

X-RAY EMISSION FROM THE WINDS OF HOT STARS¹

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ABSTRACT

A phenomenological theory is proposed for the structure of the unstable line-driven winds of early-type stars. These winds are conjectured to break up into a population of blobs that are being radiatively driven through, and confined by ram pressure of, an ambient gas that, because of shadowing by the blobs, is not itself being radiatively driven. With plausible choices for the theory's two free parameters, radiation from the bow shocks preceding the blobs can account for the X-ray luminosity of ζ Puppis. The theory breaks down, however, when used to model the much lower density wind of τ Scorpii, for then the blobs are destroyed by heat conduction from shocked gas. This effect appears to explain why the profiles of this star's UV resonance lines depart from classical P Cygni form.

Subject headings: line profiles — stars: early-type — stars: emission-line — stars: winds — X-rays: sources

1. INTRODUCTION

Harnden *et al.* (1979) have reported the discovery of X-ray emission from O and B stars and given observational arguments for believing it to originate in their expanding circumstellar envelopes. Accepting this conclusion, we interpret the X-ray emission as further evidence for nonradiative heating of these stars' radiatively driven winds. The existence of N v and O vi ions in the outflowing gas had previously been recognized as implying a source of mechanical heating, and this led to two strikingly different empirical models for the temperature stratification in these winds. In one model (Lucy and Solomon 1970; Lamers and Morton 1976), $T \approx 2 \times 10^5$ K is postulated throughout the flow in order to produce these ions by collisional ionization. In a later model (Hearn 1975; Cassinelli, Olson and Stalio 1978), $T \approx 5 \times 10^6$ K is postulated in a thin corona; the high-ionization stages in the cool wind are then produced by the Auger mechanism following K-shell absorptions of X-rays emitted by the corona.

Because the first of these models does not predict X-ray emission, while the second one does, the detection of X-rays would seem to be a decisive confirmation of the hot corona-cool wind model. However, the severe attenuation of soft X-rays ($E < 1$ keV) predicted for this model (Olson 1978; Cassinelli and Olson 1979) is contradicted by the observed X-ray spectra (Rosner and Vaiana 1979; Long and White 1980), and this failure cannot plausibly be removed by reducing mass-loss rates. For example, for ζ Puppis (O4f), for which Cassinelli and Olson (1979) predict ~ 30 optical depths of attenuation in the cool wind at $E \approx 0.5$ keV, and for which Long and White (1980)

report no deficiency of soft X-rays, a reduction of mass-loss rate to obtain $\tau \lesssim 3$ would seem to be firmly excluded by the agreement to within a factor of 3 of observational mass-loss estimates obtained with different techniques (Lamers and Morton 1976; Hutchings 1976; Barlow and Cohen 1977; Klein and Castor 1978; Abbott *et al.* 1980). We conclude therefore that the hot corona-cool wind model does not provide an acceptable description of the stratification of mechanically heated gas.

The above discussion suggests that the X-ray-emitting gas and the cool gas observed in UV spectra somehow coexist throughout these winds. A phenomenological model for which such a thermal structure is indeed expected will be developed in this paper. In obtaining this model, we suppose that the nonradiative heating is fundamentally a consequence of the instability that radiatively driven winds are subject to when *line* absorption is the mechanism of momentum transfer between radiation and matter. This instability was suggested by Lucy and Solomon (1970) and has been partially analyzed by Nelson and Hearn (1978), Martens (1979), and MacGregor, Hartmann, and Raymond (1979), all of whom speculate that the nonradiative heating of these winds is due to the dissipation mechanism that limits the growth of this instability. To investigate this speculation further, we now conjecture that the stellar wind, having been broken up by this rapidly growing instability (Nelson and Hearn 1978), has a final state that may be approximated by a simple two-component structure: a population of radiatively driven blobs ploughing through, and being confined by the ram pressure of, an ambient nonradiatively driven gas. We suppose this second component to be not radiatively driven because it is shadowed by the blobs—a beam of radiation

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from the star always interacts first with the blobs. Nevertheless, the drag force between the blobs and the ambient gas usually ensures that this gas also flows away from the star.

