STAR FORMATION IN MOLECULAR CLOUD CORES

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ABSTRACT. The problem of gravitational collapse and star formation is entirely different when the ratio of the mass of a molecular cloud $M_{\rm cl}$ to its magnetic flux Φ is high than when it is low. Magnetically-diluted overall collapse of a large dense core and the formation of an OB association or a bound cluster are the likely outcomes in the former case; quasi-static contraction of many small cores and their ultimate collapse to form a T association, in the latter. In our picture, the birth of a T association in a dark cloud like Taurus proceeds by ambipolar diffusion on a time-scale of $\sim 10^7$ years. As magnetic and turbulent support is gradually lost from a small condensing core, it approaches a state resembling a slowly rotating singular isothermal sphere which, when it passes the brink of instability, collapses from "inside-out," building up a central protostar and nebular disk. The emergent spectral energy distributions of theoretical models in this stage of protostellar evolution resemble closely those of recently found sources with steep spectra in the infrared. The protostellar phase is ended by the reversal of the infall by an intense stellar wind, whose ultimate source of energy derives from the differential rotation of the star. We argue that the initial breakout is likely to occur along the rotational poles, leading to collimated jets and bipolar outflows. The stellar jet eventually widens to sweep out gas in nearly all 4π steradian, revealing at the center a T Tauri star and a remnant nebular disk. We give rough scaling relations which must apply if an analogous process is to succeed for producing high mass stars.

1. INTRODUCTION

Our paper adopts a scenario for star formation which is complementary to that reviewed by Professor Hayashi in this volume. We begin with the assumption that magnetic fields play an important dynamical role in the interstellar medium (see, e.g., Mestel and Spitzer 1956, Mestel 1965; Mouschovias 1981), and we concentrate on clouds where, on the collapse scale of protostellar masses, the outer boundary conditions can essentially be considered to lie at infinity. The latter approximation is motivated by the observation that a typical molecular cloud has a mass which much exceeds the totality of the masses of the stars currently forming from it.

The strength of the magnetic fields needed to provide appreciable mechanical support of molecular clouds against their self-gravity can be ascertained by virial theorem arguments (Strittmatter 1966), or more accurately, by detailed model calculations (Mouschovias and Spitzer 1976). A magnetic flux Φ can support a

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cloud provided its mass $M_{\rm cl}$ does not exceed a critical value approximately given by

$$M_{\rm cr} = 0.15 \frac{\Phi}{G^{1/2}}. \tag{1}$$

This formula implies that a 30 μG field could support a clump of mass $10^3\ M_\odot$ if its radius were 2 pc; or a giant molecular cloud of mass $10^5\ M_\odot$ if its radius were 20 pc. Direct evidence that magnetic fields of such magnitude are present in (self-gravitating) molecular clouds has been obtained in the Zeeman splitting measurements of Cruchter and Kazes (1985), Troland and Heiles (1985) and Kazes and his colleagues (this conference). Indirect evidence that magnetic fields are dynamically important can be deduced from the interstellar polarization measurements of Vrba, Strom, and Strom (1976) and Tamura (this conference), which show that the direction of the magnetic field is well correlated over the entire dimension of a dark cloud and therefore is strong enough to resist bad tangling by any chaotic velocity field ("turbulence") that may be present. We regard these twin lines of evidence, *magnitude* and *direction* of $\mathbf{B}$, as powerful indicators that magnetic fields play a central role in the structure and evolution of star-forming regions.

As Mestel has emphasized for many years (see Mestel 1985), the evolution of supercritical clouds (or pieces of them) with $M_{\rm cl} > M_{\rm cr}$ is very different from subcritical clouds with $M_{\rm cl} < M_{\rm cr}$. In the supercritical case, when the mass to flux ratio is high, we have magnetically-diluted collapse of the cloud as a whole (e.g., Scott and Black 1980); in the subcritical case, when the mass to flux ratio is low, we have quasi-static contraction by ambipolar diffusion (e.g., Nakano 1979).

The gravitational collapse and fragmentation of an extremely supercritical cloud, with $M_{\rm cl} \gg M_{\rm cr}$, must resemble the field-free calculations summarized by Professor Hayashi at this conference (see also Bodenheimer 1980 and Tohline 1982). It is therefore tempting to speculate that supercritical regions generally have higher star-formation efficiencies than subcritical ones: in the simplest examples, supercritical clouds would form bound clusters (e.g., Rho Ophiuchus); subcritical ones, T associations (e.g., Taurus). An exceptional circumstance may arise if the rate of ambipolar diffusion becomes very enhanced when $M_{\rm cl}$ is sufficiently greater than $M_{\rm cr}$ in a large and dense molecular cloud core (which may itself be fragmented into many small cores). The extra frictional heating may so elevate the temperature T_0 of the large core as to favor only the formation of high mass stars (for reasons which will become clear later in our discussion). The appearance of massive luminous stars may then prove so disruptive to the cloud as a whole that an unbound OB association results rather than an open cluster (cf. Lada, this symposium).

In the present paper, we shall focus our attention on the evolution of a molecular cloud which is initially subcritical, but which has a mass (supported by a combination of magnetic fields and subalfvenic "turbulence") that much exceeds the Jeans mass M_J computed on the basis of the *average* density and temperature. We now argue that the magnetic support of such an object will automatically lead to the production of many small cores.

