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Jets from accretion discs

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THE ROYAL SOCIETY

Jets from accretion discs

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Highly collimated jets are associated with young stellar objects of all masses. They are also prevalent in microquasars and a variety of active galactic nuclei. Common to all of these systems is the presence of accretion discs around central objects, be they young stars, compact objects, or black holes. Current observations of protostellar systems support a model that posits accretion-powered, centrifugally accelerated, hydromagnetic disc winds as the underlying mechanism for protostellar jets. This model also has the attractive possibility of unifying protostellar, microquasar and extragalactic jets within a single conceptual framework. We summarize the observations, theory and time-dependent numerical simulations of disc winds as the origin of protostellar jets. We also address the relevance of these winds for the theory of star formation, as well as the more general issue of a unified theory for astrophysical jets.

Keywords: accretion discs; jets; MHD; star formation; microquasars

1. Introduction

Astrophysical jets are among nature's most beautiful and interesting phenomena. Studied by extragalactic radio astronomers since the 1960s, jets acquired an association with relativistic processes and accretion discs in the neighbourhood of black holes in galactic nuclei following the seminal ideas and work by Martin Rees and his collaborators. The most enduring models for extragalactic jets have all emphasized the extraction, by hydromagnetic torques, of gravitational binding energy released during the accretion of gas onto the central black hole through an accretion disc, or the rotational energy of a spinning black hole (see Blandford, this issue). The microquasar phenomenon gives us a wonderful laboratory in which to study these processes on shorter time-scales in nearby, accreting, stellar mass black-hole systems (see Mirabel, this issue).

The discovery of bipolar outflows associated with young stellar objects in 1980, followed in 1983 by the discovery of the highly collimated jets that now appear to be their drivers, led to the realization that the jet phenomenon was far more universal. Hydromagnetic processes were quickly identified as important components of the central engines in protostellar systems because of the high thrusts that are observed in these outflows. The thrusts are orders of magnitude greater than anything that could be provided by radiative drives from their central stars. The mechanism of outflow cannot directly involve the spin of central protostars either, because they are observed to rotate at far less than their break-up speeds. If there is a universal mechanism for astrophysical jets, therefore, it must involve their magnetized accretion discs.

The wonderful advantage that studies of protostellar jets enjoy over their extragalactic counterparts is their proximity, universality and abundance of forbidden line

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emission. There is now good evidence that all protostars are associated with jets and outflows. They commence in the earliest phases of the star-formation process, and continue for as long as gaseous accretion discs are present. Protostellar jets are also nearby (the closest region of star formation in the galaxy is 140 pc away in the Taurus clouds) and plentiful. Optical protostellar jets are characterized by a variety of shock-induced forbidden lines, making it possible to directly measure many of their important physical properties. This contribution focuses upon protostellar jets with the hope that the basic principles learned in such studies are applicable to more general problems in astrophysics than in star formation. Readers may find recent reviews in Königl & Pudritz (2000) and Shu *et al*. (2000) for a comprehensive discussion of the literature, and may consult other related reviews in *Protostars and Planets IV*.

2. Observations of protostellar jets

Bipolar molecular outflows and narrow atomic jets are ubiquitous phenomena in protostars. There are now more than 200 bipolar CO sources known. They typically appear as comparatively low-velocity (*ca.* 25 km s^{-1}) and moderately collimated (length-to-width ratios of about 3:10) lobes, although several highly collimated CO outflows that exhibit high velocities (above 40 km s^{-1}) near the flow axis have now been detected. The mass outflow rate exhibits a continuous increase with the bolometric luminosity of the driving source for L_{bol} in the range *ca.* 1–$10^6 L_\odot$. Molecular outflows are present through much of the embedded phase of protostars, and, in fact, appear to be most powerful and best collimated during the earliest (class 0) protostellar evolutionary phase (Bontemps *et al*. 1996).

The bipolar lobes are generally understood to represent ambient molecular material that has been swept up by the much faster, highly supersonic jets that emanate from the central-star–disc system. Jets associated with low-luminosity ($L_{\mathrm{bol}} < 10^3 L_\odot$) young stellar objects (YSOs) have velocities in the range *ca.* 150–400 km s^{-1}, large Mach numbers (above 20), and can have opening angles as small as *ca.* 3–5° on scales of 10^3 to 10^4 AU. The inferred mass outflow rates are *ca.* 10^{-10} to $10^{-8} M_\odot\ \mathrm{yr}^{-1}$.

The total momentum delivered by the jets, taking into account both the density corrections implied by their partial ionization state and the long lifetimes indicated by the detection of pc-scale outflows, appears to be consistent with that measured in the associated CO outflows (e.g. Hartigan *et al*. 1994; Eislöffel & Mundt 1997). A critical review of the physical mechanisms of coupling the jets and the surrounding gas is given in Cabrit *et al*. (1997).

The momentum discharge deduced from the bipolar-outflow observations is, typically, a factor of approximately 10^2 higher than the radiation-pressure thrust L_{bol}/c produced by the central source (e.g. Lada 1985), which rules out radiative acceleration of the jets. Since the bolometric luminosity of protostars is, by and large, due to accretion, and since the ratio of jet kinetic luminosity to thrust is of the order of the outflow speed (*ca.* $10^{-3}c$), it follows that the jet kinetic luminosity is on average *ca.* 10% of the rate at which gravitational energy is liberated by accretion. This high ejection efficiency is most naturally understood if the jets are driven magnetically.

Magnetic fields have also been implicated in jet collimation. A particularly instructive case is provided by Hubble Space Telescope observations of the prototypical disc/jet system HH 30 (Burrows *et al*. 1996). The jet in this source can be traced

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to within *ca.* 30 AU of the star, and appears as a cone with an opening angle of 3° between 70 and 700 AU. The narrowness of the jet indicates some form of intrinsic collimation, since external density gradients would not act effectively on these small scales. Magnetic collimation is a likely candidate, made even more plausible by the fact that the jet appears to recollimate: its apparent opening angle decreases to 1.9° between 350 and 10^4 AU. Similar indications of recollimation have also been found in other jets. Given that any inertial confinement would be expected to diminish with distance from the source, this points to the likely role of intrinsic magnetic collimation.