In this phenomenological model, we picture each blob as a coherent, long-living entity made up of a fixed parcel of gas. It is entirely conceivable, however, that disturbances, even at limiting amplitude, are always propagating relative to the gas; every parcel of gas would then have a checkered history of being radiatively driven at one moment and shadowed at the next and would thereby suffer many episodes of shock heating. The observable consequences of these alternatives are likely to be qualitatively similar, if not indistinguishable.

In the remainder of this paper, we give perhaps the simplest possible version of this two-component model, having the aim of rapidly bringing it to the point where X-ray luminosities and spectra can be calculated. Comparison with the data presently being accumulated with the *Einstein* Observatory should soon tell us whether further elaboration is worthwhile.

II. ASSUMPTIONS AND EQUATIONS

In this section, we derive the equations governing the structure and the X-ray emission characteristics of a simple version of our phenomenological model for the nonlaminar flow of line-driven winds. In doing so, we consider only that part of the flow that is sufficiently far downstream from the onset of instability for its further growth to have been halted by nonlinear effects. In the context of the two-component model, we conjecture that this is achieved when the radiative driving of the ambient gas, which causes the instability and therefore also blob production, has been rendered negligible as a result of shadowing by previously created blobs.

In addition to no further blob production, we will assume that blobs are not destroyed, that they do not shed mass, and that they all have the same mass, m . We also take the outflow to be isothermal, with the same temperature for both components.

a) Basic Equations

If g_R , g_D , and g denote the forces per gram due to radiation, drag, and gravity, respectively, then

$$v_b \frac{dv_b}{dr} = g_R - g_D - g \quad (1)$$

is the equation of motion for a blob. The corresponding equation for the nonradiatively driven ambient gas is

$$\bar{\rho}_a v_a \frac{dv_a}{dr} = -\frac{d\bar{P}_a}{dr} + \bar{\rho}_b g_D - \bar{\rho}_a g, \quad (2)$$

where $\bar{\rho}_a$ and $\bar{\rho}_b$ are the smoothed-out densities of the two components, and $\bar{P}_a = a_a^2 \bar{\rho}_a$ with $a_a^2 = kT_a/\mu m_H$.

With neither creation nor destruction of blobs and no mass shedding, the equations of continuity integrate to give

$$\Phi_{a,b} = 4\pi r^2 \bar{\rho}_{a,b} v_{a,b}, \quad (3)$$

where Φ_a and Φ_b are constants whose sum Φ is the star's mass-loss rate.

The filling factors for the two components are $f_a = \bar{\rho}_a/\rho_a$ and $f_b = \bar{\rho}_b/\langle\rho_b\rangle$, where $\langle\rho_b\rangle$ is the mean density within a blob. These satisfy the condition $f_a + f_b = 1$.

b) Properties of Blobs

We follow De Young and Axford's (1967) treatment of inertially confined plasma clouds. For an isothermal blob in mechanical equilibrium in its comoving frame, we have $P_b(z) = P_b(0) \exp(-z/h)$, where z measures distance along the axis of the blob from its apex and $h = \zeta a_b^2/g_D$, with $\zeta = 14/9$ being a correction factor allowing for the presence of sonic turbulence within the blob. (N.B. In taking the effective gravity g_D to be constant within the blob, we are implicitly making the dubious assumption that g_R is likewise constant.) To obtain the shape of the blob, we equate $\zeta P_b(z)$ to the confining ram pressure, for which we adopt the De Young-Axford approximation $\rho_a U^2 \cos^2 \xi$, where $U = v_b - v_a$ and ξ is the angle between the incident stream of ambient gas and the outward normal to the blob's surface; the resulting radius of cross-section is $r(z) = 2h \cos^{-1} [\exp(-z/2h)]$, which $\rightarrow \pi h$ as $z \rightarrow \infty$. With the shape known, the blob's mass can be readily calculated and is found to be $m = C_m \rho_b(0) h^3$, with $C_m = 9.22$.

