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STAR FORMATION IN MOLECULAR CLOUDS: OBSERVATION AND THEORY

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1. OVERVIEW

Stars are the fundamental objects of astronomy; thus, the formation of stars constitutes one of the basic problems of astrophysics. Star formation in galaxies is a complex process, which spans roughly 12 orders of magnitude in both mass and linear scale (10^{11} to 10^{-1} $M_{\odot}$, 10^{23} to 10^{11} cm) and involves many diverse physical phenomena. This review, which was written in October and November of 1986, concentrates on those aspects of the problem that occur on the scale of a giant molecular cloud (roughly, 10^{6} $M_{\odot}$ and 10^{20} cm) and smaller. Moreover, we concern ourselves with the problems of present-day star formation, as posed by observations of star-forming regions in our own Galaxy and in other spiral galaxies. We do not discuss the implications of star formation in dwarf galaxies, the possibility of star formation in the cooling flows of hot gas in galaxy clusters and elliptical galaxies, the nature of primordial star formation, the origin of starburst galaxies, and many similarly intriguing topics. Even with such truncation, our subject is a large one, and we are forced to give cursory treatments of many interesting and controversial ideas.

1.1 *Origin of Stellar Masses*

The first issue confronting any fundamental discussion of star formation is the question of the origin of stellar masses. Normal stars are self-gravitating balls of gas whose central temperatures are high enough to sustain fusion reactions. Most of the stars that form in the cosmos have masses of a few tenths of a solar mass and, hence, are only marginally

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capable of burning hydrogen (see also the review chapter by Liebert & Probst in this volume). *Why should the Universe ubiquitously produce objects that have just the right mass range to stably transform the primary product of the Big Bang—hydrogen—into heavier elements?*

There are at least two different approaches to answering this question. The first is to suppose that the result—objects with masses sufficient for thermonuclear fusion—occurs somewhat by chance. Here, the interstellar medium, through some process of hierarchical fragmentation spanning many orders of magnitude in mass scale, happens to preferentially produce objects with masses between 10^{-1} and 10^{2} $M_{\odot}$. The second approach is to suppose that star formation is intrinsically an accretion process, whereby protostellar masses are built up from initially small values, and that this accumulation is eventually halted by some process, possibly a stellar wind, which is itself somehow triggered by the onset of thermonuclear fusion. In what follows, we review the various lines of evidence that bear on the problem of star formation in molecular clouds before returning to evaluate the merits and disadvantages of these two schools of thought.

1.2 *Bimodal Star Formation*

Bimodal star formation, the notion that the birth of low- and high-mass stars may involve separate mechanisms, originated in its modern guise with Herbig (1962) and Mezger & Smith (1977; see also Elmegreen & Lada 1977). As an empirical idea, it has been invoked in an increasing number of contexts—from helping to resolve difficulties with models of the chemical evolution of the Galaxy (Gusten & Mezger 1982) to solving the problem of Oort's (1960) limit through the placement of the missing disk mass in the form of stellar remnants (Larson 1986). In his original work, Herbig noted the existence of star-forming regions like Taurus that contain no stars more massive than ~ 2 $M_{\odot}$, and he contrasted this situation with Orion, which contains young stellar objects (YSOs) of both high and low mass. Citing also the discrepancy in the turnoff and contraction time scales of the Pleiades star cluster, Herbig speculated that star births in a given cloud occur over a period of (at least) a few tens of millions of years, and that stars of low mass are born before those of high mass. Mezger & Smith gave various lines of argument that indicated that the sites of high-mass and low-mass stars may be spatially distinct on a global scale, with the formation of OB stars taking place primarily in large cloud complexes situated in the spiral arms of the Galaxy. The formation of T Tauri stars was envisaged to occur in both small and large clouds distributed throughout the disk, but little direct evidence was available on the issue. Mezger & Smith also left open the question whether one type of star-forming cloud could evolve into the other.

More recent investigations have reinforced these basic findings (see, e.g., Elias 1978a, Cohen & Kuhi 1979, Adams et al. 1983, Stauffer 1984, Dame et al. 1986, Solomon & Sanders 1985, Montmerle 1987, Scoville 1987). In particular, Solomon et al. (1985) divide molecular clouds into a cold population (<10 K) and a warm one (including many clouds with cores >20 K). The cold clouds, which do not contain stars earlier in spectral type than late B, are found to be smoothly distributed throughout the galactic disk. The warm clouds, which tend to be among the largest and most massive, are associated with radio H II regions and appear to be a spiral arm population (see also Stark 1985, Waller et al. 1987, Lo et al. 1987).

Thus, although the idea of a *temporal* sequence in the formation of stars of different masses remains controversial (see, e.g., Stahler 1985), the *spatial* distinction between regions of high-mass and low-mass star formation is well founded. In particular, stars are now known to form only from the densest portions of a molecular cloud, and the properties of such molecular cloud "cores" differ for regions of high- and low-mass star formation (Evans 1978, Rowan-Robinson 1980, Myers 1986; see Section 3.3 below).

1.3 *Initial Mass Function*

The first empirical derivation of the initial-mass function from the observed stellar luminosity function in the solar neighborhood was made by Salpeter (1955), who found that the number of stars born with masses between m and $m+dm$ satisfied

$$f(m)\,dm \propto \xi\,(\log m)\,d\log m \propto m^{-2.35}\,dm, \qquad 1.$$

for stellar masses m between 0.4 and 10 $M_\odot$, with considerable uncertainties in the shape of the mass spectrum outside of this range. The problem was reanalyzed by Miller & Scalo (1979) and by Scalo (1986). Figures 16 and 17 in the latter paper suggest that the function ξ may have *two* local maxima, one at $m \approx 0.3\ M_\odot$ and another at 1.2 $M_\odot$, with the existence and location of the 1.2 $M_\odot$ maximum dependent on the exact assumptions made for the history of stellar birthrates. The importance of a primary maximum at 0.3 $M_\odot$ is the apparent existence, unlike the Salpeter law (Equation 1), of a preferred mass scale around a few tenths of a solar mass; the importance of a second maximum at 1–2 $M_\odot$ is that it is a possible indicator of the process of bimodal star formation. Given the uncertainties present in transforming luminosity functions into initial-mass functions, however, the subtle presence of a second maximum should be regarded with some caution.

1.4 *Spontaneous and Induced Star Formation*

The idea that OB star formation is somehow *triggered* by external events has gained wide currency, and the evidence that environmental effects on a galactic scale do play a role seems fairly compelling (see, e.g., de Jong & Maeder 1977, Elmegreen & Elmegreen 1983, Persson 1987). What remains controversial is whether phenomena such as spiral structure, merging galaxies, etc., merely yield conditions that are conducive for thc agglomeration of giant complexes of molecular clouds, which subsequently undergo spontaneous star formation (see, e.g., Shu 1985, 1987, Scoville et al. 1986), or whether the larger scale structure is itself (sometimes) produced by induced star formation (e.g. Seiden 1983, Elmegreen 1985a).

It has been proposed that star formation might spread like an infectious disease through an individual cloud (and, perhaps, even through an entire galaxy) as a wave of gravitational collapse is propagated by the compression associated with supernova shocks (e.g. Mueller & Arnett 1976, Cameron & Truran 1977, Gerola & Seiden 1978, Herbst & Assousa 1978) or ionization fronts (e.g. Elmegreen & Lada 1977, Klein et al. 1983, Ho et al. 1986) or stellar winds (e.g. Norman & Silk 1980, Cameron 1985). These mechanisms have various degrees of a priori plausibility; however, they also share a few common shortcomings. First, except for that of Elmegreen & Lada, they predict the collapse of protostars to begin from "outside-in," whereas the currently available evidence suggests that the process occurs from "inside-out" (Ho & Haschick 1986, Walker et al. 1986, Keto et al. 1987, Reid et al. 1987). Second, they have limited ability to induce gravitational collapse if magnetic fields play a major role in the support of the parent clouds (see Section 3.1). And third, they require the prior presence of mature stars; thus, they implicitly invoke the existence of some *spontaneous* mechanism, in any case, to provide a primary generation of stars.

1.5 *Binary Stars, Bound Clusters, and Hierarchical Fragmentation*

The majority of stars in the Galaxy are found in binary or multiple systems (Batten 1973). A particularly intriguing aspect of binary stars is the apparent dichotomy of primary and secondary masses when the periods P are longer than or shorter than ~ 100 yr (Abt & Levy 1976; see also Abt 1983). For $P > 100$ yr, the mass of the secondary is unrelated to that of the primary; hence, there are many more companions of low mass rather than of high mass (comparable to that of the primary), as if the two stars were formed by independent processes. Thus, the probability distribution for the mass of the secondary is the same as that for a field star. On the other hand, for binaries with $P < 100$ yr, there are disproportionate numbers of high mass ratios (≈ 1) compared to low mass ratios ($\ll 1$), as if the formation of the companion was influenced by the formation of the primary, so that the outcome is weighted toward obtaining stars of more nearly equal mass.

Abt & Levy (1976) interpreted their finding to imply that systems with $P > 100$ yr formed by independent condensation (or capture), whereas systems with $P < 100$ yr formed by fission. Given the observational selections that enter at intervals of 10^2–10^3 yr, however, the nominal dividing point of 100 yr probably should not be taken too literally.

Theorists generally distinguish between fission and fragmentation on the basis of whether the splitting of the two (or more) components takes place in a state of (unstable) equilibrium or against a dynamically collapsing background. The stability of rapidly rotating and self-gravitating equilibria against fission has been studied for both incompressible and compressible fluids (for reviews, see Chandrasekhar 1969, Tassoul 1978). Evolutionary simulations of fairly coarse numerical resolution even seemed to suggest that the outcome is two bodies of nearly equal masses (Lucy & Ricco 1979). However, more recent studies (Durisen et al. 1986) demonstrate that rotationally unstable polytropes tend to shed rings and disks (via the gravitational torques of bars and spiral structure) that absorb the excess angular momenta of the systems and prevent the central bodies from fissioning in the manner envisaged by earlier workers. Equally devastating for the classical fission hypothesis are the observational findings that T Tauri stars start their pre-main-sequence contractions with fairly small radii (Cohen & Kuhi 1979) and with fairly low rotation speeds (Vogel & Kuhi 1981, Bouvier et al. 1986, Hartmann et al. 1986), so that fission from quasi-static configurations into two bodies with orbital separations of AU scales would seem to be impossible. A more attractive scenario for the origin of binaries with periods $< 10^2$–10^3 yr would now seem to be formation from the dust and gas disks of $\sim 10^2$ AU that are being discovered prolifically around young stars (see Section 6.2).

The origin of binaries with periods between a few hundred and a few million years, which constitute about 25–30% of the total population, may require the (single-stage) fragmentation of a rapidly rotating and collapsing cloud core (e.g. Boss 1986) on a scale of 10^{17} cm or smaller. Tidal capture in a dense stellar cluster is another possibility (e.g. Press & Teukolsky 1977), but the small fraction ($\sim 10\%$) of all stars likely to be formed in bound clusters makes this an unlikely origin for all except the widest binaries [which are then subject to disruption within the same cluster (Heggie 1975)].

Hoyle's (1953) concept of hierarchical fragmentation—the isothermal collapse of a single Jeans mass of (nonrotating and nonmagnetic) gas followed by successive rounds of dynamical fragmentation—has not been realized in detailed numerical simulations (Tohline 1982, and references therein), which tend to produce a single accreting, centrally condensed object. Including the effects of rotation alleviates some of the difficulties (see the reviews by Bodenheimer 1980, Hayashi 1986, Boss 1987). Also, rapid decoupling of the principal source of mechanical support against self-gravity by shock compression (Woodward 1976, Klein et al. 1985) or catastrophic cooling (D. N. C. Lin, private communication) would produce, effectively, many Jeans masses in the pre-gravitational-collapse configuration and thereby lead to fragmentation. The value of the minimum fragment mass, variously estimated at 0.003 to 0.01 $M_\odot$ (Low & Lynden-Bell 1976, Rees 1976, Silk 1977, Boss 1986), is somewhat embarrassing. Appeal could be made to accretion or agglomeration to increase the final mass, but then there are no obvious reasons why stellar values should result from this scenario.

When applied to the problem of present-day star formation, the concept of hierarchical fragmentation suffers from the observational criticism that large-scale free-fall motions in molecular clouds have never been found (see, however, Section 3.2), and from the theoretical criticism that magnetic fields can strongly inhibit fragmentation for subcritical clouds (see Section 3.1). On the other hand, flattening of a supercritical cloud (along field lines) may allow it to fragment into quasi-spherical pieces with sizes given approximately by the vertical thickness of the cloud (Mestel 1965). This latter process may explain the formation of open clusters on a time scale given by magnetically diluted collapse (see Sections 3.2). The field-free hierarchical fragmentation scenario may apply to gaseous configurations (e.g. proto-globular clusters) where magnetic fields are not a factor.

1.6 *Efficiency of Star Formation*

In order to form a bound system, the minimum efficiency of star formation has been estimated to be between 20–50%, depending on the rate of gas dispersal and the degree of central concentration of the system (Lada et al. 1984, Elmegreen & Clemens 1985). The transformation of a large fraction of the gas of a region into stars in a time comparable to the intrinsic dynamical scale must be a relatively rare occurrence, as only $\sim$10% of all stars today are believed to have been born in bound clusters (Roberts 1957, Wilking & Lada 1985). From an analysis of the total rate of star formation in the first Galactic quadrant, Myers et al. (1986) estimate that the overall efficiency of star formation is $\sim$2% (see also Blitz 1978). The total production rate of stars in the Galaxy is often cited as $\sim$3–5 $M_\odot$ yr^{-1}. This value is to be compared with a rate of gas return back to the interstellar medium (primarily in the form of mass loss from evolved stars) of ~ 1–$2\ M_\odot\ \mathrm{yr}^{-1}$ (Knapp & Morris 1985, Jura 1987). Therefore, despite the low overall efficiency of star formation in molecular clouds, an irreversible conversion of interstellar gas into stars seems to be occurring in the Galaxy with a characteristic depletion time of $\sim 10^9$ yr. Whether such a relatively short time scale should be of concern and whether appeal should be made to mechanisms such as galactic infall of fresh gas or different histories for the production of high- and low-mass stars (bimodal star formation) is a matter of some debate (see, e.g., the discussion of Wyse & Silk 1987).

2. PROPERTIES OF MOLECULAR CLOUDS

2.1 *Observed Physical Characteristics*

It is well established that molecular clouds are the principal sites of active star formation in our Galaxy (Zuckerman & Palmer 1974, Burton 1976). A substantial part of the CO emission, and most of the total molecular mass, is contained in giant clouds with masses between 10^5 and $3 \times 10^6\ M_\odot$ and with sizes of several tens of parsecs (Solomon et al. 1979). The precise values are controversial (see, e.g., Blitz & Thaddeus 1980) and depend on the conversion factor for transforming integrated CO intensity into H_2 column density (see the review of van Dishoeck & Black 1987).

Observations of CO emission with moderately high spatial resolution (see Figure 1) indicate that giant molecular clouds (GMCs) are actually cloud complexes, composed of smaller units ("clumps") with masses $M_{\mathrm{cl}} \sim 10^3$–$10^4\ M_\odot$, sizes $R \sim 2$–5 pc, densities $n_{\mathrm{H}_2} \sim 10^{2.5}\ \mathrm{cm}^{-3}$, and temperatures $T \sim 10$ K (Sargent 1977, Evans 1978, Blitz 1978, Rowan-Robinson 1979). It is unclear to what extent the observed clump sizes, masses, and densities are the result of observational selection, since the excitation of ^{12}CO, even with radiative trapping, requires H_2 densities of a few hundred per cm^3 (Kutner & Leung 1985). Subregions within the clumps that contain gas of much higher density—cloud cores—can be mapped in H_2CO (e.g. Evans & Kutner 1976), in CS (e.g. Linke & Goldsmith 1980), in HC_3N (e.g. Avery 1980), in NH_3 (reviewed by Ho & Townes 1983), and even in CO if they are appreciably hotter than their surroundings.