There now exist measurements of magnetic fields in the flows themselves at large distances from the origin. In particular, the strong circular polarization detected in T Tau S in two oppositely directed non-thermal emission knots separated by 20 AU indicates a field strength of at least several gauss (Ray *et al*. 1997). This high value can be attributed to a magnetic field that is advected from the origin by the associated stellar outflow and that dominates the internal energy of the jet. This observation thus provides direct evidence for the essentially hydromagnetic character of jets.

Good evidence for a disc origin of jets is available for the energetic outflows associated with FU Orionis outbursts. The outbursts have been inferred to arise in young YSOs that are still rapidly accreting, although it is possible that they last into the class II phase (the infrared spectral energy distribution of protostars can be classified into four categories: from the most embedded systems, the class zero objects, to those with virtually no accretion discs left, the class III objects (see Pringle, this issue)). The duration of a typical outburst is *ca.* 10^2 yr, and during that time the mass accretion rate (as inferred from the bolometric luminosity) is *ca.* $10^{-4} M_{\odot}\ \mathrm{yr}^{-1}$, with the deduced mass outflow rate $\dot{M}_{\mathrm{wind}}$ (at least in the most powerful sources like FU Ori and Z CMa) being a tenth as large. The ratio $\dot{M}_{\mathrm{wind}}/\dot{M}_{\mathrm{acc}} \approx 0.1$ is similar to that inferred in class II YSOs, and again points to a rather efficient outflow mechanism.

Detailed spectral modelling demonstrates that virtually all the emission during an outburst is produced in a rotating disc. Furthermore, the correlation found in the prototype FU Ori between the strength and the velocity shift of various photospheric absorption lines can be naturally interpreted in terms of a wind accelerating from the disc surface (Calvet *et al*. 1993; Hartmann & Calvet 1995). Because of the comparatively low temperatures (*ca.* 6000 K) in the wind acceleration zone, thermal-pressure and radiative driving are unimportant (the latter due to the lack of a high-temperature source): magnetic driving is, thus, strongly indicated. The recurrence time of outbursts has been estimated to lie in the range *ca.* 10^3–10^4 yr, and if these outbursts are associated with the large-scale bow shocks detected in pc-scale jets (e.g. Reipurth 1991), then a value near the lower end of the range is implied. In that case, most of the stellar mass would be accumulated through this process, and, correspondingly, most of the mass and momentum ejected over the lifetime of the YSO would originate in a disc-driven outflow during the outburst phases (Hartmann 1997).

Hydromagnetic disc wind models require that large-scale, open magnetic field lines thread at least some portion of an accretion disc. There is now good evidence from submillimetre polarimetric observations of star-forming cores within molecular clouds that ordered fields exist on 0.01–0.1 pc scales (e.g. Schleuning 1998). It is still far from clear, however, that these poloidal fields continue down to the scale of the

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circumstellar discs that are two decades smaller in scale size. Recent observations of H30α, recombination line maser emission that is associated with the corona of the circumstellar disc around the massive peculiar emission line star MWC 349, may shed new light on this whole issue however (Thum & Morris 1999). Zeeman observations of these lines reveal the existence of a magnetic field of strength 22 mG at a distance of 40 AU from the central star. Any stellar field would be quite insignificant at these distances, and it appears that the only reasonable explanation is that of a disc magnetic field. The energy density in this field is 70% of the thermal energy of the disc. This, as we shall see in the next section, is more than enough to provide an excellent hydromagnetic drive for protostellar outflows. These observations hold great promise for detecting significant ordered fields in the discs surrounding at least massive YSOs.

3. Theory of centrifugally driven disc winds

The theory of centrifugally driven winds was first formulated in the context of rotating, magnetized stars (Schatzman 1962; Weber & Davis 1967; Mestel 1968). Using one-dimensional axisymmetric models, it was shown that such stars could lose angular momentum by driving winds of this type. This idea was applied to magnetized accretion discs in the seminal paper of Blandford & Payne (1982). Every annulus of a Keplerian disc may be regarded as rotating close to its 'break-up' speed, so discs are ideal drivers of outflow when sufficiently well magnetized. The removal of disc angular momentum allows matter to move inward, which produces an accretion flow.

Consider, for the moment, the simplest possible description of a magnetized rotating gas threaded by a large-scale open field (characterized by an even symmetry about the midplane $z = 0$). The equations of stationary, axisymmetric, ideal MHD are the conservation of mass (continuity equation); the equation of motion with conducting gas of density ρ subject to forces associated with the pressure P, the gravitational field (from the central object whose gravitational potential is ϕ), and the magnetic field $\mathbf{B}$; the induction equation for the evolution of the magnetic field in the moving fluid; and the solenoidal condition on $\mathbf{B}$:

$$\nabla \cdot (\rho \mathbf{V}) = 0, \tag{3.1}$$

$$\rho \mathbf{V} \cdot \nabla \mathbf{V} = -\nabla P - \rho \nabla \phi + \frac{1}{4\pi} (\nabla \times \mathbf{B}) \times \mathbf{B}, \tag{3.2}$$

$$\nabla \times (\mathbf{V} \times \mathbf{B}) = 0, \tag{3.3}$$

$$\nabla \cdot \mathbf{B} = 0. \tag{3.4}$$

Consider the *angular momentum equation* for axisymmetric flows. This is described by the ϕ component of equation (3.2). Ignoring stresses that would arise due to turbulence, and noting that neither the pressure nor the gravitational term contributes, one finds that

$$\rho \mathbf{V}_{\mathrm{p}} \cdot \nabla (r V_{\phi}) = \frac{\mathbf{B}_{\mathrm{p}}}{4\pi} \cdot \nabla (r B_{\phi}), \tag{3.5}$$

where we have broken the magnetic and velocity fields into poloidal and toroidal components: $\mathbf{B} = \mathbf{B}_{\mathrm{p}} + B_{\phi} \hat{\mathbf{e}}_{\phi}$ and $\mathbf{V} = \mathbf{V}_{\mathrm{p}} + V_{\phi} \hat{\mathbf{e}}_{\phi}$.