Formally, the mean density within a De Young-Axford blob is zero, since its volume is infinite. Realistically, however, a blob's volume will be finite and must be estimated for the calculation of filling factors. Accordingly, taking account of the strong concentration of a blob's mass to its apex, we define its effective mean density to be $\langle\rho_b\rangle = m^{-1} \int_0^\infty \pi r^2 \rho_b^2 dz$, which gives $\langle\rho_b\rangle = C_\rho \rho_b(0)$ with $C_\rho = 0.290$. Correspondingly, we define a blob's effective volume to be $V = m/\langle\rho_b\rangle$ and its effective projected area to be $A = \pi \sigma^2 = (3\sqrt{\pi V/4})^{2/3}$. Now, from the formulae $g_D = \zeta a_b^2/h$, $a_b^2 = P_b(0)/\rho_b(0)$, and $\rho_b(0) = m/C_m h^2$, it follows that $mg_D = C_m \zeta P_b(0) h^2$. Substituting $\zeta P_b(0) = \rho_a U^2$ and eliminating h^2 in favor of A , we readily find that the drag force on a hypersonically moving blob is

$$mg_D = \frac{1}{2} C_D \rho_a U^2 A, \quad (4)$$

with $C_D = 1.519$.

Although the hypersonic limit (i.e., $U > a_a$) is appropriate when X-ray emission is occurring, it is not so near the star where we naturally wish to start the integrations. A simple heuristic procedure is therefore adopted to extend the analysis to slowly moving blobs. From the above formulae, we may show that $V = \zeta m a_b^2 / (C_\rho \rho_a U^2)$, which displays the confining effect

of ram pressure and which therefore suggests the generalization

$$V_{\dagger} = \frac{\zeta m a_b^2}{P_a + C_{\rho} \rho_a U^2}. \quad (5)$$

This modified formula is of course exact when $U = 0$.

The drag force is now obtained from equation (4) with A replaced by the correspondingly generalized projected area $A_{\dagger} = \pi \sigma_{\dagger}^2$; and the filling factors are derived from equations (3) and (5) by insisting that $V_{\dagger} = f_b m / \bar{\rho}_b$.

c) Radiative Transfer

From equations (1)–(3), one may readily show that the rate at which the pressure and radiation forces do work in a unit volume exceeds the rate at which the two-component medium gains mechanical energy in this same volume. The difference in these rates,

$$\dot{Q} = g_D \bar{\rho}_b U, \quad (6)$$

gives the rate at which energy is being dissipated per unit volume. Now, if the length scale for the radiative cooling of shocked ambient gas is small compared to the typical distance between blobs, this dissipated energy is radiated instead of being used to raise the mean temperature of the ambient gas. In this circumstance, ϵ , the effective thermal emissivity integrated over frequency, is given by

$$4\pi \bar{\rho} \epsilon = \dot{Q}, \quad (7)$$

where $\bar{\rho} = \bar{\rho}_a + \bar{\rho}_b$.

If T_* denotes the post-shock temperature, then gas at all temperatures $T < T_*$ contributes to ϵ , with the contribution at each temperature being, to a good approximation, inversely proportional to the total rate of radiative cooling at that temperature. Accordingly, we take the effective thermal emissivity of the two-component medium at frequency ν to be

$$\epsilon_{\nu} = \epsilon \int_0^{T_*} \frac{j_{\nu}}{j} \frac{dT}{T_*} \quad (8)$$

where j_{ν} is the conventional thermal emissivity and $j = \int_0^{\infty} j_{\nu} d\nu$. For the postshock temperature, we take $T_* = (9/64)(\mu m_H/k)U^2$, which corresponds, in the strong shock limit and with $\gamma = 5/3$, to the point on the bow shock where the angle of the incident flow is 30° . With this procedure, we are in effect approximating an integral over incident angle by a one-point quadrature formula. Higher accuracy is not warranted and would, in any case, require knowledge of the shape of the bow shock.

