Feedback from Protostellar Outflows in Star and Star Cluster Formation

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Abstract. Magnetic stresses collimate protostellar winds into a common distribution of force with angle. Sweeping into the ambient medium, such winds drive bipolar molecular outflows whose properties are insensitive to the distribution of ambient gas and to the details of how the wind is launched, and how its intensity varies over time. Moreover, these properties are in accord with the commonly observed features of outflows.

This model is simple enough to permit a quantitative study of the feedback effects from low-mass star formation. It predicts the rate at which star-forming gas is ejected by winds, and hence the efficiency with which stars form. Applied to individual star formation, it relates the stellar initial mass function to the distribution of pre-stellar cores. Applied to cluster formation, it indicates whether the resulting stellar system will remain gravitationally bound.

Using the energy injection and mass ejection implied by this model, we investigate the dynamical evolution of a molecular clump as a stellar cluster forms within it. This depends critically on the rate at which turbulence decays: it may involve equilibrium star formation (slow decay), overstable oscillations, or collapse (fast decay).

1. Introduction

Stars, even low-mass stars, are not born quietly. In the process of accreting gas through an accretion disk, each star returns a fraction of the inspiralling material to the interstellar medium at speeds comparable to the escape velocity from the stellar surface. These powerful winds, observable as jets of ionized gas, sweep the ambient material into opposed streams known as bipolar molecular outflows. Because of their intensity, protostellar winds pose severe challenges to those who wish to model star forming regions. On a small scale, the wind from a forming star is capable of stripping gas from its parent molecular core even as that core feeds material to the star. On a larger scale, the outflow driven by a star's wind can punch through a massive molecular clump that is forming a stellar cluster. There is now plentiful observational evidence that these interactions are common, even ubiquitous, within star-forming clouds. Moreover, there is evidence to support the argument (Norman & Silk 1980) that the feedback from protostellar winds is the primary driver of turbulence within molecular clouds, and may be responsible for establishing Larson's (1981) line width-size relation in these clouds (McKee 1989).
Because these phenomena have profound consequences for the formation of stars and star clusters and for the evolution of molecular clouds, it is important to quantify the ways in which protostellar winds impact their molecular environment. In the past, models of feedback in star formation have either ignored the effects of low-mass stars (e.g., Elmegreen 1983) or considered only spherically symmetric protostellar winds (e.g., Nakano, Hasegawa, & Norman 1995), despite the fact that these winds are clearly jet-like. More sophisticated models of protostellar outflows (such as those of Shu et al. 1991; Masson & Chernin 1993; Li & Shu 1996) have not been applied to the basic questions of feedback. Partly, this has been caused by a lack of consensus (see Richer et al. 2000) as to which of the several classes of models for outflows best describes their dynamics.

We shall argue that there is a single model for protostellar outflows that is both motivated by theories of protostellar winds and validated by its ability to explain some of the common, otherwise mysterious features of protostellar outflows. Although this model cannot explain all the details of these flows, it provides a stable platform from which to launch an exploration wind-cloud interactions. We shall then use this model to address basic questions of star formation: the stellar initial mass function (IMF); the efficiency of star formation in clusters; and the dynamical evolution of star-forming gas during stellar cluster formation.

The work presented here has been conducted in collaboration with Chris McKee and appears or is intended to appear in: Matzner (1999), Matzner & McKee (1999a), Matzner & McKee (2000), and Matzner, Bertoldi, & McKee (2000).

2. Common Properties of Bipolar Molecular Outflows

Protostellar molecular outflows typically share several common features, as discussed by Masson & Chernin (1993) and Lada & Fich (1996). These are: a roughly linear position-velocity (PV) diagram (a “Hubble law”); a lack of receding material in the approaching lobe, and vice versa; and a power law distribution of mass with velocity:

\[ \frac{dM}{v_{\text{obs}}} \propto v_{\text{obs}}^{\Gamma}, \]

where \( v_{\text{obs}} \) is the line-of-sight velocity relative to the systemic velocity and \( \Gamma \) is typically about -1.8. Of these, the power law distribution of mass with velocity is the most difficult to explain, as it often holds quite well over a factor of five or ten in \( v_{\text{obs}} \) (e.g., Lada & Fich 1996).

