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STAR FORMATION AND THE NATURE OF BIPOLAR OUTFLOWS

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ABSTRACT

We present a simple physical model for the bipolar molecular outflows that frequently accompany star formation. Our model forges an intrinsic link between the bipolar flow phenomenon and the process of star formation, and it helps to explain many of the systematics known for existing sources.

Subject headings: hydrodynamics — nebulae: internal motions — stars: formation — stars: mass loss — stars: pre–main-sequence

1. INTRODUCTION

Since the discovery of bipolar CO outflows a decade ago (Snell, Loren, & Plambeck 1980), it has become clear that the generation of energetic, cold outflows of molecular gas around young stellar objects in molecular clouds is intimately related to the processes of star formation and early stellar evolution (e.g., Bally & Lada 1983; see also the reviews of Lada 1985; Welch et al. 1985; Bally 1987; Snell 1987; Fukui 1989; Rodríguez 1991). Yet, the nature of the *physical* connection between the bipolar outflow phenomenon and the star formation process has remained a mystery. Indeed, it has not yet been established when in the formation and early evolution of a star the onset of the bipolar outflow phase occurs. Recently, however, observational studies have begun to provide evidence indicating that outflow activity may be initiated during the protostellar (infall) phase of evolution (Lada 1988; Margulis, Lada, & Snell 1989; Snell et al. 1988; Berrilli et al. 1989; Wouterloot, Henkel, & Walmsey 1989). Theoretical considerations suggest that bipolar outflows represent a phase of star formation in which inflow (along the equatorial regions) and outflow (along the poles) occur *simultaneously* (see the review of Shu, Adams, & Lizano 1987). According to this view, the stellar wind that drives the molecular outflow arises as a natural and fundamental consequence of the accretion of mass from a centrifugally supported disk to a rapidly rotating protostar (Shu, et al. 1988, hereafter SLRN; for the basic idea that a protostar rotating at break-up can drive very heavy mass loss, see also Hartmann & MacGregor 1982). In this *Letter* we demonstrate that this hypothesis, coupled with the knowledge that bipolar flows are made up of swept-up shells of ambient molecular cloud material, both explains many of the systematics known about outflows and provides the critical physical link between these fascinating objects and the process of star formation.

2. A SIMPLE DYNAMICAL MODEL

We adopt spherical polar coordinates (r, θ, ϕ) with origin centered on the source (see Fig. 1), and we assume symmetry with respect to the rotation axis of the system (no ϕ dependence) as well as reflection symmetry about the equatorial plane ($\theta = \pi/2$). We suppose that the molecular cloud core giving birth to the protostar has the density distribution of a (modified) singular isothermal sphere,

$$\rho(r, \theta) = \frac{a^2}{2\pi G r^2} Q(\mu) , \tag{1}$$

where $Q(\mu)$, with $\mu \equiv \cos\theta$, yields the flattening of density contours that arises either because of partial support against self-gravitation by rotation or by magnetic fields (cf. Lizano & Shu 1988). In equation (1), G is the universal gravitational constant; a represents the (angle-averaged) effective sound speed of the molecular cloud core; and the angular distribution function $Q(\mu)$ has been normalized so that

$$\int_0^1 Q(\mu)d\mu = 1 . \tag{2}$$

The molecular cloud-core models of Lizano & Shu typically show isodensity contours that have a flattening ratio (pole to equator) of about 1:2, which agrees reasonably well with observations of isolated cloud cores and Bok globules (e.g., Clemens & Barvainis 1988; cf. Myers et al. 1990). With the parametrization (1), this would correspond to a choice $Q(1)/Q(0) = \frac{1}{4}$.