The effective absorption coefficient is given by

$$k_{\nu} \bar{\rho} = (k_{\nu})_b \bar{\rho}_b + (k_{\nu})_a \bar{\rho}_a. \quad (9)$$

This follows from the result that, for a line of length l , the summed lengths of the segments where blobs are being traversed has expectation value $f_b l$.

With the effective absorption and emission coefficients thus defined, the normal transfer equation now applies. At X-ray energies, in addition to the usual boundary condition of no incoming radiation at infinity, we also assume no outgoing radiation at the star's surface—i.e., no emission by the photosphere at these energies and no reflection of incoming X-rays.

In the calculation of ϵ_{ν} , we use values of j_{ν} generated by a computer code kindly supplied by J. C. Raymond (see Raymond and Smith 1977). In the calculation of k_{ν} , we use the results of Brown and Gould (1970) and Daltabuit and Cox (1972). For X-rays with $E \gtrsim 0.3$ keV, the dominant sources of absorption are K-shell ionizations of C and O, the cross-sections of which are insensitive to ionization conditions when only L-shell electrons are stripped, as in these winds. For softer X-rays, the dominant sources of absorption are L-shell ionizations, the cross-sections of which we estimate by reducing the measured values from Brown and Gould (1970) by the fraction of L-shell electrons missing in the dominant stage of ionization, viz., C III, N III, Ne IV, Si V. In addition, we assume that He is nearly completely ionized and therefore does not contribute significantly to the absorption of X-rays. Allen's (1973) cosmic abundances were used in these calculations.

In obtaining models in § III for individual stars, we adopt a velocity law and then derive the implied g_R instead of computing it from first principles. Nevertheless, the transfer of the radiation responsible for the radiative driving must still be investigated in order to impose our conjectured condition of effective shadowing of the ambient gas by the blobs. A measure of this shadowing is the factor, $e^{-\tau_1}$, by which the blobs attenuate the star's continuum flux. Now, if we suppose that

$$v_b = v_{\infty} [1 - (1 - \alpha)x - \alpha x^2]^{1/2}, \quad (10)$$

where $x = R/r$ and $-1 < \alpha < 1$, and that the blobs are optically thick for beams of radiation with impact parameters less than $\sigma_{\dagger}$ and with frequencies $\in (\nu_0 \pm \frac{1}{2}\Delta\nu)$, where $\Delta\nu = \nu_0 w/c$, then we find from equations (30) and (33) of Lucy (1971) that

$$\tau_1 = \frac{\sigma_{\dagger}^2 \Phi_b w}{m R v_{\infty}^2} \chi(x; \alpha), \quad (11)$$

where $\chi^{-1} = 2(1 - \alpha) + 2(1 + 2\alpha)x - 3(1 - \alpha)x^2 - 4\alpha x^3$. To be consistent with our earlier assumption of sonic turbulence within the blobs, we set $w/a_b = 4\sqrt{\gamma}$.

III. SOLUTIONS

We now use this phenomenological theory to obtain models of the winds of two well-observed early-type stars, ζ Puppis (O4f) and τ Scorpii (B0 V).