In order to construct a theory for the feedback effects of protostellar winds it is necessary to develop a model for these outflows, which represent the interaction between winds and the ambient gas. To be tenable, such a model must confront the common properties of outflows listed above. Moreover, it must also be compatible with the ways in which protostellar winds are launched and collimated, a topic we now address.

3. Structure and Intensity of Protostellar Winds

Protostellar winds are launched centrifugally from the inner regions of protostellar accretion disks (Blandford & Payne 1982). This is made possible by the presence of intense poloidal magnetic fields, which act as rigid tubes anchored in
the disk, along which gas is flung away. Once the wind has traveled sufficiently far, the field weakens and gas begins to travel ballistically rather than being forced to corotate with the disk. Since it is still stuck to the same field line as when it was launched, and since that field line stretches back to the rotating disk, the magnetic field is wound into a tight spiral. By the Biot-Savart law, the tightly-wound field implies that the wind encloses a current. Magnetic stresses associated with current within the wind can generate forces that push the wind toward or away from its axis. These forces are important for the collimation of the wind on scales ($\lesssim 1$ AU) comparable to the launching region of the accretion disk, but on the larger scales of molecular outflows ($\sim 0.1$ pc), the wind relaxes into transverse pressure balance. Therefore, the wind’s magnetic flux lines tend toward a force-free state in which there is no current within the wind: the enclosed current lies entirely along the axis, rather than threading the wind itself. The field thus settles into the force-free state $B_\phi = 2I(r)/(c\varpi)$, where $\varpi = r \sin \theta$ is the cylindrical radius and the current $I(r)$ is a very slowly varying function of $r$. (There must be currents, and hence forces, in the wind for $I(r)$ to vary at all: these effect a continuing gradual collimation of the wind.)

This force-free distribution of toroidal magnetic field is sufficient to determine the ram pressure of the wind at large distances from its source. For, the magnetic field lines must corotate with the disk in a steady state. This, along with the fact that the wind will expand to fill the entire solid angle available to it, allows one to trace the field lines from their origin on the disk to their destination on the sky (Ostriker 1997; Matzner 1999). Then, the conservation of mass and energy along each streamline determines the ram pressure of the wind in each direction. In the case of a wind whose density just above the disk varies with disk radius $\varpi_0$ as $\rho_{w0} \propto \varpi_0^{-q}$, and whose field lines expand significantly ($\varpi \gg \varpi_0$) according to the force-free distribution of the field, the wind ram pressure varies as $\rho_{w}v_{w}^2 \propto C(r)\varpi_0^{(1-q)/2}\varpi^{-2}$ (Ostriker 1997). Now, the disk radius $\varpi_0$ is generally restricted within a very narrow range compared to the wind radius $\varpi$: for instance, in the model of Shu et al. (1995), $\varpi_0$ takes only a single value whereas $\varpi$ extends from a few AU to $\sim 0.1$ pc. For this reason, the ram pressure of the wind is very close to $\rho_{w}v_{w} \propto \varpi^{-2}$ in practice (Matzner & McKee 1999a). This force distribution, which is characteristic of the x-wind model (Shu et al. 1995), is thus common among hydromagnetic winds.

In reality, variability and instabilities are likely to broaden the wind force inside some angle $\theta_0$; therefore, we consider the smoothed distribution

$$\rho_{w}v_{w} = \frac{\dot{m}_{w}}{4\pi r^2} \frac{1}{\ln(2/\theta_0)} \frac{1}{(\sin \theta)^2 + (\sin \theta_0)^2}. \quad (1)$$

In the model for outflows presented below, observations imply $\theta_0 \simeq 0.01$.

The fact that hydromagnetic winds tend to $\rho_{w}v_{w} \propto (\sin \theta)^{-2}$ is consistent with the theoretical expectation (Matzner & McKee 1999a) that the fast Alfvénic Mach number, $B^2/(4\pi \rho_{w}v_{w}^2)$, is roughly constant. This expectation follows from the fact that this ratio is unity near the source, and varies logarithmically with distance along each streamline. The $(\sin \theta)^{-2}$ force distribution implies that the cumulative wind momentum within some angle, $p_w(< \theta) \equiv \int^\theta 2\pi r^2 \rho_{w}v_{w} \sin \theta' d\theta'$, varies as $\ln(\theta/\theta_0)$ for a broad range of angles, and this property will be useful for understanding outflows and their effects.
4. Bipolar Molecular Outflows from Hydromagnetic Protostellar Winds