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Important links between the velocity field and the magnetic field are contained in the induction equation (3.3), whose solution is

$$\mathbf{V} \times \mathbf{B} = \nabla \psi, \tag{3.6}$$

where ψ is some scalar potential. This shows that the electric field due to the bulk motion of conducting gas in the magnetic field is derivable from an electrostatic potential. This has two important ramifications. The first is that, because of axisymmetry, the toroidal component of this equation must vanish ($\partial\psi/\partial\phi = 0$). This forces the poloidal velocity vector to be parallel to the poloidal component of the magnetic field, $\mathbf{V}_\mathrm{p} \| \mathbf{B}_\mathrm{p}$. This, in turn, implies that there is a function k, the mass load of the wind, such that

$$\rho \mathbf{V}_\mathrm{p} = k \mathbf{B}_\mathrm{p}. \tag{3.7}$$

Substitution of this result into the continuity equation (3.1), and then use of the solenoidal condition (equation (3.4)), reveals that k is a constant along a surface of constant magnetic flux, i.e. that it is conserved along field lines. This function can be more revealingly cast by noting that the wind mass loss rate passing through an annular section of the flow of area $\mathrm{d}A$ through the flow is $\mathrm{d}\dot{M}_\mathrm{wind} = \rho \mathbf{V}_\mathrm{p} \cdot \mathrm{d}A$, while the amount of poloidal magnetic flux through this same annulus is $\mathrm{d}\Phi = \mathbf{B}_\mathrm{p} \cdot \mathrm{d}A$. Thus, the mass load per unit time and per unit magnetic flux, which is preserved along each streamline emanating from the rotor (a disc in this case), is

$$k = \frac{\rho V_\mathrm{p}}{B_\mathrm{p}} = \frac{\mathrm{d}\dot{M}_\mathrm{wind}}{\mathrm{d}\Phi}. \tag{3.8}$$

The mass load is determined by the physics of the underlying rotor, which is its source.

A second major consequence of the induction equation follows from the poloidal part of equation (3.6). Taking the dot product of it with $\mathbf{B}_\mathrm{p}$ and using equation (3.7), one easily proves that ψ is also a constant along a magnetic flux surface and that it must take the form $\psi = \Omega - (kB_\phi/\rho r)$. In order to evaluate ψ, note that $B_\phi = 0$ at the disc midplane by the assumed even symmetry. Thus ψ equals Ω_0, the angular velocity of the disc at the midplane. One thus has a relation between the toroidal field in a rotating flow and the rotation of that flow:

$$B_\phi = \frac{\rho r}{k}(\Omega - \Omega_0). \tag{3.9}$$

Let us now examine the angular momentum equation. Returning to the full equation (3.5) and applying equation (3.7) and the constancy of k along a field line, one obtains

$$\mathbf{B}_\mathrm{p} \cdot \nabla \left(rV_\phi - \frac{rB_\phi}{4\pi k} \right) = 0. \tag{3.10}$$

Hence, the angular momentum per unit mass,

$$l = rV_\phi - \frac{rB_\phi}{4\pi k}, \tag{3.11}$$

is constant along a streamline. This shows that the specific angular momentum of a magnetized flow is carried by both the rotating gas (first term) and the twisted field

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(second term). The value of l may be found by eliminating the toroidal field between equations (3.9) and (3.11) and solving for the rotation speed of the flow,

$$rV_\phi = \frac{lm^2 - r^2\Omega_0^2}{m^2 - 1}, \tag{3.12}$$

where the Alfvén Mach number m of the flow is defined as $m^2 = V_\mathrm{p}^2/V_\mathrm{A}^2$, with $V_\mathrm{A} = B_\mathrm{p}/(4\pi\rho)^{1/2}$ being the Alfvén speed of the flow. The Alfvén surface is the locus of the points $r = r_\mathrm{A}$ on the outflow field lines where $m = 1$. The flow along any field line essentially corotates with the rotor until this point is reached.

From the regularity condition at the Alfvén critical point (where the denominator of equation (3.12) vanishes), one infers that the conserved specific angular momentum satisfies

$$l = \Omega_0 r_\mathrm{A}^2. \tag{3.13}$$

If we imagine following a field line from its footpoint at a radius r_0, the Alfvén radius is at a distance $r_\mathrm{A}(r_0)$ from the rotation axis and constitutes the lever arm for the torque that this flow exerts on the disc. The other critical points of the outflow are where the outflow speed, V_p, equals the speed of the slow and fast magnetosonic modes in the flow (at the so-called SM and FM surfaces).

(*a*) *Connection with underlying accretion disc*

We now apply the angular momentum conservation relation (3.5) to calculate the torque exerted on a thin accretion disc by the external magnetic field. The vertical flow speed in the disc is negligible, so only the radial inflow speed V_r and the rotation speed V_ϕ (Keplerian for thin discs) contribute. On the right-hand side, both the radial and vertical magnetic contributions come into play, so

$$\frac{\rho V_r}{r_0}\frac{\partial(r_0 V_\phi)}{\partial r_0} = \frac{B_r}{4\pi r_0}\frac{\partial(r_0 B_\phi)}{\partial r_0} + \frac{B_z}{4\pi}\frac{\partial B_\phi}{\partial z}. \tag{3.14}$$

One sees that specific angular momentum is removed from the inward accretion flow by the action of two types of magnetic torque. The first term on the right-hand side represents radial angular momentum associated with the radial shear of the toroidal field, while the second term is vertical transport due to the vertical shear of the toroidal field. In a thin disc, and for typical field inclinations, the second term will dominate. Note that the first term vanishes at the disc midplane because $B_r = 0$ there.