In these simple calculations, we take $v_b(x)$ to be given by equation (10); our system of equations

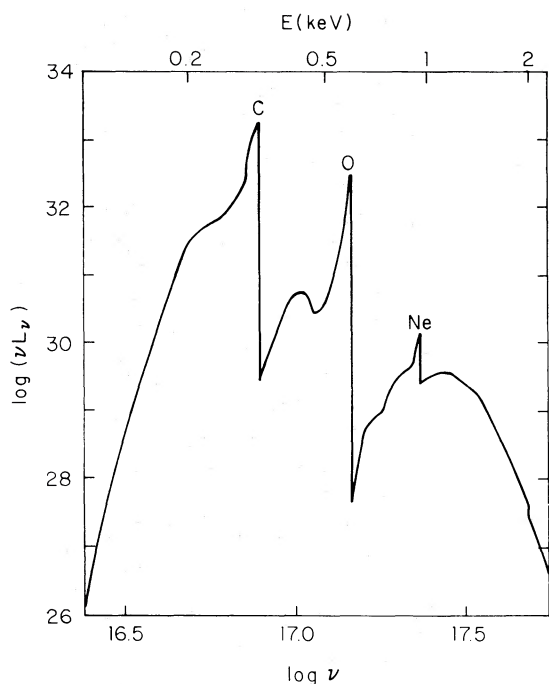


FIG. 1.—Theoretical X-ray spectrum for ζ Puppis. K-shell absorption edges of C, O, and Ne are indicated.

(§ IIa) then reduces to a first-order initial value problem, which we solve numerically starting at the point where $v_b = 150 \text{ km s}^{-1}$ with the assumption that $v_a = 100 \text{ km s}^{-1}$. Having obtained a solution, we can derive $g_R(x)$, the implied radiative acceleration of the blobs, from equation (1). In adopting this procedure, we are in effect accepting that radiative driving does indeed account for the observed mass-loss rates and terminal velocities of these and similar stars (Lucy and Solomon 1970; Castor, Abbott, and Klein 1976; Lamers and Morton 1976).

a) ζ Puppis

On the basis of previous analyses of observational data (Lamers and Morton 1976; Abbott 1978; Abbott *et al.* 1980), we adopt the following parameters for ζ Puppis: $M = 69 M_\odot$, $R = 17 R_\odot$, $v_\infty = 2660 \text{ km s}^{-1}$, and $\Phi = 3.5 \times 10^{-6} M_\odot \text{ yr}^{-1}$. We also set $T_a = T_b = 5 \times 10^4 \text{ K}$, $\mu = 0.62$, and $\alpha = 0$ (eq. [10]).

The parameters that are specific to our phenomenological model are m , the blob mass, and $\eta = \Phi_b/\Phi$, the fraction of the mass loss carried by the blobs. If we suppose that the onset of instability occurs in the neighborhood of the sonic point and that the blobs then have diameters comparable to the local scale height, we obtain the crude estimate that $m \approx 10^{15} \text{ g}$ for ζ Puppis. With m specified, we determine η by imposing the condition that $\tau_1 = \tau_m$ when $v_b = v_\infty/3$, which is roughly the velocity of maximum ram-pressure confinement and therefore minimum shadowing. (Note that $m \approx 10^{15} \text{ g}$ implies $\sim 10^9$ blobs in

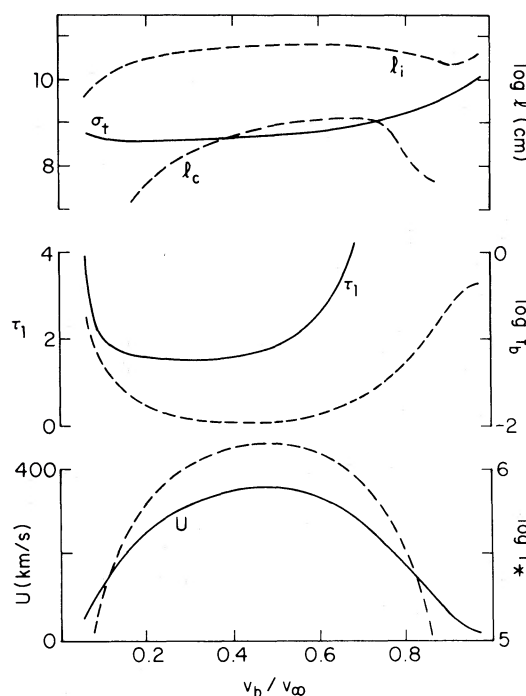


FIG. 2.—Properties of two-component wind model for ζ Puppis as functions of v_b/v_∞ . Upper diagram gives comparison of length scales defined in text. Middle diagram shows variation of τ_1 , the effective optical thickness of the blob component, and of f_b , its filling factor. Lower diagram shows variation of U , the relative velocity of the two components, and of T_* , the post-shock temperature.