The protostellar winds described above are essentially radial, both because they expand to fill the available angle, and because the wind coasts after achieving force balance. At wind speeds typical of low-mass protostars, the shocks that separate the wind from the surrounding gas are radiative (Koo & McKee 1992),
so the swept-up shell conserves momentum in each direction. For a steady wind and an ambient medium whose density varies as $\rho_a \propto Q(\theta) r^{-2}$, Shu et al. (1991) showed that these ingredients reproduce the Hubble law for outflows. However, Matzner & McKee (1999a) demonstrated that neither a steady wind nor $\rho_a \propto r^{-2}$ is necessary for this conclusion: as long as the ambient medium is relatively featureless (e.g., a power law) on the scale of the outflow, the outflow motion is self-similar even if the wind intensity varies. For radial motion this implies the Hubble law, since $v(\theta, t) \propto r(\theta, t)/t$. The collimated wind force distribution $\rho_w v_w^2 \propto (\sin \theta)^{-2}$ leads to outflows that are highly elongated and have bowshock-shaped tips of width $r_\theta$. This morphology is thus consistent with a “wide-angle” driving wind. A collimated outflow with $v \propto r$ lacks blue material in its red lobe and vice versa, consistent with observations.

That leaves only the power-law distribution of mass with velocity, $dm/dv_{\text{abs}} \propto v_{\text{obs}}^{-1.8}$, to be explained. Masson & Chernin (1992) argued that the radial outflow model of Shu et al. (1991) was inconsistent with this relation, but they only considered angular distributions $\rho_w v_w^2 \propto (\cos \theta)^{-\beta}$ for some $\beta$. Their results do not apply to hydromagnetic winds, for which $\rho_w v_w^2 \propto (\sin \theta)^{-2}$. Indeed, Li & Shu (1996) demonstrated that the outflow driven by a steady x-wind into a magnetically-flattened core, with $\rho_a \propto Q(\theta) r^{-2}$, could produce $dm/dv_{\text{abs}} \propto v_{\text{obs}}^2$.

I shall now show that this conclusion is much more general (see also Matzner & McKee 1999a), and is essentially independent of the ambient medium. The outflow shell expands at the rate dictated by momentum conservation; since $\rho_w v_w^2 \propto \theta^{-2}$ for small $\theta$, and since one expects the ambient medium also to be a power law of $\theta$, the rate of expansion obeys a power-law relation $v(\theta) \propto \theta^{-x}$ quite generally. The power $x$ can be computed, but it is not important for the present argument. Now, hydromagnetic winds satisfy the cumulative momentum distribution $p_w(< \theta) \propto \ln(\theta/\theta_0)$ for a wide range of angles, as described in §3. But since $v \propto \theta^{-x}$, momentum conservation [$p_{\text{shell}}(< \theta) = p_w(< \theta)$] implies a cumulative momentum distribution with outflow velocity $p_{\text{shell}}(> v) \propto \ln(v/v_{\text{max}})$, where $v_{\text{max}}$ is the rate of expansion along the wind axis. This, in turn, implies that $dm/dv = d^2 p_{\text{shell}}(> v)/dv^2 \propto v^{-2}$; and, since $v_{\text{obs}} \propto v$ for an elongated outflow, $dm/dv_{\text{obs}} \propto v_{\text{obs}}^{-2}$, i.e., $\Gamma = -2$. Notice that essentially nothing about the ambient medium entered into this argument; therefore, $\Gamma \simeq -2$ is generic.

Figure 1 demonstrates that models with this scaling can reproduce the mass-velocity distributions of real outflows.

The commonly-observed features of molecular outflows are thus the natural products of momentum-conserving shells driven by hydromagnetic protostellar winds: they are insensitive to the details of the wind’s launching region, of variations in the wind’s intensity, and of the ambient density distribution. Although this model remains to be tested in detail, it matches observation well enough to deserve attention as a mechanism for feedback.

