Winds from Hot Stars

S.P. Owocki

Bartol Research Institute, University of Delaware,
Newark, DE 19716, USA

1 Introduction

Stellar winds from hot stars (i.e., spectral types O, B, and WR) have been known and studied for about 20 years now. An excellent, succinct summary of both theory and observations was given recently by Abbott [1]. (This was one of Abbott's last active contributions before giving up astrophysical research, a decision which, given his many contributions to our understanding of hot stars, was most unfortunate for science.) More extensive discussions of both theory and observations are contained in the monograph edited by Conti and Underhill [2], and there have been numerous additional reviews of theory [3 – 8] or observations [9 – 11]. In nearly all of these the emphasis has been on steady-state models and time-averaged spectra. The general overview of Kudritzki in this volume emphasizes the now quite impressive sophistication of such models and their quantitative agreement with observations. Given this already rich collection of reviews, I will not attempt here to give any kind of exhaustive list of the literature. The next section gives my own general overview of the subject, but the bulk of the remaining discussion will concentrate on examining the physical driving mechanism – line-scattering in the wind of the star’s continuum radiation field – and what it implies for both the basic wind properties (e.g., mass loss rate, $\dot{M}$ and terminal flow speed, $v_\infty$) (Sect. 3), as well as for the wind structure (e.g., variability, clumpiness, embedded shocks, etc.) (Sect. 4). The emphasis on structure and variability admittedly reflects a personal bias, but it is also intended to complement the above-mentioned emphasis on steady-state properties in previous reviews. In particular, one goal will be to place in some perspective the successes of the steady-state approach.
2 Overview

Let us first consider what makes hot-star winds such an interesting subject for astrophysical research. First, there is the fact that, just as hot stars themselves are extreme in such basic properties as mass ($M \approx 10 - 50M_\odot$), temperature ($T \approx 10,000 - 50,000$ K), and luminosity ($L_* \approx 10^5 - 10^6 L_\odot$), so are their winds also extreme, with mass loss rates ($\dot{M} \approx 10^{-7} - 10^{-5}M_\odot$/yr) that exceed that of the sun by some nine orders of magnitude, and flow speeds ($v_\infty \approx 1000 - 3000$ km/s) of the order of 1% of the speed of light. Indeed, even though such stars only live $\sim 10^7$ years on the main sequence, the large mass loss rate means that a substantial fraction ($\sim 50\%$) of the star's initial mass can be lost during its main sequence lifetime, implying important effects on the star's evolution. Such winds contribute an important component to the mass and energy balance of the interstellar medium [12]. In high-mass X-ray binary systems they also provide the source material accreting onto the collapsed companion [13 - 15], a topic which, of course, provides another focus of this conference and of the accompanying reviews in this volume. Finally, as is emphasized by the general review of Kudritzki (also this volume), the detailed agreement between theory and observations makes it possible to use observations of lines formed in the wind to infer stellar properties like mass, luminosity, and temperature, and so this opens the way to using the luminous O and B stars as standard candles to infer distances in other galaxies.

In my own view, however, the most interesting aspect of hot-star winds is the opportunity to study in detail aspects of the relatively unexplored field of radiation hydrodynamics. Often this term is used in a relatively broad sense to refer to the quite common case when radiation plays an important role in the energy balance of a plasma (e.g., [16]); but here it applies in the stricter sense that the stellar radiation imparts momentum (as well as energy) to the plasma, and so drives a supersonic outflow. The former situation is more common because photons are, in a sense, better carriers of energy than momentum, since their momentum is given by dividing their energy by the largest speed possible, namely c. But the latter typically applies to the most luminous objects – those near their Eddington limit – which are also often the most energetic and hence most interesting objects in the universe, e.g., quasars, active galactic nuclei, and accretion disks. The relative proximity of hot stars makes it possible to obtain high spectral and/or temporal resolution observations, and these both guide and test a detailed wind theory that includes radiation-hydrodynamical processes likely to have relevance in the more-distant, and hence less well-observed, luminous objects.

In the case of hot-star winds the specific coupling mechanism between matter and radiation – namely line opacity – might at first seem surprising, since line transitions can only occur for a very narrow range of photon frequencies. However, what makes this mechanism in fact quite efficient here is that, once a flow is initiated, the line frequency becomes Doppler shifted, so that a fresh supply of relatively unattenuated flux from the star's continuum can then interact with the wind material in the spectral line. In this way a single line can
"sweep up" the momentum from a fraction \( \frac{v_{\infty}}{c} \sim 1\% \) of the star's spectrum, vs. only \( \sim \frac{v_{\text{th}}}{c} \sim 10^{-4} \) for the static case in which the scattering ion has a thermal velocity \( v_{\text{th}} \). Thus if there are \( \sim 100 \) strong lines spread across the spectrum, then nearly every photon will impart momentum to the wind, and so typically there is rough equivalence between the wind and radiative momentum loss, i.e.,

\[ \dot{M} v_{\infty} \sim L_*/c. \]

Historically, the basic idea of line-driving dates back to early work by Milne [17] and the pioneering work by Sobolev [18, 19]; but the specific application to hot-star winds was first proposed by Lucy and Solomon [20] shortly after the original rocket-based observations by Morton [21] of asymmetric "P-Cygni" profiles [22] in UV lines had indicated the existence of continuous mass outflows from such hot, luminous stars. The original Lucy and Solomon model included only a few lines and so yielded mass loss rates significantly smaller (\( \sim 1\% \)) than had been inferred observationally. The landmark work of Castor, Abbott, and Klein (hereafter CAK) [23] developed a simple but powerful formalism, still in use in models today, for including the ensemble effect of many lines, and thereby obtained good general agreement with the observationally inferred \( \dot{M} \).

The original CAK model had assumed radial streaming of radiation, as from a point source, and correct inclusion of the angular size of the stellar disk was found independently by Friend and Abbott [24] and by Pauldrach, Puls, and Kudritzki (hereafter PPK) [25] to yield substantially improved agreement with the observationally inferred terminal flow speed \( v_{\infty} \). A more fundamental approximation of CAK was to use Sobolev theory to compute the line force [18, 19, 26, 27], but this too has since been relaxed with wind models that compute the radiative force from full comoving frame solutions for the line transfer [28, 25]. PPK showed that the CAK/Sobolev models, modified to include the finite disk effects, yield wind parameters nearly identical to the full comoving frame results. Recent extensions of the comoving frame models (see, e.g., review by Kudritzki and Hummer [8], and references therein) now treat unified wind/atmosphere models with self-consistent solution of level populations in thousands of transitions, and the theoretical line profiles generated thereby are in impressive quantitative agreement with observed line spectra.

