Outflow-Confined H II Regions and the Formation of Massive Stars by Accretion

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Abstract. If massive stars form by disk accretion, then bipolar outflows should be generated as in the case of low-mass star formation. High accretion rates lead to high outflow rates and make the wind density very large near the protostar. We therefore predict that massive protostars have very small, jet-like H II regions confined by their outflows; we identify these H II regions with “hypercompact” (≤ 0.01 pc) H II regions. Their lifetime is approximately equal to the accretion timescale (∼ 10⁵ yr), much longer than the sound-crossing and dynamical timescales of the ionized region. We present an analytic description of the density distribution of the outflow, relate the overall mass loss rate to the accretion rate, and normalize to values appropriate to massive protostars. For a given ionizing luminosity, we calculate the extent of the H II region and its radio spectrum. A detailed comparison is made with observations of radio source “I” in the Orion Hot Core. The spectra and morphologies of many other sources are broadly consistent with this model. We argue that confinement of ionization allows disk and equatorial accretion to continue up to high masses, potentially overcoming a major difficulty of standard accretion models.

1. Introduction

The basic mode of massive star formation is a matter of debate. Extensions of core collapse models (e.g. Shu, Adams, & Lizano 1987) from low to high masses have been criticized on the grounds that radiative feedback is likely to disrupt accretion and that massive stars are known to form in the crowded environments of star clusters, where dynamical interactions between forming stars may be important. These concerns have stimulated the development of models that involve direct stellar collisions as well as competitive accretion of material in proto star clusters (Bonnell, Bate, & Zinnecker 1998; Bonnell & Bate 2002). However, many of the difficulties facing standard accretion models may be resolved once the appropriate initial conditions of massive star formation are allowed for (McKee & Tan 2003). In particular, the high pressures of these regions mean that the initial cores, if close to equilibrium, will be very dense with small volume filling factor, short collapse times, and high accretion rates.
Here we investigate the consequences of high accretion rates on hydromagnetic outflows, and the ability of these outflows to confine ionizing radiation from the protostar. The radio emission from these H II regions is a useful diagnostic of the massive star formation mechanism. The outflow may shield the collapsing gas core from the brunt of the protostar’s radiative feedback.

2. Protostellar Outflow Density

We wish to determine the density distribution in the winds (we shall use the terms winds and outflows interchangeably) from protostars and their accretion disks. Such winds are believed to be hydromagnetic in origin, at least for low-mass stars. Two classes of models have been developed: disk winds, in which the wind flows along field lines threading the accretion disk (e.g. Königl & Pudritz 2000), and X-winds, in which the wind is launched from the interaction region between the stellar magnetic field and the inner edge of the accretion disk (e.g. Shu et al. 2000). Here we shall develop an approximate formulation for the density that applies to both models. We shall not attempt to describe the region where the wind accelerates, since that depends on the details of the acceleration mechanism. Instead, we shall determine a characteristic density near the disk surface, and a general expression for the wind density far from the disk.

Consider a thin, Keplerian accretion disk (Fig. 1). We assume that the wind extends from $\varpi_{c0}$ to $\varpi_{\text{max0}}$, where $\varpi \equiv r \sin \theta$ is the cylindrical radius and $\varpi_0$ is measured in the plane of the disk. Let $\rho_0(\varpi_0)$ be the density for the streamline originating at $\varpi_0$ at the point at which the vertical velocity $v_z$ reaches the Keplerian velocity $v_K(\varpi_0)$. We assume that $v_z$ reaches $v_K(\varpi_0)$ at...
\(\varpi \simeq \varpi_0\), and that the density there varies as a power-law in radius,

\[
\rho_0(\varpi_0) = \rho_0 \left(\frac{\varpi_0}{\varpi_c}\right)^{-q} = \rho_0 \varpi_0^{-q}.
\]

(1)