The clumps resemble the nearby dark clouds surveyed by Lynds (1962). A well-known complex of such dark clouds exists in Taurus and contains $\sim 10^4\ M_\odot$ of material actively forming an unbound association of low-mass stars (Cohen & Kuhi 1979, Kleiner & Dickman 1984, Cernicharo et al. 1985). Such clumps have small cloud cores (e.g. Myers & Benson 1983) whose observed sizes and densities depend on the molecular line used for

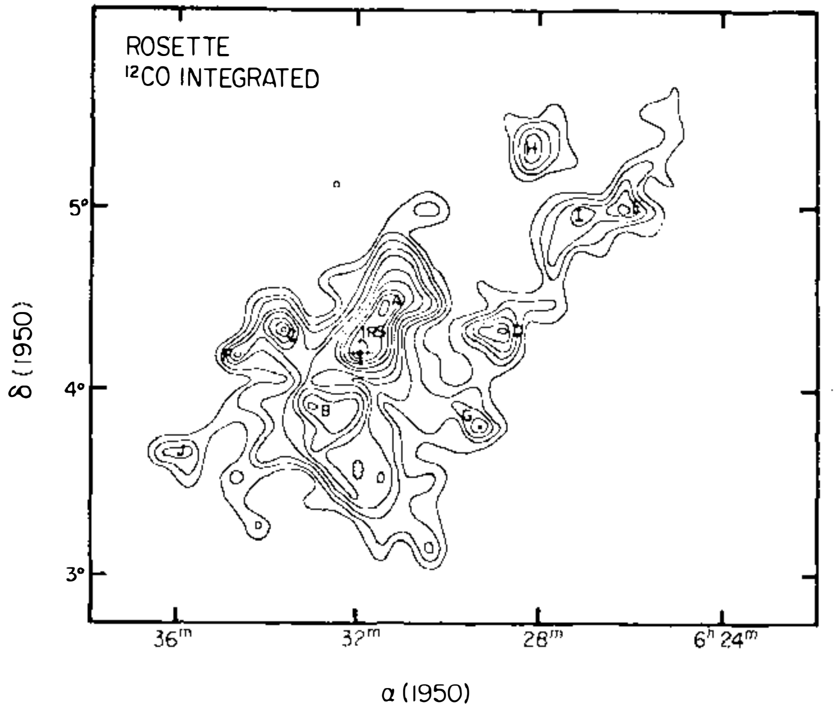

Figure 1 Clumps observed in ^{12}CO within the molecular cloud complex associated with the Rosette Nebula. A hot core has been illuminated by a luminous infrared source at the position of the cross (from Blitz & Thaddeus 1980).

the mapping (Myers 1985, Walmsley 1987). The masses enclosed within the density contours to which ammonia is sensitive (roughly $n_{H_2} > 10^4$ cm^{-3}) are $M_{core} \sim 10^0\ M_\odot$. The cores defined in this manner have sizes $R \sim 10^{-1}$ pc, temperatures $T \sim 10$ K, and NH_3 linewidths that are almost thermal. Perault et al. (1985) find from ^{13}CO observations that the dense cores are embedded in massive envelopes of several hundred $M_\odot$ (see also Falgarone 1987). The contention (Myers & Benson 1983) that the dense, quiet NH_3 cores are the sites of low-mass star formation is supported by their close association with known T Tauri stars (see Figure 2), by the correlation of bipolar outflow sources with them (Fuller & Myers 1987), and by the fact that IRAS detected infrared sources in approximately half of the NH_3 cores (Beichman et al. 1986). The deeply embedded sources can be identified as protostars by the characteristics of their infrared spectral energy distributions (Adams & Shu 1986, Adams et al. 1987, Myers et al. 1987a).

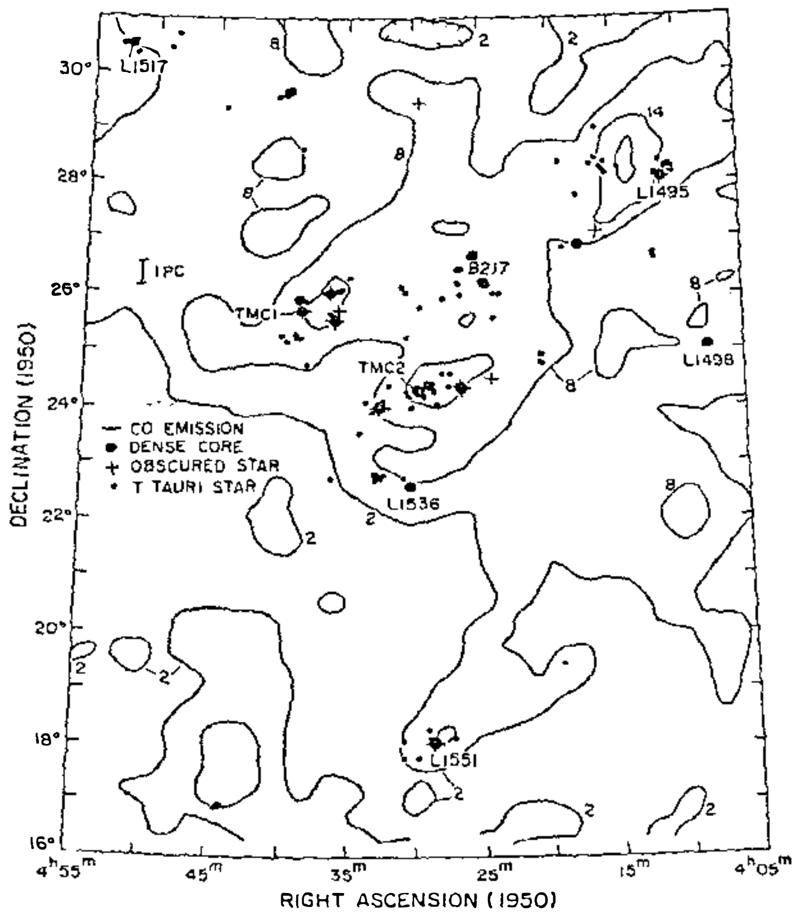

Figure 2 CO contour map of the Taurus molecular cloud with positions of dense NH_3 cores, embedded infrared sources, and visible T Tauri stars (from Myers 1986).

There are also denser clumps, like the R Coronae Australis and ρ Ophiuchi regions, which have a rather high efficiency of star formation (between 25–50%; see Wilking & Lada 1983, Lada & Wilking 1984, Wilking et al. 1986) and seem to be forming bound clusters. These clumps have radii $\sim$0.3–0.6 pc, masses $\sim$19–110 $M_{\odot}$, average densities $\sim 10^3$ cm^{-3} (Loren et al. 1983), and may be closely packed collections of cores with internal densities rising above 10^6 cm^{-3} (Snell et al. 1984).

2.2 *Lifetimes of Giant and Dwarf Molecular Clouds*

The issue of the lifetimes of molecular clouds attracted considerable controversy at the turn of the decade (Solomon et al. 1979, Blitz & Shu 1980, Cohen et al. 1980), but further analysis of the observational data has led to a greater consensus of opinion. A substantial fraction of the total molecular material is now thought to be contained in dwarf molecular clouds (DMCs), which are not restricted to the spiral arms. These clouds probably survive one or more rotation periods of the Galaxy and thus have lifetimes of 10^8 yr or more. In contrast, molecular clouds with hot cores (presumably GMCs containing OB stars) are probably assembled and dispersed as they cross the spiral arms on a time scale of $\sim 10^7$ yr (e.g. Bash & Peters 1976). In principle, the determination of cloud ages (at least, those parts that are relatively young) could be obtained more directly through the use of molecular clocks (see the review by Glassgold 1985). Unfortunately, large uncertainties dominate present discussions, and definitive results await future developments.

In any case, the existing evidence is consistent with the interpretation that giant cloud complexes are not much older than the crossing times of their individual clumps. Their heterogeneous internal structures pose no severe dynamical problems if they are transient objects (Blitz & Shu 1980); the main theoretical issue for GMCs is then how the individual clumps of gas are gathered. In high-angular-resolution observations of M83 and M51 (Allen et al. 1986, Lo et al. 1987), the ridge of H I emission is observed to be displaced *downstream* from the dust lane and peak molecular and nonthermal emission (presumably the region of maximum density-wave compression; see, e.g., Roberts & Yuan 1970). This finding indicates that GMCs in the regions being observed are assembled from molecular rather than atomic material (e.g. Kwan & Valdes 1983, 1986, Roberts & Stewart 1987), and that the H I gas comes from the partial dissociation of the H_2 gas after the formation of OB stars. The possible role of various instabilities in helping to gather the complex is left unsettled (Mouschovias et al. 1974, Cowie 1981, Elmegreen 1982a, 1987, Balbus & Cowie 1985, Lubow et al. 1986), but there is little doubt that theoretical models must grapple with the basic clumpy nature of the interstellar gas (see, e.g., Bash & Kaufman 1986).

2.3 *Mechanical Balance of Molecular Clumps*

In contrast to giant molecular clouds, DMCs and individual molecular clumps must be near mechanical equilibrium. The Jeans mass M_J associated with the average conditions of a clump is only a few $M_\odot$, much smaller than the mass of the clump M_{cl}, yet molecular clouds cannot be collapsing on a free-fall time scale or else the rate of star formation in the Galaxy would be far too high (Zuckerman & Palmer 1974). Since $M_J \ll M_{cl}$, thermal pressure could be important for molecular cloud cores, but not for their envelopes. Various alternative mechanisms of support for self-gravitating interstellar clouds have been invoked at one time or another: magnetic fields (Chandrasekhar & Fermi 1953, Mestel 1965, Spitzer 1968, Mouschovias 1976), rotation (Field 1978), and turbulence (Norman & Silk 1980, Larson 1981). Of these possibilities, magnetic fields are probably the crucial ingredient, because without them the observed levels of rotation and turbulence in molecular clouds would be very difficult to understand.

2.3.1 FLUID TURBULENCE Since the discovery of CO lines in molecular clouds, it has been known that the CO linewidths correspond to very supersonic fluid motions. When interpreted as turbulence, the velocities approach virial values $\Delta v \sim (2GM/R)^{1/2}$, where R is the radius of the cloud and Δv is the FWHM of the molecular line. The CO linewidths have also been empirically determined to follow a power-law correlation with the size of the region being observed ($\Delta v \propto R^{\alpha}$). From a compilation of data from different sources, Larson (1981) derives $\alpha \approx 0.3$; Solomon & Sanders (1985) and Dame et al. (1986) prefer a value of $\alpha \sim 0.5$–0.6. Myers (1987) attributes the difference to Larson's not having subtracted out the thermal contribution to the linewidth before drawing his correlation diagrams.

Theoretically, the difficulty with supersonic turbulence is that in the absence of magnetic fields, it should be highly dissipative (see, e.g., the discussion of Goldreich & Kwan 1974). Scalo & Pumphrey (1982) have performed numerical simulations, in which the turbulence was modeled with moving blobs; such a model might apply to the dissipation of the relative motions of clumps within a cloud complex. In any case, Scalo & Pumphrey claim that the turbulent dissipation rate is usually overestimated because the relevant interaction cross section is (in effect) πR^2 instead of the gas kinetic formula $4\pi R^2$. In addition, the collisions are generally oblique, and coalescence reduces the number of fragments with time. However, most estimates in the literature do not include the extraneous factor of 4 [see, e.g., the discussion of Spitzer (1968) or Blitz & Shu (1980), which discuss what happens to a given *small* parcel of gas], and the other effects discussed by Scalo & Pumphrey are highly model dependent (e.g. gravitational focusing will enhance the collision rate).

For individual clumps the situation is worse, since the turbulent dissipation of energy per unit mass in a nonmagnetic continuous medium (compressible or not) will always occur at a rate $\sim v^3/l$, where v is the typical turbulent velocity and l is the characteristic interaction scale. (Basically, the specific energy $\sim v^2$ in macroscopic motions is ultimately converted to heat in a time $\sim l/v$ for both cascading eddies and radiative shocks. Any acoustic radiation transported away in the process only hastens the drain of energy.) In order for turbulence to provide mechanical support for a cloud, the turbulent velocity v must be of the order of the virial speed, and l should be much smaller than the cloud size R. Consequently, the turbulent dissipation time l/v must be much smaller than the crossing time R/v. Thus, turbulent support of a cloud is not possible unless there is a mechanism to replenish the fluctuations on a sufficiently short time scale. Fleck (1981) has proposed the tapping of shear energy in galactic differential rotation as a possible mechanism, but this has long been considered an unpromising source of interstellar turbulence (Spitzer 1968). A more likely possibility—winds (in a snowplow phase) from young stars (Norman & Silk 1980)—may take too long to pump up the general velocity field (Bally & Lada 1983; P. Goldreich, private communication). Moreover, direct conversion of wind energy into gas motions would most likely deposit the largest velocities at the smallest scales, contrary to the observed correlations (Fleck 1981). Finally, turbulent "pressure" leads to forces only if it has a gradient, yet empirically the combination ρv^2 appears to be nearly a constant.

Observationally, the nature of turbulent velocity fields in dark clouds is further constrained by interstellar polarization maps (Vrba et al. 1976), which show that the directions of the embedded magnetic fields are well ordered over the dimensions of the clouds (see also Monetti et al. 1984). This demonstrates that turbulence cannot be totally dominant over the magnetic fields. If the energy of a turbulent velocity field, e.g. in the form of a cascade of eddies (see Landau & Lifshitz 1959), were much greater than the energy of the magnetic field, a more tangled magnetic field configuration would result. Since this is not the case, the observed "turbulent" motions are probably more wavelike than eddylike; thus, although fluctuations of the field are present (say, in the form of Alfvén waves), the field still has a well-defined mean direction.

2.3.2 MAGNETIC SUPPORT AND ALFVÉNIC TURBULENCE Since, unlike turbulence, magnetic fields have the virtue that they are not easily dissipated, they deserve consideration as the major agent for molecular cloud support. Their longevity makes them a natural candidate as a resilient obstacle to rapid star formation. 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 1976).

A magnetic flux Φ can support a cloud provided its mass does not exceed a critical value given by the following relation [see the extrapolation in Figure 1 of Mouschovias & Spitzer (1976) to zero $(p_{\mathrm{T}}/p_{\mathrm{m}})^{1/3}$]:

$$M_{\mathrm{cr}} = 0.13\frac{\Phi}{G^{1/2}} \approx 10^3\ M_{\odot}\left(\frac{B}{30\ \mu\mathrm{G}}\right)\left(\frac{R}{2\ \mathrm{pc}}\right)^2. \qquad 2.$$

Zeeman splitting of thermal OH (the excitation of which requires densities $> 10^3\ \mathrm{cm}^{-3}$) allows the measurement of the component of **B** along the line of sight for Orion A (125 μG), Orion B (38 μG), W22A (18 μG), W22B (32 μG), W40 (14 μG), W49B (21 μG), S88B (69 μG), and S106 (130 μG) (Troland et al. 1986, Crutcher & Kazes 1983, Kazes & Crutcher 1986, Heiles & Stevens 1986, Troland & Heiles 1986, Crutcher et al. 1987, Kazes et al. 1987; see also the review of Heiles 1987). It should be noted that there are many other cases with only upper limits; on the other hand, for random orientations, the total field strengths should be larger, on average, by a factor of 2 than the values deduced from a Zeeman analysis for differential circular polarization in the line wings. Except for W22 and W49, all the measured sources have nearby H II regions; the last two sources have bipolar flows. In OH maser sources associated with collapsed regions near young stars (with gas densities of the order of 10^5–$10^8\ \mathrm{cm}^{-3}$), the derived *total* field strengths (from direct observation of the separate Zeeman components) are typically several milligauss (see the review of Reid & Moran 1981).

If magnetic fields do provide the dominant mechanism for the support of self-gravitating clouds, then the properties of the observed turbulence become explicable (Myers & Goodman 1987, Shu 1987). The energy for the turbulence may originate with many sources—stellar winds, cloud collisions, expanding H II regions, supernovae, etc.—but ultimately the disturbances will excite a spectrum of MHD waves (Arons & Max 1975, Zweibel & Josafatsson 1983, Elmegreen 1985b, Falgarone & Puget 1986). Waves with super-Alfvénic fluid motions will generate compressive shocks and will dissipate rapidly. Hence, the fluctuating part of the fluid velocities will generally become sub-Alfvénic (but still supersonic), i.e. $\delta v \lesssim v_{\mathrm{A}}$. In the ^{12}CO line, there is an observational selection bias toward seeing the highest velocities, $\delta v \sim v_{\mathrm{A}}$, because photon trapping tends to shield regions of common (low) velocities (P. Goldreich, private communication; Kutner & Leung 1985). However, from Equation 2, it is easy to show that for clouds near the critical state $M_{\mathrm{cl}} \approx M_{\mathrm{cr}}$, the mean Alfvén speed is automatically of the magnitude needed for virial equilibrium, i.e.

$$v_{\mathrm{A}}^2 \equiv \frac{B^2}{4\pi\rho} \sim v_{\mathrm{V.T.}}^2 \equiv \frac{GM_{\mathrm{cl}}}{R}. \qquad 3.$$

Thus, $\delta v \sim v_A$ implies that $\delta v \sim v_{V.T.}$ (i.e. cloud "turbulence" automatically has a tendency to appear sufficient for virial equilibrium). Moreover, if B does not vary strongly from region to region (in CO envelopes), Equation 2 implies that clouds near the critical state, $M_{cl} \approx M_{cr}$ (i.e. clouds for which magnetic fields provide the dominant means of support), should have nearly the same mean column densities, $M_{cl}/\pi R^2 \propto \rho R \approx$ constant, implying $v_A \propto \rho^{-1/2} \propto R^{1/2}$. Thus, this line of argument provides a simple mechanistic explanation of the observed correlation $\Delta v \propto R^\alpha$, with $\alpha \approx 0.5$. The assumption that B does not vary from region to region may be relaxed if the observed relation, $\rho R \sim$ constant, arises (partially) from selection effects, because the observations may be sensitive to only a relatively narrow range of column densities (Larson 1981).

