The Dynamical Role of Radiative Driving in the Sources and Sinks of Circumstellar Matter in Massive Stars

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Abstract. The high luminosity of massive stars drives strong stellar wind outflows. In magnetic massive stars, the channeling and trapping of wind material can feed a circumstellar magnetosphere, characterized either by transient suspension and dynamical infall in slow rotators, or by long-term centrifugal support in moderately fast rotators. In the non-magnetic but rapidly rotating Be stars, direct centrifugal ejection of material from the equatorial surface can feed a Keplerian decretion disk, with radiative forces now playing a potential key role in disk dissipation, through line-driven ablation from the disk surface. This contribution reviews these dynamical roles of radiative driving in the sources and sinks of circumstellar matter, within the context of using high resolution observations to test and constrain circumstellar models.

1. Introduction

Stellar spectra are commonly dominated by absorption lines formed in a hydrostatically bound, narrow, photospheric layer of thickness much less than the stellar radius. By contrast, many massive stars also show distinctive emission lines, formed from the larger effective surface area of circumstellar matter. A key issue is how such matter is lifted from strong gravitational binding of the stellar surface, which in massive stars can occur by a combination of centrifugal, radiative, and/or magnetic forces.

For example, the characteristic Balmer emission of Be stars is understood to arise from a circumstellar decretion disk, fed by rapid, near-critical surface rotation that places material into a nearly Keplerian orbit, in which centrifugal forces nearly balance the local gravitational acceleration. Alternatively, the high luminosity of O and WR stars can drive strong, outflowing stellar winds, with mass loss rates sufficiently high that even in the relatively short, dynamical flow time, the steady-state wind density is sufficient to result in recombination emission in H-alpha, He, or other spectral lines. Finally, in the subpopulation of massive stars with strong, large-scale magnetic fields, a combination of magnetic trapping and/or magnetic torquing of the radiative driven outflow can feed a circumstellar magnetosphere, again characterized by emission in Balmer and other lines.

The discussion here reviews this dynamical interplay of radiative driving with both rotation and magnetic fields, with emphasis on its role in the sources and sinks of these three specific forms of massive-star circumstellar matter.
2. Mass Loss in Wind vs. Disk

2.1. Basic Scalings for Line-Driven Wind

Let us begin by reviewing the basic scalings for stellar winds driven by line-scattering of the star’s continuum radiation. For the simple case of continuum electron scattering with opacity $\kappa_e$, the ratio of the radiative acceleration to gravity is just a fixed constant, set by the ratio of the stellar luminosity $L$ to mass $M$,

$$\Gamma_e \equiv \frac{\kappa_e L}{4\pi GMc} = 2 \times 10^{-5} \frac{L/L_\odot}{M/M_\odot}.$$  

(1)

But the resonant nature of line (bound-bound) scattering from metal ions leads to an opacity that is inherently much stronger than from free electrons, by a factor set by the “Quality” of the resonance (Gayley 1995). In the somewhat idealized, optically thin limit that all the line opacity could be illuminated with a flat, unattenuated continuum from the full stellar luminosity, the total line-force would exceed the free-electron force by a factor $Q \approx 2000$. In contrast, in the dense, nearly static layers of the atmosphere and interior the lines become saturated, greatly reducing the line-force and so keeping these regions gravitationally bound. But within the accelerating wind, the Doppler shift of the line-resonance out of the absorption shadow of underlying material partly desaturates the lines, exposing the line opacity to a less attenuated flux, and making the line-force scale with the velocity gradient divided by the density, $(1/\rho)dv/dr$ (figure 1).

The maximal wind mass loss rate $\dot{M}$ is thus set at a level for which the net line driving is just sufficient to overcome gravity. Since the net wind acceleration $v dv/dr$ thus also scales with the local gravity $GM/r^2$, direct integration gives a velocity law of the form,

$$v(r) = v_\infty \left(1 - \frac{R}{r}\right)^\beta,$$  

(2)
where the wind terminal speed scales in proportion to the effective surface escape speed, 
\( v_{\text{esc}} = \sqrt{GM(1 - \Gamma_e)/R} \). For the classical CAK (for Castor, Abbott & Klein 1975)
model that assumes a point-star radiation source and a power-law ensemble of driving
lines, we obtain \( \beta = 1/2 \). The associated mass loss rate is given by the simple scaling,

\[
\dot{M}_{\text{CAK}} = \frac{L}{c^2} \frac{\alpha}{1 - \alpha} \left[ \frac{\mathcal{Q}}{1 - \Gamma_e} \right]^{(1-\alpha)/\alpha}, \tag{3}
\]

with the power index \( \alpha \) measuring the relative contribution of optically thick vs. thin
lines. If one takes into account the finite angular extent of the stellar disk, then near the
stellar surface the radiative force is reduced by a factor \( f_d \approx 1/(1 + \alpha) \), leading to a
reduced mass loss rate (Friend & Abbott 1986; Pauldrach, Puls & Kudritzki 1986).