The classic Blandford \& Payne (1982) self-similar wind has \(q = 3/2\); the cylindrical self-similar wind models of Ostriker (1997) have \(1/2 \leq q \leq 1\). (Note that the Ostriker winds are prevented from expanding laterally by an external pressure, and have magnetosonic Mach numbers \(M_F < 1\).) Since the Keplerian velocity can be expressed as \(v_K = v_{Kc} \varpi_0^{-1/2}\), the mass loss rate in the wind from both sides of the disk is

\[
\dot{m}_w = 4\pi \int_{\varpi_c}^{\varpi_{\text{max}}} \varpi_0 d\varpi_0 \rho_0(\varpi_0) v_K(\varpi_0) = 4\pi \rho_0 v_{Kc} \varpi_c^2 I_w \varpi_0,
\]

(2)

where \(I_w \varpi_0 \equiv (x_{0\text{max}}^{3/2-q} - 1)/(3/2 - q)\). For an X-wind, assume that the wind emerges from a narrow band of width \(\Delta x_{0\text{max}}\), so that \(x_{0\text{max}} = 1 + \Delta x_{0\text{max}}\); then \(I_w \varpi_0 = \Delta x_{0\text{max}}\). For both disk winds and X-winds, the density in the wind near the disk is then

\[
\rho_0(\varpi_0) = \left(\frac{\dot{m}_w}{4\pi v_{Kc} \varpi_c I_w \varpi_0}\right) \varpi_0^{-q}.
\]

(3)

Since the flow near the disk is approximately one-dimensional, the density at a point where the vertical velocity differs from the local Keplerian velocity is \(\rho = \rho_0 [v_K(\varpi_0)/v_z]\).

Next, consider the wind far from the disk. The flux-to-mass ratio is constant along streamlines, so \(d\Phi/d\dot{m}_w = B_z/(\rho v_z) = B_{z0}/(\rho_0 v_K)\). To simplify the discussion, we assume that the wind is self-similar, with the Alfvén velocity \(v_A\) proportional to the Kepler velocity \(v_K\). In that case, \(B_{z0} \propto x_0^{(q+1)/2}\) and

\[
\frac{B_z}{\rho v_z} = \left(\frac{B_{zc}}{\rho_c v_{zc}}\right) x_0^{q/2}\]

(4)

(e.g., Ostriker 1997), where \(\varpi_c\) is the inner edge of the wind. We assume that the wind has expanded so that \(\varpi \gg \varpi_0\). In that case, the azimuthal velocity is small, and the toroidal wrapping of the field gives \(B_z/B_\phi = v_z/(\Omega_0 \varpi)\), where \(\Omega_0\) is the angular velocity of the footpoint. The field is approximately force free, and since the azimuthal field dominates, we have

\[
B_\phi = B_{\phi c}/x,
\]

(5)

where \(x \equiv \varpi/\varpi_c\). We then obtain

\[
B_z = \left(\frac{B_{\phi c} v_{zc}}{\Omega_0 \varpi_c}\right) \frac{x_0}{x^2} \equiv B_{zc} \frac{x_0}{x^2},
\]

(6)

so that from equation (4)

\[
\rho = \rho_c \frac{x_0^{3-q}/2}{x^2}.
\]

(7)
The result that $\rho \propto \varpi^{-2}$ was first found by Shu et al. (1995) for the particular case of X-winds, and was derived more generally by Ostriker (1997) and by Matzner & McKee (1999).

Since the wind is self-similar, the velocity scales with the footpoint velocity, $v_z = v_z c_{sc}^{-1/2}$. Evaluating the mass loss rate in the wind far from the disk determines the density at the inner edge of the wind, $\rho_c$. The density distribution far from the disk is then

$$\rho = \frac{\dot{m}_w}{4\pi v_z c_{sc}^2 I_w \infty} \frac{x_0^{(3-q)/2}}{x^2},$$

(8)

where $I_w \infty = \int_1^{x_{\text{max}}} dx x_0^{1-q/2}/x$. We assume that $x_0 = x^\beta$ for some value of $\beta$; the coefficient in this expression is unity because the streamline that originates at $x_0 = 1$ has $x = 1$. Evaluating the integral, we find

$$I_w \infty = \left(\frac{1-x_0^{1-q}}{1-\frac{q}{2}x_0^{1-q}}\right) \ln x_{\text{max}} - \ln x_0 \max,$$

(9)

Since X-winds emanate from a narrow range of radii, we set $x_0 \approx 1$ for them; X-winds therefore have $I_w \infty = \ln x_{\text{max}}$.