The presence of nonlinear Alfvén waves circumvents a primary difficulty associated with magnetic support—magnetic fields cannot provide support along the direction of the field, but the cloud clumps are not observed to be highly flattened along field lines. If thermal pressure and static magnetic fields were the only means of support for a self-gravitating clump, the cloud would have a characteristic scale height in the direction parallel to the field equal to (Spitzer 1942)

$$z_0 = \frac{kT/m}{\pi G\sigma}, \tag{4}$$

where σ is the surface mass density projected along $\mathbf{B}$. For $T \sim 10$ K, a mean molecular weight $m = 2.3\ m_H$, and $\sigma = 0.01$ g cm^{-2} (corresponding to 2.5 mag of visual extinction), Equation 4 yields $z_0 \sim 0.06$ pc. Since the lateral dimensions of molecular cloud clumps are typically a parsec or more, some theorists (e.g. Pudritz 1986) have suggested that the cloud clumps should essentially be sheets. Indeed, the L204 cloud (McCutcheon et al. 1985) and the denser condensations in the Taurus complex (Monetti et al. 1984) appear quite flattened and have often been interpreted in terms of either sheets or filaments. However, the flattening is probably not as severe as indicated by Equation 4 because bipolar flows in the Taurus region tend to be aligned with the ambient field (Strom et al. 1985), and some of these flows have apparently swept up ambient molecular gas out to a parsec and are still within the cloud (see also Figure 18 of Bally & Lada 1983). Moreover, in Ophiuchus and R Coronae Australis, dust filaments are found that are aligned *parallel* to the mean field (Vrba et al. 1976, Heyer et al. 1986).

Alfvén waves may provide the requisite support along $\mathbf{B}$ (Dewar 1970). Consider a field $\mathbf{B}$ made up of a static (mean) part $\mathbf{B}_0$ pointing in the z direction, and a fluctuating (Alfvén-wave) part $\delta\mathbf{B}$ with magnitude

$\approx(\delta v/v_A)B_0$ and with direction perpendicular to $\mathbf{B}_0$. For simplicity, assume $\delta\mathbf{B}$ has the form of a spatially damped traveling wave,

$$\delta\mathbf{B} = \mathbf{b}\sin(k_R z - \omega t)\exp(-k_I z), \tag{5}$$

with $\mathbf{b}$ lying in the x-y plane. This wave provides a secular transfer of momentum to the ambient medium, which is given by the time average of the quadratic term in the Lorentz force per unit volume:

$$\frac{1}{4\pi}(\nabla \times \delta\mathbf{B}) \times \delta\mathbf{B}. \tag{6}$$

The substitution of Equation 5 into Equation 6, averaged over one cycle, gives

$$\frac{1}{8\pi}|\mathbf{b}\exp(-k_I z)|^2 k_I \hat{\mathbf{z}}, \tag{7}$$

which equals the average value of $-\nabla(|\delta\mathbf{B}|^2/8\pi)$. The force (Equation 7) is directed along the direction of wave propagation (parallel to $\mathbf{B}_0$) and has magnitude comparable to the Lorentz force associated with the mean field,

$$\frac{1}{4\pi}(\nabla \times \mathbf{B}_0) \times \mathbf{B}_0 \tag{8}$$

(which supports the cloud in the lateral directions), provided that $\delta\mathbf{B} \sim \mathbf{B}_0$ (i.e. if $\delta v \sim v_A$) and that the characteristic damping length k_I^{-1} is comparable to the size R of the cloud clump. We have already commented upon the evidence from the observed CO linewidths that suggests that δv is comparable to v_A; standard formulae (see the references cited by Zweibel & Josafatsson 1983) for the damping of Alfvén waves by ion-neutral collisions or by nonlinear steepening suggest that the condition $k_I^{-1} \sim R$ is also likely to occur in molecular clouds. (Here, the real part of the wave number k_R may be associated with wave generation by molecular outflows.) Presumably, an observed cloud clump with a given distribution of sources has contracted to a size Z in the direction parallel to the mean field, so that the available set of waves can provide support against gravity for that configuration.

Notice that the "Reynolds stress" in the transverse fluid motions associated with an Alfvén wave provides no momentum transfer along the direction of the field, $\mathbf{B}_0 \cdot \nabla \cdot (\rho\delta\mathbf{v}\delta\mathbf{v}) = 0$, and hence does not support the cloud in the direction of the mean field. In addition, only outwardly propagating waves support the cloud against self-gravity; inwardly propagating waves would tend to *compress* the cloud. Thus, cloud support by this mechanism requires that there be a negative gradient in wave sources from inside to outside the cloud.[1]

2.3.3 MAGNETIC BRAKING The generation of (torsional) Alfvén waves may also explain the small values of the angular velocity Ω commonly deduced for molecular clouds (see below). There are exceptional cases with relatively large rotation rates (see, e.g., Vogel & Welch 1983), but these clouds probably represent collapsed objects in later stages of evolution (Vogel 1984). Usually, when the envelopes of molecular clumps have measurable rotation rates, the rates are of order $\Omega \sim 1$ km s^{-1} pc^{-1} = 3×10^{-14} rad s^{-1}. An angular velocity of the same order (1 km s^{-1} pc^{-1}) is often observed for the ammonia cores in the Taurus complex (Myers 1987), which suggests that the cores of clumps may be rotationally coupled to their envelopes by magnetic braking (Gillis et al. 1974, 1979, Mouschovias & Paleologou 1979, 1980). Since the rotation rates of NH_3 cores in Taurus are always less than 5 km s^{-1} pc^{-1}, rotation cannot be important for the support of these objects (Myers et al. 1987c). In addition, there appears to be no clear correlation between the projected clouds' rotation axes and the axis of Galactic rotation (Goldsmith & Sernyak 1984).

The process of magnetic braking of cloud cores by their envelopes is important because it produces a potential reservoir of low-angular-momentum material for the formation of stars, planets, and binary systems (Mouschovias 1978). In general, Alfvén waves generated by the spin of the core relative to its envelope will tend to bring the two into corotation; significant braking occurs when the amount of envelope material affected by the outwardly propagating Alfvén wave has a moment of inertia equal to that of the core. In an approximation that models the core as a cylinder of half-height Z and radius R with uniform density $\rho_{\rm core}$ embedded within an envelope of uniform density $\rho_{\rm env}$, Mouschovias & Paleologou (1980) show that the characteristic time scale for braking the component of Ω parallel to $\mathbf{B}_0$ is given by

$$\tau_{\parallel} = \frac{\rho_{\rm core}}{\rho_{\rm env}} \frac{Z}{v_{\rm A}}. \qquad 9.$$

On the other hand, the corresponding time scale for braking the component of Ω perpendicular to $\mathbf{B}_0$ is given by

$$\tau_{\perp} = \frac{1}{2}\left[\left(1+\frac{\rho_{\rm core}}{\rho_{\rm env}}\right)^{1/2} - 1\right]\frac{R}{v_{\rm A}}. \qquad 10.$$

[1] The same comment applies to proposals that invoke stellar outflows as a direct mechanism for cloud support (P. Goldreich, private communication).

In Equations 9 and 10, the Alfvén speed v_A in the envelope is to be evaluated just outside the surface of the core under the assumption that farther out it is independent of z but declines as $\tilde{\omega}^{-1}$.

For the case $\rho_{core} \gg \rho_{env}$, the time scale $\tau_{\perp}$ for braking the perpendicular component will be much less than that for the parallel component $\tau_{\parallel}$; this situation occurs because the greater lever arm in the perpendicular directions offsets the assumed decrease of propagation speed with increasing $\tilde{\omega}$. The lower efficiency of parallel magnetic braking will tend to leave the spin axes of molecular cloud cores (and the objects to which they collapse) aligned parallel to $\mathbf{B}_0$; this effect may explain why the bipolar flow axes of several sources in Taurus line up with the mean direction of the ambient magnetic field (Strom 1985, Strom et al. 1985, Heyer 1986).

2.3.4 MEASURED ROTATION RATES Goldsmith & Arquilla (1985; see also Fuller & Myers 1987) have compiled a list of rotation rates in 16 dark cloud regions, with masses in the range 6×10^3 to 0.7 $M_{\odot}$, sizes from 17 to 0.1 pc, and velocity gradients from 0.2 to 6 km s^{-1} pc^{-1}; they find a tendency for the smallest and densest regions to have the shortest rotation periods (see their Figure 6).[2] Unfortunately, the various correlations derived on the basis of these data may not be extremely meaningful, since they may represent limits below which rotation in dark clouds cannot be detected. What is significant about the rotation rates for ordinary dark clouds is that they are uniformly low or undetectable (see also Table 3 of Arquilla & Goldsmith 1986), which suggests that the clouds are magnetically controlled, i.e. they are subcritical regions (see Section 3.3) and are not contracting rapidly as a whole.

Observations of rotation in molecular cloud cores in regions other than Taurus are fairly sparse. Two cases are known in ρ Ophiuchi (Wadiak et al. 1985), where H_2CO measurements show two fragments with diameters $\sim$0.01–0.02 pc and masses $\sim$1–2 $M_{\odot}$ enclosed within density contours $\sim 5 \times 10^6$ cm^{-3} and having rotation rates $\sim$40 km s^{-1} pc^{-1}. In Orion, five cases have been found and measured (Harris et al. 1983; L. Mundy, private communication); NH_3 observations show cores with diameters $\sim$0.05 pc and masses $\sim$10 $M_{\odot}$ enclosed within density contours $\sim 3 \times 10^4$ cm^{-3} and having rotation rates 25–50 km s^{-1} pc^{-1}. The faster rotation rates and

[2] It is tempting to interpret the apparent spin-up as being due to contraction and to conclude that it has proceeded with considerable braking because the specific angular momentum J/M of the various regions decreases with decreasing size (see their Figure 5). This conclusion may well be correct; however, it is not warranted on the basis of the evidence presented. The specific angular momentum in any stably stratified medium will generally increase outward, so that averaging over progressively larger volumes will naturally produce larger values of J/M, independent of whether there is any evolutionary connection between the material on the inside and that on the outside.

larger masses within a given radius are all consistent with a picture in which these cores have formed from a supercritical background that is itself contracting (see Section 3.2), thereby lowering the available time for magnetic braking.

Vogel & Welch (1983), Zheng et al. (1985), and Ho & Haschick (1986) have studied the kinematics of the molecular gas in the neighborhood of the H II regions W58, ON1, and G10.6−0.4; these regions presumably represent the core environments in which high-mass stars form. The gas masses within $\sim$0.5–1 pc contours and mean densities $\sim 10^5$ cm^{-3} are large (of order 10^2–10^3 $M_\odot$). The measured angular velocities are all of order 10 km s^{-1} pc^{-1} and may again represent the spin-up associated with the overall contraction of a supercritical region.

The central regions of G10.6−0.4 studied by Ho & Haschick (1986; see also Keto et al. 1987) seem to be collapsing toward a compact group of OB stars, with the inflow velocities increasing inward. The infalling gas appears to spin faster than the radially stationary material farther out; these are characteristics naturally produced from an "inside-out" collapse such as that discussed in Section 4.1.2.

2.3.5 IONIZATION FRACTION The ionization fraction in typical regions of molecular clouds is believed to be very low, although modern attempts to measure it precisely (see the review of Langer 1985) have been foiled by the revision of the dissociative recombination coefficient for H_3^+ (Smith & Adams 1984). The lower limit on the ionization number fraction, as determined from the abundance of $HC^{18}O^+$, is still 10^{-8}. An upper limit of $\sim 10^{-6}$ can be set from deuterium fractionization reactions (Dalgarno & Lepp 1984), and a value of $\sim 10^{-7}$ can be derived from the abundance of HCO^+ if the cosmic-ray ionization rate is assumed to be $\zeta \sim 10^{-17}$ s^{-1} (Smith & Adams 1984). It may be possible to determine the cosmic-ray flux in dense cores by measuring (in self-absorption) the amount of atomic hydrogen formed by the dissociation of H_2 (M. Jura, private communication).

Until better empirical results are available, the most accessible determinations at present are theoretical. When cosmic-ray ionization at a rate $\zeta = 1 \times 10^{-17}$ s^{-1} is balanced by two-body recombinations of charged particles and on charged grains, the ion mass density ρ_i depends on the neutral mass density ρ_n according to the law (Elmegreen 1979)

$$\rho_i = C\rho_n^{1/2}, \qquad 11.$$

where C is a weak function of gas temperature and is proportional to the square root of the metal depletion. For average metal depletions of 0.1 and gas temperatures of 10–30 K, Elmegreen (1979) obtains results that correspond to $C = 3 \times 10^{-16}$ cm$^{-3/2}$ g$^{1/2}$. Thus, with a typical ion mass $m_i \approx 30\ m_H$ and a typical neutral mass of $m_n \approx 2.3\ m_H$ (de Jong et al. 1980), the ionization number fraction is about 10^{-7} at a density of 10^4 cm^{-3}. At densities much higher than 10^8 cm^{-3}, natural radioactivity becomes dominant over cosmic rays as the primary ionizing agent (Nakano & Tademaru 1972), grains begin to carry a major fraction of the total charge, and ρ_i approaches a constant value (Umebayashi & Nakano 1980).

2.3.6 AMBIPOLAR DIFFUSION AND MAGNETIC FLUX LOSS When the ionization fraction is very low, magnetic support of molecular clouds becomes synonymous with ambipolar diffusion—the process by which magnetic fields (and the charged plasma to which they are coupled) drift relative to a background of neutrals (Mestel & Spitzer 1956). The magnetic forces are felt directly only by the charged particles; the neutrals can be supported against their self-gravity only through the frictional drag that they experience as they slip relative to the ions. To see this quantitatively, note that with the low ionized mass fractions implied by Equation 11, the pressure and gravitational forces acting on the fluid of charged species are always very small in comparison with the Lorentz force,

$$\frac{\mathbf{j}}{c} \times \mathbf{B} = \frac{1}{4\pi}(\nabla \times \mathbf{B}) \times \mathbf{B}, \qquad 12.$$

where we have used Ampere's law, $\mathbf{j} = (c/4\pi)\nabla \times \mathbf{B}$. The mean Lorentz force will drive the ions through a sea of neutrals at a (mean) relative terminal velocity given by a balance with the drag force (per unit volume):

$$\mathbf{f}_d = \rho_n \rho_i \gamma (\mathbf{u}_i - \mathbf{u}_n), \qquad 13.$$

where $\mathbf{u}_i$ and $\mathbf{u}_n$ are the fluid velocities of the ions and neutrals, respectively, and γ is the drag coefficient associated with momentum exchange in ion-neutral collisions. When the collision process is dominated by the dipole moment induced in the neutral by the passing ion, we have $\gamma \approx 3.5 \times 10^{13}$ cm^3 g^{-1} s^{-1} (Draine et al. 1983). For drift speeds in excess of ~ 10 km s^{-1}, the Langevin approximation breaks down, and γ should approach the geometric value $\sim \pi(r_i + r_n)^2 |\mathbf{v}_d|/(m_i + m_n)$ (Mouschovias & Paleologou 1981). A balance between Equations 12 and 13 leads to a drift velocity of the ions (which in turn nearly comove with the electrons) relative to the neutrals given by (Spitzer 1958)

$$\mathbf{v}_d \equiv \mathbf{u}_i - \mathbf{u}_n = \frac{1}{4\pi\gamma\rho_i\rho_n}(\nabla \times \mathbf{B}) \times \mathbf{B}. \qquad 14.$$

Because the cyclotron frequency $eB/m_i c$ of the typical ion is much greater than the mean collision frequency of ions with neutrals, $\rho_n\gamma$, ions gyrate about a given field line many times before they are knocked off by collisions; thus, the ions are effectively tied to **B**. The time evolution of the magnetic field itself may then be obtained from the approximation that it is frozen in the plasma of ions and electrons:[3]

$$\frac{\partial \mathbf{B}}{\partial t} + \nabla \times (\mathbf{B} \times \mathbf{u}_i) = 0. \qquad 15.$$

The substitution of Equation 14 into Equation 15 yields the nonlinear diffusion equation

$$\frac{\partial \mathbf{B}}{\partial t} + \nabla \times (\mathbf{B} \times \mathbf{u}_n) = \nabla \times \left\{ \frac{\mathbf{B}}{4\pi\gamma\rho_i\rho_n} \times \left[\mathbf{B} \times (\nabla \times \mathbf{B}) \right] \right\}. \qquad 16.$$

If the right-hand side were zero, the field would be well coupled to the motion of the neutrals. As it is, Equation 16 states that the effective "diffusion coefficient" $\mathscr{D}$ associated with the drift of the field relative to the neutrals has order of magnitude $v_A^2 t_{ni}$, where $v_A \equiv (B^2/4\pi\rho_n)^{1/2}$ is the Alfvén speed in the (combined) medium, and $t_{ni} \equiv (\rho_i\gamma)^{-1}$ is the mean collision time of a neutral molecule in a sea of ions. If the magnetic field has characteristic length R, the time scale for ambipolar diffusion is $t_{AD} \sim R^2/\mathscr{D} \sim R/v_d$. For a magnetically supported self-gravitating cloud that satisfies Equations 3 and 11, the ratio of t_{AD} to the dynamic time scale $t_{dyn} \sim R/v_A \sim (G\rho_n)^{-1/2}$ is approximately given by the dimensionless combination

$$\frac{t_{AD}}{t_{dyn}} \sim \frac{\gamma C}{2(2\pi G)^{1/2}}, \qquad 17.$$

where the numerical coefficient has been adjusted to agree with a detailed analysis in a slab geometry (Shu 1983, Lizano & Shu 1987). Notice that the ratio (Equation 17) is independent of M_{cl} or R (and therefore the mean ρ_{cl}). For values of γ and C quoted above, the coupling constant on the right-hand side of Equation 17 has the value 8; thus, ambipolar diffusion can be generally expected to take place at a relatively slow rate compared with dynamical events until the ionization fraction begins to depart appreciably from the law (Equation 11).