the wind. Statistical fluctuations will therefore not give detectable time variations.)

Setting $m = 10^{15} \text{ g}$ and $\tau_m = 1.5$, we obtain a model with $\eta = 0.36$ and for which the 0.15–4.5 keV X-ray luminosity is $L_x = 2.1 \times 10^{32} \text{ ergs s}^{-1}$, in reasonable agreement with the observed value of $1.6 \times 10^{32} \text{ ergs s}^{-1}$ (Long and White 1980). The computed X-ray spectrum is shown in Figure 1, from which it will be seen that the emergent flux of soft X-rays ($E < 1 \text{ keV}$) is not negligible, as it is for the hot corona-cool wind model (Cassinelli and Olson 1979, Fig. 5); the two-component model is therefore not similarly in immediate contradiction with observation. A possibly critical test of models of this kind would be provided by the modest resolution X-ray spectra that could be obtained if the *Einstein* Observatory's imaging proportional counter (IPC) were used in conjunction with its $1000/\text{mm}^{-1}$ objective grating spectrometer (OGS). In particular, the pronounced K-shell absorption edges would be detectable.

Figure 2 illustrates various characteristics of this model. It shows, for example, that at its maximum $U = 363 \text{ km s}^{-1}$, corresponding to $T_* = 1.4 \times 10^6 \text{ K}$, and this occurs at $r/R = 1.31$, where $v_b = 1290 \text{ km s}^{-1}$. Moreover, at this point, $f_b = 1.1\%$, close to its minimum value.

The upper diagram in Figure 2 provides consistency

checks on some of our assumptions. The plotted quantities are: σ_t , a blob's characteristic radius of cross-section; l_c , the cooling length in a blob's frame for the shocked gas, which is taken to be $l_c = 0.035m_H^2 U^3 / \rho_a \Lambda(T_*)$, where Λ is the radiative cooling function; and $l_i = 1.33\sigma_t/f_b$, the typical distance in a blob's frame that a parcel of ambient gas travels between encounters with bow shocks. This plot reveals that throughout the flow $l_c \ll l_i$, so that dissipated energy is indeed radiated.

The plot of length scales also shows that $l_c > \sigma_t$ for $0.39 < v_b/v_\infty < 0.73$; consequently, in this part of the flow, a blob finds itself immersed in shock-heated gas at $T \approx 10^6$ K, and its survival is therefore in question. Taking $t_h = 1.75 < \rho_b > \sigma_t^2 a_b^2 / T_* K(T_*)$ with $K = 1.0 \times 10^{-6} T^{5/2}$ ergs cm⁻¹ s⁻¹ K⁻¹ (Spitzer 1962) as an estimate of a blob's heating time scale due to conduction and $t_c = 1.5a_b^2 m_H^2 / \langle \rho_b \rangle \Lambda(T_b)$ evaluated at $T_b = 2 \times 10^5$ K as an estimate of a blob's minimum radiative cooling time scale, we find that $t_c \ll t_h$ throughout the region where $l_c > \sigma_t$. For example, at the point where $v_b/v_\infty = 0.5$, $t_h = 1.7 \times 10^3$ s, but $t_c = 0.1$ s. Accordingly, we conclude that heat conduction from surrounding shocked gas is not able to raise the blobs' temperatures significantly; they therefore retain their ability to be radiatively driven.