Despite this great success of steady-state, line-driven-wind models in quantitatively matching many time-averaged spectral observations, there are strong reasons, both observational and theoretical, for believing that hot-star winds are not at all smooth, time steady flows, but rather are both spatially and temporally highly variable. The observational evidence for this will be discussed in greater detail below (Sect. 4.1), but listed briefly, it includes Superionization [29], X-ray emission [30, 31], Nonthermal radio emission [32], Black profiles in saturated UV lines [33], and Discrete absorption components in unsaturated UV lines [34 - 36]. The theoretical evidence stems from the expectation, which dates back to Lucy and Solomon [20] (and even to Milne [17]), that such a line-driven flow would be dynamically unstable. A small-scale increase in radial flow speed Doppler-shifts the local line-frequency out of the absorption shadow of intervening material, leading to an increased line-force which then tends to further increase the flow speed. Formal stability analyses [37 - 40]

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(See review by Rybicki [41] and Sect. 4.2 below.) yield a growth rate nearly a hundred times the wind expansion rate. This implies that even very small base perturbations should quickly become nonlinear, perhaps giving rise to a wind embedded with multiple shocks and dense clumps [42, 43] (Sect. 4.3). Current efforts carry out numerical hydrodynamics simulations of the nonlinear wind structure that results from applied perturbations [44] (Sect. 4.4), or that arises intrinsically (Sect. 4.5). A important question is how much this structure affects the mean wind properties inferred by steady-state models (Sect. 4.6). The simulations carried out so far are intriguing, but are still based on several simplifying assumptions (e.g., pure absorption, isothermal flow, spherical symmetry) that will have to be relaxed in future work (Sect. 5).

3 Physics of Line Driving

3.1 The Radiative Force from a Single Line

The radiative acceleration $g_{\text{rad}}$ (i.e., force-per-unit mass) resulting from the interaction of a radiative flux $F_\nu$ at a photon frequency $\nu$ with material of opacity $\kappa_\nu$ is given quite generally by

$$g_{\text{rad}} = \int_0^\infty d\nu \, \kappa_\nu \frac{F_\nu}{c}. \quad (1)$$

A simple but useful limiting case is that of an optically thin line, for which one may neglect the attenuation of an externally imposed flux, so that (1) becomes

$$g_{\text{thin}} = \frac{F_\nu}{c} \Delta \nu \kappa_L$$
$$= \frac{\nu_L L_\nu}{L_*} \frac{v_{\text{th}} \kappa_L}{4\pi r^2 c^2}, \quad (2)$$

where $L_\nu$ is the source luminosity spectrum, $\nu_L$ is the line frequency, and $\kappa_L$ is the total line opacity; $\Delta \nu = \nu_L v_{\text{th}} / c$ is the line width, assumed here to be set by Doppler broadening from thermal motion at speed $v_{\text{th}}$. The latter expression applies at a distance $r$ from a spherically symmetric radiative source with total luminosity $L_*$, and emphasizes that, just like gravity, the optically thin force falls off with an inverse square law. The force from an optically thin continuum opacity $\kappa_c$ (e.g., electron scattering) also follows this scaling and so is often accounted for by simply considering an effective gravity that is reduced by a factor $1 - \Gamma_c$, where $\Gamma_c \equiv \kappa_c L_* / 4\pi G M_\ast c$.

For an optically thick opacity source, the simplicity of equation (1) is actually somewhat misleading, since the flux $F_\nu$ depends sensitively on the opacity $\kappa_\nu$, and in general must be determined from solution of the full nonlocal radiative transport problem. For the winds from O and B stars, the optically thick
Line Scattering in an Expanding Wind

Fig. 1. Radial variation of photon wavelength as viewed in the local comoving frame of material in the wind. The shaded region represents a thermally broadened spectral line with which the outward propagating radiation interacts.

Opacity is from numerous (~100) strong scattering lines. (For Wolf-Rayet stars the continuum can also be optically thick in much of the wind; see below.) Fortunately, for this case of line transfer in a supersonically expanding flow, methods first developed by Sobolev [18, 19] make it possible to obtain a good approximation to the force that depends only on local quantities.

The basic ideas of this Sobolev approximation are illustrated graphically in Fig. 1, which plots the radial variation of a photon’s wavelength as viewed in a frame comoving with the mean flow in the expanding wind. The shaded band represents a line that is thermally broadened by a Doppler width \( \Delta \lambda_D \) centered about the rest wavelength \( \lambda_1 \). A photon emitted in the stellar continuum at a wavelength \( \lambda < \lambda_1 \) propagates freely until red-shifted into the blue edge of the line, whereupon it is scattered. In addition to changing direction, the photon will in general be redistributed in frequency, most likely into the line core, but for purposes of this illustration, let us assume the scattering is coherent. The opacity on the blue edge is relatively weak, and so the scattered photon then travels some mean free path on the order of the Sobolev length \( L \equiv v_{th}/(dv/dr) \), which in this context may be viewed as a local “opacity scale height”.

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A crucial point is that, because the wind is an expanding medium, distant material viewed in any direction from a given location always appears to be receding. Hence, whether a photon is scattered into the forward or backward directions, the subsequent distance traveled always causes its comoving wavelength to be red-shifted. After many scatterings this systematic drift toward the red brings the photon to the line’s red edge, from which it’s last scattering can, with roughly equal probability, be either into a fore or aft direction. Since the resulting diffuse radiation field thus has rough fore-aft symmetry, the net force associated with this diffuse radiation is small, i.e., $g_{\text{diffuse}} \approx 0$. Thus the overall force associated with such line scattering is very nearly equal to what it would be for pure line absorption.

For simplicity let us assume (as did CAK) a point source of radially streaming radiation. The radiative force associated with an isolated line can then be written as

$$g_{\text{rad}}(r) \approx g_{\text{abs}}(r) = g_{\text{thin}} \int_{-\infty}^{\infty} dx \, \phi(x - v(r)/v_{th}) \, e^{-t(x,r)},$$

(3)

where $\phi$ is the line-profile function and $x \equiv (\nu - \nu_L)/\Delta \nu_D$ is the frequency displacement from line center, measured in Doppler-units in the star’s rest frame. The optical depth $t(x,r)$ at radius $r$ is given by

$$t(x,r) = \int_{R_*}^{r} \kappa_L \rho(r') \phi(x - v(r')/v_{th}) dr'.$$

(4)