\[
\dot{M}_{f_d} = f_d^{1/\alpha} \dot{M}_{\text{CAK}} = \frac{\dot{M}_{\text{CAK}}}{(1 + \alpha)^{1/\alpha}} \approx \dot{M}_{\text{CAK}}/2. \tag{4}
\]

Away from the star, the correction factor increases back toward unity, which for the re-
duced base mass flux implies a stronger, more extended acceleration, giving a terminal
speed, \( v_\infty \approx 2.5v_{\text{esc}} \), and a flatter velocity law, approximated by replacing the exponent
in eqn. (2) by \( \beta \approx 0.8 \).

In practice such line-driven winds are subject to a strong line-deshadowing in-
stability that leads to extensive wind structure characterized by embedded shocks and
clumps. See the proceedings contribution by Sundqvist for basic references, along with
an update on ongoing efforts to account for unstable wind structure in quantitative mod-
ing of both UV resonance lines like PV, as well as recombination lines like H\( \alpha \).

### 2.2. Effect of Rotation

Initial investigations (Friend & Abbott 1986; Pauldrach et al. 1986) of the effect of
rotation on radiatively driven winds derived 1D models based on the standard CAK
line-driving formalism, but now adding the effect of an outward centrifugal acceleration
in the equatorial plane, \( g_{\text{cen}}(r) = v_\phi^2/r = V_{\text{rot}}^2 R^2/r^3 \). In the latter equality, \( V_{\text{rot}} \)
the equatorial rotation speed at the stellar surface radius \( r = R \), with a fixed specific
angular momentum \( v_\phi r \) in the outflowing wind. In terms of the critical rotation fraction
\( \Omega = V_{\text{rot}}/\sqrt{GM/R} \), the reduction in the effective surface gravity by a factor \( 1 - \Omega^2 \)
leads to an increased mass loss rate, but since the centrifugal acceleration falls off faster than
the \( 1/r^2 \) of gravity, the wind flow speed is lower. Indeed, for \( \Omega > 0.75 \), the flow can
even stagnate (Madura et al. 2007), and/or switch to a slow-acceleration solution (Curé
2004). Curé et al. 2005 proposed that the lower speed and thus higher density of these
low-acceleration, 1D solutions could explain enhanced equatorial density inferred for
Be or B[e] stars.

But more general 2D analyses show the importance of accounting for latitudinal
effects. At co-latitude \( \theta \), the centrifugally reduced, effective surface gravity is

\[
g_{\text{eff}}(\theta) = \frac{GM}{R^2} \left( 1 - \Omega^2 \sin^2 \theta \right), \tag{5}
\]

where for simplicity we have ignored any rotational distortion of the surface radius \( R \).
This allows one to rewrite the standard CAK mass loss rate scaling law (3) in terms
Figure 2. Contours of density in stellar wind from a star rotating at 75% of critical rate, plotted vs. colatitude $\theta$ and radius $r$, spaced logarithmically with two contours per decade. The superposed vectors represent the latitudinal velocity, with the maximum length corresponding to a magnitude of $v_0 = 100$ km/s. The three panels show the cases (a) without nonradial forces or gravity darkening, (b) with nonradial forces but no gravity darkening, and (c) with both nonradial forces and gravity darkening.

