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Modeling Ejecta of Massive Stars

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Abstract. In this review I describe the various ways in which massive stars can shed matter. At first take, these stars seem to share little in common with the planetary nebula phenomenon. Their mass loss typically takes the form of periods of continuous outflow, rather than the discrete outbursts that form PNe shells. Nevertheless, the ejecta of massive stars tend to be highly structured in space and in time, and it seems likely that some of the mechanisms responsible for this may also explain the diverse morphologies seen in PNe.

1. Foreword

Throughout my discussion, I'll be placing deliberate emphasis on four key questions. What causes the structure we see in the ejecta of massive stars? What tools can we use to model this structure? How have these tools so far been applied? And, what problems do we still face?

Due to limitations on page space, and moreover on the boundaries of my own knowledge, I shall cravenly pass over the most extreme forms of massive-star mass loss: binary interactions, giant eruptions and supernovae.

2. Massive-Star Winds

Although making up a small fraction of the mass in galaxies, massive stars[1] are responsible for a preponderance of the radiant energy generation. This is a reflection of their immense luminosities, which — with a main-sequence scaling $L_*/\mathrm{L}_\odot \sim (M_*/\mathrm{M}_\odot)^3$ — can reach up to a dizzying million-times solar, for the stars at the top end of the mass range ($M_* \sim 100\,\mathrm{M}_\odot$).

These luminosities give rise to strong wind outflows. In contrast to the wimpy ($\dot{M} \sim 10^{-14}\,\mathrm{M}_\odot\,\mathrm{yr}^{-1}$) winds of the Sun and other cool main-sequence stars, which are accelerated by the gas pressure of a hot corona, massive-star winds are driven radiatively by the direct conversion of photon momentum into flow momentum. Mass-loss rates range up to $10^{-5}\,\mathrm{M}_\odot\,\mathrm{yr}^{-1}$, and typical terminal velocities $v_\infty \sim 1,000$ – $2,000\,\mathrm{km\,s^{-1}}$ are factors of a hundred times the sound speed — that is, the winds are *hypersonic*.

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[1]Here, defined loosely as those having spectral types O and B on the main sequence.

A necessary condition for a radiation-driven wind to arise is that the outward radiative force exceed the inward pull of gravity. This familiar breaching of the Eddington limit is usually framed as $\Gamma_e > 1$, where

$$\Gamma_e = \frac{\kappa_e L_*}{4\pi G M_* c} \tag{1}$$

is the ratio of radiative and gravitational accelerations. However, the presumption in this expression is that the radiative force arises from Thomson scattering of continuum photons (as represented by the opacity κ_e). While such *continuum-driven* winds can occur in extreme cases (for instance, the Great Eruption of η Car; see Smith & Owocki 2006), massive-star winds are more typically driven by line opacity associated with resonance-scattering transitions.

Within the Castor et al. (1975, hereafter CAK) formalism for these *line-driven* winds, the Eddington parameter above is replaced by

$$\Gamma_\ell = \frac{1}{1-\alpha}\Gamma_e \bar{Q} \left(\frac{\mathrm{d}v/\mathrm{d}r}{\rho c \bar{Q} \kappa_e} \right)^\alpha \tag{2}$$

where the notation follows Owocki (2004), and in particular $\bar{Q}$ is a measure of the total line opacity in the wind, in units of κ_e (see Gayley 1995). The term in parentheses is the reciprocal of the Sobolev optical depth; through the appearance of the spatial velocity gradient $\mathrm{d}v/\mathrm{d}r$, it represents the degree to which the wind is able to Doppler-shift lines out of their own shadow.