A crucial element of the Sobolev approximation is that the variation in the integrand of Eq. (4) is dominated by the Doppler-shift associated with changes in the velocity $v(r')$; by switching the variable of integration to the comoving frame frequency $x' \equiv x - v(r')/v_{th}$, we thus find

$$t(x,r) \approx \tau \int_{x}^{\infty} \phi(x') dx',$$

(5)

where the Sobolev optical depth

$$\tau \equiv \frac{\kappa_L \rho v_{th}}{dv/dr}$$

(6)

represents a collection of spatial variables that are assumed to be approximately constant over a Sobolev length $L \equiv v_{th}/(dv/dr)$. For example, in the case of a smooth, steady flow this means that the density scale length

$$H \equiv \frac{\rho}{d\rho/dr} \approx \frac{v}{dv/dr} \gg \frac{v_{th}}{dv/dr} \equiv L,$$

(7)

and hence that $v \gg v_{th}$. Since the ion thermal speed is on the order of the sound speed $a$, we can expect that the approximation will be well satisfied in the supersonic portions of a smooth flow. Applying Eq. (5) to Eq. (3) we see that both required frequency integrations can now be done analytically, yielding

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\[ g_{\text{rad}}(r) = g_{\text{thin}} \frac{1 - e^{-\tau}}{\tau}. \] (8)

Note that for a weak line \( \tau \ll 1 \) we recover the optically thin expression (2), while for a strong line \( \tau \gg 1 \) we obtain the optically thick form

\[ g_{\text{rad}}(r) = g_{\text{thin}} \frac{\nu_{\lambda} L_{\nu}}{L_*} \frac{L_*}{4\pi \rho r^2 c^2} \frac{dv}{dr}. \] (9)

The latter is independent of the opacity \( \kappa_{\lambda} \), depending instead on the velocity gradient \( dv/dr \). This can be understood physically from Fig. 1. Once a line is optically thick, all photons shifted into resonance with it are scattered, and further increasing the opacity only pushes the first scattering further into the blue wing. The number of photons scattered per unit length is thus independent of the opacity, and depends instead on the slope of the photon red-shift, which depends on the velocity gradient. Normally, such a velocity gradient is part of an inertial term that represents the acceleration that results from a driving force, but here we see that it helps determine the driving force. In general terms, just as the optically thin line force helps reduce the effect of gravity on the fluid, so does the optically thick line force reduce the effect of inertia.

A simple consideration of just the inertial force balance enables us to estimate the mass loss rate that results from such line-driving. The active acceleration resulting from such driving is just

\[ \nu \frac{dv}{dr} = g_{\text{rad}} = \frac{\nu_{\lambda} L_{\nu}}{L_*} \frac{L_*}{4\pi \rho r^2 c^2} \frac{dv}{dr}, \] (10)

which implies that the mass loss rate \( \dot{M}_l \) associated with driving by a single, isolated line is

\[ \dot{M}_l = 4\pi \rho vr^2 = \frac{\nu_{\lambda} L_{\nu}}{L_*} \frac{L_*}{c^2}. \] (11)

If the line has a frequency near the peak of the stellar luminosity, the factor \( \nu_{\lambda} L_{\nu}/L_* \) is of order unity, implying that the wind mass loss rate driven by a single line is only about the photon mass loss rate! It is essentially for this reason that the original Lucy and Solomon [20] model, which considered the effect of driving in only a few strong lines, yielded such low mass loss rates.

### 3.2 The CAK Model: Driving by an Ensemble of Nonoverlapping Lines

A major advance of the CAK [23] model was to develop a formalism for efficiently including the cumulative effect of a large ensemble of lines, which they effectively assumed to have a flux-weighted number distribution that was a power law in opacity, i.e., \( N(\kappa) \sim \kappa^{\alpha-2} \), where \( 0 < \alpha < 1 \). Assuming the lines do not significantly overlap, the cumulative force can then be computed by integrating the expression (3) over this number distribution [44, 75]. Applying, as before, the Sobolev approximation, one obtains the CAK line-force
\[ g_{\text{CAK}} = \frac{KL_*}{r^2} \left( \frac{1}{\rho} \frac{dv}{dr} \right)^\alpha, \]  

(12)

where \( K \) is a normalization constant for the line-distribution (related to the CAK constant \( k \) by \( K = k(\kappa_e^{1-\alpha}/4\pi c v_\text{th}^\alpha) \), where \( \kappa_e \) is the electron scattering opacity.) Since \( 0 < \alpha < 1 \), this force can be viewed as kind of a “geometric-mean” mixture of the optically thin and thick expressions (2) and (9). But note that the presence of the density in this expression means that whether any given line in the distribution is optically thick or thin is not determined \textit{a priori}, but rather results from a self-consistent solution of the steady-state equations of motion including this force.

CAK derived their wind solutions by applying this expression in both the subsonic and supersonic regions of a steady wind with a finite temperature. However, strictly speaking, this Sobolev limit can only be attained for real driving ions in the limit of vanishing temperature, implying also a vanishing sound speed \( a \). Furthermore, since such gas pressure forces play a relatively minor role in driving the wind, we can derive the basic wind characteristics by considering the balance among inertial, gravitational, and radiative forces in this limit \( a \to 0 \) of negligible gas pressure terms. Applying Eq. (12) for the radiative force, the requirement of momentum balance reduces to the condition,

\[ F \equiv w' + 1 - C w'^\alpha = 0, \]  

(13)

where \( w' \equiv r^2 uv'/GM_* \) and

\[ C \equiv \frac{K}{GM_*} \left( \frac{4\pi GM_*}{Mc} \right)^\alpha. \]  

(14)

Fig. 2 illustrates that, depending on the value of \( C \), equation (13) has either 0, 1, or 2 solutions. The critical case with one solution has the property that \( F = \frac{\partial F}{\partial w'} = 0 \), which implies the critical values \( w'_c = \alpha/(1 - \alpha) \) and \( C_c = \alpha^{-\alpha}(1 - \alpha)^{1-\alpha} \). Since Eq. (13) has no explicit spatial dependence, this critical value \( w'_c \) must hold throughout the flow. Integrating \( w'_c \) from the stellar radius \( R_* \), we thereby obtain the usual “CAK velocity law”,

\[ v_{\text{CAK}}(r) = v_\infty \sqrt{1 - \frac{R_*}{r}}, \]  

(15)

where the asymptotic wind speed is given by

\[ v_\infty \equiv \sqrt{\frac{2\alpha GM_*}{(1 - \alpha)r_o}} = \sqrt{\frac{\alpha}{1 - \alpha}} v_\text{esc}. \]  

(16)

Likewise, if we solve the definition (14) for the mass loss rate and apply the critical value \( C_c \), we obtain the CAK mass loss rate formula,

\[ \dot{M}_{\text{CAK}} = 4\pi \alpha \left( \frac{GM_*}{1 - \alpha} \right)^{1 - 1/\alpha} (KL_*)^{1/\alpha}. \]  

(17)

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Zero-Sound-Speed Wind Solutions

Fig. 2. Graphical solution of a CAK/Sobolev line-driven wind model in the limit of zero sound speed. The CAK mass loss rate represents the critical case with one solution; for \( \dot{M} > M_{CAK} \) there are no solutions, but for \( \dot{M} < M_{CAK} \) there are two solutions.