[3] Because grains are negatively charged in opaque clouds, Elmegreen (1979, 1986) and Nakano & Umebayashi (1980) considered the extent to which they are also tied to field lines and can impede the slip of the neutrals. Large grains have small gyrofrequencies and are not well coupled to **B**, while small grains present cross-sectional areas for collisions with neutrals that are not much larger than the Langevin cross sections of ions with neutrals. Consequently, the effects of charged grains on the rate of ambipolar diffusion are important only at high densities, when the charges are mostly carried by grains.

The equation of motion for the neutrals reads as

$$\rho_n\left[\frac{\partial \mathbf{u}_n}{\partial t}+(\mathbf{u}_n\cdot\nabla)\mathbf{u}_n\right] = -\nabla P_n - \rho_n\nabla\phi+\mathbf{f}_d, \qquad 18.$$

where P_n is the neutral gas pressure, and ϕ is the gravitational potential satisfying

$$\nabla^2\phi = 4\pi G\rho_n \qquad 19.$$

in the approximation $\rho_i \ll \rho_n$. In a heuristic treatment that does not follow fluctuating motions (Alfvén waves), e.g. the adoption of Equation 14, one might include in P_n the "turbulent" contributions discussed in Section 2.3.2. In any case, if we now substitute Equations 13 and 14 into Equation 18, we see that the magnetic forces provide indirect support to the neutrals against their self-gravity via the intermediary of frictional drag. Thus, in a lightly ionized gas like a molecular cloud, the processes of magnetic support and ambipolar diffusion are synonymous; one cannot have one without having the other. As is shown in Section 3.4, this automatically leads to a mechanism for the formation of substructures (cores and envelopes) in molecular clouds.

Ambipolar diffusion is believed to be the process by which the classic flux problem is resolved for forming stars (see the review of Mestel 1985). Moreover, the issue of when the magnetic field decouples from the matter may determine the range of internal densities present in a molecular cloud core at the point when it becomes unstable to dynamical collapse (see Section 4.1.1). The value of the cloud density required in order for dynamical decoupling of the field to occur is controversial. Nakano (1984, 1985) asserts that it happens when the density becomes greater than 10^{11} cm^{-3}; Mouschovias et al. (1985, and references therein) claim that decoupling will occur for 10^5 cm^{-3}. The difference in opinion arises because of an unorthodox definition that leads the latter authors (see their Equation C1a) to label relatively slow motions as being "dynamical" even when the evolution proceeds quasi-hydrostatically in the core by ambipolar diffusion. (The large infall velocities in the outer layers are artifacts of the slab geometry, in which the gravitational field approaches a constant at infinity instead of falling off as $1/r^2$.) In any case, Equation 17 favors a high density (that for which Equation 11 begins to break down) for decoupling; however, the argument depends sensitively on the precise values adopted for γ, C, and the numerical coefficient in the denominator. Observations of systematic velocity differences between neutral and ionized species with similar critical densities for excitation can help to clarify this issue (M. Walmsley, private communication).

2.4 *Thermal Balance*

Temperatures of about 10 K are generally inferred from CO observations of molecular cloud envelopes and can be understood in terms of a balance between cosmic-ray heating and CO cooling (e.g. de Jong et al. 1980). In the outer parts, heating by photoelectrons from dust grains is an important additional contribution (for a review of several possible mechanisms for heating molecular clouds, see Goldsmith & Langer 1978). The cloud cores associated with regions of low-mass star formation also have low temperatures (Myers & Benson 1983, Menten & Walmsley 1985), but those near OB stars are considerably hotter (for a review, see Wynn-Williams 1982). The conventional explanation for this latter observation has been that luminous embedded stars heat the surrounding dust grains, which in turn heat the gas (Goldreich & Kwan 1974). This interpretation has been questioned in a few well-studied cases where the gas appears to be warmer than the dust (e.g. Wilking & Lada 1983) or where the densities are insufficient to provide good thermal coupling between gas and dust (e.g. Evans et al. 1981); also, there now appears to be some reason to believe that a general elevation of the gas temperature may have taken place in such regions before the appearance of massive young stars (see Sections 3.2 and 3.4).

3. ORIGIN OF SUBSTRUCTURE IN MOLECULAR CLUMPS

3.1 *Two Modes of Cloud Contraction*

If we accept the importance of magnetic fields for molecular cloud support, there are logically two regimes of interest in the problem of star formation (see, e.g., the discussion of Mestel 1985):

1. In the *supercritical* regime, $M_{cl} > M_{cr}$, the clump's self-gravity can overwhelm the magnetic support *even if the fields remain frozen in the fluid*. Cloud evolution in this state is characterized by magnetically diluted collapse (e.g. Scott & Black 1980). The flow of gas along field lines (cloud flattening) would eventually provide subregions, of size comparable to the vertical dimension, with the same mass-to-flux ratio (on average) as the entire cloud; thus, these regions would themselves be supercritical and may be able to fragment from the contracting background (Mestel 1965).

2. In the *subcritical* regime, $M_{cl} < M_{cr}$, indefinite gravitational collapse (star formation) cannot be induced by *any* amount of increased external load (external pressure) if Φ is conserved (field freezing) because the mass-to-flux ratio M_{cl}/Φ remains fixed and subcritical. Neither can *dynamical*

fragmentation of the cloud take place. Cloud evolution in this state is most likely driven by ambipolar diffusion (e.g. Nakano 1979).

3.2 *Supercritical Case: High-Mass Star Formation or High Star Formation Efficiency*

The two cases discussed in Section 3.1 may provide a natural basis for the phenomenon of bimodal star formation (Shu et al. 1986). In particular, supercritical clouds may arise naturally when clumps are agglomerated in large cloud complexes (Shu 1987). The condition $M_{cl} > M_{cr}$ is equivalent to the existence of a critical surface density (see, e.g., Field 1970):

$$\frac{M_{cl}}{\pi R^2} > 80 \frac{M_\odot}{\text{pc}^2}\left(\frac{B}{30\ \mu\text{G}}\right). \qquad 20.$$

In the absence of other means of support, a cloud with a supercritical mass will suffer relatively rapid contraction and compress the embedded magnetic fields well above the starting values. Cloud contraction under these circumstances may be expected to efficiently form stars (see Section 1.5). If significant core heating also takes place (see Section 3.4), the stars that form could have relatively high masses (see Section 4.5). Although the contraction process can be expected to increase the mean surface density well beyond the starting value, we note that if one were to put even a healthy fraction of 80 $M_\odot$ pc^{-2} into O and B stars, the resulting areal luminosity densities would be $\sim 10^4$–10^5 $L_\odot$ pc^{-2}. This is close to the value seen in the region of the Trapezium stars in Orion.

For a canonical gas-to-dust ratio, Equation 20 is equivalent to a critical mean visual extinction:

$$A_V > 4\ \text{mag}\left(\frac{B}{30\ \mu\text{G}}\right). \qquad 21.$$

The value 30 μG may typify the average conditions only in the envelopes of small dark clouds; after gravitational contraction, dense cores and protostellar regions may have considerably larger values. For example, Simonetti & Cordes (1986) observe significant variations in the rotation measure toward L1551 that they interpret as being due to a magnetic field of strength 300 μG. In any case, it is interesting to note the following observed progression:

1. The visual extinction through the envelope of the Taurus molecular cloud is $\sim$1–2 mag (e.g. Dickman 1978, Cernicharo et al. 1985). Its cores have $A_V \sim 10^1$ mag (Myers & Benson 1983, Cernicharo et al. 1984) and probably condensed from clumps that are subcritical (see below). The gas

temperatures in the cores are generally 10–11 K. Taurus is, of course, a region with low star formation efficiency and seems to be forming an unbound association of low-mass stars (see the review of T Tauri stars by Cohen 1984).

2. The visual extinction through the envelope of the ρ Ophiuchi molecular cloud is ~ 6 mag (e.g. Encrenaz et al. 1975, Elias 1978b, Frerking et al. 1982); the cores in its densest portion have $A_V \sim 10^2$ mag (Wilking & Lada 1983). The gas temperatures in the cores may be higher than in Taurus (Zeng et al. 1984), and ^{12}CO and ^{13}CO measurements indicate temperatures of 30–35 K for the general ρ Ophiuchi region (Loren et al. 1983, Wilking & Lada 1983). This portion of the cloud has a high star formation efficiency and may be forming a bound cluster containing mostly low-mass stars but also a few B stars (Grasdalen et al. 1973, Elias 1978b, Lada & Wilking 1984). Wilking & Lada (1983) observed CO line profiles that were asymmetrically self-reversed over the entire core region of the ρ Ophiuchi cloud (C. Lada, private communication; see also Encrenaz et al. 1975). The sense of this asymmetry corresponds to overall contraction at 1–2 km s^{-1}, consistent with the entire dense region being supercritical.

3. The average visual extinction through GMCs is controversial; estimates range from 4 mag (Blitz & Shu 1980) to 12 mag (Sanders et al. 1984, Solomon et al. 1987). High-angular-resolution studies near the water masers in W3(OH) suggest a young stellar object embedded in a dense core with $A_V \sim 10^3$ mag (Turner & Welch 1984). GMCs are very inhomogeneous, and portions of them are very likely to be supercritical; hence, it is informative that an inverse P Cygni profile has been found (after the submission of this review) in HCO^+ toward the brightest compact H II region in W49, which indicates that about 10^5 $M_\odot$ is collapsing on itself at 10–20 km s^{-1} to form the most luminous association of OB stars in the Galaxy (W. J. Welch, private communication). Since HCO^+ is a molecular ion, embedded magnetic fields must be dragged in by the infall. The gas kinetic temperature in the region is in excess of 50 K, which is typical of hot-core sites containing OB stars.

3.3 *Subcritical Case: Low-Mass Star Formation and Low Star Formation Efficiency*

When a cloud has less mass than M_{cr}, it can attain a stable quasi-equilibrium state by adjusting its size and shape in accordance with an external medium of finite pressure and/or finite magnetic field (Mouschovias 1976). However, such a cloud cannot remain forever in the same state because the neutral particles slip relative to the ions (their means of magnetic support), resulting in the quasi-static condensation of a dense pocket of gas and dust (Nakano 1979). If $M_{cl} > M_{cr}$, ambipolar diffusion will also take place, but against the backdrop of a dynamically shrinking envelope (Black & Scott 1982).

The physical principle of ambipolar diffusion is illustrated by the idealized problem of the quasi-static evolution of a plane-parallel, self-gravitating slab of isothermal gas that is lightly ionized (Shu 1983). As time increases, the field tends to decay to its background value, and the gas evolves toward 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 one-dimensional problem is unrealistic because the gravitational field of the slab saturates at a finite value proportional to the total surface density. Thus, one-dimensional slabs with finite thermal pressures can never undergo true gravitational collapse. In a more realistic three-dimensional problem, where the original cloud clump contains many Jeans masses but M_{cl} is still less than M_{cr}, we may expect the clump to fragment *quasi-statically* into many "cores," each asymptotically evolving toward a configuration in which the field lines become almost uniform and straight (the background condition) and where the total velocity dispersions approach purely thermal values as the Alfvén velocity drops because of the increase in density. This picture provides an attractive evolutionary explanation for the small quiet cores that are observed in the Taurus dark cloud.

However, *stable* asymptotic support by thermal pressure alone in a quasi-spherical geometry is impossible if the potential reservoir of matter is very large in comparison with the Jeans mass in the initial state. Thus, a three-dimensional isothermal core must ultimately undergo gravitational collapse when the core becomes sufficiently centrally condensed. This has apparently happened to roughly half of the NH_3 cores in Taurus (Beichman et al. 1986).

The relatively long time scale for core formation by ambipolar diffusion (see Equation 17) probably explains why star formation is generally an inefficient process in subcritical molecular clouds. On the time scale of the dynamical collapse of an unstable molecular cloud core, star formation is badly synchronized. Newly formed T Tauri stars have ample time to turn on their winds and disrupt the incipient condensation of neighboring cores, which makes star formation in T associations loosely aggregated. There is no guarantee, however, that the coupling coefficient (Equation 17) is large compared with unity in all circumstances. Severe metal depletions or high initial mass-to-flux ratios (which, in a three-dimensional geometry, can lead to cloud contraction even in the absence of ambipolar diffusion) may well lead to relatively high star formation efficiencies.

3.4 *Heating of Cloud Cores by Ambipolar Diffusion*

The friction between ions and neutrals due to the ambipolar diffusion of the field is a source of heating for the cloud (cf. Scalo 1977, Mouschovias 1978, Elmegreen 1982b). Quantitative estimates of this heating (Lizano & Shu 1987) indicate that significant elevation of gas temperatures can occur in cores with high column density (see Figure 3). In particular, this heating mechanism provides a natural means by which the gas might become hotter than the surrounding dust, a characteristic of the ρ Ophiuchi region (Loren et al. 1983, Wilking & Lada 1983). In contrast, ambipolar diffusion yields little extra heating above that provided by cosmic rays ($T \sim 10$–11 K) for cases that simulate the ammonia cores in Taurus.

In Ophiuchus, the higher gas temperature cannot be due to heating of gas by warm dust grains (which are in turn heated by luminous stars) because the entire region only has three embedded B stars (Elias 1978b). In addition, the presence of large quantities of warm dust would produce far-infrared emission greatly exceeding that measured by balloon-borne observations (Fazio et al. 1976; see also Wilking 1985) and by IRAS (Young et al. 1986). The situation is more ambiguous for regions of true

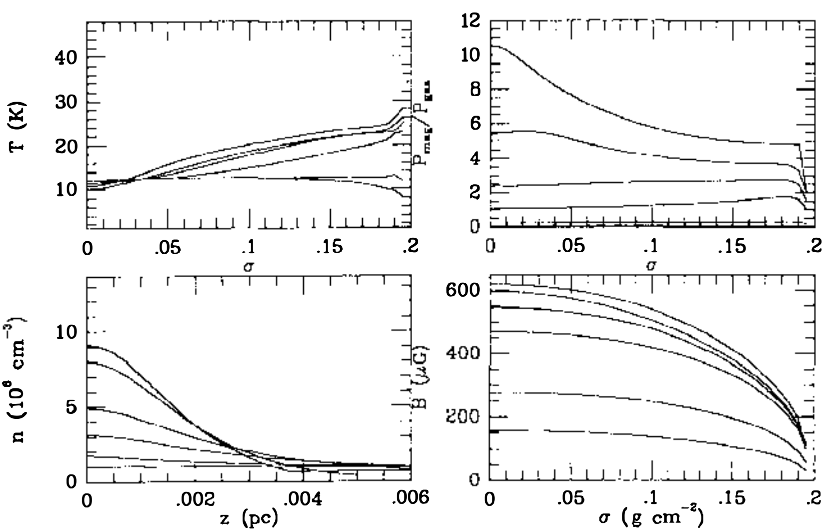

Figure 3 Time evolution of the temperature, magnetic to gas pressure, density, and magnetic field for a condensing cloud core with a column density corresponding to a (double-sided) visual extinction of 100 mag. The profiles shown occur at times 0, 0.06, 0.2, 0.3, 0.6, and 1.2 Myr (from Lizano & Shu 1987).

high-mass star formation. Additional observations made with high angular resolution are needed to help to sort out this problem.