2. THE PRODUCTION OF SMALL CLOUD CORES

If magnetic fields do provide support for a molecular cloud, then it is automatic that such a cloud (or clump) will develop subclumping, *i.e.*, small cloud cores. The reason is that magnetic support of self-gravitating matter of low fractional ionization necessarily involves slip of the neutral component of that matter relative to the field (ambipolar diffusion; Mestel and Spitzer 1956; Nakano 1979, 1982; Mouschovias and Paleologou 1981; Scott and Black 1980), leading to the gradual concentration of the neutrals into subcondensations where the magnetic field plays increasingly less of a part in the total support.

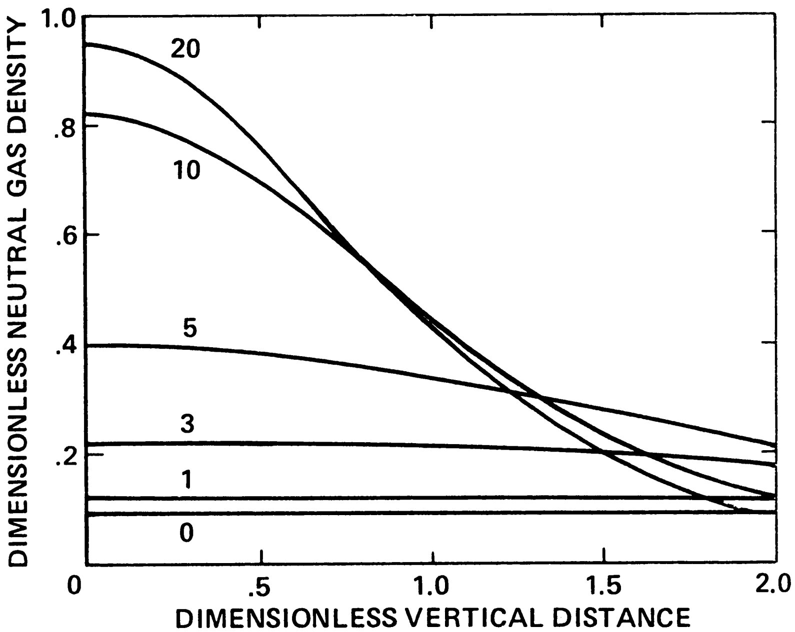

Figure 1. A slab of lightly ionized gas initially supported by magnetic fields against its self-gravity relaxes by ambipolar diffusion to a state supported only by thermal pressure (Shu 1983).

Shu (1983) has solved a simple idealized problem which illustrates the physical principle (see Fig. 1): the quasi-magnetostatic support of a plane-parallel self-gravitating slab of isothermal gas which is lightly ionized. As time increases, the magnetic field wants to decay to its background value (here, 0), and the gas wants to acquire the configuration that applies in the absence of a magnetic field, *i.e.*, support against self-gravity by thermal pressure alone. With a large initial ratio of magnetic to thermal pressure, the outcome is the production of a small dense core in the background of a more extended envelope.

The 1-D problem is unrealistic because the gravitational field of a slab saturates at a finite value proportional to the total surface density. In a more realistic 3-D problem, where the original cloud clump contains many Jeans masses, but M_{cl} is still less than the magnetic mass M_{cr}, we may expect the clump to fragment quasi-statically into many "cores," each core asymptotically trying to reach a configuration where the field lines become almost uniform and straight (the background condition) and where the total velocity dispersions approach purely thermal values. This picture provides an attractive evolutionary explanation for the small quiet cores which are observed in the Taurus dark cloud by Myers and Benson (1983).

Stable asymptotic support by thermal pressure alone in quasi-spherical symmetry is impossible if the potential reservoir of matter is very large in comparison with a Jeans mass in the initial state; thus, 3-D isothermal cores must ultimately undergo gravitational collapse when they become sufficiently centrally condensed. This has apparently happened in the Taurus region to roughly half of Myers and Benson's cores, which are known to be the sites of T Tauri star formation. The theoretical elucidation of how and when quasi-static evolution is tranformed into dynamical collapse is an ongoing project of Lizano and Shu (in progress). It should also be noted that although the core formation process may be quasi-static, the epoch of major flux loss probably occurs during the dynamical collapse phase (Mouschovias, Paleologou, and Fiedler 1985).

A rough estimate for the ambipolar diffusion time t_{AD} required for the quasi-static phase of cloud core formation is provided by the 1-D calculations; t_{AD} equals the sound travel time across the scale height of the core in the final state, multiplied by a coefficient which depends, among other things, the degree of ionization of the medium. Thus, ambipolar diffusion can be considered a slow evolutionary process only to the extent that this coefficient is a pure number large compared to unity. For the conditions which apply typically in the envelopes of molecular clouds, the coefficient *is* large in comparison with unity, of the order of 10 or 100, and this probably explains why star formation is generally an inefficient process in present-day clouds. On the time scale of the dynamical collapse of unstable molecular cloud cores (the sound travel time across the core, $\sim 10^5$ y), star formation is a badly synchronized affair in molecular cloud clumps because the time it takes to form the many individual cores takes $\sim$ 10-100 times longer. This means that newly-formed T Tauri stars have ample time to turn on their winds and disrupt the incipient condensation of neighboring cores, making star formation in T associations a loosely aggregated affair. *There is no guarantee, however, that the coefficient is large compared to unity in all circumstances.* Low fractional ionizations or high initial mass to flux ratios may well lead to relatively high-star formation efficiencies.

