Aust. J. Phys., 1992, **45**, 513–29

Mechanisms for the Production of Bipolar Flows and Jets in Star Formation Regions*

G. V. Bicknell

Mount Stromlo and Siding Spring Observatories, Australian National University, Private Bag, Weston Creek, A.C.T. 2611, Australia.

Abstract

Two types of mass ejection are associated with the formation of young stars: poorly collimated bipolar flows and well collimated jets. Some mechanisms which have been suggested for the driving of these flows are reviewed. These include centrifugally driven magnetised winds, magnetic pressure driven winds and bubbles driven by ionised or neutral winds. The idea that the bipolar flows are bubbles driven by a neutral wind seems attractive on both theoretical and observational grounds but the source of the neutral wind—disk or star—is uncertain. It is possible that the jets are driven by magnetic pressure or by a rapidly rotating magnetic field close to the star. However, no definitive theory exists at the present time.

1. Introduction

The discovery of energetic (total energy $\sim 10^{43}$–10^{47} erg) large scale ($\sim$1 pc) bipolar flows of molecular gas in the vicinity of star formation regions (Snell *et al.* 1980; Rodriguez *et al.* 1980; Lada and Harvey 1981; Bally 1982) ushered in a new era in the study of star formation. Bipolar flows were initially discovered through the existence of high velocity wings on the CO profiles of such regions; subsequent maps revealed the existence of poorly collimated (but nevertheless bipolar) regions of red- and blue-shifted emission (see Fig. 1). These flows, whose velocities are $\sim$10–50 km s^{-1}, are associated with protostars and young stellar objects (YSOs) and pre-main-sequence stars (T-Tauri stars) and represent an important phase of the life of these objects before they come onto the main sequence.

Jet flows are frequently (but not always) associated with these pre-main-sequence objects. Usually they are detected through $H\alpha$ imaging; they are well collimated with velocities $\sim$200–400 km s^{-1} and a mass flux 10^{-10}–10^{-7} $M_\odot$ yr^{-1} (Mundt 1986). Herbig–Haro objects are almost always associated with these jets and apparently arise either through internal shocks in the jet or the bow shocks enveloping the head of the jet as it forces its way out through the surrounding molecular cloud. The internal shocks may occur either through non-linear development of the Kelvin–Helmholtz instability or through variations of the jet power (Raga *et al.* 1990).

* Paper presented at the Workshop on Star Formation in Different Environments, held at the University of Sydney, 9–11 October 1991.

0004-9506/92/040513$05.00

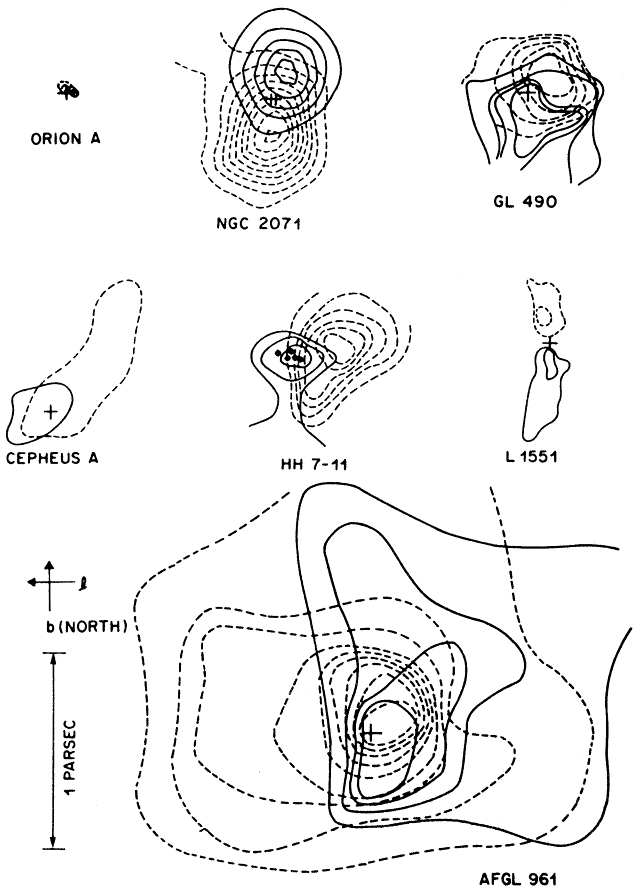

Fig. 1. Maps of several bipolar flows reproduced from Bally and Lada (1983). The solid contours represent blue-shifted emission; the dashed contours represent red-shifted emission.

Bipolar flows and jets may not merely be a by-product of star formation, they may regulate it significantly. Bipolar flows may carry away the angular momentum from the accreting gas, thereby enabling the star formation process itself. The energy in jets may be sufficient to disrupt the dense environment of a new star, and an understanding of their energetics may be necessary if we are to fully understand the determination of the initial mass function. The study of jets in star formation regions may also be instrumental in a clearer understanding of the processes which lead to the formation of jets in the environments of radio galaxies and quasars.

This review is not concerned with *all* the details of jets and bipolar flows in star formation regions. In particular I shall not consider the interesting physics concerning the large scale structure of these objects, for example, the production of Herbig–Haro objects. Rather, I shall focus on the mechanisms which have been considered for the production of jets and bipolar flows. As we shall see, there is no real consensus on this at the present time. However, I have written this review with the intention of making clear the salient physical characteristics of each particular theory so that observers and theoreticians alike may be able to take it as a starting point for their own particular endeavours. This review extends the excellent review of Padman *et al.* (1991) in that it includes a more detailed review of centrifugally driven flows, as well as more recent work on magnetic pressure driven and wind driven flows. However, readers should consult Padman *et al.* for a more comprehensive compilation of some of the observational papers.

2. Morphology, Energy and Momentum of Bipolar Flows and Jets

Bipolar flows and jets obviously represent two quite different types of mass loss activity: the massive, poorly collimated molecular outflows are associated with the early stages of star formation whereas the jets may be associated with the later stages. It is not clear however whether the detection of jets in the late stages of star formation is due to the smaller optical depth. In some objects (e.g. the much-studied source L1551; Moriarty-Schieven *et al.* 1987) the two coexist and observations such as this have promoted the development of theories which give rise to the simultaneous production of jets and bipolar flows.