The inverse dependence of the radiative acceleration on the density ρ naturally introduces a negative feedback loop that helps to self-regulate the wind. If too much mass is launched from the surface, the acceleration declines to a point where Γ_ℓ drops below unity; the wind then stalls and falls back to the stellar surface. Accordingly, the mass-loss rate $\dot{M}$ of a radiatively driven wind is not a free parameter, but instead established by this self-regulation process — that is, $\dot{M}$ is an *eigenvalue* of the system. In the idealized point-star case, this eigenvalue is given by CAK theory as

$$\dot{M}_{\mathrm{CAK}} = \frac{L_*}{c^2} \frac{\alpha}{\alpha - 1} \left(\frac{\bar{Q}\Gamma_e}{1-\Gamma_e} \right)^{(1-\alpha)/\alpha}, \tag{3}$$

and solution of the equation-of-motion for $\dot{M} = \dot{M}_{\mathrm{CAK}}$ gives a wind terminal velocity that scales with the escape velocity,

$$v_\infty = \sqrt{\frac{\alpha}{1-\alpha}} v_{\mathrm{esc}} = \sqrt{\frac{\alpha}{1-\alpha}\frac{2GM_*}{R_*}}. \tag{4}$$

Owocki (2004) provides an excellent review of the formalism leading to these various results.

3. Line-Deshadowing Instability

Although the CAK formalism envisages steady, smooth outflows, in reality line-driven winds exhibit significant structure and variability, on scales both small and large, and with morphologies that can be regular/symmetric or stochastic. As I shall discuss in subsequent sections, the structure in many cases arises via imprinting from an external agent such as pulsation or rotation. However, there is also an intrinsic instability in radiative line driving that leads to the spontaneous generation of structure. The instability reveals itself in a linear perturbation analysis (e.g., Lucy & Solomon 1970), indicating a perturbed radiative force that is directly proportional to the perturbed velocity; this simply reflects the amount of additional opacity that is Doppler shifted out of its own shadow at the blue edge of line profiles. The growth timescale of this so-called *line-deshadowing instability* (LDI) is $\tau_{\rm grow} \sim 10^{-2}\,\tau_{\rm wind}$, which — with wind flow timescales $\tau_{\rm wind} = R_*/v_\infty \sim 0.5\,\rm d$ — is very short.

The Lucy & Solomon (1970) analysis suggests instability at all length scales; but a more-careful investigation (see Owocki & Rybicki 1984) shows that the LDI operates only for velocity perturbations whose physical length scale is shorter than the Sobolev length $\ell_{\rm Sob} = v_{\rm th}/({\rm d}v/{\rm d}r)$. At larger scales, the perturbed line force is proportional to the perturbed velocity *gradient* (as one might expect from eqn. 2), and the wind is stable. Accordingly, the LDI is primarily responsible for the generation of *small-scale* structure in line-driven winds.

Numerical simulation of the LDI can be computationally expensive, as the Sobolev approximation (which assumes $\ell_{\rm Sob}$ is small compared to any structure in the wind) cannot be used. 1-D hydro simulations by Feldmeier & Owocki (1998) indicate that the instability breaks up a smooth CAK wind solution into a sequence of reverse shocks, where fast, low-density wind material runs into the back of slower-moving, high-density material. These wind shocks are considered a likely source for the soft, broad-lined X-ray emission observed in many single OB stars (e.g., Owocki & Cohen 2006, and references therein).

Extending the simulations to 2-D, Dessart & Owocki (2005) find that the shell-like shocks produced by the LDI are fragmented by Rayleigh-Taylor instabilities, and the wind structure rapidly becomes incoherent down to angular scales approaching the grid scale. These results are difficult to reconcile with observations of clump-induced stochastic variability in massive-star winds (e.g., Townsend & Mast 2010), suggesting that there's something still missing from the simulations. Lateral forces from side-scattered radiation are a likely candidate, since they will tend to retard the shell fragmentation process (e.g., Owocki 2004). However, inclusion of these forces will require multi-D radiation hydrodynamics, which for the time being lies beyond our computational grasp.

4. Cyclical Wind Structure

In addition to the stochastic structure caused by the LDI, massive-star winds show evidence for larger-scale structures that recur periodically or episodically. These structures usually manifest themselves as *discrete absorption components* (DACs) that migrate blueward through the absorption troughs of ultraviolet P Cygni line profiles (e.g., Kaper & Henrichs 1994). Puzzlingly, the typical lifetime of DACs is significantly longer than wind flow times, suggesting that they are caused not by embedded clumps, but instead by patterns in the wind that remain coherent over large spatial scales and long time scales.

