THE COLLAPSE OF CLOUDS AND THE FORMATION AND EVOLUTION OF STARS AND DISKS

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We consider the interrelationships among the structure of molecular clouds; the collapse of rotating cloud cores; the formation of stars and disks; the origin of molecular outflows, protostellar winds, and highly collimated jets; the birth of planetary and binary systems; and the dynamics of star/disk/satellite interactions. Our discussion interweaves theory with the results of observations that span from millimeter wavelengths to X-rays.

I. OVERVIEW

This chapter gives a status report on the current astrophysical problems that confront the theory and observation of the collapse of clouds and the formation and evolution of stars and disks. The assignment contains too vast a topic to review in any detail in the allotted pages, and we have interpreted our task accordingly as setting the context for some of the following chapters by others. Parts of this review in essence can also be found in Shu (1991).

We start in Sec. II with the notion that two different modes seem to account for the birth of most stars in the Galaxy: a "closely packed" mode characterized by the more or less simultaneous formation of a tight group of many stars from large dense clumps of molecular gas and dust (see Chapters by Blitz and by Lada et al.), and a "loosely-aggregated" mode in which an unbound association of individual systems (some of which may be binaries) forms sporadically from well-separated, small, dense, cloud cores embedded within a more rarefied common envelope. We introduce the working hypothesis, adopted currently by many workers in the field, that the formation of sunlike stars by the second mode occurs in nearby dark clouds like the Taurus region in four conceptually distinct stages (Fig. 1).

Of the four stages *a–d* outlined in Fig. 1, the most surprising—the one totally unanticipated by prior theoretical developments—is *c*, the *bipolar outflow phase*. In Sec. III, we review the observational discoveries and

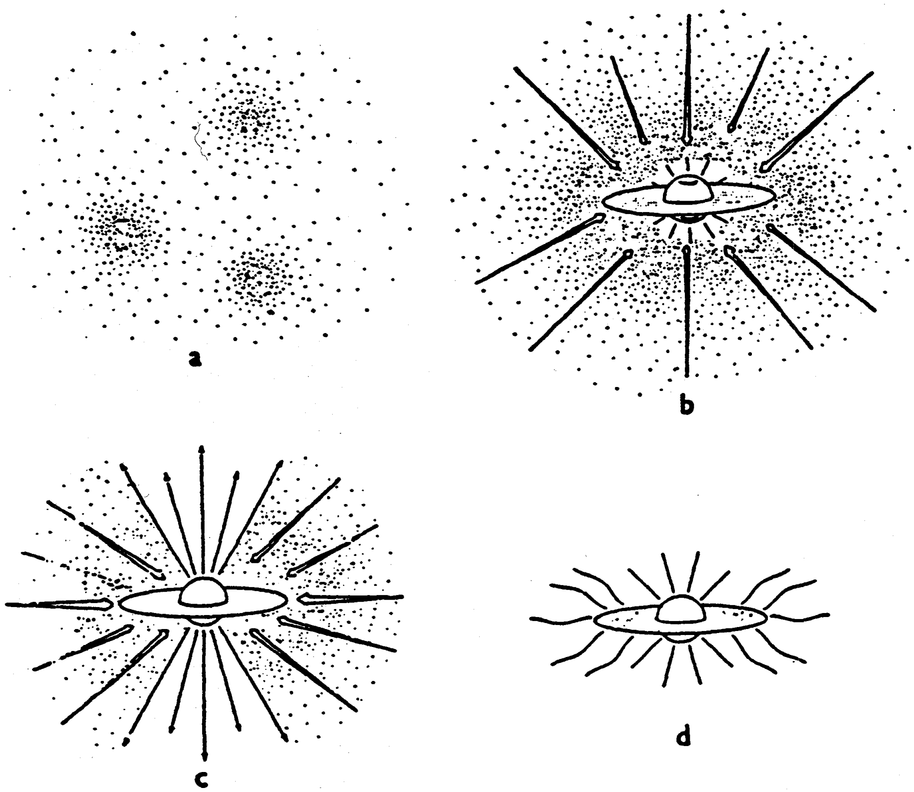

Figure 1. The four stages of star formation. (a) Cores form within molecular cloud envelopes as magnetic and turbulent support is lost through ambipolar diffusion. (b) Protostar with a surrounding nebular disk forms at the center of a cloud core collapsing from inside-out. (c) A stellar wind breaks out along the rotational axis of the system, creating a bipolar flow. (d) The infall terminates, revealing a newly formed star with a circumstellar disk (figure from Shu et al. 1987*a*).

interpretations that have led up to our current empirical understanding of this fascinating phenomenon. We pay particular attention to those aspects that set tight constraints on possible theories. In Sec. IV, we begin a theoretical discussion that indicates why, in retrospect, we should perhaps have known all along that the formation of stars would necessarily involve the heavy *loss* of mass. Given the angular momentum difficulty likely to be faced by any object which forms by contraction through many orders of magnitude from an initially extended, rotating state, the birth of stars (and, perhaps, galactic nuclei) through the accretion of matter from the surroundings may simply not be possible without the *simultaneous* accompaniment of powerful outflows.

Under the working hypothesis, gas and dust in the first stage (Fig. 1*a*) slowly contract under their own gravitation against the frictional support provided by a background of ions and magnetic fields via the process of ambipolar diffusion (see also Chapters by Heiles et al. and McKee et al.). The principal feature in this stage involves a *quasi-static evolution toward a* $1/r^2$ *density configuration* appropriate for a singular isothermal sphere. The relative statistics of cores with and without embedded infrared sources, as well as the theory of ambipolar diffusion, suggests that the time scale over which the configuration is observable as a quiescent ammonia core before it enters dynamical collapse (i.e., before it contains an embedded infrared source) should span about 10^6 yr.

When the contracting configuration becomes sufficiently centrally concentrated, it enters stage 2 (Sec. V; cf. Chapter by Tscharnuter and Boss), wherein the cloud core gravitationally collapses from *inside out*. In such a situation, the inner regions form an accreting but otherwise secularly evolving protostar plus nebular disk. An infalling envelope of gas and dust that rains down from the overlying (slowly rotating) molecular cloud core covers the growing star plus disk. The visual extinction to the central star measures from several tens to a thousand or more, so the embedded source during this stage does not appear as an optically visible object, but must be studied principally by means of the infrared, submillimeter, and millimeter radiation produced by dust reprocessing in the surrounding envelope (see Chapter by Zinnecker et al.).

At some point during this phase of the evolution (see Sec. VI), a powerful wind breaks out along the rotational poles of the system, reversing the infall and sweeping up the material over the poles into two outwardly expanding shells of gas and dust. This stage (stage c) corresponds to the *bipolar outflow* phase observed spectroscopically at radio wavelengths by CO observers (cf. Chapter by Fukui et al.). Theory suggests that this stage features *combined* inflow (in the equatorial regions) and outflow (over the poles). Current consensus in the field holds that magnetohydrodynamic forces drive the wind sweeping up the molecular outflow. The main debate concerns whether the wind originates from the *star*, or from the *disk*, or from their *interface* (see Chapter by Königl and Ruden). From the mass infall rate as well as the statistics of numbers of embedded sources compared to revealed ones (T Tauri stars), we can estimate the combined time spent in stages *b* and *c* (when the system still gains mass in net) as roughly 10^5 yr, almost independent of mass.