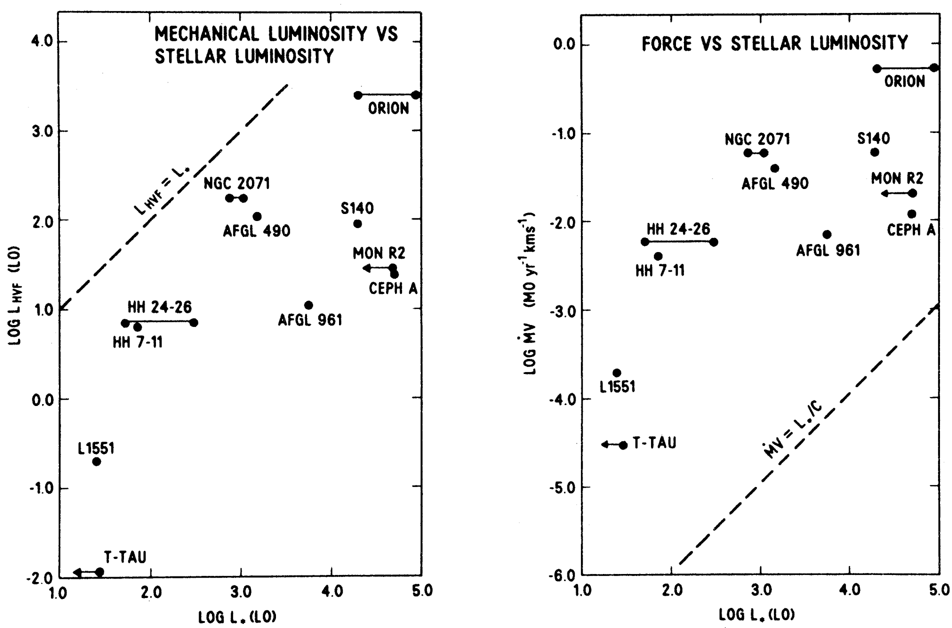

Fig. 2. Plots of the flow mechanical luminosity (left panel) and momentum (right panel) against the bolometric luminosity of the embedded infrared source. [Reproduced from Bally and Lada (1983).]

The pioneering work of Bally and Lada (1983) clearly presented one of the major dilemmas of bipolar flows: what drives them? Plots of mechanical luminosity and momentum against stellar luminosity, L_*, respectively (Fig. 2) show that it is *energetically* possible for the embedded stars to radiatively drive the outflow via single photon scattering. However, there is a momentum deficit of up to a factor of 10^3. Moreover, multiple photon scattering does not save the day because the required optical depths are far too large (see Padman *et al.* 1991). The Bally–Lada diagrams are not immediately pertinent when assessing the energy and momentum requirements of driving by hydrodynamic winds, a point which was emphasised by Dyson (1984) and which is discussed in Section 5.

The most prominent mechanisms which have been invoked to explain the observed features of bipolar flows are: (1) centrifugally driven magnetised winds from disks; (2) magnetic pressure driven winds and (3) driving by neutral winds which themselves are centrifugally driven from the star–accretion disk interface. These are the mechanisms which are reviewed in this paper.

3. Centrifugally Driven Winds, Star Formation and Bipolar Flows

A theory for bipolar flows which is intimately related to the star formation process has been developed in a number of papers (Pudritz and Norman 1983, 1986; Pudritz 1985). Some of the initial ideas for this theory come from the papers by Lovelace (1976), Blandford (1976), Hartmann and MacGregor (1982) and Blandford and Payne (1982). The Blandford and Payne paper was especially important in that it showed, through self-similar solutions, that a steady state wind could be produced above a magnetised accretion disc. Magnetic self-pinching of this wind led to the formation of a collimated jet. In this section the Pudritz–Norman theory is reviewed in some detail. This has the added advantage of exhibiting some of the features of centrifugally driven winds in detail. These are relevant to a number of other theories.

(3a) The Pudritz Star Formation Cycle

The essential ideas encompassed in the Pudritz–Norman theory can be summarised by the following cycle (from Pudritz 1985):

(1) Molecular gas in the core of a molecular cloud collapses preferentially along the local field lines forming a rotating disc. The molecular gas at this stage is only loosely coupled to the magnetic field through neutral–ion collisions.

(2) The accretion of mass from the inner disk forms an accretion shock, the UV luminosity of which ionises the disk envelope, and a hydromagnetic wind with a mass-loss rate $\sim 10^{-6}\ M_\odot\ \mathrm{yr}^{-1}$ is initiated. During the same epoch the increased heating of the disk due to the absorption of photons from the accretion region switches the disk into the tightly coupled regime. (In the inner region of the disk a photoionised Strongren region $\sim 10^{15}$ cm in size is produced by photons shortward of 912 Å; in the outer region photons longward of 912 Å heat the disk through absorption by grains.) When the disk material is in the regime where the neutral–ion collision time is much less than the orbital period, the relative drift between the ions and molecules can be ignored in the first approximation and the hydromagnetic wind can be treated as a single component wind.

(3) The magnetised wind of molecular and ionised gas (the bipolar flow) effectively removes angular momentum from the disk promoting the inwards accretion and maintaining the luminosity of the protostar which heats and photoionises the disk in a self-consistent fashion.

(3b) MHD Wind Solutions

Many of the papers in this field make use of the simplification of the MHD equations that are possible for axisymmetric flow (Chandrasekhar 1956; Mestel 1961). The magnetic field and flow velocity are composed of poloidal (subscript p) and toroidal (subscript ϕ) components. The fluid velocity can be written in the form

$$\rho \mathbf{v} = \alpha \mathbf{B} + \rho \mathbf{\Omega} \times \mathbf{R},$$

which consists of a component of the momentum density along the field lines ($\alpha\mathbf{B}$), together with a rotating component ($\rho\mathbf{\Omega} \times \mathbf{R}$). (Here $\mathbf{R}$ is the three-dimensional radius vector.) Qualitatively, therefore, one can see that a hydromagnetic wind driven from a disk which has drawn the magnetic field into an hour-glass pattern may adopt a morphology that is similar to the hour-glass itself. There also exists the possibility of collimation from the increasingly important toroidal field as the inertia of the fluid winds up the toroidal field beyond the Alfven radius (Blandford and Payne 1982; Heyvaerts and Norman 1989; see Section 3g).

In the mathematical treatment of centrifugally driven winds one can make use of three streamline constants; these are

(1) The specific energy:

$$E = \tfrac{1}{2}(v_p^2 + v_\phi^2) + h + \psi - \frac{\Omega r B_\phi}{4\pi\, \alpha},$$

where h is the specific enthalpy and ψ is the gravitational potential.