Combining these results, we express the density in the wind as

$$\rho = \frac{\dot{m}_w f(x_0, x)}{4\pi v_z c_{sc}^2 I_w},$$

(10)

where

$$f(x_0, x) = \begin{cases} \frac{x_0^{-q}}{x_0^2} & x \simeq x_0, \\ \frac{x_0^{(3-q)/2}}{x^2} & x \gg x_0. \end{cases}$$

(11)

At the disk surface, $v_z c = v_{Kc}$ and $I_w = I_w 0$, whereas far from the disk $v_z c > v_{Kc}$ and $I_w = I_w \infty$. Note that for $q = 1$, we have $f = x_0/x^2$ in both cases, so that in this case

$$I_w = 2(x_0^1/2 - 1) \left(\frac{\ln x_{\text{max}}}{\ln x_0 \max}\right) \quad (q = 1),$$

(12)

both at the disk and far from the disk. For X-winds, $I_w \infty = \ln x_{\text{max}}$ reduces to $I_w 0 = \Delta x_0 \max$ at the disk surface, so we can take $I_w \approx \ln x_{\text{max}}$ everywhere for X-winds.

We evaluate equation (10) numerically in terms of the density of hydrogen nuclei, $n_w \equiv \rho/\mu_H$, where $\mu_H = 2.34 \times 10^{-24}$ g is the mean mass per hydrogen for an assumed helium abundance 10% of hydrogen. We also write the mass loss rate in the wind as $\dot{m}_w = f_w \dot{m}_*$, where $\dot{m}_*$ is the accretion rate onto the star. For example, Matzner & McKee (1999) estimate $f_w \approx \frac{1}{2}$. We then find

$$n_w = \frac{\rho}{\mu_H} = 9.52 \times 10^{10} \left[\frac{f_w f(x_0, x)}{v_{z sc}^7 I_w}\right] \left(\frac{\dot{m}_*}{10^{-4} \text{ M}_\odot \text{ yr}^{-1}}\right) \left(\frac{1 \text{ AU}}{c_{sc}}\right)^2 \text{ cm}^{-3},$$

(13)

where $v_{z sc}^7 \equiv v_{z sc}/(10^7 \text{ cm s}^{-1})$. 

2.1. Axial Wind Cavity

A question of critical importance for H II regions associated with massive protostars is the shape of the cavity that forms along the axis, since ionizing photons can propagate freely through this cavity. There are several effects that govern its shape: For disk winds, the flow is launched at an angle greater than 30° from the rotation axis, and the hoop stresses of the toroidal field must overcome the inertia of the outflow in order to begin to close the cavity. For X-winds, poloidal field lines from the protostar determine the shape of the cavity (Shu et al. 1995), and the same appears to be true for disk winds (E. Ostriker 2002, private comm.). In either case, once the protostar becomes massive enough to generate a strong main sequence wind, the pressure due to this wind will tend to open the cavity further.

For simplicity, we shall focus on X-winds, for which \( x_0 = 1 \) as discussed above. The magnetic flux in the cavity is \( \Phi_c = \pi B_c \omega_c^2 \), whereas that in the wind far from the disk can be evaluated with the aid of equation (6),

\[
\Phi_w = \frac{2\pi \omega_c v_{zc} B_{\phi c}}{\Omega_0} \ln \frac{\omega_{\text{max}}}{\omega_c} \, .
\]  

(14)