As time proceeds, we envisage the angle occupied by the outflow to open up (like an umbrella) from the rotation axes and to spread, halting even the rain of infalling matter over the equator. At this point, stage d, the system becomes visible, even at ultraviolet, optical, and near-infrared wavelengths as a star plus disk to all outside observers (see Sec. VII; Chapter by Basri and Bertout). The location of the star in the Hertzsprung-Russell diagram yields a constraint on the accretion time scale (see Chapter by Stahler and Walter), and the numbers estimated by this method agree well with the dynamical calculations based on the concept of inside-out collapse from a molecular cloud core modeled as a (rotating) singular isothermal sphere. The principal questions during this stage then concern the mechanisms by which mass, angular momentum, and energy transport take place within the disk (see Chapter by Adams and Lin), and the nature of any companions (stellar or planetary) that may condense from such a disk (see Sec. VIII; Chapters by Bodenheimer et al., Weidenschilling and Cuzzi, Lissauer and Stewart). Intimately tied to these issues are the estimates of disk masses that can be obtained from observational measurements at submillimeter and millimeter wavelengths (see Chapter by Beckwith and Sargent). Another important question concerns whether disk accretion mainly occurs episodically during FU Orionis outbursts or during the relatively quiescent (normal) states as well (see Chapter by Hartmann et al.). FU Orionis outbursts may have important connections with the meteoritic evidence for transiently high nebular temperatures (see Chapter by Palme and Boynton); however, the empirical need to appeal to such outbursts for geochemical peculiarities becomes less clear if nebular lightning truly offers a solution for the origin of chondrules (see Chapter by Morfill et al.).

For the purposes of this book, we may append a fifth stage to the above four: an epoch of disk clearing. The chapter by Strom et al. identifies this phase with planet making and sets disk lifetimes at 10^6 to 10^7 yr. We can usefully compare this astrophysical constraint with the models of giant planet formation that do not begin to accrete large amounts of gas until a critical core mass of solids has been reached (see Chapter by Podolak et al.). Edwards et al. in their chapter argue persuasively that the presence of an inner disk constitutes the crucial distinction between those T Tauri systems that drive extraordinary winds (classical T Tauri stars) and those that do not (weak-line or "naked" T Tauri stars). The absence of an absorbing wind (and inner disk) makes the magnetic activity on the surfaces of the latter objects observable as X-ray sources (cf. Chapter by Montmerle et al.). Spallation reactions driven by energetic particle fluxes associated with the enhanced flaring activity observed for these relatively gas-poor systems may help to explain some of the isotopic anomalies seen in the meteoritic record. Surviving dust grains from the protosolar core may have brought other anomalies intact into the solar nebula (cf. Chapters by Cameron, Swindle, and Ott).

In Sec. IX, we recapitulate by discussing the question of the origin of stellar and planetary masses. In particular, we emphasize the emerging view that in environments where there exist more than enough material to form the final condensed objects, stars and giant planets help, in part, to determine their own masses: stars, by blowing powerful winds that shut off the continuing infall from a molecular cloud; and giant planets, by opening up gaps that turn off the continuing accumulation from a nebular disk (see also Chapter by Lin and Papaloizou). A worrisome issue in the latter context concerns how the solar system managed to clear the debris left over in the disk from making the Sun and the planets. The Chapter by Duncan and Quinn makes clear that within the region of the giant planets the solid debris could have been cleared by gravitational interactions with the planetary bodies. The observations of remnant particulate disks around main-sequence stars (see Chapter by Backman and Paresce) demonstrate empirically, however, that planetesimal dispersal mechanisms become inefficient at large distances from the central star; such regions probably become the reservoir for (an inner belt of) comets (see Chapter by Mumma et al.). The gas appears to pose a greater difficulty. The old view envisages nebular gas left over from planet building to be dispersed by a T Tauri wind. The new view suggests that T Tauri winds do not exist in the absence of (the inner portions of) nebular disks; thus, Edwards et al. argue that the former cannot get rid of the latter (however, see Sec. IX).

The above summary gives a sample of the rich diversity of physical processes that confronts the research worker in the field of star and planet formation. This variety arises for a simple and basic reason: the problem spans physical conditions ranging from the depths of interstellar space to the interiors of stars and planets, involving all the known states of matter and forces of nature, with observational diagnostics available across practically the entire electromagnetic spectrum, and experimental access to relevant primitive materials that has no parallel in any other branch of astronomy and astrophysics. To attack its problems, the field has developed a broad range of technical tools—observational, theoretical, and experimental; and the full exploitation of this range of tools during the past decade has led to rapid and impressive progress.

II. BIMODAL STAR FORMATION

The empirical notion that the birth of low- and high-mass stars may involve separate mechanisms has a long and controversial history (Herbig 1962*b*; Mezger and Smith 1977; Elmegreen and Lada 1977; Gusten and Mezger 1982; Larson 1986; Scalo 1986; Walter and Boyd 1991). The name "bimodal star formation" usually attaches to this concept, but, more recently, the emphasis has changed from "low mass versus high mass" to "loosely aggregated versus closely packed" (see Chapter by Lada et al.). The latter phrasing roughly coincides with the older one of "associations versus clusters," except that it need not carry the connotation of "gravitationally unbound versus gravitationally bound."

The theoretical distinction between "loosely aggregated" and "closely packed" refers to whether gravitational collapse occurs independently for individual small cores to form single stars (or binaries); or whether it involves a large piece of a giant molecular cloud to produce a tight group of stars created more or less simultaneously. This tight group need not form a bound cluster if the winds or other violent events that accompany the formation of stars expel a large fraction of the gas not directly incorporated into stars (Lada et al. 1984; Elmegreen and Clemens 1985).

The occurrence of two separate modes of star formation has a natural theoretical explanation (Mestel 1985; Shu et al. 1987*b*) if we adopt the point of view that magnetic fields provide the primary agent of support of molecular clouds against their self-gravity (Mestel 1965*a*; Mouschovias 1976*b*; Nakano 1979; see also Chapters by Heiles et al. and McKee et al.). The inclusion of a conserved magnetic flux Φ threaded by an electrically conducting cloud introduces a natural mass scale (cf. Mouschovias and Spitzer 1976; Tomisaka et al. 1988*a*,1989),

$$M_\Phi \equiv 0.13\, G^{-1/2}\Phi \tag{1}$$

that is analogous to Chandrasekhar's limit $M_{\rm Ch}$ in the theory of white dwarfs. A critical mass arises whenever we try to balance Newtonian self-gravity by the internal pressure of a fluid which varies as the 4/3 power of the density (because of ultrarelativistic electron degeneracy pressure in the case of a white dwarf; because of the pressure of a frozen-in magnetic field in the case of a molecular cloud). White dwarfs with masses $M > M_{\rm Ch}$ cannot be held up by electron degeneracy pressure alone, but must suffer overall collapse to neutron stars or black holes. Molecular clouds with masses $M > M_\Phi$ (supercritical case) cannot be held up by magnetic fields alone (even if perfectly frozen into the matter), but, in the absence of other substantial means of support, must collapse as a whole to form a closely packed group of stars. The only trick in this case concerns how to get a supercritical cloud or clump from an initially subcritical assemblage (or they would have all collapsed by now). Theory and observation both suggest a natural evolutionary course: the agglomeration, with an increase of the mass-to-flux ratio, of the discrete cloud clumps that comprise giant molecular complexes (Blitz and Shu 1980; Blitz 1987*b*,1990; Shu 1987; see also Chapters by Blitz and by Lada et al.).