(2) The angular velocity:

$$\Omega = \frac{1}{r}\left(v_\phi - \frac{\alpha B_\phi}{\rho}\right).$$

(3) The specific angular momentum:

$$L = r\left(v_\phi - \frac{B_\phi}{4\pi\alpha}\right).$$

(In these equations r is the cylindrical radius.)

These streamline constants can be used to derive a single equation (the transfield equation) to determine the flow. However, this equation is extremely complex (albeit with some appealing symmetries) and its solution involves various critical surfaces (the Alfven surface and the fast and slow magnetoacoustic surfaces). Thus, to date, analytical work and order of magnitude estimates have been based upon self-similar solutions and special solutions involving a monopolar magnetic field (as well as the asymptotic analysis of Heyvaerts and Norman 1989). One general feature which is relevant to all flows however is the relation between the Alfven radius, r_A, and the specific angular momentum, namely

$$L = \Omega r_A^2 .$$

This illustrates the consistency of a strong magnetic field (large 'lever-arm' r_A) with a large angular momentum of the hydromagnetic wind. This in turn is consistent with an efficient removal of angular momentum from the disk when the magnetic field is strong. One can visualise this through the bead-on-a-wire analogy (Henriksen and Rayburn 1971; see also Fig. 3). The magnetised fluid is driven along the field lines by the rotating field gaining angular momentum as it does so; eventually the inertia of the fluid becomes important and it 'holds back' the field counter to the direction of rotation. The stronger the magnetic field (the stiffer the 'wire') the larger the distance from the disk at which this happens.

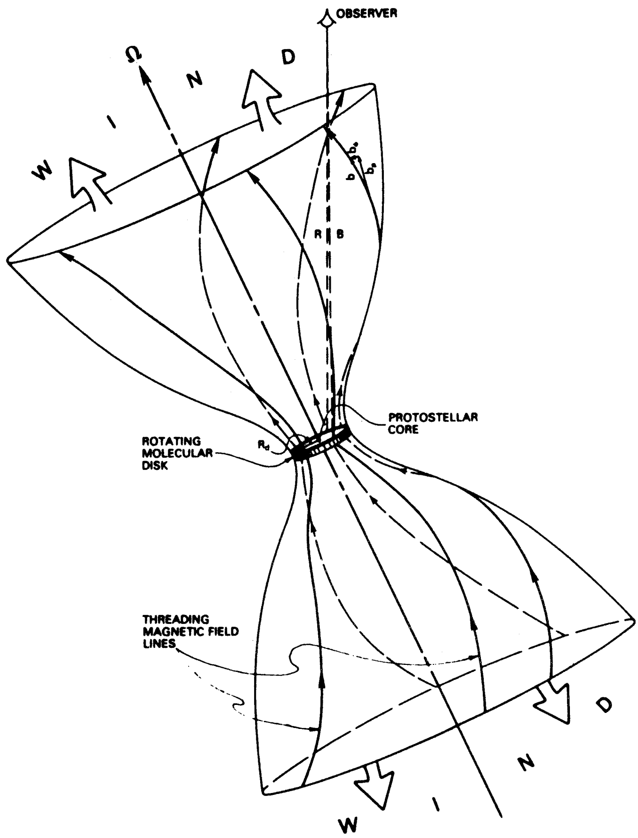

Fig. 3. The conceptualisation by Pudritz and Norman of a bipolar flow driven off a circumstellar accretion disk and centrifugally accelerated by a strong magnetic field. [Reproduced from Pudritz and Norman (1986).]

(3c) Michel's Monopolar Minimum-torque Solution

Because of the difficulties involved in formulating exact analytical solutions, Pudritz and Norman used a special (equatorial) solution for a monopolar magnetically dominated wind, derived by Michel (1969), together with approximate geometrical factors to make allowance for the application of this solution to a polar wind. The following conservation laws apply in this solution:

$$
\begin{aligned}
\rho v_{\mathrm{r}} r^2 &= \frac{\dot{M}_{\mathrm{w}}}{4\pi} = \text{mass loss rate per steradian} = \text{constant}, \\
r^2 B_{\mathrm{r}} &= \Phi = \text{magnetic flux per steradian} = \text{constant}.
\end{aligned}
$$

The wind accelerates through the Alfven point to a terminal velocity given by

$$
\begin{aligned}
V_\infty &= \left(\frac{\Omega^2 \Phi^2}{\dot{M}_{\mathrm{w}}}\right)^{1/3} \\
&= 50 \text{ km s}^{-1} \left(\frac{\Omega}{10^{-11}\mathrm{s}^{-1}}\right)^{2/3} \left(\frac{R_{\mathrm{disk}}}{10^{17}\,\mathrm{cm}}\right)^{4/3} \left(\frac{B}{10^{-2}\mathrm{G}}\right)^{2/3} \left(\frac{\dot{M}_{\mathrm{w}}}{10^{-4} M_\odot \mathrm{y}^{-1}}\right)^{-1/3} .
\end{aligned}
$$

The velocity at the Alfven point is $2V_\infty/3$ and the Alfven point is at

$$
\begin{aligned}
R_A &= \tfrac{3}{2}\left(\frac{\Phi^2}{\Omega \dot{M}_{\mathrm{w}}}\right)^{1/3} \\
&= 2{\cdot}5 \times 10^{18} \text{ cm} \left(\frac{R_{\mathrm{disk}}}{10^{17}\mathrm{cm}}\right)^{4/3} \left(\frac{B}{10^{-2}\mathrm{G}}\right)^{2/3} \left(\frac{\Omega}{10^{-11}\mathrm{s}^{-1}}\right)^{-1/3} \\
&\quad \times \left(\frac{\dot{M}_{\mathrm{w}}}{10^{-6} M_\odot \mathrm{yr}^{-1}}\right)^{-1/3} ,
\end{aligned}
$$

where R_{disk} is the radius of the disk. For these fiducial parameters, which Pudritz and Norman take to be appropriate for star formation regions, it can be seen that the Alfven radius is an order of magnitude larger than the disk radius, as it should be, if the magnetic field is to act as a lever for extracting angular momentum from the disk. The flow field accelerates out to and beyond the Alfven radius and that is the Pudritz–Norman explanation for the observed gradient of the velocity in bipolar flows.

The argument for applying this solution to a disk magnetosphere is that the accretion of material through the disk should pull the magnetic field into a characteristic hour-glass configuration resembling to some extent, in the neighbourhood of the disk, a monopolar field. It should also be realised that the Michel solution is a *singular* solution out of an infinite class of solutions. However, this solution has the property that it minimises the angular momentum for given magnetic flux and mass-loss rate and therefore provides a minimum estimate for the rate of angular momentum loss. An implicit assumption in the use of this solution is that it provides estimates of wind velocities etc. which are indicative and provides the correct scaling of the various parameters.