Since the maximum extent of the wind \( \omega_{\text{max}} \) enters only logarithmically, we approximate it as \( \omega_{\text{max}} \simeq r_c \). Let \( v_{zc*} \equiv v_{zc}/v_{Kc} = v_{zc}/\Omega_0 \omega_0 \) be the normalized velocity; in the numerical example worked out by Shu et al. (1995), \( v_{zc*} = 2.1 \). Magnetic pressure balance requires that the field in the cavity, \( B_c \), equal the field at the inner edge of the wind, \( B_{\phi c} \) (which is much greater than \( B_z \) there). Furthermore, in the X-wind model, the cavity flux and the wind flux are comparable, \( \Phi_c \simeq \Phi_w \). These relations give an equation for the radius of the cavity,

\[
\omega_c = 2v_{zc*} \omega_c \ln \frac{r_c}{\omega_c} \, .
\]  

(15)

which is valid far from the disk. We develop an approximate solution of equation (15) following the approach of Matzner & McKee (1999), and then add a term in the logarithm to make it valid at the disk surface:

\[
\omega_c \simeq 1.4v_{zc*} \omega_c \ln \left[ \frac{r_c}{1.4v_{zc*} \omega_c} + e^{1/(1.4v_{zc*})} - \frac{1}{1.4v_{zc*}} \right] \, .
\]

(16)

This form remains valid for regions near the disk even when \( v_{zc*} \neq 1 \).

We now evaluate the factor \( \ln x_{\text{max}} \) in \( I_{\infty} \), which enters the expression for the density far from the disk:

\[
\ln x_{\text{max}} = \ln \frac{\omega_{\text{max}}}{\omega_c} \simeq \ln \frac{r_c}{\omega_c} \simeq 0.7 \ln \left( \frac{r_c}{1.4v_{zc*} \omega_c} \right) \, ,
\]

(17)

where we have assumed \( r_c \gg \omega_c \). We generalize to arbitrary values of \( r_c \) by adding a term in the logarithm,

\[
\ln x_{\text{max}} \simeq 0.7 \ln \left( \frac{r_c}{1.4v_{zc*} \omega_c} + x_{\text{max}}^{1/0.7} - \frac{1}{1.4v_{zc*}} \right) \, .
\]

(18)

Our result for the density agrees with that obtained by Shu et al. (1995) to within about 20% for \( 10 \lesssim r_c/\omega_c \lesssim 10^5 \).
Figure 2. Left: Outflow isodensity contours from a $20M_\odot$ protostar, with $r_\ast = 16R_\odot$, $\gamma_c = 2r_\ast$, $\dot{m}_\ast = 10^{-4}M_\odot\,\text{yr}^{-1}$, $\dot{m}_w = \dot{m}_\ast/3$, $v_{z,c} = 2.1$, and $q = 1$. Right: H II region growth with ionizing luminosity, $S$.

3. Outflow-Confined H II Regions and Application in Orion

Consider a beam of ionizing radiation from the protostar that intercepts the cavity boundary at a distance $r_c(\theta)$. Let $S$ be the ionizing photon luminosity, so that the rate at which ionizing photons cross an element of area $dA$ is $S dA/4\pi r_c^2$. This flux of ionizing photons is able to balance the recombinations of the ions in the wind out a radius $r_i$ given by

$$ \frac{S dA}{4\pi r_c^2} = \frac{dA}{r_c^2} \int_{r_c}^{r_i} \alpha^{(2)} n_w r^2 dr $$

(19)

where $\alpha^{(2)}$ is the radiative recombination rate to the excited states of hydrogen and the gas is assumed to be fully ionized inside $r_i$. This equation ignores the influx of neutrals from the wind and the absorption of ionizing photons by dust, which both make the H II region smaller. However, the effect of neutral influx is negligible for the specific numerical models we consider for the Orion source (below). If the outflow originates from close to the protostar then most dust grains are likely to have been destroyed. We present an analytic description of the H II region elsewhere. For now we solve equation (19) numerically (Fig. 2).

For small ionizing luminosities only a thin skin of material around the cavity is ionized. As the luminosity increases the H II region expands out to large angles from the rotation axis and greater distances from the protostar. At a critical ionizing luminosity the H II region breaks out at angles $\sim 30^\circ$ from the rotation axis. Soon the entire solid angle interior to this becomes ionized, while in the equatorial directions the gas can remain neutral for a much longer time.