The analogy between white dwarfs and molecular clumps breaks down for the subcritical case. White dwarfs with $M < M_{\rm Ch}$ can last forever as uncollapsed degenerate objects (if nucleons and electrons are stable forms of matter) because quantum principles never weaken. Magnetic clouds with $M < M_\Phi$ initially cannot last forever because magnetic fields do weaken, and the local loss of magnetic flux from unrelated dense regions allows loosely aggregated star formation. In particular, in a lightly ionized gas, the fields can leak out of the neutral fraction by the process of *ambipolar diffusion* (Mestel and Spitzer 1956; Nakano 1979; Mouschovias 1978; Shu 1983). In Sec. IV, we shall examine in more detail the production of individual molecular cloud cores by this process. For the present, we merely note that theory predicts that self-gravitating cores sustained by magnetic fields will inevitably evolve to a state of spontaneous gravitational collapse. This prediction seemingly flies in the face of observations, which almost always find young stellar objects associated with outflows rather than inflows.

III. THE BIPOLAR OUTFLOW PHASE: OBSERVATIONS

At the heart of the central paradox concerning star formation lies the problem of *bipolar outflows* (for reviews, see Lada 1985; Welch et al. 1985; Bally 1987; Snell 1987; Fukui 1989; Rodriguez 1990). Astronomical visionaries (see, e.g., the reminiscences of Ambartsumian 1980) have long worried that forming astronomical systems frequently exhibit *expansion*, rather than the *contraction* that would be naively predicted by gravitational theories. Astronomical conservatives have long persisted in ignoring such warnings, doggedly pursuing the theoretical holy grail that bound objects, such as stars and planetary systems, should form from more rarefied precursors by a process of gravitational contraction. The reconciliation of this fundamental dichotomy remains the central challenge facing theorists and observers alike (see the review by Lada and Shu 1990).

In 1979, Cudworth and Herbig discovered that two H-H objects in a nearby dark cloud L1551 (Lynds 1962) exhibit very high proper motions (corresponding to $\sim$150 km s^{-1}) that trace back to an apparent point of origin near the location of an embedded infrared source IRS 5 found by Strom et al. (1976). Knapp et al. (1976) had earlier found that CO millimeter-wave spectra in this region possess linewidths of the order 10 km s^{-1}, too large according to Strom et al. to correspond to gravitational collapse. The latter authors suggested instead that the disturbance in the ambient molecular cloud material might be produced by a powerful outflow from IRS 5. Snell et al. (1980) mapped the CO emission in the L1551 region and verified this conjecture in a dramatic fashion. They found that the high-velocity CO surrounds the tracks of the fast H-H objects, taking the form of two lobes of gas moving in diametrically opposed directions from IRS 5 (see Fig. 2). Putting together all of the empirical clues, Snell et al. made the prescient proposal (Fig. 3) that a stellar wind must blow at 100 to 200 km s^{-1} from IRS 5, in directions parallel and anti-parallel to the rotation axis of a surrounding accretion disk, and that this collimated wind sweeps up ambient molecular cloud material into two thin shells, which manifest themselves as the observed bipolar lobes of CO emission. The thin-shell nature of the CO lobes of L1551 has since received empirical validation in the investigations of Snell and Schloerb (1985) and Moriarty-Schieven and Snell (1988). Except for the further interpretation that IRS 5 represents a *protostar*, which, apart from suffering mass loss in the polar directions, accretes matter from *infall* occurring in the equatorial regions (see panel *c* in Fig. 1), the proposal of Snell et al. (1980) corresponds in every detail to the theoretical model that we shall pursue in Sec. VI.

Shortly after the original suggestion, however, events moved to a different interim conclusion. In common with the central source of several other bipolar outflows (see reviews of Cohen 1984; Schwartz 1983), IRS 5 lies along a chain of HH objects (Strom et al. 1974; Mundt and Fried 1983) that may just constitute the brighter spots of a more or less continuous optical jet (Mundt 1985; Reipurth 1989*c*).

The emission knots probably arise as a result of the interactions of a highly collimated ionized stellar wind with the surrounding medium (possibly another [neutral] wind; see Stocke et al. [1988] and Shu et al. [1988]). The free-free thermal emission from the ionized stellar wind resolves itself in VLA radio-continuum measurements as a two-sided jet centered on IRS 5 (Bieging et al. 1984). It was natural to suppose that this ionized wind represents the driver for the bipolar molecular outflow. Unfortunately, the momentum input provided by the ionized wind, integrated over the likely lifetime of the system, falls below the value required to explain the moving CO lobes by one to two

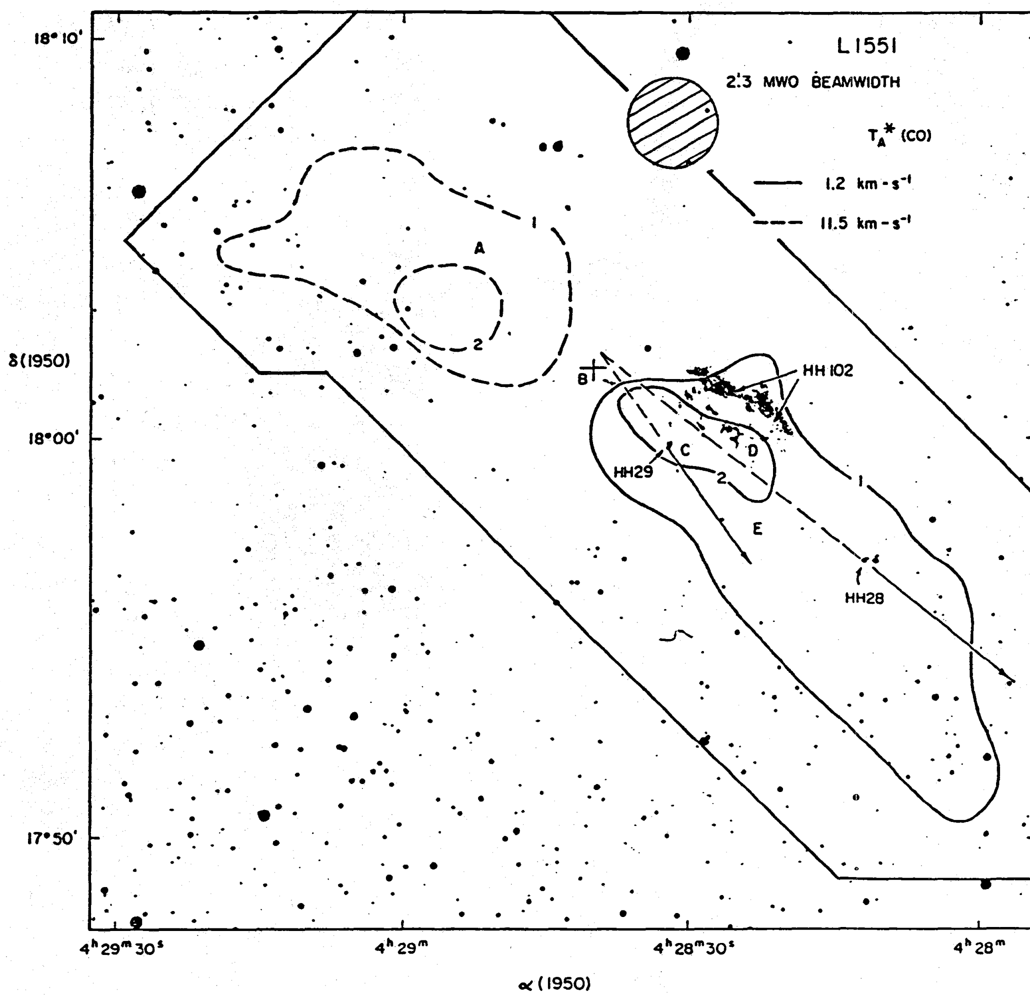

Figure 2. Contour map of the antenna temperature of the $J = 1 - 0$ transition of ^{12}CO at high velocities, superposed on an optical photograph of the L1551 dark cloud. The cross indicates the position of IRS 5; also shown are the directions of the proper motions of two Herbig-Haro objects, HH 28 and HH 29 (figure from Snell et al. 1980).