Now, following standard thin-disc theory, vertical integration of the resulting equation gives a relation between the disc accretion rate, $\dot{M}_\mathrm{acc} = -2\pi\Sigma V_r r_0$, and the magnetic torques acting on its surfaces (subscript s):

$$\dot{M}_\mathrm{acc}\frac{\mathrm{d}(r_0 V_\phi)}{\mathrm{d}r_0} = -r_0^2 B_{\phi,\mathrm{s}} B_z. \tag{3.15}$$

Angular momentum is thus extracted out of discs threaded by open magnetic fields. The angular momentum can be carried away either by torsional Alfvén waves or, when the magnetic field lines are inclined by more than 30° from the vertical, by a centrifugally driven wind. By rewriting equation (3.11) as $rB_\phi = 4\pi k(rV_\phi - l)$ and

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using the derived relations for k and l, the disc angular momentum equation can be cast into its most fundamental form:

$$\dot{M}_{\mathrm{acc}} \frac{\mathrm{d}(\Omega_0 r_0^2)}{\mathrm{d}r_0} = \frac{\mathrm{d}\dot{M}_{\mathrm{wind}}}{\mathrm{d}r_0} \Omega_0 r_{\mathrm{A}}^2 \left(1 - \left(\frac{r_0}{r_{\mathrm{A}}}\right)^2\right). \tag{3.16}$$

This equation shows that there is a crucial link between the mass outflow in the wind and the mass accretion rate through the disc:

$$\dot{M}_{\mathrm{acc}} \simeq (r_{\mathrm{A}}/r_0)^2 \dot{M}_{\mathrm{wind}}. \tag{3.17}$$

One has arrived at the profoundly useful expression of the idea that, if viscous torques in the disc are relatively unimportant, the rate at which the disc loses angular momentum ($\dot{j}_{\mathrm{d}} = \dot{M}_{\mathrm{acc}} \Omega_0 r_0^2$) is exactly the rate at which it is carried away by the wind ($\dot{j}_{\mathrm{wind}} = \dot{M}_{\mathrm{wind}} \Omega_0 r_{\mathrm{A}}^2$).

The value of the Alfvén lever arm $r_{\mathrm{A}}/r_0 \simeq 3$ for most theoretical and numerical calculations that we are aware of, so that one finds $\dot{M}_{\mathrm{wind}}/\dot{M}_{\mathrm{acc}} \simeq 0.1$. Thus, the observed link between accretion and outflow has a natural explanation in the context of disc wind theory. Its origin lies in the efficient manner by which the magnetic torques extract angular momentum from discs. This efficiency can be traced to the fact that magnetic wind torques have lever arms that are *external* to the disc, and, therefore, potentially very large. This is in obvious contrast with viscous torque mechanisms that have lever arms that are no more than αH, where α is the SS parameter for the torque ($\alpha \simeq 10^{-2}$ in the MHD simulations of Balbus–Hawley turbulence; see Hawley, this issue). The ratio of an MHD disc wind torque to a viscous disc torque may be written as $(B_z^2/4\pi P) \cdot (r_{\mathrm{A}}/\alpha H)$ (e.g. Pelletier & Pudritz 1992), which shows that even if the magnetic energy density within a disc is far less than the gas pressure P, its lever arm is still sufficiently large to beat viscous stresses in the competition to transport disc angular momentum.

4. Numerical simulations of jets from accretion discs

The advent of numerical simulations of disc winds is allowing us to explore the rich behaviour of time-dependent MHD jets unhindered by the simplifications inherent in steady-state models of the last decade. The published simulations may be grouped in two classes:

(1) dynamic MHD discs, in which the structure and evolution of the magnetized disc is also a part of the simulation; and

(2) stationary MHD discs, in which the underlying accretion discs do not change and provide fixed boundary conditions for the outflow problem.

(*a*) *Dynamic discs*

The first disc wind numerical calculations were published by Uchida & Shibata (1985) and Shibata & Uchida (1986). These simulations were carried out using magnetized discs with sub-Keplerian rotation. The models developed a rapid radial collapse in which the initially poloidal field threading the disc is wound up due to the differential rotation of the collapsing disc. The vertical Alfvén speed being smaller

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than the free-fall collapse speed implies that a strong, vertical toroidal field pressure gradient, $\partial B_z^2/\partial z$, must build up rapidly. The resulting vertical pressure gradient results in the transient ejection of coronal material above and below the disc as the spring uncoils. The work by Uchida & Shibata (1985) was confirmed by Stone & Norman (1994) using their ideal MHD, ZEUS two-dimensional code (Stone & Norman 1992). These are single, transient events however, and it is difficult to see how such a process would give rise to jets operating on many disc dynamical time-scales.

In two dimensions, a strong competitor for the removal of angular momentum by an external MHD torque is the Balbus–Hawley instability, which generates a vigorous channel flow and rapid outward transport of angular momentum (Balbus & Hawley 1991). The Balbus–Hawley instability grows on the orbital time-scale in sheared discs with weak magnetic fields. It is stabilized when fields reach strengths such that their magnetic and thermal energy densities are comparable ($v_{\mathrm{A}} \simeq c_{\mathrm{s}}$). Stone & Norman (1994) ran a series of two-dimensional simulations for a uniform magnetic field threading a wedge-shaped disc of constant opening angle. The vertical structure of the disc was set to be independent of latitude in spherical polars. The four parameters of the simulations were the ratio of the Keplerian velocity to the rotational velocity, the square of the ratio of the sound speed to the rotational speed, the square of the ratio of Alfvén speed to the rotational speed, and finally the ratio of the gas pressure to the magnetic pressure. Their simulations covered three cases: (A) sub-Keplerian rotation; (B) a Kepler disc with strong disc field; and (C) a Kepler disc with weak field. In case A, rapid collapse immediately ensues with an expanding, transient outflow appearing in 2.5 orbits (reproducing Uchida & Shibata (1985)). The collapse occurs until the centrifugal barrier is obtained, at which point a shock develops and the infall stagnates. In case B, the disc is Balbus–Hawley stable, but collapse occurs anyway because of the very strong braking that occurs for the disc due to the external wind torque. The disc collapses in a couple of orbits. In case C, one again sees rapid radial collapse of the disc, this time because of the Balbus–Hawley instability. This case differs from the other two because one does not encounter a rapidly growing toroidal field.

(*b*) *Stationary discs*

Given the highly dynamic behaviour of a dynamic disc and wind in two dimensions, it is difficult to compare such calculations with the predictions of steady-state wind models. One expects that in three dimensions, after the initial transients have died away, a Keplerian disc with a magnetic energy comparable with the gas pressure would be achieved. The turbulence in such a disc would lead to slight slippage of matter across field lines, allowing for a slowly evolving nearly equilibrium disc to be produced. Thus it is natural to consider simulations in which the disc provides fixed boundary conditions for the disc wind problem.