The dynamical collapse of a core with the density profile (1) produces an unperturbed mass infall rate equal to (cf. Shu 1977)

$$\dot{M} \approx \frac{a^3}{G} , \tag{3}$$

with a substantial fraction of the inflow occurring near the equatorial regions if an appreciable degree of rotation accompanies the collapse (Terebey, Shu, & Cassen 1984). A geometrically thin, centrifugally supported disk soon forms from this infall. Accretion of this rapidly rotating material by the central protostar will eventually cause it to blow out a wind at a rate that is some fraction f of the infall rate:

$$\dot{M}_w = f\dot{M} , \tag{4}$$

where $f \sim 0.3$ for HH 7–11 (cf. SLRN; Lizano et al. 1988; Koo 1989; Masson, Mundy, & Keene 1990). We assume that the momentum input of such a wind per steradian has the form,

$$\frac{\dot{M}_w v_w}{4\pi} P(\mu) , \tag{5}$$

where v_w is the average velocity (≈ 150 km s^{-1} for HH 7–11).

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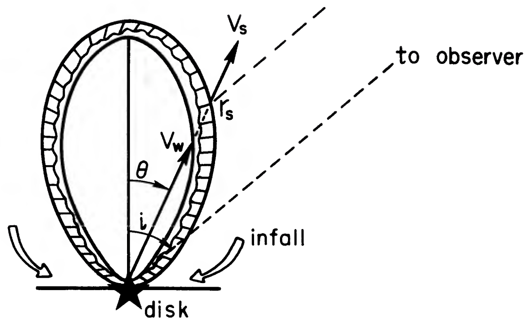

FIG. 1.—Geometry for shell of molecular gas, moving at speed $v_s(\theta)$ at position $r_s(\theta)$, as swept up by a protostellar wind moving at speed $v_w(\theta)$ into a molecular cloud core with the density profile given by eq. (1). The observer views the system at an inclination angle i with respect to the rotation axis of the protostellar disk, and residual infall of cloud material onto the disk occurs in those directions near the equatorial plane not affected by the outflow.

In equation (5), the angular distribution function $P(\mu)$ also satisfies the normalization condition,

$$\int_0^1 P(\mu)d\mu = 1 . \tag{6}$$

The value of $P(\mu)$ toward the poles ($\mu = 1$) will exceed its value (perhaps greatly) toward the equator ($\mu = 0$) for two reasons: (1) the stellar wind will be intrinsically collimated by the magnetohydrodynamic acceleration mechanism (Nerney & Suess 1975; Sakurai 1985; SLRN; for the problem of the collimation of a disk wind, see Lovelace 1976; Blandford & Payne 1982; Pudritz & Norman 1983; and Königl 1989), and (2) breakout of this wind to the molecular cloud core occurs more easily toward the rotational poles than toward the equator because of the greater amount of material that needs to be traversed in the latter direction (including the rotating infall as well as a disk; see Chevalier 1983 for a detailed discussion of this problem). With breakout having occurred from $\mu = 1$ to $\mu = \mu_{\min}$, it would not surprise us to find a ratio, say, $P(1)/P(\mu_{\min}) \sim 2$, or even more. The average momentum input rate given by equation (5) is independent of time if $\dot{M}_w$ is constant (eqs. [3]–[4] with f independent of time) and since $v_w \propto (GM_*/R_*)^{1/2}$ is constant as long as the protostar exists on the deuterium main sequence, where the stellar radius R_* is proportional to the stellar mass M_* (Shu & Terebey 1984; SLRN; Stahler 1988).

We now adopt a "snowplow" model to compute the properties of a swept-up shell when the wind interacts with the ambient molecular cloud core. Cantó (1980) and Barral and Cantó (1981) have performed related calculations for the problem of an isotropic wind expanding into an anisotropic medium. In assuming strongly radiative shocks and good mixing on both sides of the separation between swept-up cloud and shocked-wind material, our model greatly simplifies the earlier considerations.