These are just the relations derived by CAK for the case of a point-star; taking account of the finite angular size of the stellar disk tends to make the mass loss rate somewhat lower and the velocity law somewhat shallower [24, 25].

It is important to realize that, unlike in the finite-sound-speed model of CAK, the mass loss rate in this zero-sound-speed case is not unique, since solutions with \( \dot{M} < M_{CAK} \) are now also allowed. Indeed, for these “subcritical” \( \dot{M} \), there are two solutions, one steep, and one shallow. In the zero-sound-speed case, the CAK mass loss rate \( M_{CAK} \) thus now only represents a maximum mass loss rate for which solutions exist. As we discuss in greater detail in Sect. 4.5, a similar kind of solution “degeneracy” exists for finite-gas-pressure models that do not use the Sobolev approximation to compute the line-force.

3.3 Multi-Line Scattering in Overlapping Lines

This CAK mass loss rate is typically \(~100~\) times the single line value, implying that there are effectively about 100 optically thick lines spread throughout the spectrum. On the other hand, note that the terminal flow speed \( v_\infty \) roughly
scales with surface speed as \( v_\infty \sim v_{\text{esc}} \) and is pretty much independent of the number of strong lines. Typically this is on the order of 1% of the speed of light, so that, quite coincidentally, in the CAK model

\[
\dot{M} v_\infty \approx \frac{L_*}{c} \left( N_{\text{thick}} \frac{v_\infty}{c} \right) \approx \frac{L_*}{c},
\]

which is often referred to as the “single scattering limit”. This means that, fortuitously, the CAK model is just barely self-consistent in ignoring line overlap, if one assumes the lines are smoothly spread throughout the stellar spectrum.

In reality, of course, there is always some line overlap, arising both from chance frequency coincidences of independent lines as well as from the tendency for lines to occur in multiplets. An early but quite elegant examination of the dynamical effects of such overlap was carried out by Friend and Castor [45], who derived an analytic extension of the CAK model to include overlap under the assumption that the line frequency spacing is Poisson distributed. Although this ignored the tendency of lines to occur in multiplets, it demonstrated quite clearly the major effect that overlapping lines can drive mass loss in excess of the single scattering limit, with \( \dot{M} v_\infty \approx (2 - 5) L_*/c \). Similar results were obtained for more realistic line lists by Abbott and Lucy [46] from a Monte-Carlo calculation, and by Puls [47], who carried out the quite formidable calculation of solving self-consistently the coupled system of statistical equilibrium equations.

Fig. 3 illustrates graphically how multi-line scattering can result in the emergent stellar radiation imparting its radial momentum more than once to the wind. As in Fig. 1, this figure shows the radial evolution of a photon’s comoving frame wavelength, but now including interaction with two scattering lines. The initially outward-propagating photons are first scattered in the bluer line, being reemitted with equal probability into the fore and aft directions and thus on the average imparting all their initial radial momentum to the wind. However, during their subsequent free flight through the spherical envelope, the component of the photon’s momentum along the local radial direction systematically increases. Thus, if they don’t strike the stellar core, the photons are eventually red-shifted into resonance with the redder line, but now again with an average net positive radial momentum, which is thus also imparted to the wind. The exact amount of extra momentum imparted depends, of course, on the line spacing and the wind velocity law.

For O and B stars such overlap effects are typically minor, but they probably play a much more fundamental role in the Wolf-Rayet stars, which are inferred observationally to have \( \dot{M} v_\infty / (L_*/c) \approx 5 - 50 \) [48]. Indeed the density in such winds is so high that the continuum remains optically thick well past the sonic point, and so models based on the usual “core-halo” treatment of the emergent continuum, which have proven so successful for O and B stars, completely fail for Wolf-Rayets. Instead, a suitable model must treat the diffuse transport of both line and continuum radiation within the expanding envelope, taking proper account of the substantial line overlap. This is sufficiently difficult that

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Multi-Line Scattering in an Expanding Wind

Fig. 3. Same as Fig. 1, except now with two spectral lines.

so far there is still no self-consistent, dynamical, Wolf-Rayet star wind model that can adequately explain the high mass loss rates inferred observationally.

3.4 Further Extensions of the Basic CAK Model

Finally, over the past decade there have been several other extensions of the original CAK wind model. These have largely been directed along three general lines: (1) Multi-dimensional effects, (2) Improved transfer theory, and (3) Instability and time-dependence.

Multi-dimensional effects have included rotation and/or magnetic fields, treated within the steady-state, CAK/Sobolev wind formalism. For typical O and B stars, MacGregor, Friend, and Gilliland [49] were able to place an upper limit $B_o < 100$ G on the open magnetic field near the equator from the apparent lack of rotational spin-down. Friend and MacGregor [50] found relatively minor effects of rotation on the basic wind parameters $\dot{M}$ and $v_\infty$ unless the rotation is at a substantial fraction (≥ 70%) of breakup. Models for such rapidly rotating stars show a substantially different wind from pole and equator, and these have been applied to Be stars [51, 52] and (more speculatively) to Wolf-Rayets [53].
Improvements in radiative transfer include the above-mentioned (Sect. 2) PPK comoving frame models and Pauldrach's [54] treatment of the NLTE statistical equilibrium. Recent further efforts (e.g., [55]) have focussed on developing unified wind-atmosphere models that abandon the artificial distinction between "core" and "halo". A major overall goal of these efforts is to carry out "quantitative spectroscopy" [8], i.e., to compare quantitatively theoretical and observed spectra in order to infer more precisely and reliably the basic parameters of both wind and star. Further details on these efforts are given by the contributions of Kudritzki, Pauldrach, and Puls in this volume.

The next section will examine in some detail the last item, the question of the wind instability and variability.

4 Wind Instability and Variability

4.1 Observational Evidence for Wind Structure

Let us now review the observational evidence that hot-star winds are highly structured and variable. As noted briefly in Sect. 2, this includes:

1. Superionization, i.e., the presence in the wind of relatively high ionization stages like OVI, NV, and CIV not normally expected in material near the stellar effective temperature. Originally it was argued [29] that such ions arose from Auger ionization from X-rays, but recently Pauldrach [54] has argued that such stages can indeed be produced near the stellar effective temperature if one correctly computes the stellar UV photoionization from excited levels.