The *IUE Megacampaign* observations of the B0.5Ib supergiant HD 64760, spanning almost 16 days, revealed a pair of DACs with lifetimes on the order of $\sim 10\,\rm d$

(Prinja et al. 1995). However, superimposed over these DACs is a 1.2 d-periodic modulation in the depth of UV line profiles, phase-locked over a wide range of blueshift velocities. Because the period is close to one-quarter of the star's rotation period, these *periodic absorption modulations* (PAMs) have been hailed as the indirect evidence of a 'photospheric connection' between surface and wind structure. Support for the existence of this connection has come from a more-recent study of HD 64760 by Kaufer et al. (2006), who find that the star's wind-sensitive Hα line is variable on a 6.8 d period, corresponding to the beat period between photospheric non-radial pulsation modes (the issue of pulsation is discussed further in §5). How this 6.8 d period is related to the 1.2 d PAM period is, however, still unclear.

What causes DACs and PAMs? Mullan (1984, 1986) first suggested that DACs are the signature of co-rotating interaction regions (CIRs) — rotating spiral structures formed by the collision between fast and slow wind streams, which trace back to some kind of inhomogeneity in the photospheric wind-launching. However, based on kinematical models, Fullerton et al. (1997) have convincingly argued that it is *PAMs* rather than DACs which are the observational manifestation of CIRs.

5. Rotational Wind Sculpting

Massive stars are systematically rapid rotators: the observational survey by Howarth et al. (1997), covering 373 O- and B-type stars, reveals a distribution of equatorial rotation velocities $v_{\rm eq}$ peaking at $\sim 100\,{\rm km\,s^{-1}}$, and extending all the way up to the critical limit $\sim 500\,{\rm km\,s^{-1}}$. What effect does this rapid rotation have on the stars' winds? Physical intuition suggests that the reduction in the effective (Newtonian plus centrifugal) gravity will make it easier to launch a wind from the equator of a rotating star, than from the poles. Thus, we should expect oblate mass loss around more-rapid rotators.

This expectation is at first glance lent support by a simple extension of eqn. (3), giving a latitude-dependent mass flux

$$\dot{m}_{\rm CAK}(\theta) = \frac{F}{c^2}\frac{\alpha}{\alpha-1}\left(\frac{\bar{Q}\Gamma_{\rm e}}{1-\Gamma_{\rm e}}\right)^{(1-\alpha)/\alpha}, \qquad (5)$$

where now

$$\Gamma_{\rm e}(\theta) = \frac{\kappa_{\rm e}F}{g_{\rm eff}c}, \qquad (6)$$

and F and $g_{\rm eff}$ are the local radiative flux and effective gravity, respectively. These expressions indeed seem to indicate that a lowering of $g_{\rm eff}$ will push Γ_e closer to unity, in turn upping the local mass flux.

However, as Cranmer & Owocki (1995) first pointed out, the von Zeipel (1924) gravity darkening law — which establishes the proportionality $F \propto g_{\rm eff}$ — means that Γ_e is *constant* over the stellar surface. The elevation of F at the stellar poles (due to the higher effective gravity) then means an enhanced mass flux there, and an overall *prolate* mass-loss morphology. Moreover, the wind terminal velocity v_∞ (cf. eqn. 4) will also be larger over the poles than at the equator, due to the higher escape velocity. Smith et al. (2003) presents evidence for both of these effects in the present-day wind of η Car.

6. Rotation and Disks

The preceding section points toward bipolar outflows from rotating massive stars. How, then, do some show the clear signature of equatorial disks — in particular, the enigmatic Be stars, whose disks are revealed in Hα emission lines (e.g., Porter & Rivinius 2003)? Bjorkman & Cassinelli (1993) proposed that wind streams launched off the surface of rapidly rotating B-type star are focused toward the equatorial plane by angular momentum conservation, where they will collide and form a disk. This kinematical *wind-compressed disk* (WCD) model was initially confirmed in hydrodynamical simulations by Owocki et al. (1994). However, when gravity darkening and stellar oblateness were incorporated in the simulations, Cranmer & Owocki (1995) found that the polar-enhanced mass loss discussed above tends to reduce the density of the equatorial disk. Moreover, when the simulations were further extended to incorporate non-radial line forces (arising from the velocity gradient between equator and pole via the dv/dr term in eqn. 2), Owocki et al. (1996) showed that the formation of a WCD is completely inhibited.