4. GRAVITATIONAL COLLAPSE AND PROTOSTARS

4.1 *Gravitational Collapse of Molecular Cloud Cores*

At some point, as a cloud core gradually loses magnetic and "turbulent" support, the growing central concentration will cause it to become unstable to gravitational collapse. For stable (nonmagnetic and nonrotating) isothermal spheres bounded by an external medium of negligible inertia and constant pressure (i.e. a very hot gas), the maximum density contrast from center to edge is approximately 14 (Ebert 1955, Bonnor 1956). This result is, however, probably irrelevant to the present problem, where the core has access to a very large reservoir of cold matter and where the residual magnetic fields may greatly affect the stability criterion without appreciably modifying the equilibrium configuration (e.g. consider the role of quasi-uniform fields). A maximum contrast of 14 also contradicts the observational evidence that molecular clouds typically have H_2 densities ranging from 10^2 cm^{-3} to well in excess of 10^6 cm^{-3}.

4.1.1 INITIAL CONDITIONS AT THE ONSET OF COLLAPSE Consider a cloud of gas that is nearly isothermal and rotating very slowly (due to efficient braking), and suppose that the residual magnetic fields become increasingly unimportant for mechanical support of a core. Under these conditions, the continued *subsonic* evolution of the gas (with any mass) will tend to produce a $1/r^2$ density distribution (Bodenheimer & Sweigart 1968), provided that the thermal pressure gradients remain nearly in balance with the gravitational field.[4] In particular, a slowly forming molecular cloud core that is not artificially separated from its surroundings by a boundary will tend to acquire the density distribution of a singular isothermal sphere (see, e.g., Chandrasekhar 1939):

$$\rho = \frac{a^2}{2\pi G r^2}, \qquad 22.$$

where $a = (kT/m)^{1/2}$ is the isothermal sound speed, and T and m are the (constant) temperature and mean molecular weight of the gas, respectively.

[4] Numerical simulations of protostellar collapse ubiquitously establish $1/r^2$ density profiles, even when the initial state is taken to be homogeneous (e.g. Larson 1969a). The analytic explanation of this phenomenon given by Larson (1969b) and by Penston (1969a,b) is probably incorrect (Shu 1977; see, however, Hunter 1977) because it relies on the notion of (constant) *supersonic* inflow velocities at large r. Such an explanation is in conflict with the quasi-equilibrium condition that different parts of the cloud remain in good acoustic contact.

The configuration (Equation 22) is subject to the criticism (e.g. Whitworth & Summers 1985) that it is singular both at the origin, where the density is infinite, and at infinity, where the enclosed mass $M_{\mathrm{core}}(r) = 2a^2r/G$ increases without limit. Nevertheless, it has some desirable properties that make it a useful initial state in modeling the collapse of real molecular cloud cores. Like the many observational claims of $1/r^2$ density distributions surrounding forming stars (e.g. Bally & Lada 1983, Myers et al. 1987a, Keto et al. 1987), Equation 22 is a power law. As a consequence, although the system has a definite temperature, it has no characteristic density and, therefore, no characteristic Jeans mass. (It has one Jeans mass at each radius r, which is the only possible equilibrium accessible to a system with a virtually unlimited supply of matter.) Regarded in this sense, the infinite total mass merely simulates a situation in which the gas potentially available for star formation greatly exceeds the mass of the object finally formed. Similarly, the infinite density at the origin may be taken to represent an idealization of the tendency for magnetic decoupling to take place only at very high central densities (see, however, the controversy cited at the end of Section 2.3.6).

4.1.2 INSIDE-OUT COLLAPSE As long as the densities in a condensing molecular cloud core span several orders of magnitude before a stage of true dynamic instability is reached, its subsequent collapse properties should closely resemble those of the singular isothermal sphere (Equation 22). The latter can be found *analytically* (Shu 1977), thereby avoiding the problems of dynamical range associated with numerical simulations of protostar formation (cf. the central "sink cell" in the calculations by Boss & Black 1982). The solution at any instant in time looks like a stretched version of the solution at previous times. The self-similarity of the wave of infall, whose head propagates outward at the speed of sound, arises because of the lack of characteristic time and length scales; hence, except for scaling factors, all quantities must be functions of the similarity variable $x \equiv r/at$.

An accreting protostar of mass $M \equiv \dot{M}t$ is formed at the center, where $\dot{M}$ is the mass infall rate and t is the elapsed time since core collapse. Notice that this solution defines no characteristic mass scale that can be associated with the forming star. Instead of a mass scale, it is the (constant) *rate* $\dot{M}$ at which the star accumulates matter that is well defined (provided that T is well defined):

$$\dot{M} = m_0 a^3/G, \qquad 23.$$

where $m_0 = 0.975$. This result (Equation 23) is more robust than its formal derivation. The mass-infall rate of any critical cloud of gas of radius R

and mass M_{cl} must be given, in order of magnitude, by the estimate $\dot{M} \sim M_{cl}/t_{dyn}$, where $t_{dyn} = R/a$ and a is now the virial speed. If this cloud of gas was originally (marginally) supported against its self-gravity, then $a^2 \sim GM_{cl}/R$, which yields (upon elimination of M_{cl}) $\dot{M} \sim a^3/G$, independent of R [i.e. independent of any assumption about a (nonexistent) characteristic density]. In particular, the strong central concentration implied by Equation 22 is a basic consequence of self-gravitation and should almost always cause a spontaneous collapse to occur from inside-out. However, $\dot{M}$ will not necessarily be a constant in all cases; its variation with time will depend on the adopted boundary conditions and other assumptions (see, e.g., Zinnecker & Tscharnuter 1984).

To summarize, in any general configuration that is initially near critical equilibrium, the mass-infall rate can be expected to be given by Equation 23, within factors of order unity, as long as the calculation of the characteristic or signal speed a appropriately accounts for *all* mechanisms of mechanical support. Thus, if we use α and β to denote the ratios of the (nonfluctuating) magnetic pressure $B_0^2/8\pi$ and the (characteristic) turbulent pressure $\rho\langle\delta v^2\rangle$ to the thermal pressure $\rho kT/m$, respectively, we are tempted to identify a with the generalized magnetosonic speed (Stahler et al. 1980a):

$$a = [(1+2\alpha+\beta)kT/m]^{1/2}. \qquad 24.$$

Given the likely anisotropic nature of the magnetic and "turbulent" support mechanisms, however (see Section 2.3.2), Equation 24 must represent an overestimate for the actual "effective speed of sound."

In this context, it is interesting to note that the one-dimensional simulations of core formation by ambipolar diffusion (see Section 3.3) tend to produce almost *spatially* constant values of α (of order unity when magnetic fields are still important for core heating). On the other hand, multi-dimensional models of clouds that flatten as they contract under a constraint of field freezing show a tendency for α (of a given fluid element) to stay *temporally* constant (Mouschovias 1976, Scott & Black 1980). If the turbulence is further constrained to be sub-Alfvénic ($\beta < 2\alpha$), the maximum correction for a in cloud cores from its purely thermal values cannot exceed a factor of ~ 2 (at least, for subcritical clouds), a conclusion that seems empirically validated by an examination of molecular linewidths (Goodman 1987). Unfortunately, certain key formulae (e.g. Equations 23 and 25) are very sensitive to the parameter a. Realistic MHD calculations of cloud core collapse are needed if we are to replace heuristic devices like Equation 24 with honest a priori estimates. Until such estimates are available, the only pragmatic approach is to regard a as a free parameter that is to be determined by best empirical fits to the observational data (see, e.g., Adams et al. 1987).

4.1.3 THE EFFECTS OF ROTATION The effects of (slow) rotation at an initial uniform rate Ω can be incorporated in a rigorous treatment using singular perturbation theory (Terebey et al. 1984). The axisymmetric time-dependent collapse, with variations in two spatial dimensions, can be followed analytically if the initial state is taken to be a slowly rotating, singular isothermal sphere. Ahead of the expanding wave of infall but where the equilibrium rotation is still subthermal, $at < r < a/\Omega$, the density distribution very nearly satisfies Equation 22. For r much less than at and much greater than a centrifugal radius,

$$R_C \equiv G^3M^3\Omega^2/16a^8 \tag{25.}$$

(where $M \equiv \dot{M}t$ is now the mass of the central star plus disk), the density law has an approximate free-fall form:

$$\rho = \frac{\dot{M}}{4\pi(2GM)^{1/2}}r^{-3/2}, \tag{26.}$$

appropriate for quasi-steady infall on nearly radial streamlines. Inside of R_C (the position where infalling matter in the equatorial plane encounters a centrifugal barrier if it conserves its initial specific angular momentum), the isodensity contours are highly flattened and the rotationally modified similarity solution asymptotically joins onto free-fall ballistic trajectories (see, e.g., Ulrich 1976), which are parabolas in the limit of a very concentrated central mass distribution. When R_C is greater than the stellar radius R_*, a nebular disk forms around the protostar, whose outer dimension extends at least to R_C (and beyond if the disk has internal mechanisms for appreciable transport of mass and angular momentum in an infall time scale $M/\dot{M}$; see Cassen & Moosman 1981). Averaged over all angles, the density outside the disk between R_C and R_* behaves approximately as $\rho \propto r^{-1/2}$ (Chevalier 1983, Adams & Shu 1986).

To fix ideas, consider a typical numerical example: a molecular cloud core having negligible magnetic field and "turbulence" with $T = 10$ K, $m = 2.3\ m_H$, and $\Omega = 1$ km s^{-1} pc^{-1}. With these values of T and m, Equation 24 implies that $a = 0.2$ km s^{-1}; Equation 22 then predicts that densities $\sim 3 \times 10^4$ cm^{-3} should be reached at $r \sim 10^{17}$ cm (in reasonably good agreement with the NH_3 maps of Taurus cores by Myers & Benson 1983); and Equation 23 yields $\dot{M} = 2 \times 10^{-6}\ M_\odot$ yr^{-1}, which would build up a central mass $M = 0.5\ M_\odot$ in $t = 8 \times 10^{12}$ s, at a time when a stellar wind may already have started to blow (see Section 4.4). Infall speeds on the order of 1 km s^{-1} would be acquired only for $r \sim 10^{16}$ cm (see Figure 3*b* of Shu 1977) and would be detectable by the current generation of single-dish millimeter-wave telescopes only through special spectroscopic techniques (see, e.g., Walker et al. 1986). With the adopted value of Ω, Equation 25 gives $R_C = 45$ AU, which would be in rough agreement with the sizes of dust disks inferred to surround T Tauri stars (see Section 6.2).

The above description gives a representative accounting of the (isothermal) collapse dynamics for situations where centrifugal support of the collapsed object arises at scales that are small in comparison with the original dimensions of the interstellar gas mass from which it formed. Thus, the results can probably be applied to the generic problems of the formation of single stars with planetary systems and (ultimately) binary stars with relatively short orbital periods (e.g. less than 10^3 yr). The fragmentation of cloud cores with large original rotation rates, or of clumps with supercritical initial mass-to-flux ratios, requires other approaches (see Section 1.5).

4.2 *Evolution of High-Mass Protostars*

As long as the relevant regions remain reasonably optically thin, the core collapse proceeds nearly isothermally (see, e.g., the discussion of Hayashi 1966). Indeed, the rapid loss of any compressional heat generated guarantees the existence of a dynamical phase in protostellar evolution (Cameron 1962, Gaustad 1963, Gould 1964). The runaway increase of the density at the center is halted only after the formation of an optically thick protostar (of mass $\sim 10^{-3}$–$10^{-2}\ M_\odot$) whose thermal time scale exceeds the sound crossing time. The protostar then grows hydrostatically (after a transient phase associated with the dissociation of molecular hydrogen; see Larson 1969a) by accreting matter from infalling gas and dust.

Certain general points concerning the structure of protostars were recognized early—for example, protostars of high mass should evolve differently from those of low mass. The reason is that the protostar changes on a Kelvin-Helmholtz time scale GM_*^2/R_*L_*—the time required to reach internal thermal equilibrium—whereas the infall takes place at a rate that is not highly variable. For a massive protostar the Kelvin-Helmholtz time scale will be shorter than the infall time scale, and the star will reach the main sequence while still gaining mass from the infalling envelope (Kahn 1974; see also Yorke & Krugel 1977). This conclusion seems in reasonable accord with the observations (see the review by Wynn-Williams 1982).

A complication arises because the radiation pressure acting on dust grains may become large enough to halt the infall. Indeed, Larson & Starrfield (1971) and Kahn (1974) proposed this as a possible mechanism (in addition to expanding H II regions) to set an upper limit on stellar masses; however, the observed upper limits (see, e.g., the review by Humphreys & Davidson 1984) disagree with the best theoretical value using modern dust opacities. Within factors of order unity, radiation pressure acting on grains can reverse spherical infall if the ratio of luminosity to mass, L/M, of the central source exceeds the critical value,

$$L/M = 6\pi c G/\kappa_P(T_d) \approx 700\ L_\odot/M_\odot, \tag{27.}$$

where $\kappa_P(T_d)$ is the Planck-mean of the opacity at the dust destruction temperature and has been taken to be 30 $cm^2\ g^{-1}$ for standard dust abundances. Main sequence stars heavier than 7 $M_\odot$ have L/M that exceed this critical value. A proposed escape from this dilemma involves the postulate that massive stars form only in regions relatively depleted of dust, especially graphite grains (Wolfire & Cassinelli 1987). Another possibility is that the infall is highly nonspherical because of rotation, which thereby allows matter to fall into a disk (which is shielded from the direct radiation of the star) at radii substantially larger than the dust destruction front. Accumulation of a high-mass star may then proceed via disk accretion. The newly born OB star may later reveal itself optically, like its low-mass counterparts, by producing a stellar wind that removes the surrounding placenta of gas and dust. Another argument against terminating the infall by radiation pressure (as contrasted with a wind) is that the radial momentum carried in a typical cold outflow (probably swept-up molecular gas) exceeds the momentum contained in all the photons that have ever left the star by $\sim$2–3 orders of magnitude (see the review of Lada 1985).[5]

4.3 *Evolution of Low-Mass Protostars*

The longer Kelvin-Helmholtz contraction times associated with lower stellar masses imply that they may end their accretion phase without having reached the main sequence. To determine where in the Hertzsprung-Russell diagram such objects might first become optically visible requires accurate and detailed computations of the ensuing radiative hydrodynamics. However, the collapse of an interstellar cloud is highly nonhomologous (McNally 1964); in order to follow the thermal and mechanical structure of the central protostar, the calculation must include a range of densities spanning more than 20 orders of magnitude. This enormous dynamical range strained the capabilities of most numerical attacks on the problem, even in its simplest spherical formulation, and it led to considerable early controversy (Larson 1969a, 1972, Narita et al. 1970, Hayashi 1970, Appenzeller & Tscharnuter 1974, 1975, Westbrook & Tarter 1975, Kondo 1978, Tscharnuter & Winkler 1979).

Improvements in numerical techniques (Winkler & Newman 1980a,b) and reformulation of the computational problem to take advantage of

[5] Multiple scatterings in a dust envelope cannot help because the photons would be degraded (by true absorption and reradiation) to the far-infrared; the photons can then easily escape before experiencing the required number of scatterings.

certain analytical insights (Stahler et al. 1980a,b, 1981) have resolved most of the controversy in the spherical case. In particular, the radius R_* of a protostar that grows in mass at a given rate $\dot{M}$ is now known with good accuracy. When the freely falling material in the envelope encounters the hydrostatic surface of the star (at R_*), it dissipates its energy in an accretion shock, producing the luminosity

$$L = GM\dot{M}/R_*. \qquad 28.$$

Notice that Equation 28 assumes that essentially all of the kinetic energy of the infall is converted into radiation; this holds to a high degree of approximation in the low-mass protostar problem. If, instead, a substantial fraction of the infall energy were to be converted into internal energy of the gas, the resulting stellar radius would be much larger than the computed value of a few solar radii.