2. X-ray Emission. Nonetheless, the X-ray emission predicted by the Auger model was in fact subsequently observed by the Einstein Observatory. Spectra taken with the Solid State Spectrometer (SSS) typically show both a soft and hard component, corresponding respectively to temperatures of $\sim 3 \times 10^6$ K and $\sim 10^7$ K [31]. The original picture of a hot, geometrically thin corona at the wind base predicted the existence of K-edge absorption features that were not observed in the limited number of SSS spectra taken. This is generally viewed as implying that the X-rays must originate well above the wind base, e.g., from shocks embedded in the wind, although Waldron [56] argues that the coronal model is still viable if one self-consistently includes the effect of X-ray ionization in reducing the strength of the absorption edges. However, for O and B stars the observed X-ray luminosity tends to scale in proportion to the star's bolometric luminosity as $L_x \approx 10^{-7}L_*$, which is quite distinct from the scaling with rotation typical of late-type stellar coronae, and which implies an originating mechanism tied, as is the wind, to the star's radiative output. Chlebowksi [57] notes that the deviation of $L_x$ from this general relation correlates well with stellar environment, and he thereby infers that some X-rays may originate in the very distant wind, i.e., at the wind terminal shock.
3. Nonthermal Radio Emission. Of the 25 or so early type stars in a volume limited sample of radio detections with the VLA, about a quarter show a spectral variation that is distinctly nonthermal [58]. White [59] has proposed that this arises from gyrosynchrotron emission from nonthermal, high-energy particles that have undergone Fermi acceleration in embedded wind shocks arising from the wind instability. (An interesting, recent variation of this idea supposes that the high-energy tail of the observed X-ray spectrum results from inverse Compton scattering of the stellar UV radiation by the same high-energy particles. [60])

4. Black Profiles. Lucy [33] first noted that the observed blackness of the blue-shifted absorption troughs of saturated UV lines are a signature of a nonmonotonic velocity field. In a smooth outflow such blue-shifted absorption troughs should always contain a residual flux from radiation forward scattered toward the observer by the wind in the forward hemisphere. Lucy reasoned that this implied that the expanding medium had to predominantly backscatter radiation, and he showed how this was a natural consequence of the systematic reduction, roughly by a factor of two, of the forward scattered radiation at each one of multiple resonance layers that can line-scatter radiation in a nonmonotonic velocity outflow.

5. Discrete absorption components. These may be the most direct observational manifestations of wind structure and variability. They consist of relatively narrow absorption features superposed upon the absorption trough of the broad, P-Cygni profile of an unsaturated UV line. Spectra taken months or even years apart sometime show such a discrete component at the same velocity [34], but IUE monitoring observations [35,36] taken about once an hour over several days show that the components typically form near a velocity \( v \approx v_\infty/2 \) and then narrow and shift to higher velocities over several days. Interestingly, the final velocity approached is typically smaller by 10-20% than the edge velocity defined in saturated lines, i.e., \( v_\infty \approx 0.8 - 0.9 v_{\text{edge}} \) [61, 62], and the recurrence and acceleration time-scales seem to correlate roughly with the stellar rotation period [63].

4.2 Linear Theory of Line-Driven Instability

As noted in Sect. 2, it was recognized quite early on [13,16] that the sensitive velocity dependence of the line-force might cause a line-driven flow to be dynamically unstable. However, the formal analysis of even the linear stability properties turned out to be quite subtle, and for a while the strength of the instability, and even whether it really existed, was not clear. The first formal perturbation analyses indicating instability [37,38] were based on the assumption that perturbations were optically thin, meaning that one could ignore any change in the optical depth, and instead focus only on changes arising from the Doppler-shift associated with the perturbed velocity. Referring to the line-force expression (3), one can see that applying such a velocity perturbation
and expanding to first order leads to a perturbed force that is proportional to
the perturbed velocity, \( \delta g_{\text{abs}} \propto \delta v \). The net work done by such a force thus
amplifies the perturbation, making it \textit{unstable}.

However, in a study aimed at clarifying the nature of wave propagation
through a CAK wind model, Abbott [64] took the quite different approach of
simply assuming that the perturbed force would scale in the same way as the
mean force, namely with the velocity \textit{gradient}, i.e., \( \delta g_{\text{abs}} \propto \delta v' \). For the usual
assumption of sinusoidal perturbations, this leads to a 90° phase difference
between \( \delta g_{\text{abs}} \) and \( \delta v \), with the consequence that \textit{no net work} is done by the
line-force, implying now that the perturbation is \textit{stable}.

Owocki and Rybicki [39] reconciled these apparently contradictory results
by determining the effect of a perturbation on the absorption-line-force (3)
without making either of these assumptions. They thereby derived a simple
"bridging" law,

\[
\frac{\delta g_{\text{abs}}}{\delta v} = \Omega_b \frac{ik}{\chi_b + ik},
\]

(19)

where the blue-edge line strength \( \chi_b \) determines a bridging length, given
roughly by half the Sobolev length, i.e., \( \chi_b^{-1} \approx L/2 \). This length bridges the sta-
ble, long-wavelength \( (k \ll \chi_b) \), Sobolev regime, for which \( \delta g_{\text{abs}} \propto ik\delta v \propto \delta v' \),
and the unstable, short-wavelength \( (k \gg \chi_b) \), optically-thin-perturbation
regime, for which \( \delta g_{\text{abs}} \propto \delta v \). The instability growth rate is \( \Omega_b \approx 2g_{\text{abs}}/v_{\text{th}} \),
\[\text{i.e., roughly twice the rate at which the mean flow is accelerated by a thermal}
\]
speed. Since \( g_{\text{abs}} \approx v dv/dr \), we see also that \( \Omega_b \approx 2v/L \), \[\text{i.e., the instability}
\]
grows at twice the flow speed. Following this through a Sobolev length. Comparison with the
wind expansion rate through a typical scale length \( H = \rho/(dp/dr) \approx v/(dv/dr) \)
shows that \( \Omega_b/\Omega_{\text{exp}} \approx 2H/L = 2v/v_{\text{th}} \), which can approach 100 in these highly
supersonic winds. An infinitesimal (i.e., \textit{linear}) perturbation advected outward
through the wind is thus predicted to have a cumulative growth on the order of \( e^{v/v_{\text{th}}} \approx 100 \)!
In practice this means, of course, that any small, but finite, amplitude perturbation would quickly become nonlinear in the wind. Efforts to
determine the likely form of this nonlinear structure are discussed below (Sects.
4.3-4.5).