orders of magnitude (for the likely case of momentum-driven rather than energy-driven flows), a conclusion that holds as well for all other well-studied bipolar flow sources (Bally and Lada 1983; Levreault 1985). By comparing the required momentum input to the photon luminosity divided by the speed of light, the same authors argued persuasively against the possible importance of radiation pressure in the outflow dynamics.

The claim by Kaifu et al. (1984) of the detection of an extended, rotating, molecular disk encircling IRS 5, contributed to the puzzle. This result, and the inability of ionized stellar winds to drive the observed molecular flows, motivated Pudritz and Norman (1983,1986) and Uchida and Shibata (1985) to suggest that bipolar outflows originate, not as shells of molecular cloud gas

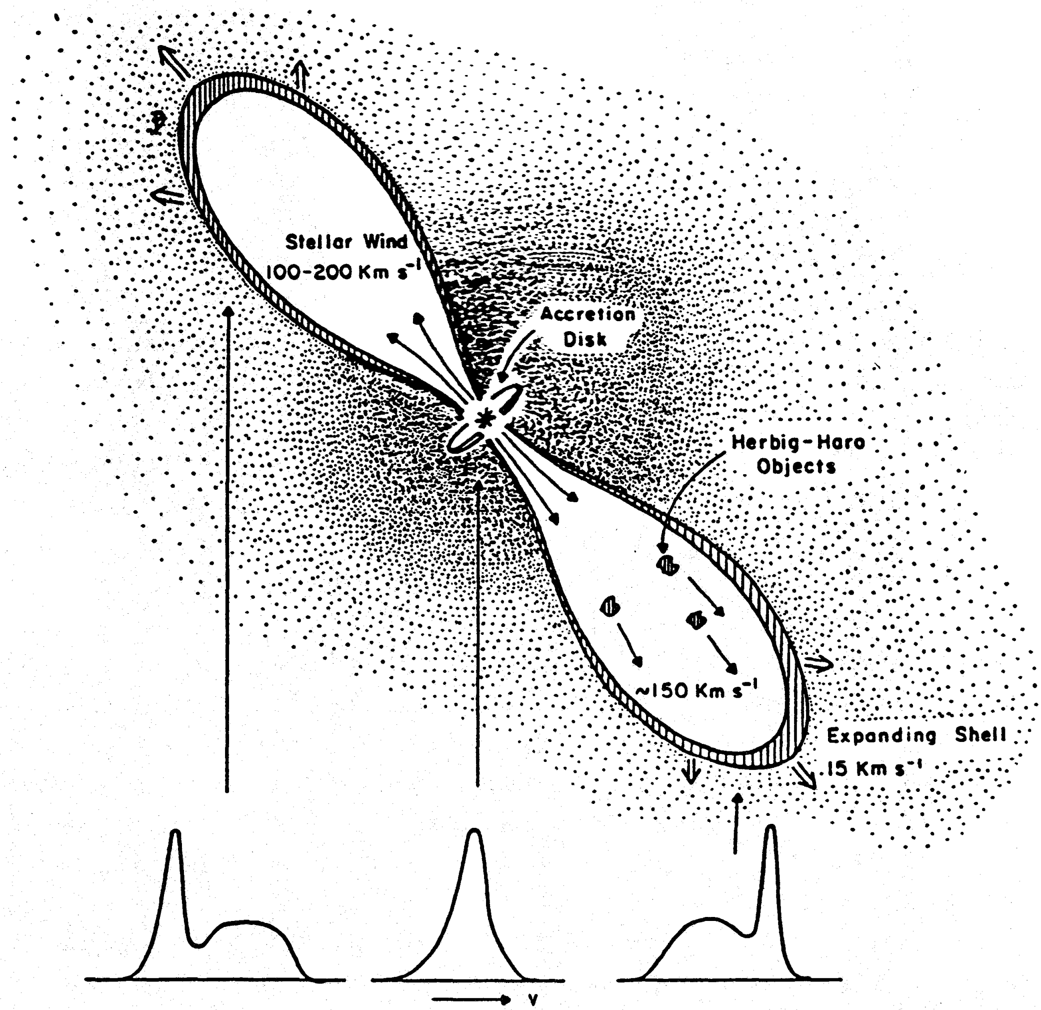

Figure 3. Schematic model for the bipolar flow in L1551 driven by a stellar wind emanating from IRS 5. At the bottom is depicted expected CO line profiles for different line of sights across the source (figure from Snell et al. 1980).

swept up by a stellar wind, but, following the work of Blandford (1976) and Lovelace (1976), as gaseous material driven magneto-centrifugally directly off large ($\sim 10^{17}$ cm) and massive ($\sim 10^2$ $M_\odot$) circumstellar disks. This suggestion seemingly received support from reports that one limb of the blueshifted lobe in L1551 exhibits rotation (Uchida et al. 1987*a*). However, subsequent observational studies have failed to confirm the presence of large and massive disks around young stellar objects with the properties required by the theoretical models (see, e.g., Batrla and Menten 1985; Menten and Walmsley 1985; Moriarty-Schieven and Snell 1988). Moreover, the claim of rotation in the blueshifted lobe may have been contaminated by confusion with another bipolar flow source in the same field of view (Moriarty-Schieven and Wannier 1991). Finally, from a theoretical point of view, the original large-disk models have severe energetic problems (Shu et al. 1987*a*; Pringle 1989*a*).

Recently, a number of different groups have attempted to salvage the disk model by postulating smaller and less massive disks, with the bulk of the wind emerging from radii much closer to the central star. This revised picture, however, suffers from the lack of a plausible injection mechanism. The paper by Blandford and Payne (1982) illuminates the basic difficulty.

