structures through its absorption features, and the many optical lines whose strengths and profiles delimit the requisite physical conditions. The time-sequence data can test the often presumption of azimuthal symmetry for either the accretion flow or radial wind and its physical conditions, and provide detailed information for a realistic modelling of both the accretion flow and radial wind geometric structures.

Further understanding of the origins and formation of spectral lines can be gained from studying the highly ionized lines, notably C IV λ1549, and optical as well as UV Fe II and Fe I lines. There are two potential sites for the highly ionized lines, the stellar wind region and the part of the accretion flow close to the impact shock. It is important to determine their origin, through analysis of their profiles and excitation mechanisms, and check for consistency/conflict with the formation and origin of the He I lines. Fe II and Fe I emission can be significant from both the H and Ca II emission regions. Their strengths and profiles, in comparison with those of H and Ca II lines, as well as the similarities/contrasts between Fe II and Fe I lines, can provide additional information on the excitation conditions in those regions.

Eventually, confirmation of a clumpy stellar wind requires an understanding of the energy generation and acceleration process. It also needs numerical simulation of the ejecta and gas flows to check if the resulting density, temperature and velocity distributions are consistent with those inferred from the line emission strengths and profiles.

In conclusion, the frequent presence in He II λ10830 of absorption features indicative of a radial flow, and our deduction of strong line emission from this flow, if correct, indicates that a significant stellar wind, in terms of mass-loss and energy-output rates, is an essential component of the star formation process. It is clearly in response to the active accretion of matter on to the star. This dynamic action–reaction between the accretion disc and star may ultimately determine the mass and angular momentum of the emerging star.

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APPENDIX A: ATOMIC PARAMETERS

A1 H1

We obtain the rate coefficients for collisional transitions between H1 levels from Anderson et al. (2002). Unfortunately, they have included only levels $n = 1–5$. The $\Delta n = 1$ transitions have the largest rates and we fit their collisional rate coefficients, that is, $C_{n+1,n}$. $1 \leq n \leq 4$, with the expression $C_{n+1,n} = C_{2,1}(n/2) \exp(\chi \ln(n/2))$, where $\kappa$ and $\chi$ are the fitted parameters, and extrapolate to $5 \leq n \leq 14$. To check the accuracy of the extrapolated rate coefficients, we compare them to those obtained from the formulae given in Johnson (1972). The differences are $\sim 30$ per cent, comparable to those for $n \leq 4$. The $\Delta n = 2$ collisional rate coefficients are the next strongest and, for the same lower level, are a factor of $\sim 2.5$ or more smaller than the $\Delta n = 1$ ones, that is, $C_{n+2,n}/C_{n+1,n} \leq 1/2.5$. They are quite different from Johnson’s values, being larger by a factor of between 1.2 and 1.8 at $T = 10^4$ K. We fit them, that is, $C_{n+2,n}$, $1 \leq n \leq 3$, with the simple power-law expression...
\[ C_{n+2,n} = C_{4,3}(n/2)^\delta, \] where \( \delta \) is the fitted parameter, and the proviso that the extrapolated \( C_{n+2,n} \) for \( 4 \leq n \leq 13 \) is no more than a factor of \( \sim 2 \) larger than Johnson’s corresponding rate coefficient. The \( \Delta n = 3 \) collisional rate coefficients are a factor of \( \sim 2.3 \) or more smaller than the \( \Delta n = 2 \) ones, that is, \( C_{n+3,n}/C_{n+2,n} \leq 1/2.3, 1 \leq n \leq 2. \) They are larger than Johnson’s values by a factor of \( \sim 2.2 \) at \( T = 10^4 \) K, so we again extrapolate them to higher values of \( n \) with a power-law fit, that is, \( C_{n+3,n} = C_{3,2}(n/2)^\gamma \), and the proviso that the extrapolated rate coefficients are no more than a factor of \( \sim 2.5 \) larger than Johnson’s. For level \( n = 6 \) only, the upper state of the Pa\( ^\gamma \) transition, we also extrapolate the rate coefficients of Anderson et al. (2002) to obtain \( C_{6,1} \) and \( C_{6,2}. \) We fit \( C_{n,1}, 2 \leq n \leq 5, \) with the expression \( C_{n,1} = C_{2,1}[(n - 1)^\eta] \) and \( C_{n,2}, 3 \leq n \leq 5, \) with the expression \( C_{n,2} = C_{3,2}[1/(n - 2)]^\epsilon, \) and extrapolate to \( n = 6. \) It turns out that in the regions responsible for the observed emission, the electron densities are quite high, \( \gtrsim 3 \times 10^{10} \) cm\(^{-3}\), so collisions dominate the population exchange among the higher energy levels. The level population then depends mostly on the scaling of the collisional rate coefficients with \( n \), which we hope to capture with our fitting procedure.