Progress in finding an alternative narrative for Be-disk formation was driven by observations constraining the disk velocity structure as Keplerian, with radial outflow speeds on the order of a couple of $\mathrm{km\,s^{-1}}$ or less (e.g., Hanuschik 1996). These observations prompted a reevaluation of a seminal paper by Lee et al. (1991) advancing a viscous decretion model. Decretion disks operate much like accretion disks, in that they transport angular momentum away from the star; the key difference is that mass is injected at the inner boundary of a decretion disk, and at the outer boundary of an accretion disk. In both cases, the disk velocity structure is Keplerian.

Smooth particle hydrodynamics simulations confirm the viability of the viscous decretion model (e.g., Okazaki 2004). However, a missing component to the model is a prescription of how stellar material is lifted into (Keplerian) orbit at the inner edge of the disk. This process requires angular momentum — and elucidating the mechanism that supplies the angular momentum can be regarded as the key to understanding the Be phenomenon. Radiative driving (as envisaged in the WCD model) is generally unsuited, because it is difficult for photons to impart an azimuthal force (see, however, Gayley & Owocki 2000, for an interesting counter-example). Magnetic fields have been considered; however, as discussed below in §7, they appear better-suited to producing *rigid* disks.

This leaves rotation itself as the most promising mechanism for supplying the necessary angular momentum. Of course, if *all* of the angular momentum is to come from rotation, then Be stars must be spinning at their critical velocity — an unlikely scenario. We therefore require an additional agent to supply the velocity boost into orbit. For agents relying on gas pressure to accelerate material, an upper limit on the magnitude of the boost is approximately the photospheric sound speed a; thus, the star's equatorial velocity must be within one to two a of the surface orbital velocity, which in massive stars translates to 90 – 95% of critical. Historically, the consensus had been that Be stars *do not* rotate this close to critical. However, Townsend et al. (2004) demonstrate that gravity darkening can cause a significant underestimation of the projected equatorial velocity $v_{\rm eq} \sin i$ measured from line widths, and that in fact observation of Be stars are consistent with (although do not *prove*) rotation in the required 90 – 95% critical window.

A specific instance of the 'rotation-plus-gas-pressure' mechanism for disk formation was suggested by Osaki (1986), who proposed that photospheric g-mode pulsation waves in a near-critical star can launch equatorial material into orbit. Owocki (2005) has supported this *pulsation-driven orbital mass ejection* (PDOME) model using 2-D (equatorial-plane) hydrodynamical simulations, in particular confirming Osaki's stipulation that *prograde*-propagating waves are required for disk formation (since they supply angular momentum in the requisite prograde direction). These theoretical developments tie in nicely with observational work by Rivinius et al. (2003) indicating that g-mode pulsation is ubiquitous in early-type Be stars. One remaining puzzle, however, is that the modes detected by these latter authors appear to be universally *retrograde* rather than prograde. My own suspicion is that so-called 'Yanai modes' (e.g., Townsend 2003), which exhibit retrograde phase propagation but transport angular momentum in the prograde direction, may therefore have some role to play.

The Be stars should not be confused with objects exhibiting the B[e] phenomenon. Although the latter show similar signatures of circumstellar material, they differ by dint of their *forbidden* emission lines and their IR excesses indicative of hot dust. Moreover, B[e] stars are a very heterogeneous group, spanning a broad range of evolutionary phases from pre-main sequence to PN. Within this group, massive stars are represented by the supergiant B[e] stars, characterized by fast polar winds and dense equatorial disks which appear to be outflowing rather than Keplerian (see, e.g., Lamers 2006, and references therein). One formation mechanism for these disks is rotation-induced bistability (Pelupessy et al. 2000): the lower effective temperature at the (gravity-darkened) equator produces $\underline{a}$ higher opacity, which in turn allows a greater line-driven mass flux there (due to the $\bar{Q}$ term in eqn. 5). Since the disks are outflowing, there is not the same angular momentum requirement as with Be-star disks.