The formula (Equation 28) yields the system luminosity only for low-mass protostars with spherical symmetry. For high-mass protostars, the interior luminosity will dominate the shock luminosity. Also, for a low-mass protostar with rotation, the (infall) luminosity will be lower (for given $\dot{M}$) if angular momentum conservation prevents material from falling all the way to the stellar surface and if part of the gravitational energy released can be stored as rotational energy in the star or in an accompanying circumstellar disk (Adams & Shu 1986).

4.4 *Termination of the Infall*

In the spherical calculations for a mass infall rate $\dot{M} = 1 \times 10^{-5}\ M_\odot\ \mathrm{yr}^{-1}$ (Stahler et al. 1980b), the protostar accumulated matter (processed through an accretion shock) of ever increasing specific entropy. Hence, the star remained radiative until the ignition of deuterium, which occurred when the stellar mass was $\sim 0.3\ M_\odot$. A convection zone then spread outward through the star until the protostar became almost entirely convective at a mass of about 0.5 $M_\odot$. Except for this event, nothing happened to distinguish any particular mass scale for the accreting protostar; in the actual calculations, the infall was terminated artificially after 1 $M_\odot$ was accumulated.[6] After a short period of readjustment, the newly formed star then descended a convective pre-main-sequence track (Hayashi et al. 1962), developed a radiative interior, and followed a radiative track to the main sequence (Henyey et al. 1955).

The above scenario resurrects the fundamental issue of star formation outlined in Section 1.1. In the present context, we may pose the question

[6] As commented upon by the authors, however, models for lower final masses can be simply recovered from the published calculations by merely terminating the accretion earlier (see, e.g., Stahler 1983).

as follows: What mechanism actually determines the end of the infall phase? The large reservoir of gas in molecular clouds and the general inefficiency of star formation make it unlikely that the star continues to gain mass until there is no available material left. A more plausible explanation is that stellar winds eventually reverse the inflow and halt the accretion (see, e.g., Shu & Terebey 1984, Lada 1985, Myers et al. 1987b). This viewpoint, however, raises another important issue—what triggers the onset of these winds?

4.5 *Deuterium Burning and the Stellar Birthline of Low-Mass Stars*

An empirical clue for low-mass stars exists in the Hertzsprung-Russell (H-R) diagrams of T Tauri stars (Cohen & Kuhi 1979). Stahler (1983) was the first to note that H-R diagrams for very different stellar groupings—Taurus-Auriga, Orion, NGC 7000/IC 5070, and Ophiuchus—all tended to show the same upper envelope (see Figure 4). It was as if pre-main-sequence stars first appeared optically visible on a universal "birthline." Stahler's explanation for this amazing fact in terms of a special accretion rate in spherical geometry was criticized by Mercer-Smith et al. (1984). Shu (1985) pointed out that the observed "birthline" could be understood

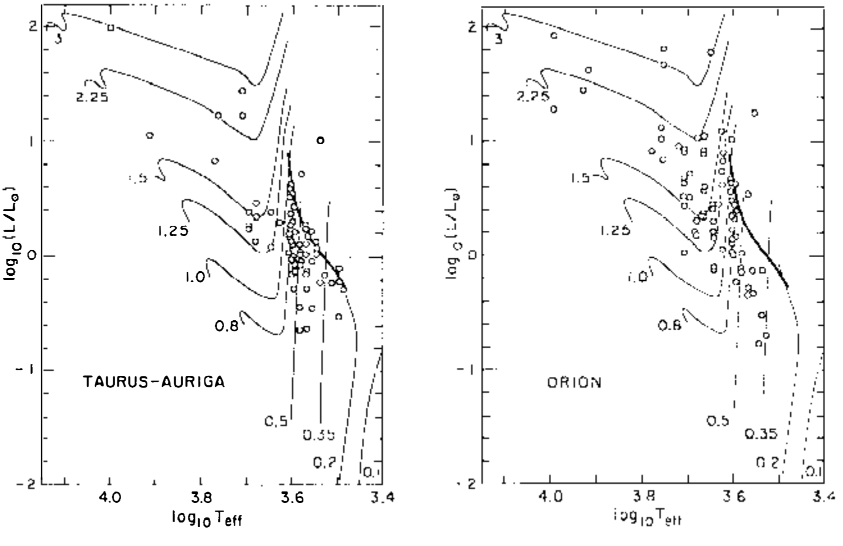

Figure 4 Hertzsprung-Russell diagrams from Cohen & Kuhi (1979) showing theoretical pre-main-sequence contraction tracks and T Tauri stars in the Taurus-Auriga and Orion cloud complexes. The heavy solid curve is the theoretical "birthline" of Stahler (1983).

as a condition that the star must be able to burn deuterium, thereby giving a connection with the fusion capabilities of bona fide stars. In a low-mass star, deuterium burning can drive convection, which, when coupled with the presence of differential rotation, can produce dynamo action and generate magnetic activity (Parker 1979). The released energy (which had been stored in rotation) may then power the intense stellar surface activity observed in YSOs. The resulting stellar wind eventually reverses the infall and thereby defines the mass of the formed star at the point when it first becomes optically visible.

For any given mass, there is a unique radius for which a completely convective star will have a central temperature (Chandrasekhar 1939),

$$T_c = 0.54 GM_* m / kR_*, \tag{29}$$

high enough for deuterium to burn. In Equation 29 we have assumed that the convective star is an ideal gas of constant mean molecular weight m (i.e. a polytrope of index 1.5). For a completely ionized gas of cosmic abundances, $m = 0.62\, m_H$, and for $0.01\, M_\odot < M_* < 2\, M_\odot$, we can obtain an approximate criterion for deuterium burning by setting $T_c = 1 \times 10^6$ K in Equation 29:

$$R_*/R_\odot \approx 0.15 + 7.6(M_*/M_\odot), \tag{30}$$

where the small correction term 0.15 (and the slight readjustment of the coefficient 7.6) on the right-hand side enters to take into account the effects of partial degeneracy at low stellar masses (cf. Nelson et al. 1986). Equation 30 produces a birthline in the H-R diagram for pre-main-sequence stars that depends only on the physics of stars and not on the details of their formation. Given the observational uncertainties and the possibility that the stellar luminosity can be contaminated by radiation from an active disk (especially at low masses; see Sections 6.4 and 6.5), the predicted birthline agrees reasonably well with the observed upper envelope for T Tauri stars.

The rate of stellar mass accumulation, $\dot{M}_*$, affects only the timing of deuterium ignition (i.e. the eventual mass of the star, M_*, on the birthline). Higher rates of stellar mass accumulation allow less time for the radiation of the stellar binding energy, which leads to larger values of R_* at every M_* and thus produces higher masses before the onset of deuterium burning. The final stellar masses will also depend on whether the infall is sufficiently intense to suppress a breakout of the stellar wind despite the convection induced by deuterium burning, and on the amount of stellar accretion that occurs from a disk. In any case, for spherical stars that accumulate $\sim 2\, M_\odot$ before deuterium ignites, the extra interior luminosity released by deuterium burning can be carried out by radiative diffusion without the inducement of convection. [See the discussions of Cassen et al. (1985) and of Stahler et al. (1986).] Therefore, the termination of stellar accumulation for high-mass stars must occur by a different process than that outlined above (possibly deuterium burning in a pseudodisk).

4.6 *Emergent Spectral Energy Distributions of Protostars*

A test of any theory of star formation is to compare the theoretical predictions of the emergent spectral energy distributions of protostars with actual observations. Historically, calculations of emergent protostellar spectra have fallen into two basic categories. The parametric approach ignores the hydrodynamic portion of the problem, assumes a known central source of luminosity, constructs a series of model dust shells, and conducts a survey of the resulting parameter space. Such a strategy has been used to study emission from dust shells around late-type stars (e.g. Jones & Merrill 1976), from stars embedded in extended dust envelopes (Rowan-Robinson 1980), and from protostellar dust clouds forming high-mass stars (Yorke & Shustov 1981). The self-consistent approach begins with a hydrodynamic model of protostellar collapse (which includes an internally determined central source of luminosity) and either solves the radiative transfer problem simultaneously or postprocesses the results of a rigorous collapse calculation. Such an approach typically has fewer adjustable parameters and has been used to study the dynamical evolution and spectral appearance of protostars of low mass (Larson 1969b, Bertout 1976), intermediate mass (Yorke 1979, 1980), and high mass (Yorke & Krugel 1977, Yorke 1977, Wolfire & Cassinelli 1986). All of these basic models have been restricted to spherical symmetry. Recently, a fast approximate technique has been developed, which is applicable to multidimensional situations as well (Adams & Shu 1985, 1986).

From the parametric studies, given a known central source, the emergent (infrared) spectra are found to be sensitive primarily to the total optical depth of the dust envelope and to the optical properties of the dust grains (usually taken to be a mixture of graphite and silicate grains), which provide the dominant contribution to the opacity κ_ν. If the dust envelope is sufficiently optically thick so that most of the original radiation from the central source is absorbed and reradiated by the grains, the resulting emergent spectra consist of a single broad hump of emission with a peak in the far-infrared ($\sim$60–100 μm), and absorption features at 10 μm (from silicate grains) and 3.1 μm (from water ice). If the column density through the dust envelope is not large, a substantial fraction of the central source radiation can escape, and a double-peaked spectrum is produced. In all cases, the emergent spectra are broader than that of a single-temperature blackbody. This is a general feature of "extended atmospheres" (Mihalas

1978) and implies limited utility to the concept of a dust "photospheric temperature."

Good infrared opacities for dust grains have become available in recent years through better observational constraints (see Hildebrand 1983, and references therein) and through better self-consistent calculations (see Draine & Lee 1984). On the other hand, our theoretical knowledge of the properties of the central source and of the total optical depth through the cloud core to the dust destruction front (along any given line of sight in rotating models) has also greatly improved (see previous section). Thus, detailed comparison of theoretical predictions with specific protostellar candidates can go substantially beyond the purely parametric studies because the adjustable parameters can be related to measurable mechanical properties of the source that are independent of the observed spectrum. In addition, far-infrared observations by the IRAS satellite (e.g. Beichman et al. 1986) and follow-up ground-based observations in the mid- and near-infrared (such as those of Myers et al. 1987a) have considerably expanded the relevant data base. Thus, it is encouraging that theoretical calculations based on the collapse models of slowly rotating cloud cores discussed in Section 4.1.3 and the protostars discussed in Section 4.4 yield spectral energy distributions that are in close agreement with observations of a class of low-luminosity infrared sources—those that have negatively steep spectra in the near- and mid-infrared and are found near the centers of dense molecular cloud cores (Adams et al. 1987; see Figure 5 for an

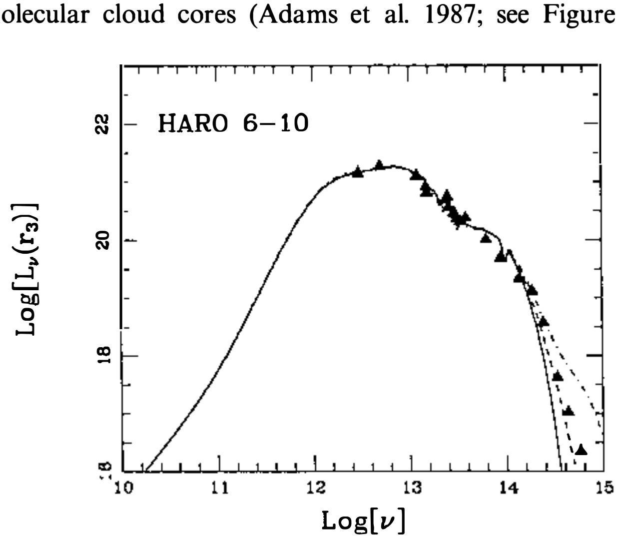

Figure 5 Model fit to spectral energy distribution (cgs units) of Haro 6–10. The solid curve assumes purely absorbing dust grains; the dashed-dotted and dashed curves make different heuristic corrections for scattering (adapted from Adams et al. 1987).

example). Major theoretical uncertainties remain concerning the appropriate values to adopt for the efficiencies η_* and η_D with which the energy of rotation in dissipated in the star and disk (see Adams & Shu 1986); however, within reasonable bounds, the inferred values of the parameters of the rotating cloud core, a and Ω, and the resulting mass M_* and radius R_* of the central star, seem consistent with current observational knowledge of these quantities (see also the discussion of Myers et al. 1987a).

5. BIPOLAR OUTFLOWS FROM YOUNG STELLAR OBJECTS

5.1 *Basic Energetics*

The growing base of observational data concerning outflows from YSOs suggests that all stars pass through an outflow stage as a fundamental part of the star formation process (see the reviews by Lada 1985, Welch et al. 1985, and Snell 1986). The outflows are highly energetic (observed kinetic energies are $\sim 10^{43}$–10^{47} erg) and exhibit a wide variety of observable phenomena. Related manifestations include Herbig-Haro (H-H) objects (see the review by Schwartz 1983), high-velocity water maser sources (see the review by Reid & Moran 1981), and "cometary" reflection nebulae (e.g. Strom et al. 1985, Goodrich 1987). Some level of collimation usually exists for the outflows; the most notable examples are massive bipolar flows of cold molecular gas (e.g. Snell et al. 1980), and long and narrow optical jets (e.g. Mundt & Fried 1983). The optical jets (often linear chains of H-H objects) can often be traced very close to the exciting (low-mass) stars and almost certainly originate from them. The molecular outflows (which occur in YSOs of all masses) may also be driven by strong stellar winds, but this suggestion is more controversial. The case for stellar winds is supported by the claim that the CO gas in the bipolar lobes of L1551 exists in spatially thin shells (Snell & Schloerb 1985, Kaifu 1986).

The best-known alternative possibility—centrifugally driven flows (Pudritz & Norman 1983) or magnetohydrodynamically driven flows (Uchida & Shibata 1985, and references therein) from large ($\sim 10^{17}$ cm) and massive ($\sim 10^2$ $M_\odot$) circumstellar disks—has severe mechanistic drawbacks. In the Pudritz & Norman model, in order to get molecular outflows with a speed $\sim 10^1$ km s^{-1} from the parts of disks where the typical rotational velocities are $\sim 10^0$ km s^{-1}, Alfvénic enforcement of corotation (constant angular velocity) must occur over a range of radii that typically spans a factor of 10. Since the moment of inertia per gram of the outflow gas is then larger (by a factor of 10^2) than that of the material in the disk that ejects it, the back reaction (i.e. the conservation of total angular momentum) produces a mass flux through the disk (leading to protostellar accretion) that is typically 10^2 times larger than the mass outflow. Since the observed *molecular* outflows (which in a more conventional interpretation would be swept-up cloud gas rather than material from a circumstellar disk) involve masses and time scales of the order of solar masses and 10^4 yr, respectively, the deduced disk accretion rates are $\sim 10^{-2}\ M_{\odot}\ yr^{-1}$ or larger. This severe requirement leads to some unusual conclusions (see, e.g., Table 1 of Pudritz & Norman 1986): The disk masses must be $\sim 10^2$ $M_{\odot}$ or larger (to be able to last 10^4 yr), the central protostars must be built up very quickly (at a rate that would make them very transient objects), and they must have very small mass-to-radius ratios ($\sim 10^{-3}$ in solar units) in order to avoid an overly large accretion luminosity from feeding material into a deep gravitational potential well. These requirements are not substantiated, to our knowledge, by any known properties of molecular cloud structures or the stars that they form.

The problem is, of course, that very big disks (size $\sim 10^{17}$ cm) have very little specific energy; thus, to use them to drive the observed molecular outflows requires making them very massive. Although spatially large disks have often been claimed in the observational literature (e.g. Kaifu et al. 1984, but see Batlra & Menten 1985), the only well-substantiated case, S106 (which is both highly flattened and rapidly rotating), has a probable mass of 2 $M_{\odot}$ (Bieging 1984), and the star at its center has a (large) mass-to-radius ratio (judging from its luminosity and ionizing flux) appropriate to a main-sequence star of $\sim 20\ M_{\odot}$. The paucity of specific energy contained in spatially big disks makes them unlikely candidates for the power behind bipolar flows.