Although the wind driving is primarily by line-scattering, the above per-
turbation analysis assumes that, just as in the mean flow (See Sect. 3.), one
may ignore the dynamical effect of the diffuse radiation field, and treat only
the direct \textit{absorption} component of the perturbed line force. However, Lucy
[65] showed that, unlike the fore-aft symmetry seen in the comoving frame of
the mean flow, the diffuse field seen by a small-scale velocity perturbation \( \delta v \)
is stronger by a factor \( \delta v/v_{\text{th}} \) against the direction of the perturbed velocity.
This diffuse field thus exerts a net \textit{drag} force on such perturbations that is on
the same order as the perturbed direct force that gives rise to the instability. In
fact, in Lucy's coherent scattering analysis, this drag effect \textit{exactly cancelled} the
instability near the wind base, where the flow is nearly plane parallel. Under
the more realistic assumption of complete redistribution, Owocki and Rybicki
[40] found complete cancellation occurs only when the flow is plane parallel
\textit{and} a non-limb-darkened stellar core fills an entire hemisphere. Away from the

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wind base, the combined effects of spherical expansion and the shrinking of the solid angle subtended by the stellar core quickly reduce the relative importance of the drag term. At $r = 1.5 R_\ast$, for example, the net growth rate is about half that obtained from a pure-absorption analysis, but this still implies a strong instability, with cumulative growth now about 50 (vs. 100) e-folds. This strong instability is also found not to be substantially reduced by any of several other potential stabilizing effects, e.g., line-overlap [40], radiative cooling, or thermal conduction [39].

In addition to reducing somewhat the growth rate of radial velocity perturbations, this line-drag effect tends to strongly dampen horizontal velocity perturbations (for which the countervailing direct term is much weaker) [66]. Hence any perturbations that are initially unpolarized at the wind base will become predominantly radially polarized not far into the wind. In this regard the wind and its variations can thus be reasonably treated within a simple 1-D, spherically symmetric model, although in general there will still be horizontal phase variations that will require a 2-D or 3-D treatment.

4.3 Heuristic Models of a Wind with Embedded Shocks

Given the supersonic nature of the wind outflow and the large instability growth rate, it seems inevitable that, whatever its detailed form, the nonlinear wind structure arising from this instability will include shocks. (A useful review of the properties of such wind shocks has recently been given by Castor [67]). Lucy [33] has proposed a heuristic model that assumes that this detailed structure consists, in fact, of a quasi-periodic train of forward shocks. As shown in Fig. 4, the assumed velocity structure has a ramped sawtooth character; each tip represents relatively fast material that sees unshadowed stellar flux and so is strongly driven, while each trough represents relatively slow material that is shadowed by the fast material and so is only weakly driven. As the fast material is pushed by the radiation force against the slower material, a forward shock forms that sweeps up and accelerates this slower material, thereby adding it to the fast material. A crucial point is that the fast material at each tip represents post-shock flow, with its associated high density and (at least initially) high temperature. To maintain the structure, this material must quickly cool and reform the driving ions that line-absorb the radiative momentum. Assuming that it does so, material at this tip then drifts backward, due to the influence of inward gravitational, inertial, and pressure forces that counter the outward radiative acceleration. At a velocity about a sound speed below the tip, the radiative term becomes sharply diminished by shadowing effects, dramatically increasing the inward drift and quickly bringing material toward the next trough, whereupon it can again be accelerated in the next shock.

The question of the ionization and energy balance behind the shock was analyzed by Krolik and Raymond [68] for a single, forward shock that propagates outward through the wind. They found that the column depth required to cool shocked material back to the ambient wind temperature scaled with the shock velocity $\Delta v$ as
Periodic Forward Shocks

\[ V(r) \]

\[ V - \Delta v \]

\[ \sim a \]

\[ I \]

\[ r \]

Fig. 4. Velocity versus height in Lucy’s model of periodic forward shocks.

\[ N_{\text{cool}} \approx 7 \times 10^{18} \left( \frac{\Delta v}{100 \text{ km/s}} \right)^4 \text{ cm}^{-2}. \] (20)

Comparing this to a characteristic value \( N_{\text{wind}}(r) = nr \) for the total wind column depth at a radius \( r \), we obtain an expression for the critical radius at which the assumption of cooling must break down,

\[ r_c \approx 10^{14} \text{ cm} \left( \frac{\dot{M}}{10^{-6} M_\odot/\text{yr}} \right) \left( \frac{v}{1000 \text{ km/s}} \right)^{-5} \left( \frac{\Delta v}{v/2} \right)^{-4}. \] (21)

Krolik and Raymond used rather elaborate arguments to obtain the approximate expression (See their Eq. (11).),

\[ \Delta v \approx v \left( \frac{N_{\text{thick}} L_*}{\dot{M} c^2} \right), \] (22)

relating the shock strength to basic stellar parameters and the number of optically thick lines (\( f \) in their notation). But recall from Sect. 3 that the factor in parenthesis in Eq. (22) is necessarily of order unity, implying that \( \Delta v \approx v \). In reality, of course, the line driving is divided between accelerating the mean flow and amplifying the instability, and so typically \( \Delta v \approx v/2 \). If one takes these points into account, then the somewhat different expression derived by Krolik and Raymond for \( r_c \) reduces to Eq. (21). Actually, since wind structure can be expected to form on a scale much smaller than the entire wind, breakdown of the cooling should occur at a radius much smaller than \( r_c \), perhaps at \( \sim r_c/10 \).

Recently, MacFarlane and Cassinelli [69] studied the evolution of wind shocks with a phenomenological numerical-hydrodynamics model, meant to
apply to the relatively weak wind from the main sequence B0 star τ Sco. The wind is assumed initially to be in a smooth and relatively slow state, but then a driving force is applied that accelerates the material to a much higher speed. The compression between the fast and slow material quickly forms a dense shell bounded on the front and back sides by a forward and reverse shock pair. Both the driving force and the initially slow flow-speed represent purely phenomenological free parameters, the former chosen to accelerate the fast wind to the observed terminal speed, and the latter chosen to give shocks of the strength necessary to match X-ray observations. Nonetheless, the temporal evolution of the temperature in the resulting shock-heated, dense shell is computed from a reasonably complete energy balance, including an optically thin, radiative cooling function. Ionization lag effects are ignored and so the radiative emission is assumed to be a known function of temperature, which is used to derive the X-ray emission spectrum from the shocks. For parameters chosen to give initial and final state terminal speeds of 500 km/s and 2500 km/s, the forward and reverse shocks each attain velocity amplitudes approaching 1000 km/s, and this turns out to be just what’s needed to reproduce the observed X-ray properties.

Another heuristic wind-structure model worth mentioning is Mullan’s application [70, 71] of the well-known solar wind phenomena of “co-rotating interaction regions” [72, 73] to the case of hot-star winds. The winds here again contain dense shells bounded by forward/reverse shock pairs, but now this is not thought of as arising from the line-driven instability, but rather from an assumed azimuthal variation in the wind flow speed and the interaction between fast and slow wind streams that results from the stellar rotation. Mullan argues that the resulting shocks and dense shells can give rise to the observed X-ray emission and discrete absorption components in UV lines. Although there is no independent evidence that O star winds have such an azimuthal stream structure, the effect of rotation in laterally extending any structure forming from the wind instability is likely to be quite similar.