Several groups have adopted this approach including Ustyugova *et al*. (1995), Ouyed *et al*. (1997), Ouyed & Pudritz (1997*a*, *b*), Romanova *et al*. (1997) and Meier *et al*. (1997). The published simulations differ in their assumed initial conditions, such as the magnetic field distribution on the disc, the plasma β ($\equiv P_{\mathrm{gas}}/P_{\mathrm{mag}}$) above the disc surfaces, the state of the initial disc corona, and the handling of the gravity of the central star. Broadly speaking, all of the existing calculations show that winds from accretion discs can indeed be launched and accelerated, much along

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the lines suggested by the theory presented above. The results differ, however, in the degree to which flow collimation occurs.

For the remainder of this paper, we describe the simulations by Ouyed *et al*. (1997) and Ouyed & Pudritz (1997*a*, *b*, 1999), who employed the ZEUS two-dimensional code to study the origin and collimation of disc winds. The accretion disc is taken to be initially surrounded by a polytropic corona that is in hydrostatic balance with the central object as well as in pressure balance with the disc surfaces. This problem has an analytic solution. After considerable effort, we eventually found the correct technique by which this solution could be loaded into the initial state and be maintained to any desired degree of accuracy for any desired amount of time. We always use an unsoftened gravitational potential from the central object. The disc has an inner edge at radius r_i, inside of which no material is rotating.

We use two initial magnetic configurations, each chosen so that no Lorentz force is exerted on the hydrostatic corona described above; more specifically, we used initially current-free configurations $\mathbf{J} = \mathbf{0}$. These are the potential field configuration of Cao & Spruit (1994), and a uniform, vertical field that is everywhere parallel to the disc rotation axis. The former configuration has field lines whose opening angles increase with increasing disc radius and is the solution for the magnetic field configuration that is expected for a conducting plate that is surrounded by a vacuum. The latter configuration is chosen because of its unfavourable opening angle for launching a wind at all radii.

We ran high-resolution simulations in which (500×200) spatial zones were used, and simulations were run up to $400t_\mathrm{i}$ (where t_i is the Kepler time for an orbit at the inner edge of the disc, r_i). Our large-scale simulations extend over a small region of disc and jet; a region $20r_\mathrm{i}$ in radius by $80r_\mathrm{i}$ in the z-direction. The model is described by five parameters, whose values describe conditions at r_i: three describe the initial corona (the coronal β_i is typically 1.0; the contrast between disc and coronal density is 100; and the turbulent disc corona is given an effective sound speed that is typically 10 times the Kepler speed of the disc), as well as two parameters to describe the disc physics (the injection speed v_inj of the material from the disc into the base of the corona, which is typically $10^{-3}v_\mathrm{K}$ at any point on the disc; and the disc toroidal field scales with its poloidal field). The start of the simulation at $t = 0$ may be regarded as the moment that the rotation of the underlying Keplerian disc is turned on.

(*c*) *Simulations: stationary jets*

A series of snapshots of the outflow for the potential configuration is shown in figure 1 (see Ouyed *et al*. 1997; Ouyed & Pudritz 1997*a*). The left-hand panels in figure 1*a*–*c* show plots of the poloidal field lines, and the corresponding right-hand panels show the density isocontours. Figure 1*a* shows the initial state, with its flaring potential field lines and the hydrostatic corona. Two snapshots of the flow at 100 and 400 inner time units are shown in figure 1*b*, *c*. At time 100, we see the bow shock separating the undisturbed coronal material from the post-shock jet gas. The outflow is seen to be highly collimated along the axis. Note that the poloidal field lines have also been collimated towards the jet axis in comparison with lines in the undisturbed corona. The toroidal field in the jet is responsible for reducing the opening angle of the flow. At time 400, the bow shock is beyond the end of our computational domain and the entire jet has achieved a cylindrically collimated, stationary state.

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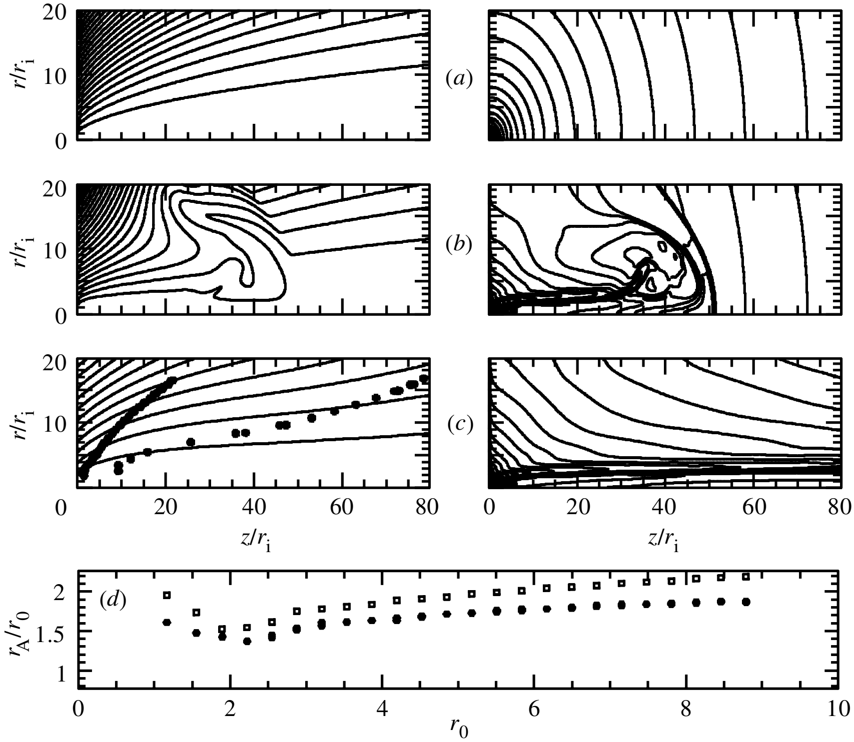

Figure 1. Potential configuration: $v_{\mathrm{inj}} = 10^{-3} v_{\mathrm{K}}$. In all figures, the axis of symmetry is plotted horizontally and the disc surface vertically. (*a*) The initial magnetic configuration, shown in the left panel, is the potential field. The right panel shows the initial isodensity contours of the corona. Evolution of the initial magnetic and density structure (the left and right panels, respectively) at 100 and 400 inner time units is shown in (*b*) and (*c*). (*c*) Location of the Alfvén critical surface (filled hexagons) and of the FM surface (stars). (*d*) Comparison of the Alfvén lever arm r_{A}/r_0 for each field line in the simulation (filled hexagons) with predictions of steady-state theory (squares). (Adapted from Ouyed *et al.* (1997).)