Slow slippage of magnetic field lines in the disk of the Galaxy provides two additionally welcome features to the theory of star formation in T associations. First, a slow initial phase of gravitational condensation almost guarantees magnetic braking (Gillis, Mestel, and Paris 1974, 1979) of the cores of molecular clouds to rotation rates characteristic of their envelopes (Mouschovias and Paleologou 1979, 1980) before the cores undergo gravitational collapse. There will also be a tendency for angular velocity vector $\mathbf{\Omega}$ to align with the direction of the clump magnetic field (Strom 1985) because the perpendicular component of $\mathbf{\Omega}$ will brake more quickly than the parallel component. Thus, these molecular cloud cores will contain a rich supply of material with relatively low angular momentum from which to form binary stars and planetary systems (Mouschovias 1978). Second, the slow evolution of these objects toward a state of gravitational instability may give well-defined initial states for dynamical collapse calculations.

3. INSIDE-OUT COLLAPSE OF UNSTABLE CLOUD CORES

In the absence of more detailed calculations, it is fruitful to speculate on what radial density profile is likely to result for a molecular cloud core as a consequence of the quasi-static process of ambipolar diffusion sketched in the last section. Nearly twenty years ago, Bodenheimer and Sweigart (1968) pointed out that subsonic evolution of a non-rotating, non-magnetic cloud would always produce a $\rho \propto r^{-2}$ density distribution (see also Larson 1969). This is interesting because Shu (1977) demonstrated later that the singular isothermal sphere,

$$\rho = \frac{a^2}{2\pi G r^2}, \tag{2}$$

has a self-similar collapse solution, whose form, apart from the integration of some simple ordinary differential equations, could be found analytically. Because the central density of the singular isothermal sphere is infinite, it might appear that starting conditions for the collapse problem based on this approximation are rather artificial (Whitworth and Summers 1985). However, it should be remembered that Nakano (1981) found that when one includes grain coupling to ambipolar diffusion calculations (Elmegreen 1979, 1985; Nakano and Umebayshi 1980), very high densities (compared to envelope values) can be reached in the central regions of a molecular cloud core before dynamical collapse is generated. As long as the power-law part of a real core solution spans several decades in density, except for initial transients, its gravitational collapse is likely to be well represented by the self-similar solution for the singular isothermal sphere (generalized, perhaps, to include the perturbational effects of the finite amounts of rotation and magnetic fields left by the process of diffusive core formation described earlier).

Note that if equation (2) is correct, there may be *no such thing* as a *typical density* or a *typical mass* for a molecular cloud core, because a power law has no typical scale. Thus, when the core undergoes gravitational collapse, the solution does not automatically define a mass scale that we may associate with the formation of ordinary stars. What is well defined instead is a *rate* at which the central object is built up by accretion; in the self-similar solution for the singular isothermal sphere, this rate equals

$$\dot{M} = m_0 a^3 / G, \tag{3}$$

where $m_0 = 0.975$. For $a = 0.2$ km s^{-1} (Taurus), $\dot{M}$ has the value $2 \times 10^{-6} M_\odot$ y^{-1}; for $a = 0.35$ km s^{-1} (Rho Ophiuchus), $1 \times 10^{-5} M_\odot$ y^{-1}.