The model's overall energy budget is of interest. From the derived function $g_R(x)$, we calculate that the rate of working by radiative forces is 1.01×10^{37} ergs s⁻¹. Of this input, 89% goes into the wind's flux of mechanical energy and the remaining 11% is dissipated into heat and then radiated. Of this radiated power, 10% or 1.10×10^{35} ergs s⁻¹ is emitted as 0.15–4.5 keV X-rays, but, because of absorption, only 0.19% of this emitted X-ray power actually emerges from the wind. The model's X-ray luminosity therefore represents only the fraction 2.1×10^{-5} of the rate of working by radiative forces or roughly 7.0×10^{-8} of the star's luminosity.

The sensitivity of L_x to the theory's parameters is of course of vital interest. Taking the above model as a standard of reference and noting that it has $\log L_x = 32.32$, we now briefly report the effects of parameter changes. Increasing m to 10^{16} g gives $\log L_x = 32.15$. Increasing τ_m to 2 gives $\eta = 0.39$ and $\log L_x = 31.77$. Decreasing α to -0.5 —i.e., slower initial acceleration—gives $\log L_x = 32.75$. Finally, decreasing Φ to $2 \times 10^{-6} \mathcal{M}_\odot \text{ yr}^{-1}$ gives $\log L_x = 32.74$. Thus L_x is gratifyingly insensitive to m but is sensitive to τ_m —i.e., to η —and also to Φ . Because of this sensitivity, our success in reproducing the X-ray power of ζ Puppis cannot be claimed as decisive support for the model.

The sensitivity of L_x to our approximations has also been tested. If the angle of incidence in the bow shock calculation (§IIc) is taken to be 45° , the standard model then has $\log L_x = 31.60$. (Our earlier choice of 30° was arrived at by matching to the result obtained by assuming the bow shock to have the shape of the De Young–Axford blob.) $\log L_x$ is not

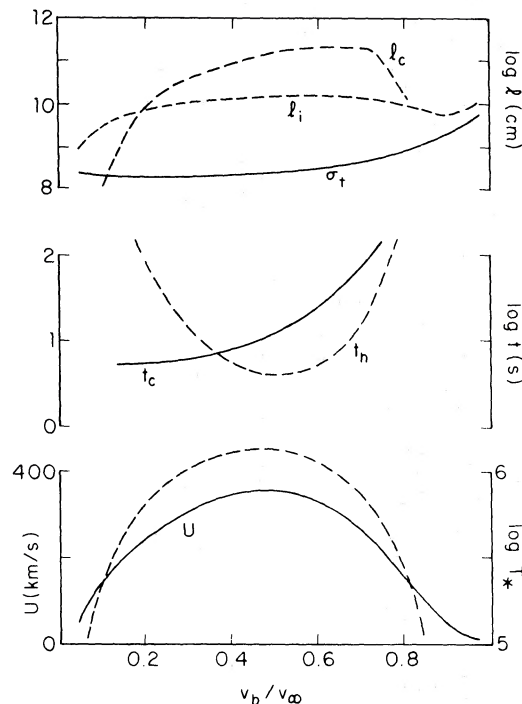


FIG. 3.—Properties of two-component wind model for τ Scorpii as functions of v_b/v_∞ . Upper diagram gives comparison of length scales defined in text. Middle diagram shows variation of heating and cooling time scales for blobs. Lower diagram shows variation of U , the relative velocity of the two components, and of T_* , the post-shock temperature.

sensitive to our choice for the velocities of the two components at the start of the integration.

b) τ Scorpii

On the basis of previous analyses of observational data (Lamers and Rogerson 1978; Abbott 1978), we adopt the following parameters for τ Scorpii: $\mathcal{M} = 23 \mathcal{M}_\odot$, $R = 6.7 R_\odot$, $v_\infty = 3440$ km s⁻¹ and $\Phi = 7.0 \times 10^{-9} \mathcal{M}_\odot \text{ yr}^{-1}$. We also set $T_a = T_b = 3 \times 10^4$ K, $\mu = 0.62$, and $\alpha = 0$. For the parameters specific to the theory, we take $m = 10^{12}$ g and $\tau_m = 1.5$.