We calculate the rate coefficients for collisional ionization from the formulae given in Johnson (1972) and determine the three-body recombination coefficients by detailed balance. In addition to UV photoionization of the ground state, we include photoionizations of the \( n = 2–4 \) levels by the stellar and veiling continua. The photoionization cross-sections are gathered from Allen (1973) and the radiative recombination cross-sections are gathered from Allen (1973) and the radiative recombination coefficients to the \( 15 \) levels from Seaton (1959). We also include radiative absorptions of the stellar and veiling continua at all permitted line transitions except the Lyman ones. In principle, the probability of stimulated absorption is the same as the escape probability calculated with the velocity gradient along the trajectory of the incident continuum photon. It is therefore different from the escape probability of a photon emitted by spontaneous emission, in general. We will ignore this distinction here, since it does not affect the order of the lines in their opacity magnitudes.

### A2 He\( ^I \)

We adopt the helium atomic parameters and many of the collision strengths from the Chianti data base (Young et al. 2003), which lists energy levels, Einstein A rates and collision strengths for transitions involving one of the lowest energy states. For the rest of the collision strengths, we obtain them from Sawey & Berrington (1993), the basis of Chianti’s compilation. We determine the collisional ionization rate coefficient of the ground state from the cross-sections measured by Montague, Harrison & Smith (1984) and that of the \( 2s^1S \) level from the theoretical calculation of cross-sections by Taylor, Kingston & Bell (1979), and gather the rest from the compilation of Benjamin, Skillman & Smits (1999). We obtain the \( 2s^1S \) and \( 2s^1S \) photoionization cross-sections from Fernley, Taylor & Seaton (1987) and the radiative recombination rate coefficients from Benjamin et al. (1999).

The absence of collision strengths involving \( n \geq 5 \) levels is the reason why our helium model atom has only 19 levels. There is the concern that, like the case of hydrogen, collisional excitation from the \( n \) levels to higher ones will likely lead to ionization, so limiting helium to \( 19 \) levels will underestimate the helium ionization rate and overestimate the helium level population. This problem may be somewhat less severe for helium, because the separations in energy of the angular momentum states with the same energy quantum number \( n \) lead to more radiative decay channels and the ionization of hydrogen by \( \lambda 584 \) photons provides a steady drainage of the \( 2p^1P \) population. Both features tend to dampen the rapid build-up of population into levels of higher \( n \) as the collisional excitations and line opacities increase with increasing density. From the results of the calculations, we find that under the physical conditions responsible for the \( \lambda 5876 \) emission, the \( 8n = 4 \) levels account for less than 20 per cent of the total helium ionization rate. Hopefully, the errors in helium line fluxes and line ratios are within 20 per cent.

### A3 O\( ^I \), Ca\( ^{II} \), Na\( ^I \)

We obtain most of the Einstein A rates of the pertinent \( O^I \) transitions from Hibbert et al. (1991), Biémont & Zeippen (1992) and Carlsson & Judge (1993). The \( \lambda 11287 \) Einstein A rate of \( 1.1 \times 10^7 \) s\(^{-1}\) is deduced from the experimental work of Christensen & Cunningham (1978). The \( 3s^3S \) level has also a radiative decay route to the \( 2p^1D \) state that is not indicated in Fig. 3. Its Einstein A rate of \( 1.83 \times 10^3 \) s\(^{-1}\) (Biémont & Zeippen 1992) is included as an additional decay rate of \( 3s^3S \) to the ground level. We gather the collisional rate coefficients for the various transitions from Barklem (2007), the radiative recombination rate coefficients from Péguynot, Petitjean & Boisson (1991), and the \( O^I - H^I \) and \( O^{II} - H^I \) charge-exchange reaction rates from Field & Steigman (1971).

We obtain the Einstein A rates and collisional rate coefficients for all relevant Ca\( ^{II} \) transitions from Burgess, Chidichimo & Tully (1995), the photoionization cross-sections from Verner et al. (1996) and Shine & Linsky (1974), the collisional ionization rate coefficients from Arnaud & Rothenflug (1985) and Shine & Linsky (1974), the total radiative recombination rate coefficient from Shull & Van Steenberg (1982), and the direct radiative recombination rate coefficients to the three levels from their photoionization cross-sections through detailed balance.

For the relevant atomic parameters of Na\( ^I \), we obtain the Einstein A rate from Sansonetti (2008), the \( 3s^3S \) and \( 3p^2P \) photoionization cross-sections from the experimental work of Hudson & Carter (1967) and Rothe (1969), respectively, the \( 3s^3S \) collisional ionization rate coefficient from Arnaud & Rothenflug (1985), the \( 3s^3S \rightarrow 3p^2P \) collisional excitation rate coefficient from Clark et al. (1982) and the radiative recombination rate coefficient from Verner & Ferland (1996).