(3d) The Angular Momentum Loss

Pudritz and Norman (1986) calculate the torque that the wind exerts on the disk from the flux of angular momentum through the Alfven surface:

$$
T_{\mathrm{w}} = -\dot{J}_{\mathrm{disk}} = \int_{\text{Alfven surface}} L\, \rho\, \mathbf{v}_{\mathrm{p,w}} \cdot d\mathbf{A} \approx f_{\mathrm{g}}\, \dot{M}_{\mathrm{w}}\, \Omega R_{\mathrm{a}}^2 ,
$$

where the geometrical factor $f_g \approx \frac{1}{3}$–$\frac{2}{3}$. Again note the importance of a large Alfven radius for the efficient removal of angular momentum from the disk.

(3e) The Relation between the Bipolar Flow and Mass Accretion

One of the features of bipolar wind driven accretion is the large accretion rate that can be produced by the efficient shedding of angular momentum. This is a consequence of the disk angular momentum equation:

$$\frac{\dot{M}_{\mathrm{a}}}{2\pi r^2}\frac{d}{dr}(rv_{\phi,\mathrm{disk}}) = T_\phi \,,$$

where $\dot{M}_{\mathrm{a}}$ is the mass accretion rate and T_ϕ is the torque on the disk due to the magnetic stress. This equation integrates to

$$\dot{M}_{\mathrm{a}} = \tfrac{1}{3}\dot{M}_{\mathrm{w}}\left(\frac{R_A}{R_{\mathrm{disk}}}\right)^2 .$$

Thus for $\dot{M}_{\mathrm{w}} \sim 10^{-4}\, M_\odot\ \mathrm{yr}^{-1}$ and $R_A/R_{\mathrm{disk}} \sim 10$ the mass accretion rate is $\sim 10^{-2}\, M_\odot\ \mathrm{yr}^{-1}$ implying an accretion luminosity

$$L_* = 6\times 10^3 L_\odot \epsilon\left(\frac{\dot{M}_{\mathrm{a}}}{10^{-2}M_\odot \mathrm{yr}^{-1}}\right),$$

where the efficiency factor $\epsilon \sim 0\cdot 1$ depends upon the details of the boundary layer between the inner accretion disc and the protostar.

Pudritz (1985) shows that on a time scale $\sim 10^5$ yr the mass accretion rate decreases to approximately the wind mass-loss rate as the disk is torqued down by the wind. Pudritz also estimates an accretion luminosity $\sim 4\times 10^{23}$ erg s^{-1} which is the accretion shock luminosity required to photoionise and heat the disk and close the loop described above.

The high accretion rate implied by the Pudritz–Norman model is one aspect of the theory which has been strongly criticised by Shibata *et al.* (1987). This problem may possibly be rectified if the dragging in of the magnetic field by the accretion flow is taken into account. Konigl (1989) has shown how an accretion disk magnetic field, subject to both advection and ambipolar diffusion, may be matched to the external Blandford–Payne solution. However, this has not been done in the case of the Pudritz–Norman theory.

(3f) The Momentum in the Hydromagnetic Wind

An appealing feature of the Pudritz–Norman theory is the large momentum flux compared with the luminosity, L, of the core protostar. Since the latter is derived from accretion,

$$L = \epsilon G \dot{M}_{\mathrm{a}}\frac{M_{\mathrm{core}}}{R_{\mathrm{core}}} \,,$$

where $\epsilon \sim 0\cdot 1$ is an efficiency factor which depends upon the details of the boundary layer between the inner edge of the accretion disk and the core. It is readily shown that

$$\frac{\dot{M}_{\rm w} V_\infty}{L_*/c} \sim \frac{(\Omega R_A)\, c}{\epsilon G M_c / R_c}$$
$$= 7000\left(\frac{\Omega R_A}{10\ \mathrm{km\ s^{-1}}}\right)\left(\frac{M_{\rm core}}{M_\odot}\right)^{-1}\left(\frac{R_{\rm core}}{10^{12\cdot 5}\mathrm{cm}}\right).$$

(3g) Collimation by the Toroidal Field

It is inevitable that the inertia of the rotating wind enhances the toroidal field to the point where it is capable of collimating the flow. The field B_ϕ may be expressed in terms of the constants of the motion Ω and L as

$$\frac{B_\phi}{\rho} = \frac{L/\alpha r}{1 - M_A^2} + \frac{r\Omega}{\alpha},$$

where M_A is the Alfvenic Mach number. Beyond the Alfven point in the wind B_ϕ/ρ increases proportionally to radius until the hoop stress becomes important. Blandford and Payne (1982) demonstrated this effect using self-similar solutions. Heyvaerts and Norman (1989), using an asymptotic analysis, have shown that it is a general property of all flows that they collimate at some point.

(3h) Massive Disks

In order for the Pudritz–Norman theory to be viable it is necessary that the disk surrounding the protostar be massive since this is the ultimate source of energy which drives the flow (via the intermediary of the magnetic field). This generally implies masses of the order of 10–100 $\mathrm{M}_\odot$ with rotational velocities of the order of a few to 10 km s^{-1}. As far as I am aware, what observations there are of disks around protostars generally support these estimates. However, the statistics are not substantial enough to make a definitive statement.

(3i) Jets

The entire description of the Pudritz–Norman theory above has been directed towards their theory of the bipolar *molecular* flows. In their theory there is also a second ionised component of the wind which originates from a region $r < 10^{15}$ cm, this being the extent of the photoionised region around the star. This wind, which they essentially identify with the jets, is also centrifugally accelerated by the magnetic field and acquires a larger speed $\sim$200 km s^{-1} through larger values of angular velocity and magnetic flux in the inner part of the disk and a smaller mass-loss rate ($\sim 10^{-6}\ M_\odot\ \mathrm{yr}^{-1}$). However, Pudritz and Norman do not describe how this ionised flow becomes better collimated than the more extended molecular flow. Moreover, the mass fluxes in the jets are generally smaller than $10^{-7}\ M_\odot\ \mathrm{yr}^{-1}$.