### 4.4 Radiation-Hydrodynamical Simulation of Non-Linear Structure

Owocki, Castor, and Rybicki (hereafter OCR) [44] have recently developed a numerical, radiation-hydrodynamics code aimed at directly simulating the dynamical evolution of the line-driven wind instability, and thereby determining the likely nature of the resulting nonlinear wind structure. Because the instability occurs for perturbations with a length scale near and below the Sobolev length, OCR had to develop a method for computation of the line force which did not use the Sobolev approximation, but which still avoided the inordinate computational expense of solving the full line-transfer problem at each time step. The crucial simplification they adopted was to ignore the diffuse, scattered radiation, and to assume that the flow is driven by an ensemble of nonoverlapping, pure-absorption lines. By integrating the pure-absorption force expression (3) over a power-law number distribution $N(\kappa) \sim \kappa^{a-2}$, we see that this combined force is of the form

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\[
g_{\text{tot}}(r) \sim \frac{L_*}{r^2} \int_{-\infty}^{\infty} dx \phi(x - v(r)/v_{\text{th}}) \eta^{-\alpha}(x, r)
\]

where \(\eta(x, r)\) is a profile-weighted mass-column-depth given by

\[
\eta(x, r) = \int_{R_*}^{r} \rho(r')\phi(x - v(r')/v_{\text{th}})dr' + \eta(x, R_*).
\]

In a smooth flow with all variations on a scale much larger than a Sobolev length \(L\), this force appropriately reduces to the CAK form (12) (See Sect. 3); but in the presence of structure on a scale smaller than \(L\), it also correctly includes the local Doppler-shift effects that give rise to the line-driven instability. Although much simpler than solving scattering-line transfer, this pure-absorption force is still much more complicated than the usual body forces assumed in hydrodynamical simulations, requiring at each time step a numerical evaluation of the double integral in \(x\) and \(r\) that dominates the code timing.

In order to focus on such dynamical terms, OCR avoided treatment of detailed energy balance, and simply assumed that radiative heating and losses would keep the wind nearly isothermal. This means they could not directly compute effects that depend heavily on the temperature, e.g., X-ray emission. In emphasizing dynamics but neglecting energy balance, the OCR calculation is in a sense the complement of that of MacFarlane and Cassinelli [69], who treated a detailed energy balance but assumed phenomenological force terms. (See Sect. 4.3.)

For a model with typical O-star parameters, Fig. 5 shows the wind spatial structure that results long \((10^5\,\text{s})\) after introduction of a 1% amplitude, 4000 s period sound wave propagating outward from the wind base. The strong correlation of the driving force with flow speed quickly amplifies this initially small perturbation, giving rise to velocity variations \(\Delta v \gtrsim 500\,\text{km/s}\) by \(r \approx 1.5R_*\). The base perturbation is initially outward propagating, and thus has velocity and density variations that are in phase, but the resulting wind structure has velocity and density with opposite phase. This turns out to be a very robust result, stemming directly from the linear instability property [39] that such opposite-phased waves are much more unstable. In fact, expected nonlinear dynamical effects like line-shadowing actually play little role, and so the waves grow at the rate predicted by linear theory up to the highly nonlinear amplitude \(\Delta v \gtrsim 500\,\text{km/s}\), at which they become steepened by the simple kinematic effect of faster flow overrunning the slower. Although advected away from the star by the supersonic flow, these unstable waves actually propagate inward relative to the fluid, and so as they steepen they naturally evolve into reverse, not forward, shocks.

It is important to emphasize that, although the velocity structure in the hydrodynamical model of Fig. 5 appears quite similar to that assumed in the Lucy periodic shock model of Fig. 4, the density structure of the two models, and hence their entire nature, is actually fundamentally different. As described in Sect. 4.3, in Lucy's model forward shocks were assumed to abruptly accelerate ambient wind material as it is rammed by a dense, strongly driven flow. In

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Fig. 5. Radial variation of radiation force, velocity, and density in OCR's numerical simulation of the nonlinear wind structure arising from introduction of a 1% amplitude, 4000 s period sound wave at the base of an unstable, line-driven stellar wind. (Note that the density is plotted increasing downward.)

In contrast, the strong reverse shocks in the OCR simulation arise to decelerate high-speed, rarefied flow as it impacts slower material that has been compressed into dense shells. Since the high-speed flow is accelerated before being shocked, there is no need to assume rapid cooling to reform driving ions. In fact, the very low density of this high-speed flow implies that, unlike Lucy's model, only a very small fraction (≈10^{-3}) of the wind material ever undergoes a shock near the maximum amplitude Δv ≈ 1000 km/s.

These overall characteristics of this dynamically computed wind structure seem in good qualitative correspondence with several observational properties. For example, the resulting shock strengths Δv ≲ 1000 km/s are roughly what's needed to obtain the 10^7 K plasma inferred from the high-energy tail observed in X-ray spectra, and the small fraction of strongly shocked material implies that a similarly small fraction ≲ 10^{-3} of the total flow kinetic energy Mv_∞^2/2 ≈ 10^{-3} L\_\odot goes into shock heating, which agrees roughly with the observed X-ray
luminosity scaling $L_x \approx 10^{-7} L_*$. The dense shells should naturally give rise to narrow line absorption features similar to the discrete absorption components observed in unsaturated lines, and the high-speed, rarefied flow should give rise to non-sharp, variable edges in saturated lines, which is again what's often observed [61]. Although it remains to be shown in detail, it also seems likely that the nonmonotonic velocity field will give rise to the observed black absorption troughs in saturated lines.

4.5 Intrinsic Nature of Wind Variability

The structure shown in Fig. 5 arises from amplification of an explicit perturbation introduced at the wind base, but subsequent work [74, 75] has shown that such winds can also often exhibit an intrinsic variability that persists even in the absence of such explicit perturbations. The incidence of this intrinsic variability was found empirically to depend on the assumed value of $v_{th}/a$, the ratio of the ion thermal speed to sound speed. For the usual case of driving by CNO ions, an appropriate value of this ratio is $v_{th}/a \approx 0.3$. However, in order to study the effect of variability arising from explicit perturbations without the complication of a background model that was itself intrinsically variable, OCR found it necessary to assume an artificially high value $v_{th}/a \gtrsim 1/2$. The situation is graphically illustrated in Fig. 6, which shows the spatial and temporal variation of velocity in two unperturbed models that differ only in the assumed $v_{th}/a$. In both cases the flow is initially disrupted because the assumed CAK/Sobolev initial condition is not an appropriate steady state for either non-Sobolev model. The model with $v_{th}/a = 1/2$, however, quickly relaxes to a somewhat steeper [75] steady solution, whereas the model with slightly smaller $v_{th}/a = 3/8$ never settles down, but exhibits a nearly periodic variability.