Blandford and Payne consider a thin disk threaded by a poloidal magnetic field **B** in which rotation balances the radial component of gravity. If **B** has sufficient strength and a proper orientation, it can fling to infinity electrically conducting gas from the top and bottom surfaces of the disk. A freely sliding bead on a rigid wire anchored at one end to a point ϖ in a disk rotating at the angular velocity $\Omega_{\rm disk}(\varpi)$ provides a useful analogy (Henriksen and Rayburn 1971). In such an analogy, the termination of the wire at the free end corresponds to the Alfvén surface, beyond which the magnetic field can no longer enforce corotation, even approximately, and the super-Alfvénic fluid motion becomes essentially ballistic. For the flinging effect to take place in a Keplerian disk (one in which all of the gravitational attraction comes from a central mass point), the poloidal field must enter or leave the disk at an angle larger than 30° from the normal. (The component of the centrifugal force parallel to the field [wire] in the meridional plane has no excess compared to gravity if the field [wire] makes too small an angle with respect to the rotation axis, e.g., if it points vertically through the disk.) However, the magnetic field on the two sides of a disk cannot *both* bend outward by a nonzero angle without generating a large kink across the midplane of the disk, one that would result in a very large $\nabla \times \mathbf{B}$, and which would yield, by Ampere's law, a very large current (infinitely large in a disk of infinitesimal thickness). The Lorentz force per unit volume,

$$\mathbf{f}_{\rm L} = \frac{1}{4\pi}(\nabla \times \mathbf{B}) \times \mathbf{B} \tag{2}$$

needed to produce order-unity departures from Keplerian motion above the disk (in order to drive a wind) must then have even bigger values inside the disk. The tendency for the field lines to want to straighten vertically (thereby shutting off the magneto-centrifugal acceleration) can be offset only by a heavy radial inflow through the disk, maintained, for example, by ambipolar diffusion in the presence of sub-Keplerian rotation inside the disk (Konigl 1989). Sub-Keplerian rotation everywhere in the disk necessitates the continuous removal of angular momentum from it (the role of the wind), but the postulated deficit of centrifugal support in the disk then works against lift-off of the gas in the first place. Indeed, no one has yet succeeded in constructing a self-consistent cool model of this type, demonstrating a smooth connection from the region of the sonic transition (near the disk surface) to the region of the Alfvénic transition (far from the disk surface).

Blandford and Payne themselves adopt another strategy. They accept the conclusion that the symmetry of a disk geometry naturally forces any poloidal magnetic field to thread vertically through the midplane of the disk, from where it *gradually* bows outward (see Fig. 4). To lift the material off the disk

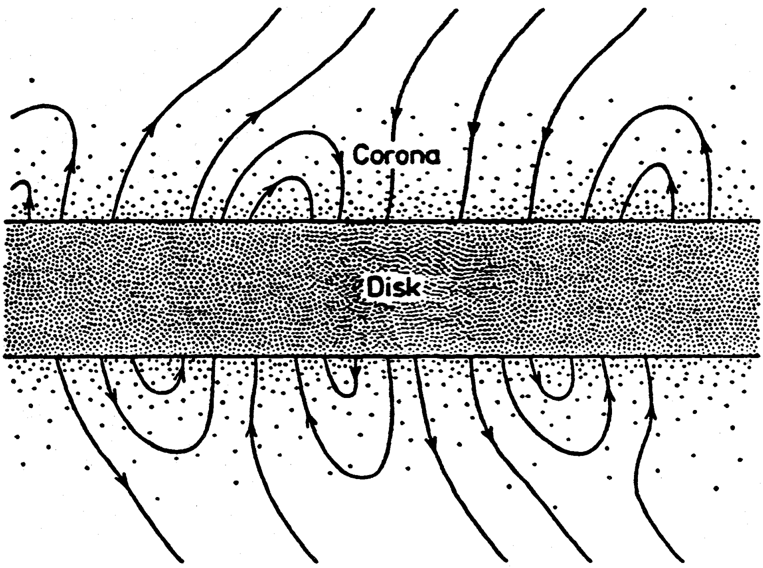

Figure 4. A schematic representation of a possible field geometry close to the disk (figure from Blandford and Payne 1982).

to heights where the magnetic field bends over sufficiently for the magneto-centrifugal mechanism to operate, they invoke the thermal pressure of a hot corona above (and below) the disk. In other words, their final model starts off as a *thermally driven wind*. A hot gas with thermal speeds comparable to the virial speeds characteristic of the depth of the gravitational potential in the inner parts of a protostellar disk, however, suffers tremendous radiative losses at the injection densities needed to supply the observed mass outflows. While astrophysical systems with highly luminous compact objects at their centers, such as active galactic nuclei or binary X-ray sources, might well provide sufficient energy to sustain analogous radiative losses, we cannot expect the same bounty from the central machine of protostellar systems (cf. the discussion of DeCampli [1981] that even the relatively low-powered T Tauri winds cannot be thermally driven).

In summary, at radii $\gtrsim 10^{17}$ cm, disks *can* maintain thermal speeds that compete with the local virial speeds, but such regions lack enough rotational energy to drive the observed flows. Greater centrifugal power exists at smaller radii, but lift-off from the disk then requires an unidentified nonthermal injection mechanism. Compounding these difficulties is the unlikelihood of strong magnetic fields in thin disks. Recall that the energy density of the field at the Alfvén surface must equal the kinetic energy density of a wind flowing close to its terminal velocity. But magnetic buoyancy effects probably limit the generation of fields by dynamo mechanisms acting within disks to energy densities comparable to or less than the prevailing thermal values (Stella and Rosner 1984). The inability of disks to retain large *toroidal* fields (Parker 1966; Matsumoto et al. 1988) limits the viability of mechanisms like those of Lovelace et al. (1991) that rely on the pressure gradient of such fields in the z-direction for the initial vertical acceleration of the gas.

Other theoretical arguments suggest that disk fields must be weak. Umebayashi and Nakano (1988) have computed the ionization state of a minimum-mass solar nebula. The large nebular densities encountered even in a minimum-mass model lead to very low ionization fractions, so low that magnetic fields would be safely decoupled by several orders of magnitude from the nebular gas throughout the disk, except for the innermost and outermost regions. Even in disks where Ohmic dissipation and ambipolar diffusion do not present severe obstacles, difficulties exist in supposing that magnetic fields acquire strengths sufficient to drive the observed outflows. For example, if accretion disks owe their viscosity to the powerful MHD instability recently discussed by Balbus and Hawley (1991) and Hawley and Balbus (1991), the *poloidal* component of the field cannot have energy densities much higher than thermal gas values. Indeed, the instability itself may serve as one of the ingredients of a self-consistent dynamo mechanism responsible for generating the poloidal component of the disk field in the first place. In addition, dynamically strong fields (energy densities comparable to that contained in the *rotation* of the part of the disk supplying the matter for the outflow) cannot have been dragged in and amplified by lightly ionized gas that collapsed from interstellar dimensions by ambipolar diffusion (see Sec. IV below), because mass infall at rates and speeds (free-fall) comparable to the inferred outflow values had to overcome the interstellar field to form the disk in the first place. Finally, even if we could surmount the above problems and drive a disk wind, we would *still* be left with an angular-momentum problem for the accreting central star (cf. the discussion in Sec. V). For all of these reasons, we believe that Snell et al. (1980) intrinsically had the right idea when they speculated that *stellar winds* must constitute the basic driver for bipolar outflows in young stellar objects.

The issue then reverts to the original question: how can stellar winds drive the observed molecular outflows when the *ionized* component lacks the requisite power? Many researchers arrived independently at the obvious answer: perhaps young stellar objects possess *neutral* winds that supply the missing momentum input. Detection of fast neutral winds would require the spectroscopic observation of broad and shallow radio emission lines, a difficult task because so little mass resides within the telescope beam at any particular time. Nevertheless, the search for *atomic* hydrogen winds has now succeeded in at least three sources: HH 7–11, L1551, and T Tau (Lizano et al. 1988; Giovanardi et al., in preparation; Ruiz et al. 1991). Fast neutral winds are also indicated through the detection of CO moving at velocities of ± 100 to 200 km s^{-1} in HH 7–11 and several other sources (Koo 1989; Margulis and Snell 1989; Masson et al. 1990; Bachiller and Cernicharo 1990; J. E. Carlstrom, personal communication). The mass-loss rate measured in HH 7–11 is especially impressive, $\dot{M}_w \approx 3 \times 10^{-6}$ $M_\odot$ yr^{-1}, more than sufficient to drive the known bipolar molecular outflow associated with this source. Thus, the discovery of massive neutral winds from (low-mass) protostars solves one of the principal remaining mysteries concerning bipolar outflows, namely, the ultimate source for their power. However, this empirical discovery pushes to the fore another puzzling question: how and why are stars born by *losing* mass at such tremendous rates?