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As a check of stationary wind theory, we show in figure 1*d* the position of the Alfvén surface on each field line, as found from our numerical data, and also plot the position predicted by our theory in § 3. One sees that the two are in excellent agreement. The outflow begins to collimate just beyond the Alfvén surface. This process is accomplished, in the main, by the time that the FM surface is reached, which is also plotted. The ratio of the toroidal to the poloidal field strength at the FM surface is between 2 and 3 for many field lines. We also find that the opening angles of field lines as they emerge from the disc, and along which an outflow is driven, obey the condition $\theta_0 \leqslant 60°$, as predicted by steady-state theory. Not all the field lines threading the disc drive an outflow, however. The size of the jet region upon a disc appears to be related to the rigidity of the magnetic field lines in the disc corona. This is measured by the plasma β, which, for values much less than unity, characterizes a field structure that is too dominant to be further shaped.

The Alfvénic Mach number achieved by the outflow at $z/r = 10$ is about 5, while the FM Mach number is quite modest, being only 1.6 at this distance along a characteristic outflow field line. The highest speeds of outflowing material are

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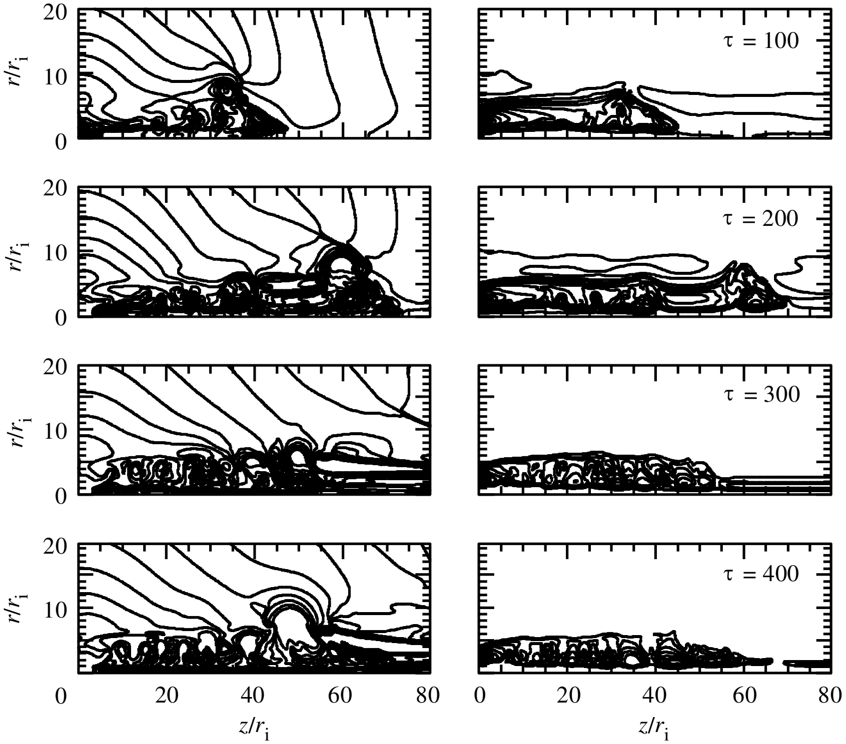

Figure 2. Vertical configuration: $v_{\rm inj} = 10^{-3} v_{\rm K}$. Jet density (left) and toroidal field (right) showing evolution of episodic outflow at four different times. Twenty logarithmically spaced contour lines are shown for density, 20 equally spaced contours for B_ϕ. (Adapted from Ouyed & Pudritz (1997*b*).)

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obtained closest to the jet axis, while lower-velocity jet material is at larger radii. This is expected on the basis of § 3 because the terminal speed of the outflow scales with the product of the lever arm and the Kepler speed of the footpoint of the field line. An additional new effect, not contained in the self-similar theory, is the appearance of a 'Hubble flow', in which the highest-speed material at any given radius is found farthest from the jet axis (see Ouyed & Pudritz 1997*a*).

(*d*) *Simulations: episodic jets*

We first discovered the presence of episodic jets in simulations of outflows in the vertical magnetic configuration (Ouyed *et al*. 1997; Ouyed & Pudritz 1997*b*). In figure 2, we show an episodic outflow in a jet with the same parameters as those for figure 1, except that a uniform vertical field is used. Figure 2 shows a large simulation region (comparable in size with figure 1). Four snapshots of the jet are shown; at 100, 200, 300 and 400 time units. The density isocontours are plotted in the left-hand panels, while contours of the toroidal magnetic field are plotted on the right-hand panels. This simulation *never* achieves steady state. The average outflow speed is *ca.* $0.5 v_{\rm K,i}$, while the highest outflow speeds achieve *ca.* $v_{\rm K,i}$. The region of knot generation is located at a distance of $z \simeq 5$–$7 r_{\rm i}$ from the central source.

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Why do outflows occur in the initially vertical magnetic field configuration? Higher-resolution simulations of the flow shown in figure 2 reveal that the initial vertical magnetic configuration has been considerably opened up. This is the effect of the pressure gradient force $\partial(B_{\phi}^2/8\pi)/\partial r$ generated in the jet. It pushes the fields lines radially outwards near to the disc surface, and thus allows the wind to be ejected. The gradient arises from the fact that the strongest toroidal magnetic field is generated in the inner region of the flow where the Kepler speed is the highest. There is also fairly good agreement between our numerical results and steady-state theory for the position of the Alfvén surface (see Ouyed & Pudritz 1997*b*). The point here is that the jet starts to develop a significant toroidal field beyond the Alfvèn surface. It is this component of the field that, given low enough mass loading of field lines, drives the episodic behaviour that we now discuss.