5. The efficiency of star and star cluster formation and the IMF

The intensity of a protostellar wind along its axis ensures that within some angle, gas will be blown away. Because nearby gas is typically accreting onto the star itself or in the process of forming other stars, this process limits the efficiency with which individual stars and multiple stellar systems can form (see Matzner
& McKee 2000, for a more detailed discussion). Gas dispersal by protostellar winds is more ubiquitous and less violent than disruption by massive stars.

To estimate the amount of gas ($M_{ej}$) lost per outflow, first note that outflows propagate by conserving momentum in each direction. To be ejected, material must travel faster than the escape velocity $v_{esc}$ of the system. The momentum of the escaping gas is at least $M_{ej}v_{esc}$; therefore, $M_{ej}v_{esc} < f_w v_w m_*$. In reality, momentum is lost in those directions where the wind is too weak to eject anything, and wasted in those where the ejected gas travels faster than $v_{esc}$. Since the wind momentum is distributed logarithmically with angle, the fraction available to eject material at around $v_{esc}$ is roughly $\ln(1/\theta_0)$; in fact, $\ln(2/\theta_0)$ turns out to be a better estimate. Moreover, an outflow’s momentum is reduced by a factor $c_g$ (typically of order unity) due to the action of gravity as it traverses the cloud. Defining the efficiency parameter $X$ and the efficiency per star $\varepsilon$,

$$X \equiv c_g \ln(2/\theta_0)v_{esc}/(f_w v_w), \quad \varepsilon \equiv m_*/(m_* + M_{ej}),$$

Matzner & McKee (2000) find that for the formation of an individual star from a collapsing molecular core,

$$\varepsilon^{-1} \simeq (2X)^{-1} + [(2X)^{-2} + (1 + f_w)^2]^{1/2};$$

and for the formation of a star within a larger cluster-forming clump,

$$\varepsilon^{-1} \simeq 1 + (2X)^{-1}.$$  

In both cases, $\varepsilon \propto X$ for $X \ll 1$ (strong winds) and $\varepsilon \rightarrow 1$ for $X \gg 1$ (weak winds). The core and clump cases differ for a number of reasons. In the formation of a single star, unlike in multiple star formation, we can safely assume that all of the core mass is either accreted or ejected. Clumps have higher values of $v_{esc}$ (which raises $X$), but their density profiles are shallower (which lowers $X$ through $c_g$); these effects roughly cancel. Finally, mass is ejected at angles much further from the wind axis in the core case; for this reason, the ejected wind mass is included in (3) but neglected in (4).

For the formation of an individual star, (3) assumes that the collapsing core is spherically symmetric. However, protostellar cores are likely to be partially supported by magnetic fields, which will cause them to flatten along the magnetic axis. Since the protostellar wind is likely to share this axis, magnetic flattening segregates material away from the wind axis, reducing the amount that will be ejected and increasing the efficiency. Employing the models of flattened, magnetized cores presented by Li & Shu (1996), Matzner & McKee (2000) find that equation (3) remains valid if $X$ is replaced by $(1 + H_0)^{3/2}X$, where $1 \lesssim (1 + H_0) \lesssim 2$ is the factor by which magnetic support increases the core’s mean density. In the limit of a completely disk-like core ($H_0 \gg 1$), only the wind mass would be lost [$\varepsilon \rightarrow 1/(1 + f_w)$]. Individual stars form at about 30% efficiency if cores are spherical ($H_0 = 0$), or $\sim 75\%$ if they are significantly flattened ($H_0 = 1$), as shown in the left panel of figure 2.

For the formation of a star within a stellar cluster-forming clump, the anisotropy of the region is not likely to correlate with the wind axis. But, (4) assumes that each star forms at the center of a spherical distribution. Comparing again simulations in which stars form in a distributed manner throughout
Figure 2. **Left panel:** Core disruption and the efficiency of single star formation. The efficiency $\varepsilon$ of protostellar core collapse (equation (3)), versus the efficiency factor $X$. Curves are plotted for spherical ($H_0 = 0$), moderately flattened ($H_0 = 0.4$) and significantly flattened ($H_0 = 1$) cores. Circles indicate the values of $X$ for thermally supported ($T = 10$ K) cores; turbulence can only increase $X$ above these values. **Right panel:** Estimates of the instantaneous star formation efficiency $\varepsilon$ in stellar cluster formation. **Dashed lines:** analytical estimate, equation (4). **Solid lines:** Numerical evaluation of the same quantity in a protostellar clump with the structure $\rho \propto r^{-1}$ and in which stars form at the local ambipolar diffusion rate. **Dotted lines** indicate the mean extinction of the clump. Above the dash-dot line, stars more massive than $20 M_\odot$ are likely to form; our estimates are therefore only rough upper limits in this region. The analytical expression relies on assumptions that break down outside of the upper and lower thin lines.