We now apply this model to the massive protostar powering the Orion hot core. At a distance of 450 pc, this is the closest example of a massive star in formation. The core is self-luminous ($L_{\text{bol}} \sim 1 - 5 \times 10^4 L_\odot$, Gezari et al. 1998; Kaufman et al. 1998). A weak radio continuum source (“I”) (e.g. Menten & Reid 1995), located within a few arcseconds of the core center, as traced by dust
and gas emission (Wright, Plambeck & Wilner 1996), almost certainly pinpoints the location of the massive protostar. SiO emission forms a “bow-tie” feature centered on “T”, that may be an inclined or flared disk (Wright et al. 1995). Perpendicular to this, a large scale, wide-angle bipolar outflow extends to the NW and SE of the core (Chernin & Wright 1996; Greenhill et al. 1998). Chernin & Wright modeled the flow axis as being inclined at 65° to our line of sight (Fig 3a). At 22 GHz, source “T” appears elongated (0″.145 by < 0″.085, Menten 2002, private comm.) parallel to the large scale outflow axis.

We model the thermal bremsstrahlung emission from source “T” assuming we are viewing the outflow at $\theta_{\text{view}} = 65^\circ$ and that the temperature of the ionized gas is $10^4$ K. As we assume negligible density in the wind cavity, the H II region is formally infinite in length. However, the emission measure rapidly decreases along the jet (as $r^{-3}$). To compare to observations at 8.4 GHz (Menten & Reid 1995), 14.9 GHz (Felli et al. 1993), 22.3 GHz (Menten 2002, private comm.), 43.1 GHz (Menten & Reid 1995), and 86 GHz (Plambeck et al. 1995), we include emission from scales about one third larger than the quoted beam sizes, i.e. 0.3″, 0.2″, 0.2″, 0.3″, and 0.5″, respectively. We extend to higher frequencies with the 0.5″ size. Upper limits at 98 GHz and 218 GHz are from Murata et al. (1992) and Blake et al. (1996), respectively. A more detailed comparison of predicted surface brightness profiles with data will be presented elsewhere. The radiative transfer calculation is accomplished by dividing the H II region, which is always axisymmetric in our modeling, into many volume elements and then calculating optical depths along ray paths to the observer. We increase the resolution to achieve convergence.

Our fiducial protostellar model is based on constraints derived from the total luminosity (McKee & Tan 2003). We set $m_\ast = 20M_\odot$ and $\dot{m}_\ast = 10^{-4}M_\odot \text{ yr}^{-1}$. We assume $\dot{m}_w = 0.33\dot{m}_\ast$ and $v_w = 2.1v_K$, which are reasonable choices for certain classes of disk winds and X-winds. Given the uncertainties in the ionizing luminosity of massive protostars, we consider a range of values (Fig 3b). A protostar with $S = 2 \times 10^{17} \text{ s}^{-1}$ provides a good match to the observations, and this luminosity is consistent with the predictions of models of massive protostellar evolution (Tan 2002).

While not unique, the model we have presented for the protostar in the Orion hot core has a bolometric luminosity consistent with infrared observations, an accretion rate that is based on theoretical expectations for massive gas cores collapsing in high pressure environments, a protostellar size that is calculated given this accretion rate, and an ionizing luminosity consistent with theoretical models and able to create an H II region with the observed radio spectrum. The geometry of the system was chosen so that this inner protostellar outflow naturally connects with the larger scale flow from OMC-1.

However, the model is nonetheless highly idealized: we have assumed that the density of material in the wind cavity is negligible, ignored the possible influence of a conventional stellar wind on the cavity geometry, treated the propagation of ionizing photons along sectors as being independent, ignored dust (although much of this may have been destroyed in the protostellar radiation field), and made simple parameterizations for the connection between accretion and outflow generation. Some of these issues will be addressed in a future paper.
4. Implications for Massive Star Formation

The class of outflow-confined H II regions is expected to be relevant to many, if not all, massive protostars. During the earlier stages of evolution the H II region is confined close to the jet axis, resulting in a relatively weak radio source, such as source ‘I’ in Orion. At somewhat later stages the H II region is able to extend to greater distances, although the peak of the emission measure is expected to remain concentrated close to the protostar. This evolutionary sequence is relevant to a host of observed sources, such as Cepheus A HW2 (Torrelles et al. 1996), IRAS 16547-4247 (Garay et al. 2003), IRAS 18162-2048 (Gómez et al. 2003), IRAS 20126+4104 (Zhang, Hunter, & Sridharan 1998), G192.16-3.82 (Shepherd, Claussen, & Kurtz 2001), AFGL 2591 (Trinidad et al. 2003), and AFGL 5142 (Zhang et al. 2002). In many of these systems jet-like morphology is clearly observed.