In all self-gravitating systems, most of the specific binding energy is contained at the smallest radii. Rapidly rotating protostars, which are $\sim 10^6$ times smaller than 10^{17} cm, potentially have 10^6 times more energy per gram that can be dissipated to power winds. On energetic grounds alone, therefore, the model of Draine (1983), which makes magnetohydrodynamic use of a rotating protostar, is preferable to that of Uchida & Shibata, which is an otherwise beautiful numerical simulation. A problem with Draine's model is that it requires a tightly wrapped magnetic field configuration for the initial state, which may be difficult to produce from a smoothly evolving scenario for the formation of the star.

In any case, supplying enough mechanical energy is not a problem for young stars, but supplying enough (radial) momentum may be. The momenta contained in the measured ionized winds from YSOs are insufficient to account for the cold molecular flows by one to two orders of magnitude (see, e.g., Levreault 1985); a heretofore undetected neutral component of the stellar wind may be necessary (Rodriguez & Canto 1983, Snell et al. 1985). In the famous bipolar flow source L1551, indirect evidence for hidden mechanical luminosity in excess of the known stellar wind (and jet) may have been found in the presence of far-infrared emission that is too strong and too extended for direct heating by the central protostar, IRS 5 (Clark & Laureijs 1986, Edwards et al. 1986, Clark et al. 1986).

5.2 *Influence of Outflows on the Surroundings*

It has been proposed that outflows from young stars represent an important input that maintains the "turbulence" in molecular clouds (see the discussion in Section 2.3.1). In further developments along such lines, feedback from such effects has been proposed as an important step in the self-regulation of star formation (e.g. Norman & Silk 1980, Franco 1984). In such scenarios, star formation will affect the subsequent evolution of the parent molecular clouds, which in turn will affect further star formation: cloud $\rightarrow$ protostar $\rightarrow$ outflow $\rightarrow$ cloud $\rightarrow$ .... However, the detailed manner in which these schematic arrows work requires more elucidation.

Stellar outflows are also useful as diagnostics of their surroundings. For example, warm H_2 molecules moving at high velocities have been amply detected in regions containing molecular outflows [e.g. Lane & Bally (1986); see Shull & Beckwith (1982) for a review]. These observations are most simply understood in terms of shock-accelerated gas produced by the interaction of powerful stellar winds with the ambient medium. However, ordinary gas-dynamic shocks of high speed should dissociate the hydrogen molecules (see the review of McKee & Hollenbach 1980). Since the molecules are left intact, warm, and radiating, they must be accelerated smoothly in magnetized "C shocks" (Draine 1980, Chernoff et al. 1982), an ungentle version of ambipolar diffusion. The required magnetic fields typically correspond to strengths of ~ 1 mG in regions where the ambient density is $\sim 10^6\ \mathrm{cm}^{-3}$.

The properties of outflows have also been used to deduce the density of the ambient gas before it was swept up into a thin shell. Dividing the total mass of swept-up gas by the volume of the lobes (as computed from the geometric mean of the major and minor diameters) in a number of flow sources, Bally & Lada (1983) show that the statistics are consistent with the radial distribution of gas density falling off as $1/r^2$ in star-forming regions.

5.3 *Collimation Mechanisms*

Proposed mechanisms for collimating stellar jets and bipolar molecular flows fall into two generic categories. One school of thought holds that outflows are produced in a well-collimated form near the star (e.g. Hartmann & MacGregor 1982, Jones & Herbig 1982). The other viewpoint holds that an initially isotropic stellar wind is channeled into a bipolar form by an anisotropic distribution of matter in the surrounding circumstellar environment (e.g. Canto et al. 1981, Konigl 1982, Torrelles et al. 1983; but see Heyer et al. 1986). Both possibilities in a very general sense may arise naturally (as a progression in time) in the scenario for star formation posed in Section 4 (Shu & Terebey 1984, Shu et al. 1986).

The basic idea is that during the infall phase, the swirling inflowing matter acts like the lid of a pressure cooker and suppresses *any* incipient stellar outflow. As the lid weakens (because the gas falls increasingly onto the disk rather than directly onto the star), breakout will occur through the channel of least resistance (the safety valve) at the rotational pole(s) of the accreting protostar, where the ram pressure of the infalling material is least during the period when R_C is not much larger than R_*. Since the system is smoothly transforming from an inflow state to an outflow state, breakout in the absence of perfect spherical symmetry should first occur at one point (on each hemisphere). Thus, the existence of collimated, dual-exhaust jets in the youngest outflow objects becomes a natural part of protostellar evolution. Later, the stellar outflow may widen and become nearly isotropic, but the swept-up shell of gas may still make less outward progress in the equatorial directions because of the larger amounts of obstructing (inflowing) gas (see Figure 6*d* of Terebey et al. 1984). Detailed modeling is needed to provide quantitative comparisons with observations.

In the interim, it is significant that the rotating infall models described in Section 4.5 correctly predict the emergent spectra for some infrared sources that are known bipolar outflow sources [see Adams et al (1987) for examples]. This finding, coupled with the observation that well-collimated sources tend to be heavily embedded objects, supports the idea that they represent systems in which inflow and outflow are taking place *simultaneously*. In this view, a bipolar flow source is the transitional phase of evolution between a purely accreting protostar and a fully revealed pre-main-sequence star. A configuration of combined inflow and outflow has been claimed on completely different grounds for the Becklin-Neugebauer object in Orion (see the review of Scoville 1985), which is probably a case of massive star formation.

If magnetic activity plays an intrinsic role in the generation of the stellar winds from low-mass protostars (e.g. Lago & Penston 1982; see also Section 4.4), the outflows will carry away material with a higher specific angular momentum than the average for the star. Intense mass loss will naturally produce intense magnetic braking. Thus, the slow rotation rates observed for (revealed) T Tauri stars (Vogel & Kuhi 1981, Bouvicr et al. 1986, Hartmann et al. 1986), which were quite surprising when first discovered, can now be attributed to an earlier epoch of heavy mass loss.

6. VISIBLE YOUNG STELLAR OBJECTS

Since high-mass stars are believed to join the main sequence while still in the accretion phase (see Section 4.2), they do not convey much information concerning their origins when they become visible as optical objects. The study of pre-main-sequence stars is thus operationally confined primarily to low-mass stars, although YSOs of intermediate mass, the Herbig Ae and Be stars (Herbig 1960, Strom et al. 1975, Finkenzeller & Mundt 1984), deserve more attention.

6.1 *T Tauri Stars*

Soon after their discovery (Joy 1942, 1945, 1949), T Tauri stars were recognized as being a very young population of (low-mass) stars, newly born within the dark clouds with which they are invariably associated both spatially and kinematically (e.g. Herbig 1977). The optical spectra of T Tauri stars are characterized by high variability and prominent emission lines, most notably the Balmer series of hydrogen and Fe II lines (see Herbig 1962). The underlying photospheres are relatively cool, usually less than 6000 K and typically $\sim$4000 K. The strength of the ultraviolet emission lines (and continuum excesses) have traditionally been interpreted to indicate heavy chromospheric activity (see, e.g., Calvet et al. 1984, Herbig & Goodrich 1986), although the boundary layer associated with an accretion disk is another possibility (Bertout 1987). T Tauri stars have strong mass loss (Kuhi 1964), and periodic components to the light variations suggest that these stars have very inhomogeneous surfaces, with an extensive network of magnetically active and quiescent zones (e.g. Herbig & Soderblom 1980, Bertout et al. 1987). Thus, T Tauri stars seem to be bubbling with surface activity, which together with appreciable surface lithium abundances (Herbig 1960) is yet another indication of their extreme youth. Since a comprehensive review of the properties of T Tauri stars is available elsewhere (Cohen 1984), we concentrate here on the characteristics of T Tauri stars that bear most directly on the theory of low-mass star formation.

T Tauri stars presumably represent the postinfall phase of young star evolution. When the object becomes optically revealed, has its bolometric luminosity measured, and can be given a spectral type, it can then be placed in the H-R diagram. The task has been completed for $\sim$450 T Tauri stars in an extensive survey by Cohen & Kuhi (1979). Their results indicate that T Tauri stars are mostly descending convective tracks, in apparent accord with the classical pre-main-sequence theory of Hayashi et al. (1962; see also Iben 1965, Ezer & Cameron 1965). The implied mass range of these objects is 0.2–3 $M_{\odot}$; the formal contraction ages range from $\sim 2 \times 10^5$ to 2×10^7 yr.

The activity of T Tauri stars may result simply from the adjustment needed to accommodate material recently acquired from the interstellar medium to stellar conditions. A finding that supports this general point of view is the observed decay, as T Tauri stars age, in the properties that earmark these objects as a distinct class: surface activity, infrared excesses, and emission lines (Cohen et al. 1987, Walter 1987).

The overall spread of ages ($\sim 10^7$ yr) in a cloud like Taurus is consistent with the theory of molecular cloud core formation by ambipolar diffusion outlined in Section 3.3. The most recently born stars require corrections to take into account the dynamical stages of protostellar evolution (e.g. Stahler 1983). The derived "ages" are further uncertain to the extent that a substantial contribution to the bolometric luminosity may not originate with the star but with an active accretion disk (see Section 6.5). Nevertheless, within a factor of 2 or so, the "contraction ages" (or, more precisely, the locations in the H-R diagram) of the youngest T Tauri stars are consistent with the mass accumulation rates implied by the core collapse scenario discussed in Section 4.1.2.

Luminosity overestimates would introduce mass overestimates if the stars are on the radiative portions of their tracks. Mass determinations for T Tauri stars may, therefore, be somewhat uncertain at the highest values. Since convective tracks are nearly vertical in the H-R diagram, variations in the stellar luminosity estimates produce little error in the mass determinations at low values. More serious for the latter are the assignment of accurate effective temperatures for objects that have strong emission lines, anomalous continuum fluxes, and (sometimes) composite spectral types. It would be important to check whether, within the uncertainties, the masses and radii of T Tauri stars really are consistent with their protostellar progenitors having had their accretion phases halted by the onset of deuterium burning (see especially the discussion of the "birthline" in Section 4.4).

6.2 *Dusty Nebular Disks*

In recent years there has been growing evidence for dusty disks of size $\sim 10^2$ AU around nearby stars. The high degree of optical linear polarization from a few sources (in excess of 10%) has been interpreted (Elsasser & Staude 1978) as scattering from an elongated or flattened dust distribution [one that partially extinguishes the starlight (i.e. is not completely spatially thin; otherwise, the direct unpolarized component coming from the star would reduce the degree of polarization)]. More recently, IRAS observations and follow-up ground-based work have produced resolved images of disklike structures around somewhat older stars (Aumann et al. 1984, Smith & Terrile 1984). Lunar occultation studies at infrared wavelengths also reveal similar elongated configurations around stars newly born from molecular clouds (Simon et al. 1985).

The star HL Tau has proven to be especially intriguing. One unique aspect of this T Tauri star is the presence of a deep water-ice feature at 3.1 μm in its spectrum; this feature has been interpreted (Cohen 1975) as evidence indicating that the object is especially young. Further observations of the unusual infrared emission and absorption features of HL Tau led Cohen (1983) to suggest that in this system we are looking through an edge-on disk, which had a radius $\sim 10^2$ AU, an (estimated) axial ratio of 100:1, and $\sim 10^{-9}$ $M_\odot$ in silicate and ice grains. Direct near-infrared imaging (Grasdalen et al. 1984) and speckle interferometry (Beckwith et al. 1984) seemed to confirm this suggestion, with the scattering optical depths yielding a revised (but still small) estimate of $\sim 10^{-7}$ $M_\odot$ for the total mass of the grains (assuming a standard interstellar size distribution). Follow-up radio-interferometric observations (Beckwith et al. 1986) in the $J = 1\text{-}0$ line of ^{12}CO and in the 2.7-mm continuum found the source to be essentially unresolved with a beam size of 6″ (corresponding to a radius of ~ 400 AU at 140 pc, the distance of the Taurus cloud). The continuum flux at 2.7 mm, interpreted as thermal emission from warm grains, implied a dust mass between 1×10^{-4} and 2×10^{-3} $M_\odot$, depending on the emissivity law adopted for the grains. Preliminary ^{13}CO measurements by Sargent & Beckwith (1987) confirm the likelihood of up to 0.2–0.3 $M_\odot$ of gas and dust lying in a rotating disk centered on the system. Thus, although Cohen's original suggestion of a dusty disk in HL Tau of roughly solar-system dimensions has proven exceptionally fertile, there is an unsettling 6 orders of magnitude difference between the smallest and largest estimates for the mass of solid material present. In Section 6.5 we offer a resolution of this discrepancy.

The above discussion summarizes the direct observational evidence concerning small disks ($\sim 10^2$ AU, as compared to the $\sim 10^4$-AU objects occasionally claimed by molecular radio astronomers) around a few YSOs. However, there is *indirect* evidence that suggests the almost universal existence of such structures around T Tauri stars; this evidence comes from the examination of the infrared excesses in T Tauri spectra.

6.3 *Infrared Excesses of T Tauri Stars*

As early as two decades ago, the spectral energy distributions of T Tauri stars were observed to exhibit excess infrared radiation (Mendoza 1966, 1968). Since then, a multitude of data have been accumulated [cf. Rydgren et al. (1984) for a catalog of T Tauri stars in Taurus-Auriga]. Using several

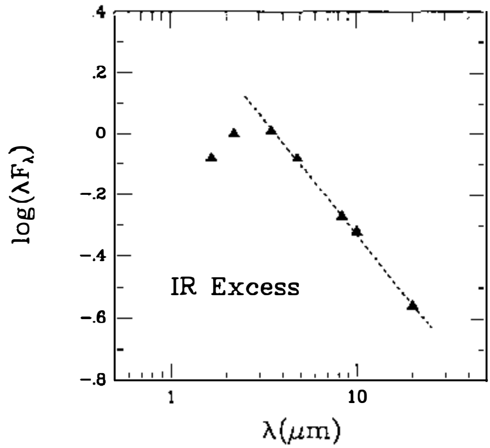

Figure 6 Composite spectral energy distribution of nine T Tauri stars. The dashed line shows that the near- and mid-infrared excess is well represented as a power-law spectrum with index $n = 0.75$. A stellar blackbody would give $n = 3$ (from Rydgren & Zak 1987).

representative sources, Rydgren & Zak (1987) have produced a composite spectrum, as shown in Figure 6, that clearly displays the nature of the infrared excess. If we define the spectral index n by

$$n \equiv \frac{d \log (\nu F_\nu)}{d \log \nu} = - \frac{d \log (\lambda F_\lambda)}{d \log \lambda}, \qquad 31.$$

then $n \approx 0.75$ for the composite spectrum at near- and mid-infrared wavelengths. Individual sources in the catalog of Rydgren et al. (1984) have a wide range of spectral indices n, varying from somewhat greater than 1 (steep spectrum sources) to $n \approx 0$ (flat spectrum sources).[7]

Mendoza (1966) suggested that the infrared excesses of T Tauri stars were caused by circumstellar dust grains, which absorb stellar photons and reradiate them at longer wavelengths. The alternative hypothesis—that the infrared excesses are given by free-free emission—cannot explain the most extreme cases (Cohen & Kuhi 1979). A wide variety of geometries for the distribution of dust have been proposed: shells, disks, and grain formation in outflows from the young stars.

Until recently, the most popular approach has been to assume that all infrared excesses are created by more or less spherical circumstellar dust shells that surround the stars. These models generally treat the properties of the star (the effective temperature T_* and the radius R_*) and those of

[7] It has not escaped our attention that an analogous situation holds for the spectral energy distributions of active galactic nuclei. Indeed, comparisons of jets and disks in YSOs with those in quasi-stellar objects have been frequently made in the literature.

the dust shell (the inner and outer radii and the density distribution, which combine to give the total optical depth) as adjustable quantities. The resulting spectra are reasonably successful in reproducing the spectral energy distributions of many observed sources (see, e.g., Rowan-Robinson et al. 1986, Wolfire & Churchwell 1987).

For evolved sources near the end point of stellar evolution, where dusty outflows may be present, the adoption of spherical symmetry is a natural first approximation. However, dust distributed in a spherical manner *close* to an optically revealed pre-main-sequence star (so that the dust has temperatures high enough to account for the near-infrared emission) is much less plausible. Since the infall has been reversed, not much dust can be infalling; but the dust cannot be outflowing either, as the stellar temperatures are generally not low enough to allow solid grains to condense. Thus, the closer and warmer dust grains must actually be in *orbit* around the star, i.e. distributed in a flat disk. Such a model is also consistent with the ideas developed in Section 5, in which stellar outflows break out at the poles and widen to the equator, so that an equatorial disk should be the last remnant circumstellar material to be cleared away.