4. Magnetic Pressure Driven Winds

The other class of models in which magnetic effects are dynamically important is magnetic pressure driven models. These originated from the work of Draine

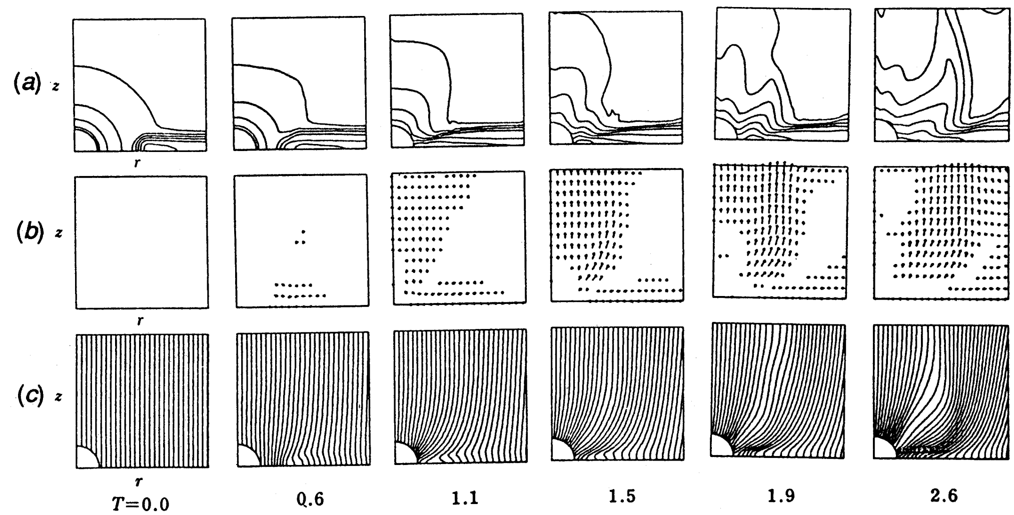

Fig. 4. Results from the simulations of Uchida and Shibata (1985). Sequences (a), (b) and (c) represent the density, the velocity and the rz projection of the magnetic field lines. The disk material is accelerated upwards due to the gradient in $B_\phi^2/8\pi$ produced by the winding up of B_ϕ by the rotational velocity.

(1983). More recent work has involved numerical calculations by Uchida and Shibata (1985) and semi-analytic approximations by Lovelace and colleagues.

(4a) The Uchida–Shibata Simulations

The Uchida–Shibata (1985) solutions involve the toroidal magnetic field which in the inner parts of the accretion disk, in particular, is wound up to the extent that the gradient of $B_\phi^2/8\pi$ becomes dynamically important. The gradient is mainly axial and therefore drives a bipolar flow perpendicular to the disk. Fig. 4 is reproduced from Uchida and Shibata (1985) and shows the density, velocity and magnetic field projected onto the rz plane from their simulation. However, it is evident that the Uchida–Shibata simulations only represent a *transient* in the possible establishment of a bipolar flow. More recent attempts to follow such solutions beyond this initial state have met with difficulties: runaway accretion results and the simulations have to be terminated. This phenomenon may be related to the assumption of ideal MHD. It seems to be necessary for the magnetic field, which continues to be wound up at the base of the transient jet, to be able to relax somewhat and this could be accomplished through the introduction of some magnetic diffusion (M. L. Norman, personal communication 1991). The runaway accretion may also be related to a recently discovered instability treated in two papers by Balbus and Hawley (1991) and Hawley and Balbus (1991). This is a shear instability within the disk itself which leads to exponentially growing angular momentum transport and corresponding runaway accretion.

As is to be expected on the general grounds presented earlier, the Uchida–Shibata solutions collimate above the disc, presumably due to the influence of the dynamically important toroidal field.

(4b) The Calculations of Lovelace et al.

Lovelace and colleagues (see Lovelace *et al.* 1991 and references therein) have also been considering the effect of dynamically important magnetic fields on the production of jets and winds. However, the physics is quite different from the Uchida–Shibata approach in that it is the *poloidal* component of the field which is important rather than the toroidal component. Moreover, rather than rely on computationally intensive numerical simulations, they derive conservation laws for the flow which are used to provide a set of *ordinary* differential equations for averaged flow variables which they solve numerically, taking into account the various critical points in the flow.

The starting point for the Lovelace *et al.* model is a strong poloidal field in the vicinity of the star which is produced by the advection of magnetic field by the flow in the accretion disk. This field is responsible for driving the flow through the slow magnetosonic point and thereafter it accelerates through the Alfven and fast magnetosonic points through the effect of the external pressure gradient. The main thrust of the paper is to examine the jet part of the flow; this is matched onto the disk (accretion) flow via conservation laws for mass, angular momentum and energy. At this stage a comprehensive exploration of parameter space has not been published but one representative solution of the flow equations is presented. This has the characteristics that the mass flux in the jets is 0·1 times the accretion rate; the jet is initially poorly collimated but collimates at greater distances from the star due to the effect of the external pressure; and there is a substantial transport of angular momentum away from the star by the jet. In general there is a toroidal field present in these solutions but it does not have a collimating effect because the net current is zero. Other consequences include the oscillation of the jet radius at large distances from the star.

Most of the Lovelace *et al.* solution is presented in dimensionless form. In order to recover numerical estimates of quantities such as velocity and magnetic field one can use the following expressions to scale the velocity and magnetic field respectively:

$$v_0 = \left(\frac{GM_*}{a_i}\right)^{1/2} \mathcal{F}_M^{1/2} \delta^{-1/2} ,$$

$$B_0 = \left(\frac{\dot{M}_{\rm jet} v_0}{a_i^2}\right) \mathcal{F}_M^{1/2} ,$$

where $a_i \sim 2\times10^{11}$ cm is the initial jet radius, $\mathcal{F}_M$ is the dimensionless mass flux in the jet and δ is a dimensionless measure of the gravitational force. For the parameters of the representative solution Lovelace *et al.* derive a terminal jet velocity $\approx$490 km s^{-1} and a magnetic field $\approx$85 G at the base of the jet, for a jet mass flux $\sim10^{-7}$ $M_\odot$ yr^{-1} and a central mass of 1 $M_\odot$. The velocity is very close to the velocity $\approx$440 km s^{-1} inferred for the ionised wind in L1551 (Cohen *et al.* 1982; Stocke *et al.* 1988). The Lovelace *et al.* approach, which is not numerically intensive and which incorporates a large part of the essential physics, is clearly a productive way of investigating the physical ideas embodied in the model.