Of surface values of the mass flux $\dot{m} = \rho v$, radiative flux $F$, and effective gravity $g_{\text{eff}}$, relative to corresponding polar ($\theta = 0$) values $\dot{m}_o$, $F_o$, and $g_o = GM/R^2$,

$$\frac{\dot{m}(\theta)}{\dot{m}_o} = \left[ \frac{F(\theta)}{F_o} \right]^{1/\alpha} \left[ \frac{g_{\text{eff}}(\theta)}{g_o} \right]^{1-1/\alpha}.$$  \hspace{0.5cm} (6)

If the radiative flux is taken to be constant in latitude, $F(\theta) = F_o$, we find

$$\frac{\dot{m}(\theta)}{\dot{m}_o} = \left[ 1 - \Omega^2 \sin^2 \theta \right]^{1-1/\alpha}; \quad F(\theta) = F_o.$$  \hspace{0.5cm} (7)

Since the exponent $1 - 1/\alpha$ is negative, the mass flux is maximum near the equator (where $\sin \theta \to 1$).

However, if we account for the von Zeipel (1924) scaling $F(\theta) \sim g_{\text{eff}}$ for equatorial gravity darkening, we find

$$\frac{\dot{m}(\theta)}{\dot{m}_o} = 1 - \Omega^2 \sin^2 \theta \quad ; \quad F(\theta) \sim g_{\text{eff}}(\theta)$$  \hspace{0.5cm} (8)

now giving the mass flux an equatorial minimum, with the maximum at the pole!

Recalling that the wind terminal speed tends to scale with the surface gravity through the escape speed, this analysis also predicts a latitudinally varying wind speed maximum is proportional to a centrifugally reduced, effective escape speed,

$$v_{\infty}(\theta) \sim v_{\text{esc}} \sqrt{1 - \Omega^2 \sin^2 \theta},$$  \hspace{0.5cm} (9)

which is also maximum near the pole. The overall latitudinal variation of wind density is then obtained from $\rho(\theta) \sim \dot{m}(\theta)/v_{\infty}(\theta)$.

More generally, the wind from a rotating star can also flow in latitude as well as radius, requiring then a fully 2D model. In this regard, a major conceptual advance was
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Figure 3. Illustration of the origin of poleward component of the radiative force from a rotating star. Left: The poleward tilt of the radiative flux arising from the oblateness of the stellar surface contributes to a poleward component of the driving force. Right: Since the wind speed scales with surface escape speed, the lower effective gravity of the equator leads to a slower equatorial speed. The associated poleward increase in speed leads to a poleward tilt in the velocity gradient, and this again contributes to a poleward component of the line force.

the development of the elegantly simple “Wind Compressed Disk” (WCD) paradigm by Bjorkman & Cassinelli (1993). They noted that, like satellites launched into earth orbit, parcels of gas gradually driven radially outward from a rapidly rotating star should remain in a tilted ‘orbital plane’ that brings them over the equator, where they collide to form compressed disk. Initial 2D hydrodynamical simulations (Owocki, Cranmer, and Blondin 1994) generally confirmed the basic tenets of the WCD model (figure 2a), with certain detailed modifications (e.g., infall of inner disk material). But later simulations (Owocki, Cranmer, and Gayley 1996) that account for a net poleward component of the line-driving force (figure 3) showed that this can effectively reverse the equatorward drift, and so completely inhibit formation of any equatorial compressed disk (figure 2b). Indeed, when equatorial gravity darkening is taken into account, the lower mass flux from the equator makes the equatorial wind have a reduced, rather than enhanced, density (figure 2c).

The net upshot then is that a radiatively driven wind from a rapidly rotating star is predicted to be both faster and denser over the poles, instead of the equator. This may help explain spectroscopic and interferometric evidence that the current-day wind of the extreme massive star η Carinae is faster and denser over the poles, leading to a prolate shape for its dense, optically thick ‘wind photosphere’ (Smith et al./ 2003; van Boekel et al./ 2003; Groh et al./ 2010). Extensions of the rotational scalings to the case of continuum-driven mass loss could also explain the bipolar shape of the Homunculus nebula (Owocki, Shaviv and Gayley 2004).

2.3. Centrifugal Ejection of Viscous Decretion Disk

An important general conclusion here is thus that the radiative driving that propels outflowing stellar winds is inherently ill-suited for producing the dense, nearly stationary
equatorial disks of Be stars. Instead the near-critical rotation of Be stars (Townsend et al. 2004) has lead to the notion that such disks are likely produced by the centrifugal ejection of material into Keplerian orbit near the equatorial surface, with then an outward viscous diffusion of mass and angular momentum leading to a radially extended **Viscous Decretion Disk** (Lee, Osaki & Saio 1991). In contrast to the well-defined radiative lifting of stellar wind from a deep gravitational potential, such decretion represents a kind of spill-over from a critical surface, somewhat akin to Roche-Lobe overflow in mass-exchange binaries. Whereas wind driving can perhaps be likened to the suction through a straw from a partly full glass, with the mass flux depending directly on the strength of the suction force, centrifugal decretion is more akin to spillage from nearly full glass, with small perturbations causing a random, chaotic mass overflow. Instead of depending on the strength of the external driving, the overall level of centrifugal mass loss depends on the internal mechanisms keeping the stellar rotation near critical.