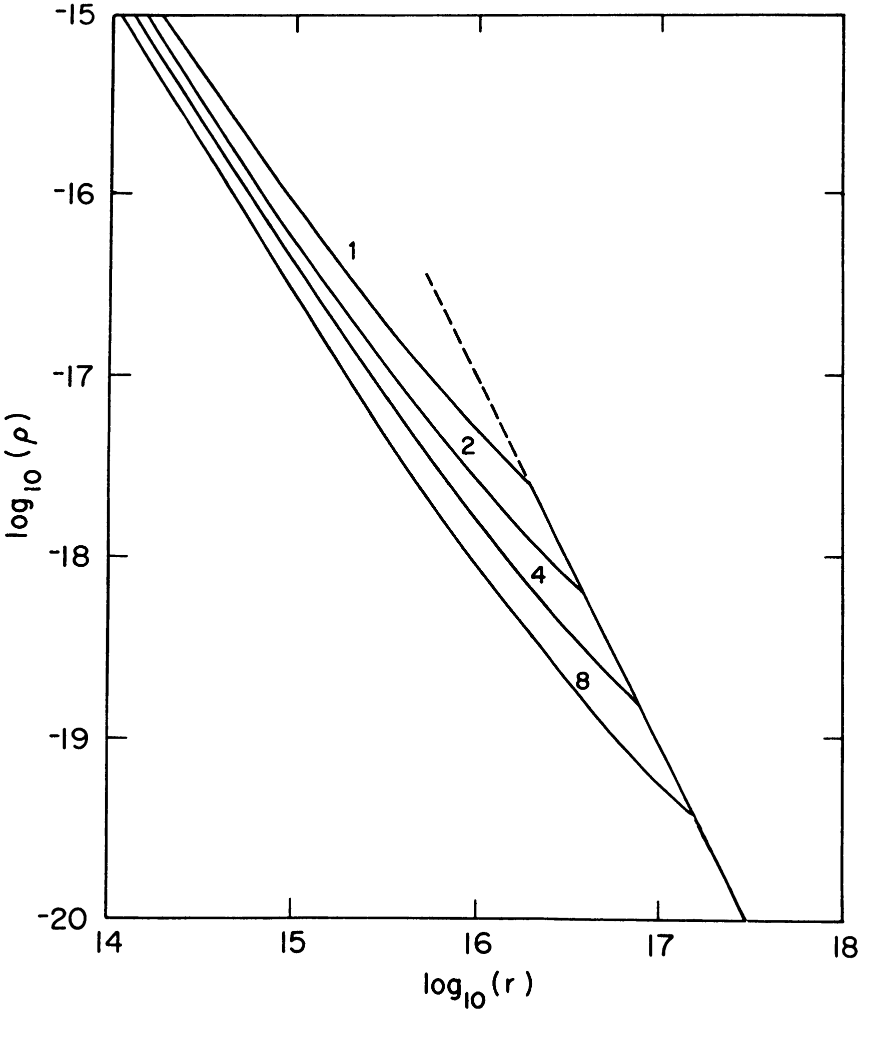

Figure 2. Similarity solution for the collapse of a cloud with an initial density profile appropriate for a singular isothermal sphere, $\rho \propto r^{-2}$ (from Shu 1977). The specific solution shown corresponds to a choice $a = 0.2$ km/s.

In the self-similar collapse solution, infall is initiated from "inside-out" by an ex-

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pansion wave which propagates outward at the speed of sound into static material with the density distribution (2). Figure 2 shows that the head of the expansion wave reaches radius $r_h = at$ in time t; and supersonic inflow velocities are generated interior to $r_s \approx 0.4\,at$, inside of which the density begins to approach the free-fall form $\rho \propto r^{-3/2}$. Identifying $t = M/\dot{M}$ as the time required to build up the central star to a mass M at the mass infall rate (3), we easily calculate that $r_s \approx 7 \times 10^{16}$ cm for $M = 0.5\ M_{\odot}$ and $a = 0.2$ km s^{-1}. Since infall velocities for $r < r_s$ scale approximately as $r^{-1/2}$, the detection of infall velocities exceeding, say, 0.5 km s^{-1}, will require linear resolutions of $\sim 10^{16}$ cm. This is only becoming possible now with a new generation of millimeter-wave telescopes, which may explain why radio astronomers have previously only found the much larger regions of high velocity *outflows* that are associated with a later phase of stellar evolution than the pure protostellar stage.

4. SPHERICAL PROTOSTELLAR EVOLUTION

The isothermal approximation begins to break down badly in low-mass protostars when the infalling dust and gas becomes optically thick to the emergent infrared radiation interior to a radius r_e of about 10^{14} cm. Fortunately, for r less than the radius of the dust photosphere, r_e, and greater than the stellar radius, $R_* \sim 10^{11.5}$ cm (bounded by an accretion shock), the material falls nearly freely toward the mass at the center, so that the crossing time for either matter or radiation is short compared to the evolutionary time, and a condition of steady-state flow holds to a high degree of approximation. The inner region of complicated radiative hydrodynamics may thus be solved (using carefully chosen closure relations for the frequency-integrated moment equations of the radiation field) as a set of ordinary differential equations, greatly alleviating the computational demands of the problem (Stahler, Shu, and Taam 1980a,b, 1981). The similarity solution discussed in the previous section provides outer boundary conditions for the inner problem, namely, inflow at free-fall speeds with mass infall rate (3).

For an adopted infall rate $\dot{M} = 1 \times 10^{-5}\ M_{\odot}\ y^{-1}$, the results showed that the protostar accumulated matter processed through an accretion shock of ever increasing specific entropy; thus, the star remained radiative until deuterium ignited near the center when the stellar mass was about 0.3 $M_{\odot}$. A convection zone then spread outward through the star until it became almost entirely convective at a mass of about 0.5 $M_{\odot}$. Apart from this event, nothing dramatic happened to distinguish a particular mass scale for the accreting protostar, and in the actual calculations, the infall was artificially shut off after 10^5 y when the star had accumulated 1 $M_{\odot}$. The surface of the star, which had been kept abnormally hot and luminous by a standing shock, then cooled in less than a day and joined a convective pre-main-sequence track (Hayashi, Hoshi, and Sugimoto 1962). The disappearance of the infall region would first make the star optically visible at this point. If the loci of such points in the Hertzsprung-Russell diagram for different shut-off times (*i.e.*, masses) are joined, we have a "birthline" for pre-main-sequence stars of low mass; Stahler (1983) showed the birthline corresponding to spherical mass accretion at the rate $\dot{M} = 1 \times 10^{-5} M_{\odot}\ y^{-1}$ gave a remarkable fit to the upper envelope for T Tauri stars in Taurus-Auriga, Orion, NGC 7000/IC 5070, and Ophiuchus.