...a cloud, Matzner & McKee (2000) find that (4) is an excellent approximation nevertheless, as shown in the right panel of figure 2. Equation (4) predicts efficiencies $30\% \lesssim \varepsilon \lesssim 50\%$ for the conditions typical of low-mass stellar cluster-forming regions, like those in Orion B (Lada 1992): each star expels one to two times its own mass from the clump. It is difficult to compare $\varepsilon$ with its observational analogue, the ratio of stellar to total mass, but the two quantities are similar in a few well-studied cluster-forming regions (Matzner & McKee 2000).

Using the efficiency of individual star formation, we can estimate the degree to which the stellar IMF differs from the mass function of pre-stellar cores (Nakano et al. 1995), a topic that has recently come under observational scrutiny (Motte, André, & Neri 1998). This comparison, discussed in detail by Matzner & McKee (2000), involves a consideration of how much $X$, and thus $\varepsilon$, differs among cores of different masses. We find that the variation of $\varepsilon$ is quite subtle, and that the slope of the IMF is thus only slightly flatter than the slope of the core mass function.

It should be noted that equations (3) and (4) predict efficiencies significantly higher than those found by Nakano et al. (1995), and that this is a direct result of the collimation of protostellar winds. By this process (in contrast to the effect of a massive star), mass is lost continuously as stars form. The efficiencies found above are consistent with the formation of bound clusters (Mathieu 1983), which raises the question of why so few actually do remain bound (Lada & Lada 1991). This could be due to a violent event associated with massive star formation;
alternatively, clusters could be unbound by motions of their parent clumps, a topic we address below.

6. Dynamical Evolution of Clumps During Cluster Formation

Embedded stellar clusters are found exclusively within the most massive clumps inside molecular clouds. Since these clumps are massive enough to be confined by their own self-gravity despite the weight of the molecular cloud above them (McKee 1999), we can gain insight into the production of a stellar cluster by approximating a clump as a distinct reservoir of gas, and examining its dynamical evolution while a stellar cluster forms within it. This is the subject of an ongoing investigation (Matzner 1999; Matzner & McKee 1999b; Matzner et al. 2000), and the results reported here should be considered preliminary.

Like any self-gravitating gaseous system, a cluster-forming clump must satisfy the virial theorem for self-gravitating gas (here in Eulerian form; McKee & Zweibel 1992):

$$\frac{1}{2} \ddot{I} = 2(T - T_0) + W + M - \frac{1}{2} \frac{d}{dt} \int_S (\rho v^2) \cdot dS,$$

where $I \approx M \Omega^2 / 2$ is the trace of the clump's moment of inertia tensor; $W$ and $M$ are its gravitational and magnetic energies; the kinetic energy of its internal motions is $T$; $T_0$ is an energy associated with its confining pressure; and the last term accounts for gas that may leave the system. Equation (5) has been used to identify equilibrium states of star-forming molecular clouds (McKee 1989; Bertoldi & McKee 1996), but it can also be taken as a dynamical equation if not taken to equal zero. In order to accomplish this, each of the energetic terms must be approximated as a function of four dynamical variables: the clump's mass $M_{cl}$ and the mass $M_*$ of its embedded cluster, its radius $R_{cl}$, and its internal velocity dispersion $v_{\text{rms}}$. To simplify matters, the results presented here assume a spherical, magnetically supercritical ($M_{cl} = 2M_\Phi$) clump with a $\rho \propto r^{-1}$ density profile and no rotation; any stars are assumed to move along with the gas.