Poe, Owocki, and Castor (POC) [75] showed that the different variability properties of these two cases reflect a difference in the nature of the corresponding steady solutions. Fig. 7 illustrates the steady-state solution topology near the critical (sonic) point. Unlike the usual saddle- or X-type solutions that apply, e.g., to the solar wind (Fig. 7a), the solution topology in this case is of the nodal type, with not one but two positive slope critical solutions (Fig. 7b). Note that along the steeper slope there is only one distinct solution passing though the critical point, whereas along the shallower slope there is a degenerate family of solutions that pass through this point. POC showed that, for reasonable boundary conditions, the distinct, steeper solution applies when $v_{th}/a \gtrsim 1/2$, whereas the degenerate, shallower solution applies when $v_{th}/a \lesssim 1/2$. Apparently, the existence of a well-defined steady solution in the former case is sufficient to enable a time-dependent model without explicit perturbations to relax to this steady state. On the other hand, the lack of a well-defined steady solution for the latter case leads in the time-dependent model to an intrinsic variability in which the flow, roughly speaking, continuously varies among this degenerate family of possible steady solutions. (Recall from Sect. 3 that a similar solution degeneracy exists for Sobolev models in the limit of vanishing sound speed $a \to 0$.)
Fig. 6. 3-D perspective plot of the height and time variation of velocity in two unperturbed wind models which differ only in the assumed values of the ratio of thermal speed to sound speed.

Fig. 6 shows that there are brief intervals when the latter flow approaches a nearly smooth, CAK-like state, but this becomes disrupted by the strong amplification of small scale structure, which is then advected away, again allowing the flow to settle temporarily into a nearly smooth state. The time scale (\(\sim 1/2\) day) for this relaxation oscillation is much longer than the insta-
Fig. 7. The velocity near the sonic radius for solution topologies of (a) the usual saddle type that applies to the solar wind and (b) a nodal type that applies to absorption-line-driven winds. Note in the nodal case how a large number of solutions converge as they approach the sonic point along the shallower of the two positive critical solutions (solid lines). The implied solution degeneracy leads to the intrinsic variability seen in the right panel of Fig. 6.

bility growth time (~1 hour) of the individual small scale instabilities, and corresponds roughly with the observed repetition time of discrete absorption features (Sect. 4.1) [61,62]. Indeed, the time-variations of synthesized absorption profiles show moving narrow absorption features quite reminiscent of such discrete absorption components [35, 36]. These arise from outward accelerating dense shells whose formation is repeatedly triggered by the quasi-periodic disruptions. It is interesting that, all by itself, such an unstable wind seems to give rise to structure with many of the qualitative features needed to explain such observational signatures of the wind structure.

4.6 Implications for Steady-State Models

If hot-star winds are indeed as temporally and spatially structured as implied by Figs. 5 and 6, how is it possible that steady-state models, which completely ignore this structure, can have been so successful in quantitatively matching most properties of time-averaged spectra?

In this regard, it is first worth noting that, despite the extensive wind structure, the gross wind properties like the terminal flow speed and time-averaged mass loss rate turn out to be in quite good agreement with those derived in steady models. Moreover, as noted above, only a small fraction of the mass in this wind is actually accelerated in a highly time-dependent way; indeed, the acceleration history of most of the mass is quite similar to what it would be in a smooth wind. From this point of view, the major effect of the instability is thus simply to spatially clump material without otherwise altering its properties. Given the intrinsic mass-weighting of spectral formation, and the extensive temporal and spatial averaging involved, it thus does not seem surprising that most spectral features would be quite similar in the smooth and structured flows.
A second point is that it is not at all clear whether the intrinsic variability characteristic of this pure-absorption model is really applicable to actual hot-star winds, for which the driving is actually by line-scattering. As noted in Sect. 3, the near fore-aft symmetry of the diffuse radiation field means that the associated net force is typically small in the supersonic regions of a smooth flow; but in the region near and below the sonic point, the radiation escape probability is rapidly increasing, leading to a diffuse field asymmetry and hence to a kind of a radiative viscous force. This can alter or even “break” the solution degeneracy found in pure-absorption models [76], and so might explain why comoving frame models that include scattering appear to have a well-defined solution [25]. Through the line-drag effect (Sect. 4.2), scattering might also stabilize the wind base enough to suppress variations that govern the intrinsic variability. The wind instability would then have an advective rather than absolute character [77, 78, 74], for which variability would persist only with some explicit driving from the underlying star, and would only become large amplitude in the supersonic wind.

5 Future Directions

Clarifying the dynamical role of scattering should thus be a primary area of focus in future work on hot-star winds. This refers to both its possible effect in regulating the wind instability, as well as its likely direct role in driving the optically thick winds from Wolf-Rayet stars, for which there is still no satisfactory dynamical theory of even the mean wind outflow. For OB stars, additional refinements in the steady-state wind models can be expected to further improve the already good quantitative agreement with time-averaged spectral observations. In my view one particularly important refinement will be to consider, if only in a phenomenological way, how some of the consequences of wind structure – e.g., X-ray heating, ionization – affect the overall properties of the time-averaged flow [79].

With regard to the direct study of the wind structure and variability itself, there is obviously a great deal of work still to be done to include the many potentially important effects neglected so far, e.g., scattering, detailed energy and ionization balance, shock X-ray emission, 2-D or 3-D structure, and stellar rotation. In addition to clarifying the role of scattering in regulating the instability, there is the question of what dynamical role it may play in the non-linear wind structure. Inclusion of a realistic energy balance in the dynamical wind models is needed to determine the temperature structure and thereby the shock X-ray emission. Likewise, inclusion of ionization balance is needed to determine whether ionization from shock heating and/or X-rays can effect the line-driving and hence the dynamics of the wind and instability. Consideration of rotation and other 2-D (and ultimately 3-D?) effects is needed to determine the likely lateral scales of the dense clumps. For example, can rotation string
out such structures in longitude in a manner similar to corotating interaction regions in the solar wind [70, 71] (Sect. 3), and thus make it possible for a given dense clump to cover a large enough fraction of the stellar disk to produce the observed discrete absorption components? Finally, perhaps one of the most urgent needs is to develop methods to compare more closely predictions from the theoretical simulations with available observational diagnostics, and thereby guide and test the further theoretical development in this complex but fascinating problem in radiation hydrodynamics.

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