The good agreement between the theoretical birthline and the observed one can be criticized on a number of points. First, it is unlikely that the infall would be purely spherical since even a small amount of initial rotation in the original molecular cloud core would have produced accretion partially through a disk. Second, we now know that while the adopted infall rate of $1 \times 10^{-5}\ M_{\odot}\ y^{-1}$ is appropriate for Ophiuchus, it is 5 times too large for Taurus. A scaling law $R_* \propto \dot{M}^{1/3}$, which can be derived heuristically and is confirmed by comparisons of the numerical results of Stahler, Shu, and Taam with those of Winkler and Newman (1980), would yield a spherical-collapse birthline for Taurus which is underluminous by a factor of about 3.

However, such disparities are relatively small, and given the ample observational evidence concerning the importance of rotation in the formation of binary stars and planetary systems, the controversy over the exact location of the "birthline" seems premature (cf. Mercer-Smith, Cameron, and Epstein 1984; Stahler 1984). At the present stage of the development, we should be grateful that the simple spherical theory works as well as it does. Its biggest failing is that it does not give an account of why the infall should stop when it does, and therefore, it provides no natural mass scale for the formation of stars. Disk accretion *per se* does no better in this regard, for as long as infall is sustained, matter continues to accumulate either in the star or in the disk.

5. INFRARED EMISSION FROM ROTATING PROTOSTARS

To test the fundamental notion that stars do form from a process of infall from a molecular cloud core, it is desirable to compute the expected emergent spectral energy distributions of protostars. Numerous groups have performed such calculations for the spherically symmetric problem (Larson 1969, Yorke and Krugel 1977, Bertout and Yorke 1978, Yorke and Shustov 1981), and Adams and Shu (1985, 1986) have discussed a fast approximate technique applicable to multi-dimensional situations as well. Provided the parameters of the low-mass protostellar models are chosen appropriately, synthetic spectral energy distributions can be constructed which yield good agreement with observations of a class of low-luminosity infrared sources, those which have steep spectra and are found near the centers of dense molecular cloud cores (see Fig. 3). The mass accretion rates needed to account for the overall luminosity and spectral shape are consistent with $\dot{M} = 2 \times 10^{-6}\ M_{\odot}\ y^{-1}$ in Taurus, and $\dot{M} = 1 \times 10^{-5}\ M_{\odot}\ y^{-1}$ in Ophiuchus. Moreover, to obtain the observed levels of emission shortward of $\lambda = 10\ \mu m$, we found it often necessary to adopt initial rates of angular rotation for the molecular cloud core of $\Omega = 1 \times 10^{-14}\ s^{-1}$ or larger. Such rates are compatible with measured values, lending support to the whole theoretical development.

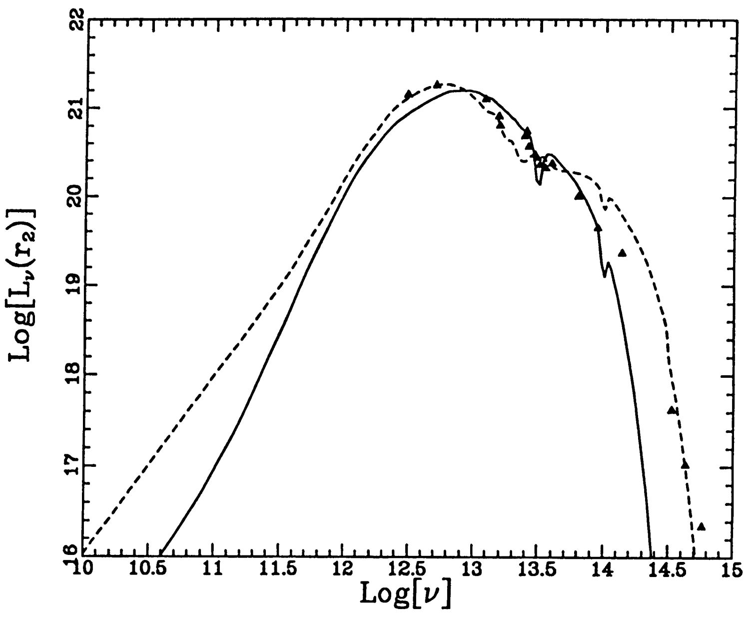

Figure 3. Comparison of emergent spectral energy distribution of Haro 6-10 (data points from Myers *et al.* 1986) and theoretical rotating models (Adams and Shu 1986) with $a = 0.2$ km/s and $\Omega = 10^{-14}$ rad/s (solid curve) or $\Omega = 10^{-13.5}$ rad/s (dashed curve).

The shape of the emergent spectral energy distributions turns out to be especially sensitive to the infallen mass M, cloud sound speed a, and cloud rotation rate Ω, through a combination defined by the centrifugal radius (Cassen and Moosman 1981; Terebey, Shu, and Cassen 1984):

$$R_C \equiv G^3 M^3 \Omega^2 / 16a^8. \tag{4}$$

The quantity R_C corresponds to the position where infalling matter in the equatorial plane encounters a centrifugal barrier if it conserves its initial specific angular momentum as it approaches the star and disk. The ratio of stellar radius R_* to centrifugal radius R_C controls the fraction of all incoming matter which impacts the star directly rather than hits the disk; thus, apart from the scale factor (3), R_*/R_C largely sets the total extinction to the central source, giving the column density of matter which reprocesses the stellar and disk photons to radiation of longer wavelengths.