For example, if stellar evolution leads to change in moment of inertia at a rate $\dot{I}$, then keeping the star at or below the critical angular rotation rate $\omega_c \equiv 2\pi \sqrt{GM/R^3}$ requires a net shedding of angular momentum $\dot{J} = \dot{I}\omega_c$. For a viscous Keplerian disk with outer radius $R_o$, the angular momentum loss associated with a steady disk mass loss is $M_d \sqrt{GM R_o}$, requiring then a mass loss rate (Krticka et al. 2011),

$$M_d = \frac{\dot{I}}{\sqrt{R^3 R_o}}. \quad (10)$$

The contribution by Granada et al. in these proceedings discusses stellar evolution models using this centrifugal scaling for disk mass loss. Other contributions (e.g. by Bjorkman, Jones, Okazaki, Haubois) discuss the detailed observational diagnostic of the extended Keplerian disk that results from viscous transport of mass outward from the source near the equatorial surface.

During quiescent periods when mass ejection is small or non-existent, viscous diffusion can also lead to infall back onto the stellar surface, and so a gradual, inside-out decay of the disk that can be tracked by detailed observations. But in addition, the inner regions of the disk can be ablated by the line-driving from the stellar radiation...
field. As illustrated in figure 4, the inherent shear in the orbital speed of the Keplerian disk means that non-radial rays through the disk see a line-of-sight velocity gradient $\frac{dv_z}{dz}$ that can be comparable to the radial velocity gradient $\frac{dv_r}{dr}$ that is essential to line-desaturation for driving a stellar wind. The resulting mass ablation flux from the inner disk follows a scaling comparable to that of the star’s line-driven wind (Gayley, Owocki & Cranmer 1999). For a disk with mass $M$, the associated disk decay timescale is thus of order $M/M_{\text{CAK}}$, with typical values measured in years. Such ablation thus could play a role in the gradual (multiyear) decline of Balmer emission in Be stars, but is too weak to explain the more rapid, month-timescale decay seen in 28 CMa (Carciofi et al. 2012).

3. Stellar Magnetospheres

3.1. Dynamical Magnetospheres from Magnetically Trapped Wind

In the subpopulation of massive stars with strong, large-scale magnetic fields, the radiatively driven wind outflow can be channeled and even trapped in closed magnetic loops. Detailed MHD simulations are discussed by the contribution by ud-Doula, so
here we just focus on some semi-analytic scaling relations for the associated mass accumulation, and how this depends on magnetic field strength and stellar rotation.

The relevant competition between field and wind can be characterized by the wind magnetic confinement parameter \( \eta_* \equiv B_{eq}^2 R^2 / M v_{\infty} \), where \( B_{eq} \) is the surface field at the magnetic equator, and \( M \) and \( v_{\infty} \) are the expected wind mass loss rate and terminal speed in the absence of a field. For a strong dipole field, the wind is magnetically confined out to an Alfvén radius \( R_A \approx \eta_*^{1/4} R \). Magnetic loops with footpoints at co-latitude \( \theta_* \) and loop apex \( R_m = R / \sin^2 \theta_* < R_A \), will thus trap the wind upflow from each footpoint, forcing a shock collision near the loop apex. In the absence of any significant rotation, this compressed, cooled material falls back toward the stellar surface on a gravitational free-fall timescale. As discussed in the contribution by Sundqvist, in slowly rotating magnetic O-stars with strong winds, the density in this Dynamical Magnetosphere (DM) can be sufficient to produce the observed Balmer emission, and its slow modulation with rotational phase.