An important point to note in comparison to the Pudritz–Norman model is that the Lovelace *et al.* jets originate much closer to the star, that is at around 10^{11} cm rather than $\sim 10^{15}$ cm. In this respect, the terminal velocity of the Lovelace *et al.* jets reflects the depth of the potential well from which they have been driven:

$$
\begin{aligned}
v_\infty &\sim \left(\frac{2GM_*}{a_i}\right)^{1/2} \\
&\approx 500 \text{ km s}^{-1}\left(\frac{M_*}{M_\odot}\right)^{1/2}\left(\frac{a}{10^{11}\text{cm}}\right)^{-1/2}
\end{aligned}
$$

5. Bipolar Flows Driven by Ionised or Neutral Winds

(5a) Energy Driven and Momentum Driven Flows

So far I have discussed theories which essentially rely upon the physics of disk magnetic fields in order to produce bipolar flows. In the case of the Pudritz–Norman theory, it is the field in the outer parts of the disk which is important; in the Uchida–Shibata picture it is the field in the inner parts of the accretion disk which drives the flow.

A different class of models involves the driving of ambient molecular material outward by an ionised or partially ionised wind from the star itself and early theories along these lines were developed by Norman and Silk (1980) and Beckwith *et al.* (1983). It was Dyson (1984) who made the important distinction between *momentum driven* and *energy driven* flows and the implication of this distinction to the interpretation of data in the Bally and Lada (1983) paper.

The distinction between momentum driven and energy driven flows can be made in the following way: when we consider a bipolar flow driven by a wind we are in essence considering the dynamics of a wind-driven bubble. [See Weaver *et al.* (1977) for the case of a bubble inflated by the wind from a blue giant star.] The bipolar flow corresponds to the shell of ambient material which is swept up by the expanding bubble. In the interior of the bubble the wind from the star is supersonic and passes through a shock since it is required to be brought almost to rest by the more slowly swept-up shell. In an *energy driven* bubble, cooling is unimportant, the wind shock is an adiabatic one and the bubble is driven outwards by the pressure of the hot, shocked gas, sandwiched between the shell and the terminal wind shock. In a *momentum driven* bubble, cooling is significant, the wind shock is radiative and the region of shocked stellar wind collapses so that the bubble is driven outwards directly by the ram pressure of the wind. The transition between momentum driven and energy driven flows is given by the criterion that the cooling of the wind is unimportant in a dynamical time scale. This leads to a critical wind velocity (Dyson 1984)

$$
v_{\text{crit}} \approx 230 \text{ km s}^{-1}\left(\frac{n_{\text{ext}}}{10^3 \text{ cm}^{-3}}\right)^{1/9}\left(\frac{\dot{M}_{\text{w}}}{10^{-6}M_\odot \text{ yr}^{-1}}\right)^{1/9},
$$

where n_{ext} is the density external to the star. Dyson's main point is that in an energy driven flow ($v_{\text{w}} > v_{\text{crit}}$), the momentum of the wind (Π_{w}) and the bipolar flow (Π_{f}) are related by

$$\Pi_{\rm w} = \left(\frac{77 + 29\beta + 2\beta^2}{9}\right)\left(\frac{v_{\rm f}}{v_{\rm w}}\right)\Pi_{\rm f} ,$$

where $n_{\rm ext} \propto r^{\beta}$ and $v_{\rm f}$ and $v_{\rm w}$ are the velocities of the bipolar flow and wind respectively. Thus, if the ratio of the flow velocity to wind velocity is small enough, the momentum required of the wind is correspondingly diminished. The physical reason for the large ratio of shell to wind momentum is that it is the *pressure* of the hot shocked wind that is driving the shell of circumstellar material and the origin of this pressure is the conversion of mechanical *energy* into thermal energy at the wind shock. Thus it may be possible for a high velocity, ionised wind to drive a bipolar flow even though it apparently has a large momentum deficit, provided that the driving is energetically feasible.

Although Dyson's proposed solution to the momentum problem is a clever one, it fails on observational grounds. As we have seen, the ionised winds have velocities which are at most a factor of 2 higher than the critical velocity and are a factor of at least a few below those inferred by Dyson from application of his model to a number of bipolar flows. In other words, the cooling of the stellar wind is important. Nevertheless, the distinction between energy and momentum driven flows is an important one and the concept pervades much of the subsequent work.

(5b) Neutral Winds

The above discussion leads to the conclusion that if a stellar wind is to be the driving agent of a bipolar flow then the flow is momentum driven and one is still left with the problem of identifying the source of the momentum. The notion that a neutral or partially ionised wind may provide the driving force receives strong support from the HI observations of Lizano *et al.* (1988) which reveal the existence (in HH7–11) of a *neutral* wind with parameters $\dot{M}_{\rm w} \approx 3 \times 10^{-6}\, M_\odot\ {\rm yr}^{-1}$ and $v_{\rm w} \approx 140\ {\rm km\ s^{-1}}$. Interpreting the core of the HI line as implying a total mass loss of $\sim 0 \cdot 06\ M_\odot$ implies that the total momentum in the associated bipolar flow can be provided by the wind. Stocke *et al.* (1988) also inferred the presence of a 'second wind' in L1551 by interpretating the spectra of Herbig–Haro objects in L1551 as arising from the interaction of the central jet with a surrounding wind. The parameters of the L1551 wind ($\dot{M}_{\rm w} > 1 \cdot 5 \times 10^{-6}\, M_\odot\ {\rm yr}^{-1}$ and $v_{\rm w} \sim 160\ {\rm km\ s^{-1}}$) imply that it can provide sufficient momentum to drive the associated bipolar flow. Stocke *et al.* note that the wind may arise from the central infrared source (IRS5), from a second star in the system or from a circumstellar disc.