We suppose that ahead of the shell with (outer) radius r_s, the core gas lies essentially at rest with respect to the star, and that a strongly radiating cloud shock sets this gas into motion at speed v_s. Similarly, a strongly radiating wind shock decelerates the flow incident at speed v_w on the back surface of the shell. If a perfect contact discontinuity were to separate the shocked wind from the shocked cloud, the nonspherical geometry of Figure 1 would require the introduction of slip velocities across the contact surface equal typically to a healthy fraction of v_w. Such a tangential discontinuity would be Kelvin-Helmholtz unstable to undulating disturbances that ruffle the surface, leading to the entrainment of material at speeds (normal to the surface) equal to some fraction of the conserved tangential velocities across the oblique shocks. Kahn (1980) calculates that the extent of mixing is small if the shocked-wind gas remains hot, but standard cooling curves readily produce the estimate that the cooling times will be short for wind speeds $\lesssim 350$ km s^{-1}. Since the radiating layers are then thin for the neutral winds that come from protostars, we assume that shocked-cloud material becomes locally well mixed with shocked-wind material, leading to the "turbulent broadening" of the line widths found by spectroscopic examination of the spatially thin shells (Moriarty-Schieven & Snell 1988), as well as making the combined medium behave on a large scale essentially as ballistic putty. On a very basic level, therefore, we agree with the conclusion by Levreault (1989) that bipolar flows are *momentum-driven* rather than *energy-driven*. In the simplest such model, no redistribution of matter occurs in the θ direction, and the net motion of each local piece of the shell is exactly radial so as to satisfy the vector conservation of momentum.

As the shell moves outward with velocity $v_s = dr_s/dt$, its mass $\mathcal{M}$ per steradian grows with time because it sweeps up new material with density $\rho(r_s, \mu)$:

$$\frac{d\mathcal{M}}{dt} = \frac{a^2}{2\pi G} Q(\mu)v_s , \tag{7}$$

where we have used equation (1) and have ignored the small amount of mass added to the shell by the shocked stellar wind. If we neglect the small external pressure of the ambient cloud core and set $v_w - v_s \approx v_w$, the momentum per steradian of the shell $\mathcal{M}v_s$ is increased by the incident stellar wind (eq. [5]),

$$\frac{d}{dt}(\mathcal{M}v_s) = \frac{\dot{M}_w v_w}{4\pi} P(\mu) . \tag{8}$$

Although the swept-up mass per steradian, $\mathcal{M}$, will generally depend on both μ and t, we see that a solution to equations (7)–(8) exists for each μ such that v_s = constant (see also Königl 1982, who assumed, however, an energy-driven bubble). For the solution relevant to our problem, the ram pressure of the outflowing stellar wind just balances the ram pressure of the incoming shocked swept-up mass such that the augmented shell continues to coast at constant speed (different for different μ). This result constitutes a special property of momentum-conserving winds blowing into $1/r^2$ density profiles (cf. Fig. 18 of Bally and Lada 1983) and gives a (CO) shell speed equal to

$$v_s = \left(\frac{\dot{M}_w}{2\dot{M}}\right)^{1/2} (av_w)^{1/2} \mathscr{B}(\mu) , \tag{9}$$

where we have made use of equation (3) to write a^2/G as $\dot{M}/a$, and where we have defined the bipolarity function:

$$\mathscr{B}(\mu) \equiv \left[\frac{P(\mu)}{Q(\mu)}\right]^{1/2} . \tag{10}$$

Notice that, despite the apparent symmetry of how $P(\mu)$ and $Q(\mu)$ enter in equation (10), these terms play quite different roles in the apparent collimation of the flow. The directional input of the original stellar wind, $P(\mu)$, represents an intrinsic focusing of the original driver; the anisotropic distribution of ambient material, $Q(\mu)$, then additionally allows the resulting flow to occur faster in some directions than in others. In this view, interstellar "toroids" (e.g., Torelles et al. 1983) arise from a passive blowing away of the polar caps of a flattened molecular cloud core, and not from an active channeling of the original stellar wind.

In any case, substitution of equation (4) now yields for the integral of equation (9):

$$r_s = (f/2)^{1/2}(av_w)^{1/2}t\mathscr{B}(\mu) , \tag{11}$$

where t is the time since wind breakout to the molecular cloud core, and we have assumed for simplicity that the fraction f is a constant independent of time. (Eq. [11] will break down near $\mu = \mu_{\min}$ if $\mu_{\min}$ itself decreases as a function of time [as it must if the wind is eventually to clear all 4π steradians around the star].)