Both the overall luminosity and the drop in infrared emission at the shorter wavelengths are exaggerated in the spherical models compared to the real objects. Nonspherical models, which include the radiation from accreting nebular disks that arise when the effects of rotation are included, give better results in these two regards: the luminosity can be lower for a given infall rate because some of the gravitational binding energy released can be stored as kinetic energy of (differential) rotation, and the disk's copious emission in the mid- and near-infrared can supplement the reprocessing of stellar optical photons by the dust envelope to radiation at these wavelengths.

The rotating infall models even give acceptable fits for some infrared sources which are known bipolar outflow sources. This finding supports the idea that well-collimated sources represent objects in which inflow and outflow are taking place *simultaneously*, indeed, that these amazing objects (Schwartz 1983, Cohen 1984, Lada 1985, Welch *et al.* 1985, Mundt 1985) represent exactly the transitional phase of evolution between a purely accreting protostar and a fully revealed pre-main-sequence star (Shu and Terebey 1984; Cassen, Shu, and Terebey 1985). Sketched below is an outline of how rotation might play the crucial role for energizing and collimating such a stellar wind (see also Konigl 1982). The accompanying conjecture of how the inflow is eventually reversed closes the logical loop concerning the physical processes which ultimately define the mass of the central star.

6. JETS AND BIPOLAR FLOWS

Although no fully self-consistent account of the formation of a rotating protostar by a combination of direct infall and disk accretion is yet available, it is easy to comment on some general features. First, as long as the star's interior remains radiative, it is likely to rotate *differentially*. (The incoming material can hardly know in advance to have a distribution of angular momentum which corresponds to uniform rotation when it is incorporated inside a star.) The store of energy represented by the differential rotation will be a healthy fraction of the total gravitational binding energy of the star, which, in turn, is about $\sim 10^3$ times larger than the kinetic energy $E_{\rm HVF}$ needed to drive CO high velocity flows (as can be calculated from Bally and Lada 1983) , *i.e.*,

$$E_{\rm HVF} = \alpha \frac{GM_*^2}{R_*}, \tag{5}$$

where $\alpha \sim 10^{-3}$. However, as long as the star has no internal circulation, there is no obvious mechanism to tap the energy contained in differential rotation. But once fluid does begin to circulate across shear layers, dynamo mechanisms should operate that amplify magnetic fields (Parker 1979) which can eventually pump energy into driving a stellar wind and coincidentally account for the high levels of surface activity that characterize, for example, the atmospheres of the T Tauri stars to be revealed at a later stage.

As long as low-angular momentum material falls in first (Terebey, Shu, and Cassen 1984), the protostar is automatically built up with a *mechanically stable* stratification – a distribution of specific angular momentum which increases outward. Any internal circulation which arises, therefore, has to be driven thermally (convection in low-mass stars, perhaps Eddington-Sweet circulation in high-mass stars). The natural time scale associated with overturning a convectively stable entropy gradient (by deuterium burning in a low-mass star or with meridional circulation currents in the radiative envelope of a rapidly rotating high-mass star) is some fraction β of the Kelvin-Helmholtz time. If we identify this with the duration $t_{\rm HVF}$ of the high-velocity flow, we obtain

$$t_{\rm HVF} = \beta \frac{GM_*^2}{R_* L_*}, \tag{6}$$

where $\beta \sim 10^{-1}$ represents a reasonable estimate both from the observational and theoretical points of view. Division of equation (5) by equation (6) then reproduces the observational correlation that the flow luminosity $L_{\rm HVF} \equiv E_{\rm HVF}/t_{\rm HVF}$ is roughly a constant fraction, $\sim 10^{-2}$, of the photon luminosity of the central source L_* over several decades of L_*:

$$\frac{L_{\rm HVF}}{L_*} = \frac{\alpha}{\beta}. \tag{7}$$

The above discussion, of course, does not represent so much a *derivation* of Figure 16 of Bally and Lada (1983; see also Snell, this volume) as an *interpretation* of the well-known striking correlation between $L_{\rm HVF}$ and L_*. For the spherical models of low-mass protostars, convection does not occur in a major way until the onset of deuterium burning. If the same statement applies to models which incoporate a combination of direct infall and disk accretion, then we would have a fundamental understanding of why stars usually form with a mass scale which is roughly marginal for thermonuclear fusion, namely, they keep on accreting mass until their interior temperatures ($\propto M_*/R_*$) rise enough to burn at least the most easily fused of all nuclear species, deuterium. And for an infall rate of 2×10^{-6} $M_\odot$ y^{-1}, this will occur typically when the star has a few tenths of a solar mass.