IV. ROTATING, MAGNETIZED, MOLECULAR CLOUD CORES

A major breakthrough occurred in the field when Myers and Benson (1983) identified the sites for the birth of individual sunlike stars in the Taurus cloud as small dense cores of dust and molecular gas that appear as especially dark regions in the Palomar sky survey or as regions of NH_3 and CS emission (cf. the references in Evans 1991). One theory for the formation of such cores begins with the view that large molecular clouds are supported by a combination of magnetic fields and turbulent motions (see Shu et al. 1987*a*; Myers and Goodman 1988*a*, *b*). Because magnetic fields directly affect only the motions of the charged particles, magnetic support of the neutral component against its self-gravity arises only through the friction generated when neutrals slip relative to the ions and field via the process of ambipolar diffusion (Mestel and Spitzer 1956). This slippage occurs at enhanced rates in localized dense pockets where the ionization fraction is especially low. As the field diffuses out relative to the neutral gas, the Alfvén velocity $B/(4\pi\rho)^{1/2}$ drops, and the turbulence must also decay if fluctuating fluid motions are to remain sub-Alfvénic. Thus, both processes—field slippage relative to neutrals and the lowering of turbulent line widths in regions of increasing mass-to-flux—reinforce the tendency for quiet dense cores to separate from a more diffuse common envelope.

Assuming axial symmetry (no φ dependence) in spherical polar coordinates (r, θ, φ) as well as reflection symmetry about the midplane $\theta = \pi/2$, Lizano and Shu (1989) performed detailed ambipolar diffusion computations, using the ionization balance calculations of Elmegreen (1979*b*) and including empirically the effects of turbulence with the scaling laws found by radio observers (cf. Chapters by Blitz, by Lada et al. and by McKee at al.). They find that the time to form NH_3 cores, $\sim 10^6$ to 10^7 yr, and the additional time to proceed to gravitational collapse, a few times 10^5 yr, agree roughly with the spread in ages of T Tauri stars in the Taurus dark cloud (Cohen and Kuhi 1979) and with the statistics of cores with and without embedded infrared sources (Fuller and Myers 1987). Toward the end of its quiescent life, a molecular cloud core tends to acquire the density configuration of a (modified) singular isothermal sphere:

$$\rho(r, \theta) = \frac{a_{\text{eff}}^2}{2\pi G r^2} Q(\theta) \tag{3}$$

where a_{eff} is the effective isothermal sound speed, including the effects of a quasi-static magnetic field $\mathbf{B}_0$ and (as modeled) an isotropic turbulent pressure. The dimensionless function $Q(\theta)$, >1 toward the magnetic equator and <1 toward the poles, yields the flattening that occurs because the magnetic field contributes no support in the direction along $\mathbf{B}_0$. We define $Q(\theta)$ so that it has a value equal to unity when averaged over all solid angles, i.e.,

$$\int_0^{\pi/2} Q(\theta) \sin\theta \, d\theta = 1 \tag{4}$$

hence, $a_{\rm eff}^2$ in Eq. (3) gives the angle-averaged contributions of thermal, turbulent and magnetic support.

Myers et al. (1991*a*) have used statistical arguments to deduce that many cloud cores have prolate, rather than oblate, shapes. They note that this result poses difficulty for the above view that cloud cores represent quasi-equilibrium structures in which quasi-static (poloidal) magnetic fields play an important part in the support against self-gravitation (see also Bonnell and Bastien 1991). Tomisaka (1991) points out that equilibrium clouds with prolate shapes become possible if they possess toroidal magnetic fields of substantial strength. Even without toroidal fields, oblate shapes represent a natural theoretical expectation only if we model the effects of molecular-cloud turbulence as an *isotropic* pressure. If "turbulent" support arises, instead, from the propagation and dissipation of nonlinear Alfvén waves in an inhomogeneous medium (see Arons and Max [1975], as well as the discussion on pp. 36–37 of Shu et al. [1987*a*]), then the largest amount of momentum transfer might well occur *parallel* to the direction of the mean field $\mathbf{B}_0$, rather than perpendicular to it. For example, if ambipolar diffusion causes $\mathbf{B}_0$ to become nearly straight and uniform, as occurs in the detailed calculations before collapse ensues, then the mean field exerts little stress on average, and the anisotropic forces associated with the fluctuating component $\delta\mathbf{B}$ in total $\mathbf{B} = \mathbf{B}_0 + \delta\mathbf{B}$ may actually cause cloud cores to assume prolate instead of oblate shapes. This tendency for condensations to stretch out along the direction of the mean field would only be enhanced by tidal forces if cloud cores tend to form as individual links on a magnetic "sausage" (cf. Elmegreen 1985*a*; Lizano and Shu 1989). Bastien et al.'s (1991) study of the fragmentation of elongated cylindrical clouds into such individual links differs from that of Lizano and Shu (1989) in that the former authors ignore magnetic effects and they assume that the core-formation process occurs by dynamical collapse rather than by quasi-static contraction.

Apart from the delicate question of the sense of elongation of the shape function $Q(\theta)$, the prediction (Eq. 3) with regard to the *radial* variation of ρ agrees reasonably well with observations of isolated cores and Bok globules (see, e.g., Zhou et al. [1990*b*], who comment particularly that the *outer* velocity fields are more representative of quasi-static contraction than dynamical collapse). The solution as written represents an asymptotic result; actual configurations become gravitationally unstable in their central regions before a singular density cusp develops there. However, because the phase of quasi-static contraction produces very high central concentrations in the numerical models, the resulting dynamical behavior would probably closely resemble the inside-out collapse known analytically for the (exact) singular isothermal sphere without rotation (Shu 1977). Such a collapse solution builds up a central protostar at a rate,

$$\dot{M}_{\rm infall} = 0.975\, a_{\rm eff}^3/G \tag{5}$$

that remains constant in time as long as the external reservoir of molecular cloud gas (with a r^{-2} density distribution) continues to last. For L1551, where $a_{\rm eff} \sim 0.35$ km s^{-1}, $\dot{M}_{\rm infall} = 1 \times 10^{-5}$ M$_\odot$ yr^{-1}.

The solution requires modification in the presence of rotation. Because ambipolar diffusion occurs relatively slowly, one might imagine that magnetic braking has time to torque down the cloud core sufficiently during the quasi-static condensation stage to allow even the direct formation of single stars, without subsequent angular momentum difficulties (see, e.g., Mouschovias 1978). Realistic estimates yield a somewhat different picture. Before the onset of dynamical collapse, magnetic braking can enforce more or less rigid rotation of the core at the angular velocity Ω of its surroundings (Mestel 1965*a*,1985; Mouschovias and Palelogou 1981). Once dynamical collapse starts, however, the infall velocities quickly become super-magnetosonic, and any initial angular momentum possessed by a fluid element carries into the interior. In the inside-out scenario presented above, dynamical collapse is initiated while much of the mass still has an *extended configuration*. For example, to contain 1 M$_\odot$, the wave of falling must typically engulf material out to $\sim 10^{17}$ cm. Even if such a core were perfectly magnetically coupled to its envelope before collapse, it would typically rotate at too large an initial angular speed Ω to allow the formation of just a single star with a dimension of $\sim 10^{11}$ cm. The simplest solution to the core's residual angular momentum problem is to form a star *plus a disk*.