Thus, in the class of theories considered in this section, the focus for bipolar flows shifts from the bipolar flow itself to the neutral (more probably partially ionised) wind from the star itself. The only existing theory at the moment is the X-celerator mechanism of Shu *et al.* (1988) (see Fig. 5). The main components of their theory are (1) a protostar rotating nearly at breakup and (2) a strong ($\sim$150 G) magnetic field rotating with the star. The reason for the maximally rotating configuration is as follows. In a normal star, the sonic point must occur close to the photosphere of the star if a large mass loss rate is to be sustained. However, Shu *et al.* argue that in a protostar the thermal speed in the atmosphere is only a few percent of the escape velocity and is therefore incapable of driving the flow. If the star is maximally rotating, then the effective potential at the

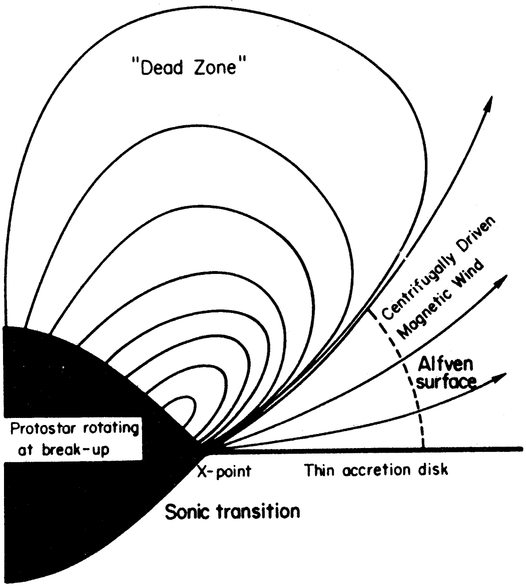

Fig. 5. A schematic representation of the driving of a strong neutral wind from a maximally rotating protostar by a strong magnetic field. [Reproduced from Shu *et al.* (1988).]

surface is reduced and a larger mass-loss rate can be sustained through the equatorial region. Essentially the process is one of Roche–Lobe overflow through the inner Lagrangian point of the potential. Shu *et al.* then invoke a centrifugally driven magnetised wind once the gas is outside the Lagrangian point. If we take the value of the terminal velocity implied by the Michel solution as indicative, then

$$\begin{aligned} v_\infty &= \left(\frac{\Omega^2\Phi^2}{\dot{M}_{\rm w}}\right)^{1/3} \\ &= \left(\frac{2GM_* R_* B^2}{\dot{M}_{\rm w}}\right) \\ &\approx 200\ {\rm km\ s^{-1}}\left(\frac{M_*}{0{\cdot}5M_\odot}\right)^{1/3}\left(\frac{R_*}{5R_\odot}\right)^{1/3}\left(\frac{B}{100\ {\rm G}}\right)^{2/3}\left(\frac{\dot{M}_{\rm w}}{10^{-6}M_\odot\ {\rm yr}^{-1}}\right)^{-1/3}, \end{aligned}$$

where expressions and parameters relevant to a maximally rotating protostar are used. Thus it can be seen that for a mass-loss rate $\sim 10^{-6}\ M_\odot\ {\rm yr}^{-1}$ a wind with the velocity observed in HH7–11 is obtained. A mass-loss rate of this magnitude is produced in the Shu *et al.* theory if the specific angular momentum ($h_{\rm out}$) of the matter in the outflowing wind is comparable with the specific angular momentum ($h_{\rm in}$) of the matter accreted from the circumstellar disk. The mass-loss rate is then

$$\dot{M}_{\rm w} \approx f\dot{M}_{\rm a}\,,$$

where

$$f = \frac{1-2b}{\bar{J}-2b}\,,$$

with

$$b = \frac{\mathcal{L}_*/\mathcal{M}_*}{(\mathcal{G}\,\mathcal{M}_*\,\mathcal{R}_*)^{1/2}}\,, \qquad \bar{J} = \frac{h_{\rm out}}{h_{\rm in}}\,,$$

where $\dot{M}_{\rm a}$ is the accretion rate from the disk, $\mathcal{L}_*$ and $\mathcal{M}_*$ are the angular momentum and mass of the protostar and $h_{\rm out}$ and $h_{\rm in}$ are the specific angular momenta of the outgoing and incoming material being accreted from the circumstellar disk. The physics described by the above equations is relatively straightforward: the excess angular momentum of the accreting material is carried off by the wind. The set of mass loss equations is closed by the expression for $\bar{J}$ in a centrifugally driven wind:

$$\bar{J} = \tfrac{3}{2} + \tfrac{1}{2}v_\infty^2\,.$$

Shu *et al.* apply their theory to the observations of HH7–11 and obtain quite reasonable consistency between the observed $\dot{M}_{\rm w} \approx 3 \times 10^{-6}\,M_\odot\,{\rm yr}^{-1}$, $\dot{M}_{\rm disk} \approx 1 \times 10^{5}\,M_\odot\,{\rm yr}^{-1}$ and the calculated value of $\bar{J} \approx 2{\cdot}5$. It should be remarked that this success appears to depend more upon the characteristics of a centrifugally driven wind than on the details of the way in which the mass loss is initiated.

Shu *et al.* state that their theory (the detailed version of which is unpublished at the time of writing) does not appear to be capable of producing jets. Magnetic focussing can produce enough collimation to produce the observed CO flows but apparently is unable to produce sufficient focussing to produce a jet. Shu *et al.* speculate that the jets may arise as an 'ordinary' wind from the 'dead zone' above the poles of the star which blows off the otherwise closed magnetic field lines and is collimated by the surrounding 'extraordinary wind' produced by the X-celerator mechanism.

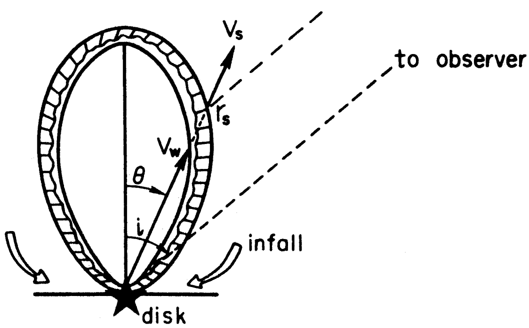

Fig. 6. A schematic representation of the dynamics of a bubble driven by a stellar wind. The apparent acceleration of the bipolar flow is a geometrical effect due to the elongation of the bubble. [Reproduced from Shu *et al.* (1991).]

(5c) The Characteristics of the Resultant Bipolar Flow

In a subsequent paper, Shu *et al.* (1991) derived the velocity field that would result from the driving of a bubble by *any* wind (including a centrifugally driven disk wind) into a surrounding partially flattened molecular cloud core. In this case the bubble is momentum driven and their explanation for the apparent acceleration of the bipolar flow as one moves away from the YSO is the velocity field produced by the expansion of the bubble (see Fig. 6). The highest velocities are observed at the extremities of the bubble and the lowest velocities are on the parts of the bubble closest to the protostar.

There is observational support for both parts of this theory. First, the theory explains the linear increase of velocity away from the driving star as observed. Second, when resolved, the CO is found to have a shell-like morphology and, third, the observed masses of bipolar flows (e.g. $\sim 1\,M_{\odot}$ in L1551) are typically close to the mass which would have been swept up by the wind, given the density of the molecular cloud and the volume of the flow. [See Snell and Schloerb (1985), Uchida *et al.* (1987), Moriarty-Schieven *et al.* (1987), Moriarty-Schieven and Snell (1988) for observational papers on these last two points.]