Knots are produced on a time-scale that is expected for the kink instability. Our simulations verify that if r_{jet} is the radial size of the jet region, then the time-scale for knot production at the generator region is, roughly $\tau_{\mathrm{knot}} \simeq r_{\mathrm{jet}}/v_{\mathrm{A},\phi}$. The knots are coherent structures that propagate down the length of the jet. They are kept coherent by the pressure of regions of strong toroidal field that both precede and follow them. That is, regions of high density (the knots) are regions with low toroidal field strength, while the inter-knot regions are characterized by low density and high toroidal-field strength. Thus, each knot experiences the confining pressure of a toroidal field gradient from ahead of and behind it, as well as the overall confinement towards the flow axis from the radially inward-directed pinch force.

Knots are formed when field lines have relatively low mass loading. Under these conditions, the toroidal pressure forces become very strong and refocus outflow towards the outflow axis. After recollimation, the outflow must accelerate, much like water going through a pinch in a garden hose. The inward motion of the gas beyond the recollimation region (located just beyond the Alfvén surface) is terminated by its reflection off a 'centrifugal barrier'; gas may not move inwards of the radius, where it comes into centrifugal balance once again. This reflected stream appears to plough into slower moving material at slightly larger radial scales, and a shock results. The material that piles up becomes the material of the knot. The outflow moves the knot away from the generating region, and the whole process repeats.

It follows from these results that the *transition from episodic to stationary jets* occurs for sufficiently well mass loaded field lines in any given magnetic configuration. Our analysis and simulations suggest that the nature of axisymmetric disc winds (stationary or episodic) is connected to the interplay between the toroidal field energy density generated in the jet and the kinetic energy in the outflow. We find (see Ouyed & Pudritz 1999) that this behaviour can be best thought of by comparing the speed of a toroidal Alfvén wave across the jet radius with the dynamical speed along the jet. This interplay can be captured in the definition of a quantity we defined as $N = B_{\phi}^2/8\pi\rho v_{\mathrm{p}}^2$. We find that, in general, stationary flows involve low values of N (high mass load), while episodic jets typically occur for higher values of N (low mass load).

5. Outflows and star formation

Why are hydromagnetic winds so prevalent and so intimately associated with accretion discs and star formation? The three most basic stellar attributes whose origins

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must be explained are their masses, spins and magnetic fields. Given that clouds are rather strongly magnetized—and that the rotation of molecular cores, while small, is not insignificant—some mechanism must be at work to rid star-forming gas of its abundant magnetic field and angular momentum. The former is suspected to proceed by the mechanism of ambipolar diffusion proposed by Mestel, Spitzer and others (see McKee *et al*. (1993) for a review). The rotation and magnetization of cores implicate the formation of discs, out of which most of the stellar material is probably accreted. Thus, the stripping of star-forming gas of its residual angular momentum is a process that is intimately associated with the transport of angular momentum in accretion discs. As long as a reasonably ordered residual magnetic field continues to thread the growing circumstellar discs as gravitational collapse from the surrounding core proceeds, one may reasonably expect that disc winds will ultimately arise. Indeed, an accretion disc has been discovered around the youngest class 0 protostar that we know of, VLA 1623, whose age cannot be more than several 10^4 years old (e.g. Pudritz *et al*. 1996, 1997).

Do winds determine stellar masses in some way? Current models of star formation suggest different views of this important question, and we briefly address these. As is well known, there are two physical pictures of individual star formation that have dominated the discussion over the last few decades (for example, see the reviews by Shu *et al*. (1987) and Pudritz *et al*. (1996, 1997)).

The first view, based on the gravitational stability of Bonner–Ebert spheres, posits that the mass of a star is, essentially, a *Jeans mass*. At a temperature of 10 K, and for reasonable core pressures, the critical Bonner–Ebert mass is indeed around $1M_\odot$ (e.g. Larson 1992). The basic problem with this approach, however, is that in a highly inhomogeneous medium, the Jeans mass becomes an ambiguous quantity. In order to explain the IMF, therefore, one would have to posit that it is 'laid-in' in the mass spectrum of clumps. Recent submillimetre observations by Motte *et al*. (1998) of clumps in the ρ Oph cloud claim to have found a core spectrum much like the Salpeter IMF, however, much more work will have to be done on this point. Winds, in this picture, would presumably play no role in the control of stellar masses.

The second basic model for star formation is the idea that it is not the mass but, rather, the *mass accretion rate*, $\dot{M}$, that is determined by a molecular cloud. The self-similar collapse picture (e.g. Shu *et al*. 1987) requires that some process *turns off* the infall. The observed outflows have been hypothesized to play this role. In this picture, winds would play a central role in determining the mass of stars, yet the reason why one star acquires a fraction of a solar mass, and another several solar masses, has never been made clear. A more basic concern, however, is that jets in class 0 objects are highly collimated and will not intercept most of the infalling envelope. In addition, there is a strong positive relation between disc accretion and jet outflow, which seems to indicate that plenty of matter can collapse onto a disc and accrete through it onto the central protostar.

A third picture of star formation is gradually being developed that incorporates elements of both of these pictures. It focuses on the observed fact that stars do not form in isolation, but rather as members of groups and clusters. The idea is that stellar mass is acquired through a process of competitive accretion for the gas that is available in a clump. In the specific numerical simulation of Bonnell *et al*. (1997), for example, the accretion rate is modelled as a Bondi–Hoyle process onto a distributed set of initial objects of mass M_0. The stellar velocity is v_∞ and the gas

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density of the clump is ρ_∞. The simulations showed that seed objects in the higher, central density region became more massive than the outliers. This can be understood analytically, since the time dependence of accretion at the Bondi accretion rate, given by $\dot{M}_{\mathrm{BH}} = \gamma M^2$, where $\gamma \equiv 4\pi\rho_\infty G^2/(v_\infty^2 + c_{\mathrm{s}}^2)^{3/2}$ is readily found. It is easy to show that the accretion time-scale in this picture is $t_{\mathrm{accr}} = (1/\gamma M_0) \propto (1/\rho_\infty M_0)$. Shorter accretion times occur in higher-density regions of a clump, which is where one might expect more massive stars to be formed. These results appear to be strongly dependent on the mass of input objects, which is not determined by the theory.