Probably as a consequence of cloud magnetic braking, Arquilla and Goldsmith (1986) and Fuller and Myers (1987) find empirically that rotation does not usually play a dynamically important role on scales of a molecular cloud core and larger. This fact suggests the possibility of a perturbational analysis for the dynamical collapse problem (Terebey et al. 1984), with core rotation at an initially uniform rate Ω treatable as a small correction, in the outer parts, to the spherical self-similar solution for the collapse of a singular isothermal configuration. Large departures from sphericity occur only at radii comparable to or smaller than a centrifugal radius,

$$R_{\rm C} \equiv \frac{G^3 M^3 \Omega^2}{16 a_{\rm eff}^8} \tag{6}$$

where $M \equiv \dot{M}_{\rm infall} t$ equals the total mass that has fallen in at time t. Typical combinations of $a_{\rm eff}$, Ω, and M yield values for $R_{\rm C} \sim 10^{15}$ cm.

For $R_* \sim 10^{11}$ cm $<< R_{\rm C} \sim 10^{15}$ cm, the bulk of the freely falling matter (on parabolic streamlines because of the nonzero values of the specific angular momentum) does not strike the star directly, but forms a disk of size

$\sim R_C$ that swirls around the protostar. Radiative transfer calculations for the emergent spectral energy distribution in an infall model of this type give quite good fits to the data for IRS 5 L1551, if we choose $a_{\rm eff} = 0.35$ km s^{-1}, $\Omega = 1 \times 10^{-13}$ rad s^{-1}, and $M = 1$ M$_\odot$ (Fig. 5).

Fermi is supposed to have remarked that with three free parameters, he could fit an elephant. Figure 5 gives an empirical proof of this claim. Indeed, one can readily discern a tail at millimeter frequencies, attached to a broad back at far- and mid-infrared frequencies, with a 10 μm silicate absorption feature separating the shoulder from the head and a 3.1 μm water-ice feature defining a tusk. The trunk of the elephant emerges from the droop to optical frequencies, although the models do not reproduce the slight raising of the trunk seen in the data points because they do not properly account for the presence of scattered near-infrared and optical light. By Fermi's standards, then, our failure to catch the trumpet call of the elephant, without the need to introduce additional free parameters, prevents us from claiming a complete victory.

However, we note that the three independent parameters that go into the fits, $a_{\rm eff}$, Ω, and M, represent an irreducible set from an *a priori* theory, which have fundamental dynamical implications apart from the spectral energy fits. For example, the value $a_{\rm eff} = 0.35$ km s^{-1} well describes the properties of the core density distribution around L1551 (cf. Eq. [3] and the review of Evans [1991]). The agreement between the size of the spatial maps at millimeter and submillimeter wavelengths (Walker et al. 1990) and at infrared wavelengths (Butner et al. 1991) with the predictions of the radiative-transfer calculations of our model also bodes well for the values of $a_{\rm eff}$ and M, derived in the spectral energy fits on the basis of providing the correct overall optical depth to the central star and the absolute luminosity scale (from protostar theory). Finally, the numerical value of Ω, chosen to fit the depth of the silicate absorption feature, yields a predicted size for the circumstellar disk, $R_C = 42$ AU, which compares well with the limits 45 ± 20 AU deduced for this source by Keene and Masson (1990) from radio interferometric measurements of the thermal emission from dust grains in the putative disk. More recently, Goodman et al. (1991; see also Menten and Walmsley 1985) report the detection of a velocity gradient $\sim$4 km s^{-1} pc^{-1}, localized to the core region of L1551, consistent in magnitude with the model value $\Omega = 1 \times 10^{-13}$ s^{-1}, although the result may have been contaminated by the velocity shear introduced by the outflow.

A similar modeling of the spectral energy distribution of HH 7–11 suggests that this source also possesses infall at a rate $\dot{M}_{\rm infall} \gtrsim 1 \times 10^{-5}$ M$_\odot$ yr^{-1} (F. C. Adams, personal communication). Our ability to fit the spectral energy distribution of famous *outflow* sources like L1551 and HH 7–11 with pure *inflow* models bears directly on the fundamental paradox of Sec. III. Many of the observed embedded systems (those in stage *c* of Fig. 1) must possess *both* outflow and inflow, with the blowing off of a small polar cap in deeply embedded sources making little difference for the problem of infrared reprocessing in the rest of the infalling envelope. The challenge to observers

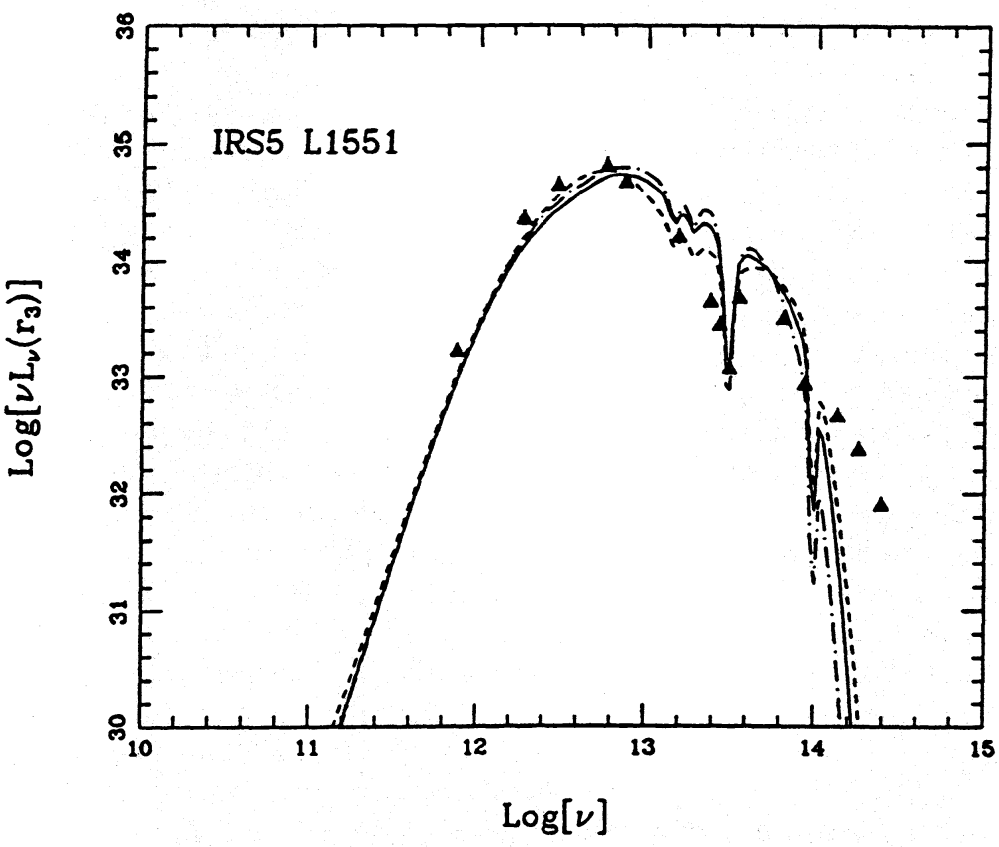

Figure 5. Spectral energy distribution for the bipolar outflow source L1551. The best fit for the data points (solid curve) results from choosing $a_{\rm eff} = 0.35$ km s^{-1}, $\Omega = 10^{-13}$ s^{-1}, and $M = 0.975\, a_{\rm eff}^3 t/G = 1$ M$_\odot$ (figure from Adams et al. 1987).