How would breakout of the stellar wind in a low-mass protostar occur in practice? When deuterium ignites (slightly off-center in the spherical models), buoyant motions across shearing layers would probably try to bring the material in the convection zone to a state of more nearly uniform rotation, releasing energy in the process. Some part of this energy will go into amplifying seed magnetic fields; another part, to overturning the stable entropy gradient and driving the outer boundary of the convection zone upwards. However, as long as there is still a substantial outer radiative envelope, the heat released (and the magnetic fields generated) cannot make its way to the surface, and the radius of the star will expand. If the clearing time scale is shorter then the thermal time scale (which seems to be the case from estimates of the age of bipolar flows), one will deduce an expanded radius for the star; perhaps it is this expansion which reconciles the observed "birthline" for T Tauri stars with the relatively low mass-accretion rate in Taurus of $\sim 2 \times 10^{-6}$ $M_\odot$ y^{-1}. In any case, when the outer radiative zone becomes thin enough, thermal equilibration of magnetic flux tubes (Parker 1984) will occur rapidly, enabling them to buoy out to the radius of the standing infall

shock. There they will pile up, like steam in a pressure cooker with the lid on, until the magnetized plasma builds up enough head to break out.

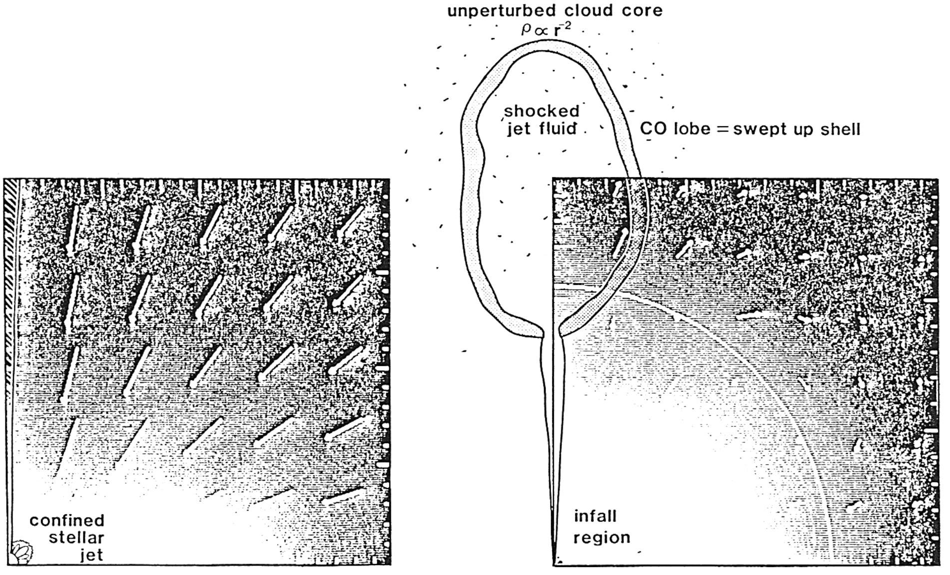

Figure 4. **Cartoon depiction of the collimation of a protostellar jet (left panel) which shocks and decollimates on much larger scales to produce a CO lobe (right panel). The cartoon is superimposed on a calculation of the collapse flow of the slowly rotating core of a molecular cloud (Terebey, Shu, and Cassen 1984).**

Just as in a pressure cooker, however, breakout will not immediately blow off the entire lid (the ram pressure of the inflow); it will occur first through the channel of least resistance (the safety valve). Figure 4 shows a cartoon depiction of the idea that the weakest point occurs at the rotational pole(s) of the accreting protostar. Thus, the gases of the incipient stellar corona should initially be channeled up and out of the rotational poles of the system, forming two highly collimated stellar jets. The heads of the jet will plow through the infalling cloud envelope, eventually sweeping up enough material to account for the CO lobes of the observed bipolar outflows. The prediction is, therefore, that highly collimated outflow sources should generally correspond to deeply embedded objects, in which outflow and inflow are occurring *simultaneously*. This prediction is in accord with the finding that bipolar outflow sources often have steep spectral energy distributions in the infrared which make them virtually indistinguishable from (protostellar) candidates for pure infall.

As time proceeds, more and more of the rotating inflowing matter falls preferentially on the disk rather than on the star. The "lid" of the pressure cooker will weaken relative to the "steam," and the stellar jet will widen, eventually to sweep outwards over all 4π steradians and bring an end to the infall phase. An outside observer will then be able to see optically a T Tauri star, newly born on a convective (Hayashi) pre-main-sequence track, still bubbling with residual surface activity (Herbig 1962; Kuhi 1964; Calvet, Basri, and Kuhi 1984), and surrounded perhaps by a remnant nebular disk (Adams, Lada, and Shu 1986). The slow rotation rates observed for T Tauri stars (Vogel and Kuhi 1981), which came as a major surprise when first discovered, may now be naturally attributed to the large magnetic braking associated with an earlier period of intense mass loss.