The cluster mode of star formation appears not to require a fundamental role of winds in determining stellar mass. A plausible picture for star formation, then, is that while stellar mass is determined by a competitive accretion process, stellar angular momentum is finally removed through processes involving jets and disc winds.

Is there a lesson here for other accreting systems that produce astrophysical jets? The fact that microquasars produce jets is a hint that the magnetic field seemingly required for jet production does not have to be inherited from a larger environment (such as molecular clouds), but could be generated locally within discs through MHD turbulence and some sort of dynamo action (see, for example, Tout & Pringle 1996; Curry *et al.* 1994). This is probably also relevant to jet production in active galactic nuclei (AGN) (see Blandford, this issue). The most puzzling problem here is to understand why all AGN do not produce jets and radio emission. Until this is better understood, there is little doubt that we still have much to learn about astrophysical jets.

I thank the organizers of this most stimulating meeting for the invitation to present this paper. I also wish to thank my collaborator and former student, Rachid Ouyed, for many interesting conversations on all of these topics. This work was supported in part by an operating grant from the Natural Science and Engineering Research Council of Canada.

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Discussion

A. Shukurov (*University of Newcastle, UK*). Models of centrifugally driven winds need a significant poloidal field. How strong is the required field? What fraction of the initial field can be retained in a collapsing protostellar cloud not withstanding magnetic (especially ambipolar) diffusion?

R. E. Pudritz. Because the lever arm of the wind torque (r_{A}), is so much larger than that for viscous torques (essentially αH), a field strength that is less than equipartition with the thermal energy of the disc can still be the dominant means of transporting away the disc angular momentum. Not much flux is lost to a collapsing region through the process of ambipolar diffusion, at least not during the collapse phase.

I. F. Mirabel (*CEA/SACLAY, France*). You said that only 50% of the accreted mass is blown into a wind. What happens with the remaining 90%? Are there observational tests for this fraction?

R. E. Pudritz. The remaining 90% flows through the disc onto the central object. Essentially, each wind particle carries off the angular momentum of 10 of its fellows left behind in the disc. Direct observations of the ratio $\dot{M}_{\mathrm{wind}}/\dot{M}_{\mathrm{acc}}$ are made using a variety of emission line diagnostics in the jet and the disc, as I mention in my paper.

Y. Uchida (*University of Tokyo, Japan*). I would say that the magnetic field in the disc can be 100 μG to 1 mG if it is frozen in the mass contraction in the star formation case, and strong enough to exert a braking effect on the disc.

The Balbus–Hawley instability that you mentioned to break up the disc does not act appreciably in the deep part of a high-density disc. Instead, it preferentially

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occurs at surfaces of the disc or torus, much like avalanches. We have been dealing with this nonlinear preferential mode of Balbus–Hawley instability earlier than they published the linear treatment of the process.

R. E. PUDRITZ. Indeed, the recent maser (H 30 α recombination line) observations of Thum & Morris suggest a 20 mG disc field. The Balbus–Hawley instability is destructive only in two dimensions, and for weakly magnetized discs. I had not realized that you had seen it in operation before they noted it.

I. W. ROXBURGH (*Queen Mary and Westfield College, London, UK*). Can you explain this magic number of 3 that you gave us for the lever arm?

R. E. PUDRITZ. It is a highly nonlinear problem, and so I'd really like to rest on the simulations here. If you take plausible field strengths that go with these accretion rates you can estimate that an Alfvén surface could be anywhere from 3 to 10; you don't really know that number very well. It can't be much larger than 10. If it's 1, then nothing's gained here. You have much too high a mass loss rate from the disc and are essentially putting as much mass into the wind as upon the central star. So reasonable numbers from a back of the envelope calculation range from 3 to 10.

K. BARKER (*Zetex Semiconductor plc*). Does the jet carry a net current in these models?

R. E. PUDRITZ. Yes it does. Ouyed and I specifically show the current distribution that is associated with our stationary jet. The current flows down the axis of the jet, outwards through the accretion disc, and returns along the flanks of the bow shock to the axis, completing the circuit. This net current may lead to the establishment of the cylindrical collimation predicted by Heyvaerts & Norman (1989) for this case.

D. LYNDEN-BELL (*Institute of Astronomy, Cambridge, UK*). Because the disc is accreting, the total magnetic flux through a ring of given radius will increase with time. Thus the disc is unsteady with the magnetic field increasing.

R. E. PUDRITZ. You are correct. Dissipation is needed in the disc if accreting matter is to pass through the magnetic field, leaving it fixed in space. Our simulation does not follow the evolution of the disc so that the disc accretion rate cannot be numerically computed. However, the radial inflow speeds are so small that it is reasonably safe to set $v_r = 0$ in our disc boundary conditions, at least over the 400 odd inner-disc Keplerian periods that characterize our simulations.

A. BRANDENBURG (*University of Newcastle upon Tyne, UK*). What determines the opening angle of the field and what prevents it from going subcritical?

R. E. PUDRITZ. The opening angle is given an initial value from the initially imposed magnetic configuration. It is then free to adjust itself as the jet dynamics develop. The strength of the coronal β, which is a function of radius in our models, appears to be important. Low β seems to make the field rather rigid. Also, the mass loss per unit area, which is controlled by our injection speed parameter, controls the opening angle. A finite region in our discs is involved in driving an outflow in all of our (non-self-similar) magnetic configurations and simulations.

R. D. BLANDFORD (*Caltech, Pasadena, USA*). There are some observations which show that approaching and receding Herbig–Haro objects can be matched. This shows that they are found at the source. Is this generally true?

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R. E. Pudritz. Herbig–Haro objects are generally believed to form in the bow shocks that accompany the episodic jets. Jets are generally two-sided, but their receding red lobes can occasionally be hidden by intervening discs. They are also often asymmetric, a consequence of the highly anisotropic medium in which protostars are formed.

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