6.4 *Passive Disks*

Given that T Tauri stars have dusty disks around them, it can be easily shown that unless the masses of the disks (gas plus dust) are much less than $\sim 10^{-1}\ M_{\odot}$, the disks will be optically thick to their own thermal radiation, from the near- to far-infrared, out to $\sim 10^2$ AU. Even a minimum-mass solar nebula ($\sim 10^{-2}\ M_{\odot}$) should be optically thick for radii within the orbit of Neptune (Lin & Papaloizou 1985), and modest gas opacities are likely to keep the disk optically thick in the dust-free environment adjacent to the star. Thus, essentially all disks will be optically thick and reprocess a significant portion of the luminosity from the central star. Such disks can be divided into two conceptual categories: those that also have appreciable radiation arising from intrinsic luminosity in the disk (an *active* disk) and those that do not (a *passive* disk).

The theory of passive disks is especially simple: An optically thick disk, radially extended and vertically thin, will intercept and reradiate 25% of the stellar luminosity and produce a spectral energy distribution that is approximately a power law in the near- and mid-infrared with a spectral index $n = 4/3$ (Adams & Shu 1986). The simple theoretical models, which contain *no* free parameters (apart from viewing geometry) when the properties of the central star are known, are in good agreement with the observed spectra of (positively) steep spectrum T Tauri stars (see Adams et al. 1987).

In a passive disk, the equilibrium temperature distribution in the disk is given approximately by $T \propto r^{-3/4}$, which is also the classical result for a Keplerian accretion disk (Lynden-Bell & Pringle 1974); hence, the corresponding spectral energy distributions for the two types of disk have the same shape. Thus, simple reprocessing of stellar photons will mask the effects of Keplerian disk accretion if the accretion rate is low enough that the intrinsic disk luminosity is less than $\sim 25\%$ of the stellar luminosity (see also Friedjung 1985). In particular, the typical accretion rate $\dot{M}_a$ is $\sim 5 \times 10^{-8}\ M_\odot\ \mathrm{yr}^{-1}$ in a minimum-mass solar nebula driven by thermal convection (Ruden & Lin 1986), and it produces a total disk luminosity $\sim G\dot{M}_a M_*/2R_*$ that is too small to compete with disk reprocessing. Such small accretion rates will be observable only if a viscous boundary layer between the rapidly rotating disk and the slowly spinning star radiates $\sim G\dot{M}_a M_*/2R_*$ in the ultraviolet, where there is competition only from the chromosphere. Ultraviolet excesses (cf. Bertout 1987) therefore may be especially valuable diagnostics in deciding whether systems like FU Orionis (for a review, see Herbig 1977) are sources in which an outburst has taken place in the star or in the disk (Hartmann & Kenyon 1985, Adams et al. 1987).

For a given $\dot{M}_a$, the accretion luminosity scales as M_*/R_*, which is constant for a star on the deuterium "birthline" and is proportional to $M_*^{0.4}$ on the upper main sequence. But the intrinsic stellar luminosity L_* varies as a relatively high power of M_*. Thus, the disks around stars of high-enough mass must be dominated by reprocessing of starlight. This fact, coupled with the fierce tendency of high-mass stars to ablate any nearby material, may explain why near-infrared excesses associated with optically revealed YSOs of high mass are not prominent.

6.5 *Active Disks*

There is another class of T Tauri stars where n is significantly less than 4/3; in particular, there are some T Tauri stars with flat spectra ($n = 0$). Adams et al. (1987) proposed that such sources could be understood if the disks were *active* and possessed intrinsic luminosity comparable to or larger than their central stars. It is easy to show that in order to produce a spectrum with $n < 4/3$, it is necessary to have a temperature distribution that falls off less steeply than $T \propto r^{-3/4}$. In particular, a spectrum with $n = 0$ will result if $T \propto r^{-1/2}$ (Elmegreen 1982b). If such a temperature distribution were produced in a viscously evolving axisymmetric system, the rotation curve would have to be flat (like in a spiral galaxy) rather than Keplerian, which would imply that the mass distributions are extended [$M(r) \propto r$].

However, if the relation $M(r) \propto r$ extends from the inner disk ($\sim$a few stellar radii) out to the outer regions of the disk (~ 100 AU), the resulting disk would be very massive (R. Thompson, private communication). Using this method to explain extreme cases such as T Tau or HL Tau would require disk masses in excess of $10^3\ M_\odot$. Such models can be ruled out on the basis of observations at millimeter wavelengths, where the thermal emission from dust grains and rare isotopes of CO is optically thin (e.g. Sargent & Beckwith 1987), and on the basis of theoretical considerations of disk stability (e.g. Lin & Bertin 1985). Indeed, the latter consideration suggests that if circumstellar disks become roughly as massive as their central stars, further growth of the disk would be limited by the onset of nonaxisymmetric gravitational instabilities (bars or spiral waves) that would redistribute the material in the disk into a more centrally condensed configuration. In other words, nonaxisymmetric processes in such systems may drive disk accretion and generate waves in the inner regions of the disk that propagate and dissipate their energies at outer locations, thereby producing a less steeply declining equilibrium temperature gradient. The *nonlocal* transport implied by such mechanisms (whose forced analogues have been found in Saturn's rings; see Borderies et al. 1985, Shu et al. 1985) would distinguish them from true viscous transport. This should be contrasted with axisymmetric instabilities that do not propagate and may therefore be amenable to a pseudoviscous formalism (Lin & Pringle 1987). A more fundamental approach to the process is needed to obtain theoretical progress on this issue. In the future, spatial maps of very high resolution may show whether the disks around flat-spectrum T Tauri stars have interesting nonaxisymmetric structures.

In the meantime, the case for moderately massive disks can be stated more robustly. The objects in question must often have intrinsic disk luminosities of several $L_\odot$ in order to explain the observed infrared excesses. Accretion onto a star (either the central source or a companion embedded in the disk) is the most likely energy source and hence must occur at a rate of 10^{-5} to $10^{-6}\ M_\odot\ \mathrm{yr}^{-1}$. Since many T Tauri stars have such infrared excesses, and their lifetimes are on the order to 10^5 to 10^6 yr, the reservoirs of material for accretion in the disks must often be 0.1–1.0 $M_\odot$, a value that compares well with Sargent & Beckwith's (1987) estimate for the disk of HL Tau.

What about the lower estimates for the mass of HL Tau's disk? In our opinion, the absorption and scattering measurements (Cohen 1983, Beckwith et al. 1984, Grasdalen et al. 1984) refer to nonplanar material that is still undergoing *residual infall* onto this very young T Tauri star. The amount of gas and dust involved in such residual infall is always very small, and it produces a characteristic signature—a hump of far-infrared radiation—that is, in fact, seen in the spectral energy distribution of HL Tau. In contrast, the 2.7-mm continuum emission measurement of

Beckwith et al. (1986) detects emission from *all* warm dust. Most of this warm dust lies in a spatially flat disk, and it can contain a large fraction of a solar mass without coming into any conflict with the determinations of absorption or scattering optical depths, provided we are *not* viewing the disk edge-on.

6.6 *Implications for Binary Star and Planetary System Formation*

Given the above discussion, it is interesting to speculate further that active disks are active precisely because they are massive, and that passive disks are passive precisely because they lack the gravitational instability mechanisms that would give them appreciable accretion rates (Shu et al. 1987). If this speculation has any foundation, then there may be a natural basis for two modes of companion formation in T Tauri disks. If a nebular disk turns out to have enough mass to make another star, it may be an excellent candidate for forming a relatively close binary system (with a period shorter than $\sim 10^2$–10^3 yr; see Abt 1983). A passive or nearly passive disk will contain a smaller amount of material, and it may be a likely candidate for forming a planetary system [see Wetherill (1980) and Pollack (1984) for reviews].

7. CONCLUSION

7.1 *Summary: Four Stages of Star Formation*

We summarize this review with a proposed outline of the various stages involved in the birth of stars from molecular clouds; this includes most of what we know theoretically and observationally about the process. We begin with a molecular cloud, after its basic constituent material has been assembled somehow and somewhere in a galaxy. The first stage of star formation is then the formation of slowly rotating cloud cores (see Figure 7*a*). In subcritical clumps this occurs through the slow leakage of magnetic (and turbulent) support by ambipolar diffusion; in supercritical clumps, the clump may also fragment as it contracts and flattens as a whole. As long as the core formation process occurs relatively slowly (i.e. with a time scale that is longer than the characteristic signal crossing time), the cores probably asymptotically approach centrally concentrated states that resemble singular isothermal spheres. However, such end states cannot actually be reached because they are unstable.

The second phase begins when a condensing cloud core passes the brink of instability and collapses dynamically from "inside-out" (see Figure 7*b*). This evolutionary phase is characterized by a central protostar and disk, deeply embedded within an infalling envelope of dust and gas. The infalling

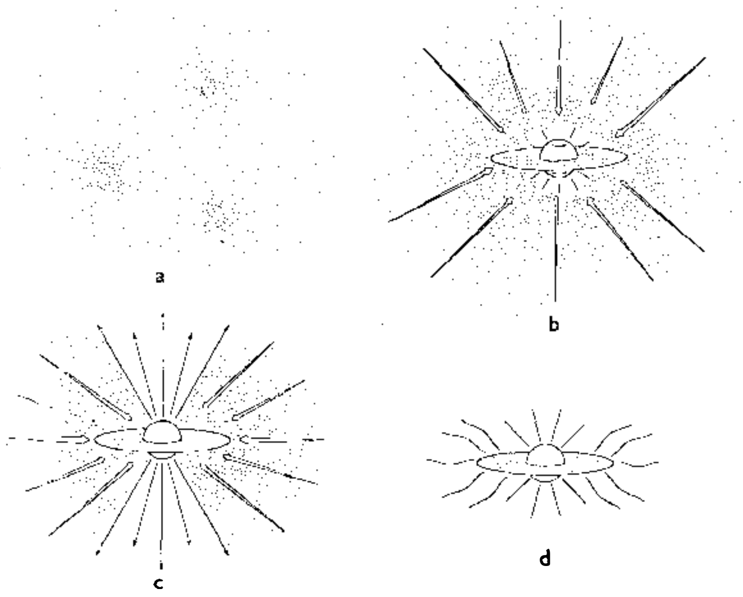

Figure 7 The four stages of star formation. (*a*) Cores form within molecular clouds as magnetic and turbulent support is lost through ambipolar diffusion. (*b*) A protostar with a surrounding nebular disk forms at the center of a cloud core collapsing from inside-out. (*c*) A stellar wind breaks out along the rotational axis of the system, creating a bipolar flow. (*d*) The infall terminates, revealing a newly formed star with a circumstellar disk.

material passes through an accretion shock as it falls onto the central star and disk, which, along with accretion within the disk, produces the main contribution to the luminosity for low-mass protostars. The emergent spectral energy distributions of theoretical models in the infall stage are in close agreement with those of recently found infrared sources with negatively steep spectra in the near- and mid-infrared. Protostars of high mass, in a pure accretion phase, have yet to be found, although the source near the water masers in W3(OH) is probably close to being such an object.

As a protostar accretes matter, deuterium will eventually ignite in the central regions and drive the star nearly completely convective if its mass is less than about 2 $M_\odot$. If the convection and the differential rotation of the star combine to produce a dynamo, the star can naturally evolve toward a state with a stellar wind. However, at first the ram pressure from material falling directly onto the stellar surface suppresses breakout. Gradually, the "lid" of direct infall will weaken as the incoming material falls preferentially onto the disk rather than onto the star. The stellar wind then rushes through the channels of weakest resistance (the rotational poles), which leads to collimated jets and bipolar outflows. Thus, the protostar enters the next step of evolution—the bipolar outflow phase (see Figure 7*c*). Examples of massive YSOs in this phase of stellar evolution (combined inflow and outflow) include the Becklin-Neugebauer object in Orion, but the trigger for the stellar wind in such cases is still theoretically obscure.

As time proceeds, more and more of the rotating inflowing matter will fall preferentially onto the disk rather than onto the star. In any reasonable picture of stellar outflows, the opening angle of the wind will widen with time, eventually sweeping outward over all 4π steradians and revealing the fourth evolutionary stage—in the case of a low-mass YSO—a T Tauri star with a surrounding remnant nebular disk (see Figure 7*d*). Radiation from the disk adds an infrared excess to the expected spectral energy distribution of the revealed source. The detailed shape of this infrared excess depends on whether the disk is largely passive and merely reprocessing stellar photons, or whether it is relatively massive and actively accreting. Both extremes of spectral shapes are observed in T Tauri stars; the amount of circumstellar material in the form of disks around nearly formed stars may be related to the dual issues of the origins of binary star and planetary systems. Disks around high-mass YSOs may also be common; in at least one well-documented case, S106, the disk is very large spatially but not extremely massive (i.e. not more massive than the central star).

A fifth phase of evolution can also be considered—namely, the final disappearance of the nebular disk as matter becomes incorporated into planets or stellar companions or is dispersed by the energetic outflow. For low-mass YSOs, such objects are probably the "naked T Tauri stars" (Mundt et al. 1983, Walter 1986, 1987). Their story more properly belongs to the subject of pre-main-sequence evolution than to star formation, although it can also be viewed as logically flowing from the developments of the four prior stages. High-mass stars apparently join the main sequence while still in the accretion phase; thus, they have no "pre-main-sequence story" to tell.

7.2 *Outstanding Problems*

There has been considerable theoretical and observational progress in recent years in our understanding of how individual stars form and evolve. In particular, the discoveries of the low general efficiency of star formation in molecular clouds and of the ubiquity of outflows from YSOs have challenged older conceptions of the ultimate origin for the masses of forming stars. Still, many gaps remain in our knowledge of how stars are born. We have only relatively rough ideas of the properties of the final system that is produced from the collapse of a cloud core with given initial conditions. Much of the uncertainty arises from ignorance about basic disk transport processes; new ideas may not be needed so much as rigorous calculations. Observations carried out with high angular resolution at millimeter, submillimeter, and infrared wavelengths, in the continuum and in spectral lines, could be of crucial help in unraveling the puzzle.

In contrast, there are few viable ideas concerning the generation of stellar jets and bipolar flows that are amenable to rigorous theoretical treatment. Quantitative work has focused, for good reason, on the easier problem of jet propagation and wind interaction with the surrounding medium. Steady progress on the latter front can be expected as the models being developed become more quantitative and can be compared fruitfully with the wealth of existing observational data, but fundamental advances on the former issue may lie in the distant future (especially if the generation of the winds requires a deep understanding of the basic processes underlying stellar activity).

Computations can be done (and are in the works) to follow in detail how a molecular cloud core condenses from a parent clump of given properties and evolves to the point of gravitational collapse. Indeed, if magnetic fields and ambipolar diffusion control the process, then in some sense theorists may be ahead of observers on this problem. The current empirical data base contains very little secure information concerning magnetic field strengths and the ionization fractions in molecular cloud cores. Resolved velocity maps at scales less than 10^{16} cm, in both neutral and ionized molecular species, would also be very useful to kinematically confirm or reject the basic concepts of the first three phases of star formation discussed in Section 7.1.

Many of the detailed comparisons between theory and observation in recent years have focused on the problem of low-mass star formation. The rcasons frequently cited are that the developed theory for low-mass stars is more complete, and that high-mass stars, once formed, do so much damage to their surroundings that it becomes difficult to sift through the remaining clues to reconstruct prior events. Yet in order to progress to the point where the initial-mass functions can be predicted (e.g. for use in models of galactic evolution), we must have an a priori theory for how high-mass stars are formed. A viable, if conservative, approach may be a greater concerted effort, both observational and theoretical, on the somewhat neglected problem of the formation of intermediate-mass stars.

Another fruitful area for future research will be the problem of the formation of star clusters. Largely as a result of the reviewers' own prejudices and expertise, this article has concentrated on the problem of star formation in clouds that are dominated by the presence of magnetic fields.

However, as discussed in Sections 3.1 and 3.2, there is a second regime logically and probably empirically as well—the supercritical regime—in which cloud evolution is gravitationally dominated. In the extreme case of a very supercritical cloud, the behavior should resemble those calculations in which the magnetic field is neglected altogether. A rich picture of fragmentation and collapse is characteristic of such calculations, but to date little success has been achieved when the theoretical results are compared with observations. This lack of agreement, we believe, arises because of a misplaced sense of where the results should apply. Ignoring magnetic fields altogether is probably not valid for most present-day molecular clouds; however, it may be an acceptable first approximation for proto-clusters (especially proto-globular clusters). In a similar vein, the concept of induced star formation gains credibility in environments that are essentially free of magnetic fields or are decoupled from them.

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