7. THE MASSES OF FORMING STARS

The above picture also yields a scaling for the conditions which might lead to the production of high mass stars. The strength of the lid of the "pressure cooker" can be measured, in some sense, by the rate $\dot{M}_*$ at which matter falls directly onto a star. When the stellar radius $R_* \ll R_C$, $\dot{M}_*$ is very nearly equal to $(R_*/2R_C)\dot{M}$ (see, e.g., eq. [24] of Adams and Shu 1986). With $\dot{M}$ and R_C given by equations (3) and (4), we now have

$$\dot{M}_* = \frac{8m_0 R_* a^{11}}{G^4 M^3 \Omega^2}. \tag{8}$$

We assume that the stellar mass M_* is proportional to M, the total mass which has collapsed to the central regions,

$$M_* = \gamma_1 M, \tag{9}$$

and we consider a stellar wind of mass loss rate $\dot{M}_w$ and with terminal velocity comparable to the escape velocity from the surface of the star. For it to reverse direct infall with mass addition rate $\dot{M}_*$, we expect $\dot{M}_w$ has to be comparable to $\dot{M}_*$ (see, e.g., the discussion of Terebey and MacGregor 1986),

$$\dot{M}_w = \gamma_2 \dot{M}_*. \tag{10}$$

If the winds roughly conserve energy while driving bipolar outflows, and if equation (7) applies,

$$\dot{M}_w G M_*/R_* = \gamma_3 \frac{\alpha}{\beta} L_*. \tag{11}$$

The combination of equations (8)-(11) implies that the final properties of the protostar, defined by when its incipient wind is capable of reversing the direct infall, satisfy

$$L_* M_*^2 = 8m_0 \gamma_1^3 \frac{\gamma_2}{\gamma_3} \left(\frac{\beta}{\alpha}\right) \frac{a^{11}}{G^3 \Omega^2}. \tag{12}$$

If the parameters γ_1, γ_2, and γ_3 are of order unity, and $\beta/\alpha \sim 10^2$, then equation (12) predicts that the Taurus cloud ($a \approx 0.20$ km/s and $\Omega \sim 10^{-14}$ to $10^{-13.5}$ rad/s) should typically form stars with a few tenths of a solar mass. The Rho Ophiuchus cloud, with $a \approx 0.35$ km/s, should form stars with comparable masses because the typical rotation rate deduced by Adams, Lada, and Shu (1986) from spectral modelling is roughly an order of magnitude larger in Ophiuchus than in Taurus. It would be important to observationally check this theoretical prediction. If the prediction holds, it is interesting to speculate that the increase in mean rotation rate arises because less time is available for magnetic braking when $M_{\rm cl} > M_{\rm cr}$. It is tempting also to associate the difference between T associations and OB associations to a substantially larger velocity dispersion in the regions which form the latter. The increase of a is due, in part perhaps, to the heating produced by rapid ambipolar diffusion in massive and dense molecular cloud cores. However, it is necessary to check observationally whether the mean cloud rotation rate Ω is higher in the latter (because $M_{\rm cl} > M_{\rm cr}$) by an amount which offsets the increase in the velocity dispersion of the star forming regions.

Clearly, the discussion of this last section has been extremely speculative, and we should not be surprised if future developments impose substantial modifications. However, even if it only contains qualitative truth, then it should be noticed that this picture supplies possible mechanistic bases for the idea of "bimodal star formation" (Gusten and Mezger 1981, Larson 1985) and for "triggers" of starburst activity (Lo *et al.*, this volume, 1986; Wynn-Williams, this volume).

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STROM: If I understand you correctly, the location of the birthline depends on the accretion rate. In turn, the accretion rate depends on the sonic velocity of the cloud. That suggests an empirical test: a correlation between the location of the birthline and the sonic velocity of the parent cloud. I should note that preliminary results comparing Lupus and Taurus suggest that the birthline in Lupus lies ~ 5 times lower in luminosity than in Taurus (if the distances of the associations are indeed nearly identical). Unfortunately, I am unaware, as yet, of molecular line observations which would permit a comparison.

SHU: If the accretion is purely spherical and if the inflow is shut off arbitrarily, then the luminosity of the "birthline" at each mass (for low mass stars) is proportional to $a^{2/3}$. This would make it difficult to get a birthline lower than the Taurus one, because the value of a there is already nearly thermal, and it is difficult to lower T below 10 K. Thus, there are, at least, two possible ways out of the difficulty you raise: (a) the influence of rotation and the turning of the ignition of deuterium burning in setting the mass of the forming star and its location in the H-R diagram when it first becomes optically visible, or (b) the distance to Lupus as you hint may have been underestimated by a factor of about 2.

ZINNECKER: In your picture, the mass of a star is determined by the onset of the wind, right? Now, there are stars of different masses, in particular there are many more stars born with $\sim 0.5\ M_\odot$ than with $\sim 1\ M_\odot$. How do you account for this in your picture? Secondly, with an r^{-2} singular isothermal sphere density profile, we agree that there is no preferred mass scale; but the r^{-2} profile cannot hold all the way out, it must eventually flatten to a constant back-ground density. Doesn't that define some mass scale?

SHU: The typical mass at which deuterium burning turns on (and therefore inducing convection) for $a = 0.2$ km s^{-1} (T = 10 K) is a few tenths of a $M_\odot$, not 1 $M_\odot$. To make higher mass stars, one probably has to increase $\dot{M}$, e.g. by increasing the temperature of the cloud core, or by implosion. The rotation rate Ω will also affect the timing of when the wind is able to reverse the inflow and, therefore, determine the mass of a forming star. 2) I agree that the turnover from r^{-2} to a shallower density profile could be used to define a mass scale, but this mass is typically greater than that of the star which ultimately forms; moreover, one still needs to involve some mechanism (e.g. winds) to prevent the outer mass from falling into the central regions.