6. Discussion

In this review I have concentrated on what are probably the major ideas surrounding the production of jets and bipolar flows in star formation regions. As we have seen, each theory has its own successes and shortcomings, and future developments will be interesting. Nevertheless, I think there is one idea which is extremely attractive and that is the notion that bipolar flows are driven by a separate wind. This appears to be entirely consistent with several pieces of observational data. If this is the case then the original idea of the Pudritz–Norman theory, that the bipolar flow may originate from the disk itself, is unattractive. Nevertheless, the notion that disks may produce centrifugally driven winds is still an appealing one in view of the way in which this can lead to angular momentum loss from the star–disk system and a scaled down version of the Pudritz–Norman theory in which, say, the 'second wind' inferred for L1551 is a centrifugally driven disk wind may be quite relevant. Indeed, one of the central issues in this field is the relation between the low mass flux, high velocity, ionised jets and the surrounding wind. Do they both originate from close to the star or is the wind more extensive? As far as the former type of theory is concerned, Lovelace *et al.* and Shu *et al.* have addressed different aspects of the same problem. However, the connection between the two is not obvious. Likewise, the Uchida–Shibata simulations, which were directed at disk winds, may be more relevant to winds driven off the central protostar. To my mind, an observational point which requires clarification is whether jets and bipolar flows always exist simultaneously (as in L1551) or whether jets appear as pre-main-sequence stars evolve. The latter seems to be the case at present but could well be a selection effect due to the larger optical depth surrounding the youngest objects.

Acknowledgment

I would like to acknowledge useful discussions with Mike Norman during the preparation of this review.

References

Balbus, S. A., and Hawley, J. F. (1991). *Astrophys. J.* **376**, 214.
Bally, J. (1982). *Astrophys. J.* **261**, 558.
Bally, J., and Lada, C. J. (1983). *Astrophys. J.* **265**, 824–47.
Beckwith, S., Natta, S., and Salpeter, E. E. (1983). *Astrophys. J.* **267**, 596.
Blandford, R. D. (1976). *Mon. Not. R. Astron. Soc.* **176**, 465.
Blandford, R. D., and Payne, D. G. (1982). *Mon. Not. R. Astron. Soc.* **199**, 883.
Chandrasekhar, S. (1956). *Astrophys. J.* **124**, 232.
Cohen, M., Bieging, J. H., and Schwarz, P. (1982). *Astrophys. J.* **253**, 707.
Draine, B. T. (1983). *Astrophys. J.* **270**, 519.

Dyson, J. E. (1984). *Astrophys. Space Sci.* **106**, 181–97.
Hartmann, L., and MacGregor, K. B. (1982). *Astrophys. J.* **259**, 180.
Hawley, J. F., and Balbus, S. A. (1991). *Astrophys. J.* **376**, 223.
Henriksen, R. N., and Rayburn, D. R. (1971). *Mon. Not. R. Astron. Soc.* **152**, 323.
Heyvaerts, J., and Norman, C. (1989). *Astrophys. J.* **347**, 1055.
Königl, A. (1989). *Astrophys. J.* **342**, 208.
Lada, C. J., and Harvey, P. M. (1981). *Astrophys. J.* **245**, 58.
Lizano, S., *et al.* (1988). *Astrophys. J.* **328**, 763.
Lovelace, R. V. E. (1976). *Nature* **262**, 649.
Lovelace, R. V. E., Berk, H. L., and Contopoulos, J. (1991). *Astrophys. J.* **379**, 696.
Mestel, L. (1961). *Mon. Not. R. Astron. Soc.* **122**, 473.
Michel, F. C. (1969). *Astrophys. J.* **158**, 727.
Moriarty-Schieven, G. H., and Snell, R. L. (1988). *Astrophys. J.* **332**, 364.
Moriarty-Schieven, G. H., Snell, R. L., Strom, S. E., Schloerb, F. P., Strom, K. M., and Grasdalen, G. L. (1987). *Astrophys. J.* **319**, 742.
Mundt, R. (1986). *Can. J. Phys.* **64**, 407.
Norman, C., and Silk, J. (1980). *Astrophys. J.* **238**, 158.
Padman, R., Lasenby, A. N., and Green, D. A. (1991). *In* 'Beams and Jets in Astrophysics'. (Ed. P. A. Hughes) (Cambridge University Press).
Pudritz, R. E. (1985). *Astrophys. J.* **293**, 216.
Pudritz, R. E., and Norman, C. A. (1983). *Astrophys. J.* **274**, 677.
Pudritz, R. E., and Norman, C. A. (1986). *Astrophys. J.* **301**, 571.
Raga, A. C., Canto, J., Binette, L., and Calvet, N. (1990). *Astrophys. J.* **364**, 601.
Rodriguez, L. F., Ho, P. T. P., and Moran, J. M. (1980). *Astrophys. J. (Lett.)* **240**, L149.
Shibata, K., Shu, F. H., Adams, F. C., and Lizano, S. (1987). *Ann. Rev. Astron. Astrophys.* **25**, 23.
Shu, F. H., Lizano, S., Ruden, S. P., and Najita, J. (1988). *Astrophys. J. (Lett.)* **328**, L19.
Shu, F. H., Ruden, S. P., Lada, C. J., and Lizano, S. (1991). *Astrophys. J. (Lett.)* **370**, L31.
Snell, R. L., Loren, R. B., and Plambeck, R. L. (1980). *Astrophys. J.* **233**, 127.
Snell, R. L., and Schloerb, P. F. (1985). *Astrophys. J.* **295**, 490.
Stocke, J. T., Hartigan, P. M., Strom, S. E., Strom, K. M., Anderson, E. R., Hartmann, L. W., and Kenyon, S. J. (1988). *Astrophys. J. (Suppl.)* **68**, 229.
Uchida, Y., and Shibata, K. (1985). *Publ. Astron. Soc. Japan* **37**, 515.
Uchida, Y., Kaifu, N., Shibata, K., Hayashi, S. S., and Hasegawa, T. (1987). *In* 'Star Forming Regions' (Eds M. Peimbent and J. Jagaku), p. 287 (Reidel: Dordrecht).
Weaver, R., McCray, R., Castor, J., Shapiro, P., and Moore, R. (1977). *Astrophys. J.* **218**, 377.

Manuscript received 19 December 1991, accepted 19 February 1992