7. Magnetically Channeled Winds

Massive stars aren't expected to harbor magnetic fields, owing to the absence of a significant outer convection zone to serve as a field-generating dynamo. Nevertheless, a subset show evidence for strong ($\sim$ kG), ordered (typically, dipole) magnetic fields which are stable over timescales of decades (see, e.g., Wade et al. 2009). Because massive-star winds are highly ionized, there is a strong coupling between wind and field, with each competing to determine the morphology of the other. ud-Doula & Owocki (2002) investigate this coupling using 2-D magnetohydrodynamical (MHD) simulations of radiation-driven winds in the presence of a dipole field. A key result from these simulations is that the global flow morphology is determined by a single parameter

$$\eta_* = \frac{B_*^2 R_*^2}{\dot{M} v_\infty} \tag{7}$$

representing the ratio between magnetic and wind-kinetic energy densities adjacent to the stellar surface. When $\eta_* \ll 1$, the wind dominates the field, and the latter is stretched out into a split-monopole configuration with radial field lines and an equatorial current sheet. Conversely, in the $\eta_* \gg 1$ limit the field remains relatively unaffected by the wind, with field lines remaining closed out to the Alfvén radius $R_{\rm Alf} \approx \eta_*^{1/4} R_*$.

In the latter 'strong confinement' case, wind streams flowing from opposite footpoints of closed magnetic loops collide near the loop summit. In their seminal *magnetically confined wind shock* (MCWS) paradigm, which foreshadowed the MHD simulations, Babel & Montmerle (1997b) proposed that the kinetic energy of the colliding streams is thermalized in shocks, heating the material to temperatures $\sim 10^6 - 10^7$ K where X-ray emission becomes significant. A follow-up paper (Babel & Montmerle 1997a) argues that the hard X-rays of the O6V star θ^1 Ori C arise in this manner (interestingly, this paper pre-dated the actual detection of the star's ~ 1.1 kG field by Donati et al. 2002). Subsequent analysis of the star's X-ray emission-line spectrum, coupled with MHD simulations of the wind shocks, has lent strong support to the basic MCWS paradigm (see Gagné et al. 2005).

More-recent MHD simulations have explored the impact of magnetic-axis-aligned rotation (ud-Doula et al. 2008, 2009). As the post-shock material cools radiatively back to photospheric temperatures, it accumulates in the equatorial plane to form a *rigidly rotating* disk supported against the inward pull of gravity by centrifugal and magnetic forces. Eventually, sufficient material builds up to overwhelm the magnetic field, which then reconnects. During such a centrifugal breakout episode (see ud-Doula et al. 2006), material in the disk is flung away from the star and escapes to infinity.

This narrative runs contrary to the *magnetically torqued disk* model introduced by Cassinelli et al. (2002), in particular showing that disks formed with the aid of magnetic fields tend to be rigid rather than Keplerian, and are thus unsuited to explaining the Be phenomenon. In fact, the MHD simulations lend support to the alternative *rigidly rotating magnetosphere* (RRM) model developed by Townsend & Owocki (2005) to describe the distribution of circumstellar material in the idealized limit $\eta_* \longrightarrow \infty$ where field lines are completely rigid. This limit is effectively realized in the chemically peculiar He-strong stars, whose ~ 10 kG fields and low mass-loss rates combine to produce confinement parameters in the $\eta_* \sim 10^5 - 10^7$ range.

Applied to the B2Vpe star σ Ori E, the RRM model can simultaneously reproduce the observed 1.2-d periodic Hα and photometric variations (Townsend et al. 2005). Building on this success, Townsend et al. (2007) have created a *rigid field hydrodynamics* (RFHD) approach for simulating the time-dependent flow along an ensemble of rigid field lines. RFHD models provide a 3-D, dynamical picture of a star's magnetosphere, including the collision shocks anticipated in the MCWS paradigm, at a tiny fraction of the computational cost of equivalent MHD simulations; however, they remain restricted to cases where $\eta_* \gg 1$. An initial application to σ Ori E (Hill et al. 2010) shows promise in reproducing the star's observed X-ray emission.

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