Apart from factors of order unity ($[f/2]^{1/2} \sim 0.4$ if $f \sim 0.3$, and $\mathscr{B}[1] \gtrsim 3$), we see from equation (9) that the characteristic velocity for bipolar CO outflows equals the geometric mean $(av_w)^{1/2}$ of the effective sound speed a of the core giving birth to the protostar and the stellar wind velocity v_w. Since v_w measures typically as hundreds of km s^{-1}, whereas a measures typically from a fraction of a km s^{-1} (for low-mass star formation) to 1 km s^{-1} or more (for high-mass star formation), we now have a fundamental explanation why bipolar outflows usually have velocities logarithmically intermediate between the characteristic velocities of stars v_w and molecular cloud cores a, i.e., $v_s \sim$ a few tens of km s^{-1}. Moreover, from equation (11), we see that, until propagation occurs out of the region of the $1/r^2$ density drop-off, this characteristic speed is maintained independent of the size (i.e., the age) of the object, from the very smallest flows (e.g., Terebey, Vogel, & Myers 1989) to the largest (e.g., Wolf, Lada, & Bally 1990). A snapshot at any given time t will show blueshifted and redshifted lobes with an (unprojected) axial ratio $\mathscr{B}(1)/[(1-\mu^2)^{1/2}\mathscr{B}(\mu)]_{\max}$, which can easily equal the typical value 3:1, if, say, $P(\mu)$ and $Q(\mu)$ have variations of the order of magnitude previously mentioned. Finally, notice that equations (9) and (11) imply that a picture at given t would tend to show a "Hubble law" for the expansion,

$$v_s = \frac{r_s}{t} , \tag{12}$$

with the expansion speed v_s being larger in those directions $\mu > \mu_{\min}$ where the shell has progressed farther (by the factor $\mathscr{B}[\mu]$) in distance r_s.

If an outside observer oriented at an angle i with respect to the bipolar-flow axis observes the ratio of line-of-sight velocity $v_{\parallel}$ to projected distance $r_{\perp}$ from the central source along the major axis, he or she will measure

$$\frac{v_{\parallel}}{r_{\perp}} = \frac{1}{t} \cot(\theta - i) . \tag{13}$$

Equation (13) will yield a linear law, $v_{\parallel}/r_{\perp} \approx$ constant, if the opening angle with respect to the polar axis θ (from 0 to arccos $\mu_{\min}$) varies little in comparison with the size of the tilt angle i. The linear velocity law, $v_{\parallel} \propto r_{\perp}$ appears to characterize the on-axis velocity fields of well resolved and highly collimated bipolar flows such as L1551 (Fridlund et al. 1984; Moriarty-Schieven et al. 1987; Fridlund et al. 1989), Mon R2 (Wolf et al. 1990; Meyers-Rice and Lada 1991), NGC 2264G (Margulis et al. 1990), NGC 2071 (Snell et al. 1984), NGC 2024 (Richer et al. 1989), B335 (Cabrit, Goldsmith, & Snell 1988) and IRAS 16293−2422 (Walker et al. 1988). Moreover, kinematic models of highly collimated outflows indicate that outwardly increasing velocity gradients are required to produce *bipolar* lobes (i.e., lobes in which little spatial mixing of redshifted and blueshifted high-velocity emission occurs) without invoking very special aspect geometries of the flows relative to Earth (Meyers-Rice & Lada 1991). The linearity of the velocity law actually observed toward highly collimated outflows cannot be due solely to acceleration of the outflow material away from the central source since a nonphysical force law would be required (Meyers-Rice & Lada 1991). While other models, e.g., freely flying shrapnel, could also give rise to the observed "Hubble law," only our present explanation links the phenomenon to the other global properties of the flow as well as to the intrinsic process of star formation.

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