To complete our description of the system, it is necessary to specify the evolution of $M_{cl}$, $M_*$, and $v_{\text{rms}}$. All three are affected by star formation and the effects of protostellar outflows on the gas. For this, we use the theory of McKee (1989): the local rate of gas conversion into stars ($t^{-1}_\text{gs}$) is identical to the local ambipolar-diffusion rate ($t^{-1}_\text{AD}$), because the formation of protostellar cores requires the assembly of magnetically supercritical condensations. Because of ionization by external FUV photons, star formation is inhibited within four visual magnitudes of the surface; interior to that, it proceeds at about a twelfth of the local free-fall rate. As stars form, their mass is added to $M_*$ and subtracted from $M_{cl}$; also subtracted is the mass of gas that is ejected according to equation (4).

For a massive molecular clump, $v_{\text{rms}}$ is dominated by non-thermal turbulent motions. The turbulent energy decays at a characteristic rate that can be parameterized by $(dv_{\text{rms}}^2/dt)_{\text{decay}} = -v_{\text{rms}}^2 / (\eta t_{\text{ff}})$, where $t_{\text{ff}}$ is the free-fall time of the clump. Therefore, turbulence must be replenished every $\eta$ free-fall times;
we expect $1 \lesssim \eta < 10$, where numerical experiments (e.g., Stone, Ostriker, & Gammie 1998) favor $\eta \lesssim 1$. Turbulent energy can derive from two sources. For one, it can be exchanged for other energies of the system, through work done during compressions or expansions of the cloud; for this, we use the analytical results of Zweibel & McKee (1995) and McKee & Zweibel (1995) to estimate the energetics of turbulent, magnetized gas. This process does not, however, prevent the cloud from collapsing. Alternatively, turbulence may be stirred up by protostellar outflows. Since these conserve momentum, it is most convenient to consider the clump’s turbulent momentum ($M_{cl} v_{rms}$) enhanced an amount $\phi_w p_w$ by a wind of momentum $p_w$; here, $\phi_w \sim 0.6$ is the fraction of the wind momentum that does not escape the clump (Matzner 1999).

Putting these elements together, it is useful to consider two approximations: one in which energy is added continuously according to the mean star formation rate within the clump, and another where stars form randomly at the current star formation rate, are chosen individually from the IMF, and affect the cloud instantaneously. These approaches differ for clumps of about five hundred solar masses or less, in which only ten to twenty stars form per $t_{ff}$. In the continuous approximation it is possible to identify equilibria (Norman & Silk 1980; McKee...
in which turbulent decay is just offset by injection from outflows. These are stable, in the sense that a compressed cloud’s turbulence is enhanced faster than it decays. But they can also be overstable, as energy in an oscillation can grow during compressive phases.

Perturbations due to individual forming stars will stimulate oscillations of the system, whose response depends critically on the rate of turbulent decay. Figure 3 shows the effect of changing $\eta$ for a clump of $500 M_\odot$: when turbulence decays slowly ($\eta = 10$), a cluster-forming cloud will remain very close to its equilibrium state, whether or not star formation is taken to be continuous. For moderate turbulent decay ($\eta = 3$), individual star formation seeds overstable oscillations about equilibrium: stars form in bursts. For fast turbulent decay ($\eta = 1$), a continuous simulation can remain near its equilibrium, but one with individual star formation cannot consistently replenish the lost turbulent energy every dynamical time; the system collapses. Although the approximations used in these simulations (e.g., spherical symmetry and no rotation) affect these outcomes, it appears that the duration and constancy of star formation are sensitive to the rate of turbulent decay.

7. Conclusions

Observations of star-forming regions and of individual accreting protostars have made it increasingly clear that stars interact violently with their parent clouds as they form. To model this phenomenon requires an understanding of protostellar winds, and of the bipolar molecular outflows they drive into the surrounding gas. We have seen that the typically observed features of molecular outflows can be explained as the natural product of magnetically-collimated winds and radiative shocks. This yields a model for outflows which can be used to estimate how much mass is ejected, either from an individual collapsing core or from a massive clump in the process of forming a stellar cluster. The estimates of star formation efficiency ($30\% - 50\%$) presented in §5 are much higher than those of Nakano et al. (1995), because they account for the collimation of protostellar winds.

Lastly, the feedback effects of protostellar outflows can be incorporated into simple dynamical models for the process of stellar cluster formation within massive molecular clumps. The investigations described in §6 show that the dynamics of stellar cluster formation are sensitive to the rate at which turbulence decays, and this will have observational consequences.

In the future, large-scale numerical simulations of star formation must include protostellar outflows, for without them such simulations are incomplete.

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