Figure 3 illustrates various characteristics of the τ Scorpii model and, in particular, exposes the model's failure to satisfy the consistency checks discussed above for ζ Puppis. From these diagrams, we see that $l_c > l_i$ throughout the region where $T_* > 10^6$ K; our assumption that dissipation does not raise the mean temperature of the ambient gas is therefore contradicted. The diagram also shows that $l_c > \sigma_t$ and $t_h < t_c$ for $v_b > 1200$ km s⁻¹. Since t_h is then much shorter than the flow time scale ($\sim 10^3$ s), the temperature of the blobs rises rapidly, and they therefore lose the ability to absorb and scatter the star's UV radiation.

The above analysis suggests that the two-component model is applicable up to the point where

$v_b \approx 1200 \text{ km s}^{-1}$, corresponding to $r/R \approx 1.14$. Beyond this point, radiative driving ceases as the blobs, following the lead of the ambient gas, heat up to temperatures $\sim 10^6 \text{ K}$. Thereafter, the relative motions of the two components dissipate and the smoothed wind coasts out to infinity. This tentative picture of τ Scorpii's wind is relevant to the interpretation of the distinctive profiles of its UV resonance lines (Lamers and Rogerson 1978). Because the outflowing gas loses its ability to scatter UV radiation while still close to the star's surface, the absence of P Cygni-type emission components is explained simply as an occultation effect. Moreover, supposing there to be a distribution of blob masses, the loss of UV-scattering ions will occur gradually over a range of velocity, an effect that can quite conceivably explain the absence of a sharply defined violet edge to the absorption troughs. A theory freed from these consistency failures is of course required to establish this interpretation of τ Scorpii's UV line profiles.

Because of the model's breakdown, we do not venture to predict the X-ray luminosity and spectrum of τ Scorpii.

IV. SUMMARY AND DISCUSSION

In effect, we have conjectured that chaotic hypersonic motions relative to the mean outflow develop within the winds of O-type stars. The grounds for this conjecture are: (1) the instability of line-driven winds; (2) the powerful differential forces existing between shadowed and nonshadowed gas; (3) the rapidity with which shocked gas cools, thus recovering the prospect of being radiatively accelerated.

The particular phenomenological model adopted here for this chaotic hypersonic motion has provided the means of estimating its properties at limiting amplitude—in particular, the rate at which the chaotic

flow is then being dissipated and the fraction of this that appears as X-ray emission. Although this model, in common with other phenomenological theories, has free parameters, it was surely not obvious *a priori* that plausible choices for these parameters would lead to a computed X-ray luminosity in reasonable agreement with that observed for ζ Puppis. Moreover, further support for the conjectured chaotic hypersonic motions is provided by our finding that τ Scorpii's failure to exhibit P Cygni line profiles in its UV spectrum can be related to the failure of the chaotic motions to be maintained beyond a short distance from the star.

A decisive test of our conjecture will be provided by X-ray spectra capable of detecting the K-shell absorption edges shown in Figure 1. We might also note that the model predicts no X-ray emission from early-type stars whose winds have low terminal velocities—e.g., P Cygni itself.

Insofar as a phenomenological model of a gas dynamic flow is incorrect, it may be expected to be unstable to modes whose growth would bring about the correct flow. For the model considered here, the blobs are known to suffer from growing surface waves due to Kelvin-Helmholtz and Rayleigh-Taylor instabilities (Blake 1972). A possible elaboration of our model, therefore, would be to take account of the destruction of blobs by the nonlinear growth of these instabilities and the consequent production of new blobs as the shadowing of the ambient medium diminishes.

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