As a complement to models based on full MHD simulations of the complex outflow/infall structure, figure 5 outlines a relatively simple Analytic Dynamical Magnetosphere (ADM) model for the mean, time-averaged density and velocity. For loop with footpoint co-latitude \( \theta_* \), the radial projection of the field has a tilt \( \mu_* \equiv \cos \theta_* \), reducing the wind mass flux fed into the loop by this same tilt factor (Owocki & ud-Doula 2004). Once mass from opposite footpoints collides near the loop apex, it free-falls back along the field lines. Since both the flow and field have zero divergence, mass conservation sets the time-averaged density in terms of this source mass flux, local free-fall speed, and field strength. The factor \( f_w \) accounts for a finite width \( w \) for mass deposition about the loop apex, avoiding the infinite density that would occur if all material started with zero speed right at the loop top. The lower panels of figure 5 show good general agreement with the time-averaged quantities from a full MHD simulation. Initial application of this ADM scaling to computation of Balmer emission also shows remarkably good agreement with both observations, and with corresponding MHD models (Sundqvist et al., in preparation).

3.2. Centrifugal Magnetospheres

The additional dynamical effect of rotation depends on the critical rotation fraction \( \Omega \). Insofar as the field is strong enough to maintain rigid-body rotation, we can define an associated Kepler co-rotation radius \( R_K = \Omega^{-2/3} R \), at which the centrifugal force for rigid-body body rotation exactly balances the gravity in the equatorial plane. Below \( R_K \), any magnetically trapped material will again fall back toward the star, as in the above DM model. But in the case of a strong field with \( R_A > R_K \), material trapped in the region \( R_K < r < R_A \) has an outward centrifugal force that exceeds the local gravity, providing support for a Centrifugal Magnetosphere (CM), and so now allowing the wind to feed a secular buildup in density.

For full 2D MHD simulations of magnetic winds with field-aligned rotation (ud-Doula et al. 2008), the mosaic of color plots in figure 6 illustrates the time evolution of radial mass distribution \( dm/dr \) within 10° of the equator, for magnetic confinement parameters ranging from \( \eta_* = 10^{1/2} \) to \( 10^3 \) (left to right), and for rotation fractions \( \Omega = 0, 1/4 \) and \( 1/2 \) (bottom to top). Within each panel, the vertical spans a radius range \( r/R \approx 1 - 5 \), while the horizontal covers a time span of 3 Msec after introduction of the field. Note that for moderately fast rotation \( \Omega = 1/2 \) and strong magnetic confinement \( \eta_* > 100 \), there is a systematic buildup of high density material in the regions between \( R_K \).
and $R_A$, whereas material in the region near and above $R_A$ streams outward, sometimes through quasi-regular intervals of "centrifugal breakout".

For the limit of arbitrarily strong field, Townsend & Owocki (2005; TO-05) developed a semi-analytic Rigidly Rotating Magnetosphere (RRM) model for accumulation of material at the minimum of effective gravitational+centrifugal potential of each magnetic loop. Applicable even for the case of non-aligned rotation, this RRM model has proven remarkably successful in reproducing the rotational phase variation of Balmer line in magnetic B-stars like $\sigma$ Ori E (Townsend et al. 2005). By analyzing the point at which the net centrifugal force of built-up material overwhelms the confining effect of magnetic tension, TO-05 were also able estimate a characteristic "breakout" time for magnetospheric mass loss,

$$t_b \approx 0.1 \eta_* t_{ff},$$

where the disk confinement parameter, $\eta_*$, is replaced with the surface escape speed $v_{esc} = \sqrt{2GM/R}$, and $t_{ff} = R/v_{esc}$ is a characteristic free-fall time. ud-Doula et al. 2008 showed this scaling matches quite well the breakout timescales seen in their MHD simulations. The associated total magnetospheric mass is then

$$M_{cm} = \dot{M} t_b \approx 3 \times 10^{-9} M_{\odot} \frac{B_3^2 R_1^{12}}{g_4}$$

where $B_3 \equiv B/10^3$ G, $R_1 \equiv R/10^{12}$ cm, and $g_4 \equiv g/10^4$ cm/s$^2$. Note that the wind mass loss rate itself cancels in this net scaling for total mass. For the prototypical magnetic B-star $\sigma$ Ori E, one finds $t_b \approx 100$ years and $M_{cm} \approx 10^{-7} M_{\odot}$. As discussed in contributions by Townsend and by Carciofi, this seems to exceed by 1-2 orders of magnitude the magnetospheric mass inferred from observed linear polarization. Current efforts are thus focused on what other MHD or plasma leakage mechanisms might be responsible for limiting the buildup of magnetospheric mass.
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References