The small observed sizes of the H II regions may be due either to the limited sensitivity of particular observations, or to a real physical confinement by the outflow. In the latter case the sound crossing time (e.g. \( t_s = \frac{r_{\text{HII}}}{c_s} = 1000[r_{\text{HII}}/0.01 \text{ pc}] [c_s/10 \text{ km s}^{-1}]^{-1} \text{ yr} \)) and flow crossing times (typically \( \sim 100 \) times shorter than \( t_s \)) are both unrelated to and much shorter than the true lifetime of the system, which is set by the accretion timescale of the protostar (\( \sim 10^5 \text{ yr} \), McKee & Tan 2002) and the timescale for evolution of the ionizing luminosity. This bears upon estimates of the expected number of observed sources, in a similar manner to that of the classic “lifetime” problem of ultracompact H II regions (Wood & Churchwell 1989).

One prediction of the outflow-confined model for compact H II regions is the presence of very broad recombination lines. The outflow velocities from massive protostars are expected to be at least of order the escape velocity, \( v_{\text{esc}} = 620(m_*/10M_\odot)^{1/2}(r_*/10R_\odot)^{-1/2} \text{ km s}^{-1} \). Our fiducial model for source ‘I’ in Orion predicts that the fastest velocities are about 1000 km s\(^{-1}\). Atomic line
profiles from HH objects much further out in the outflow from OMC-1 show line-of-sight widths of several hundred km s\(^{-1}\) (Taylor et al. 1986), which would be consistent with the presence of a \(\sim\) 1000 km s\(^{-1}\) flow from the protostar. Of course as this flow interacts with nearby dense molecular gas, it will create outflows with a wide range of lower velocities, as are observed (Chernin & Wright 1996; Stolovy et al. 1998).

Outflow-confined H II regions have broader implications for the massive star formation paradigm. First of all, their properties are a quantitative diagnostic of the accretion scenario for massive star formation: radio spectra allow estimation of the outflow density near the protostar. Theoretical models of outflows from accretion disks and protostars then allow a direct connection between these observations and the accretion physics. If the morphologies of outflows very close to the protostar align with larger scale features, then this would suggest that massive stars form from accretion disks that are relatively stable in terms of their orientation. Stellar collisions would be expected to disrupt such configurations. We would argue that the current observational evidence, particularly in the Orion hot core, supports formation models based on accretion rather than stellar collisions.

High accretion rates to massive protostars naturally result in high outflow rates, which create dense and collimated gas structures close to the star. These in turn should help to shield the bulk of the accreting material from radiative protostellar feedback. In particular, ionizing photons are degraded in the H II region, never reaching the disk in the equatorial plane. Models of disk photoevaporation (Hollenbach et al. 1994) are therefore inhibited during the main accretion phase. By the time the protostar has reached a stage where the ionizing luminosity is important \((m_\ast \gtrsim 10M_\odot)\) we expect the outflow to have significantly altered the density distribution of the initial gas core in its polar regions (Matzner & McKee 2000). Radiant energy from the protostar will thus tend to escape along these directions, rather that in the equatorial plane. This is similar to the so-called “flashlight effect” discussed by Yorke & Bodenheimer (1999), although they considered that arising only from the presence of a disk. Thus radiation pressure feedback on infall may be substantially weakened, to an extent greater than that resulting from rotation alone (Nakano 1989; Jijina & Adams 1996). The evolution of the size of the outflow-enhanced flashlight effect may play an important role in determining the maximum stellar mass.

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