is to find spectroscopic evidence, in the low-luminosity sources, for the (rotating) collapse along the equatorial directions in an analogous manner that this task has been accomplished for high-luminosity protostars (see, e.g., Ho and Haschick 1986; Keto et al. 1987; Welch et al. 1987; Rudolph et al. 1990). In any case, here then lies a possible answer to the question: how can a star (SVS 13, the driving source in HH 7–11) form if it loses mass (through a neutral wind emerging from the poles) at the enormous rate $\dot{M}_{\rm w} = 3 \times 10^{-6}$ M$_\odot$ yr^{-1}? Clearly, a net gain can still take place if the star is being "force-fed" (through the equatorial regions) at a rate $\dot{M}_{\rm infall}$ which is $\gtrsim 3$ times larger yet. This empirical finding answers the "how" of our original question, but it does not address the question "why" a new-born star should be losing mass. A clue to the resolution of the latter issue lies in an examination of the very different way in which a star, in contrast to a disk, may deal with its angular momentum legacy.

V. PROTOSTAR FORMATION BY DISK ACCRETION

Notice that no mass scale that we can identify with stars emerges naturally in the theoretical picture of infall summarized in the previous section. To obtain a stellar mass $M_* \sim \dot{M}_{\rm infall} t_{\rm infall}$ from an object (a giant molecular cloud typically), whose characteristic mass scale much exceeds anything that we normally associate with stars, requires us to choose a small total duration $t_{\rm infall}$ (say, 10^5 yr) over which infall occurs. This duration might be set, for example, by the time required for an incipient outflow from the star to completely reverse the inflow. In spherical simulations of protostar formation and evolution (Stahler et al. 1980; Palla and Stahler 1990), the only event of significance to take place on such a time scale concerns the onset of deuterium burning and the establishment of an outer convection zone at the bottom of which the deuterium burns at a rate equal to that which accretion brings in a fresh supply. In low-mass protostars ($<2\ M_\odot$, i.e., precursors to T Tauri stars), deuterium burning occurs only slightly off-center; in intermediate-mass stars (2 to 8 $M_\odot$, i.e., precursors to Herbig Ae and Be stars), it occurs in a thin shell. In either case, the important feature is that the process induces an outer convection zone—a condition, when combined with rapid stellar rotation, that Shu and Terebey (1984) speculated would be conducive to dynamo action and to the appearance of strong magnetic fields on the surface of the star.

Calculations by Picklum and Shu (in preparation) demonstrate that accretion through a disk rather than by spherical infall does not modify the above conclusions appreciably, except to lower by a significant factor the emergent photon luminosity for a given rate of mass accretion. The latter effect may alleviate the discrepancies in time scales inferred for embedded protostars and revealed pre-main-sequence stars noted by Kenyon et al. (1990; see also Hartmann et al. 1991).

For star formation via the collapse of a cloud core modeled initially as a uniformly rotating singular isothermal sphere, the fraction of mass that suffers direct infall onto the star compared to that which first lands in the disk equals 1.29 $(R_*/R_C)^{1/3}$ (Adams and Shu 1986). This expression yields a small number (several percent) for stars with radii $\sim$ a few $R_\odot$ and disks with radii $\sim$10 to 100 AU. The small cross-sectional area of stars in comparison with their disks leads us to expect generically that most of the mass from a collapsing molecular cloud core, which eventually ends up inside the star, must first make its way through the disk (an original point of view argued first by Cameron [1962]; see also Mercer-Smith et al. [1984]).

The process of disk accretion is not well understood in any astrophysical system. The most frequently invoked mechanism involves the inward transport of mass and the outward transport of angular momentum by the friction associated with some form of anomalous viscosity in a differentially rotating disk (see, e.g., Lynden-Bell and Pringle 1974). In the case of the nebular disks that surround young stellar objects, Lin and Papaloizou (1985; see also Ruden and Lin 1986) identified thermal convection in the disk as a plausible source of the turbulent viscosity, and this may well hold for the optically thick dusty disks believed to surround all classical T Tauri stars. More recently, Balbus and Hawley (1991; see also Hawley and Balbus 1991) have rediscovered a powerful magnetohydrodynamic instability (Chandrasekhar 1960,1965; Fricke 1969) that afflicts weakly magnetized differentially rotating systems in which the angular velocity of rotation decreases outwards (as occurs in all known astrophysical disks). These developments hold great promise for providing a physical foundation for the central assumption of standard accretion-disk theory that the viscosity has an anomalous magnitude much higher than molecular values (for a review, see Pringle 1981).

However, in the most extreme cases (so-called "flat spectrum sources"; see Sec. VII; the Chapter by Beckwith and Sargent), the far-infrared radiation (relative to the near-infrared values) exceeds by a factor approaching 10^3, the value one would have naively predicted by the steady-state viscous model (Adams et al. 1988). The observed disks also have masses, as estimated from their submillimeter and millimeter emission, that compare favorably with those contained in the central stars (Adams et al. 1990; Beckwith et al. 1990; Keene and Masson 1990).

The above considerations suggest an interesting scenario for protostellar disk accretion (see also Chapter by Adams and Lin). The overall time scale for viscous transport of mass and angular momentum scales roughly as

$$t_{\rm vis} \sim \alpha^{-1}(R/H)^2 t_{\rm rot} \tag{7}$$

where (R/H) equals the aspect ratio of the disk (radius to vertical scale height), $t_{\rm rot}$ is the rotation period at the outer edge of the disk, and α is a dimensionless parameter in the so-called "alpha" prescription for the anomalous viscosity. For application to nebular disks, typically, $R/H \gtrsim 10$, $t_{\rm rot} \sim 10^3$ yr, and $\alpha < 1$ (e.g., $\alpha \sim 10^{-2}$ in the convection models of Ruden and Lin [1986]); thus, $t_{\rm vis} > 10^5$ yr. The statistics of disks around classical T Tauri stars sets a lifetime of $\sim 10^6$ to 10^7 yr (see Chapter by Strom et al.). This represents only a lower limit for $t_{\rm vis}$ because the disappearance of near-infrared radiation only implies the disappearance of small dust particles surrounding the central star, and not necessarily the viscous transport of all of the material, gas and dust, from the entire disk.

Suppose rotating infall piles matter into the disk faster than viscous accretion can transfer it to the central star. (Models for the main accretion phase of low-mass protostars suggest that $M/\dot{M} \sim 10^5$ yr.) The disk would then build up mass relative to the star until the disk becomes comparably massive. At this point, strong gravitational instabilities (nonaxisymmetric density waves) may develop in the disk that result, in the nonlinear regime, in the inward transport of mass and the outward transport of angular momentum. The energy dissipation associated with this "wave dredging" differs, however, from the viscous mechanism, and we might expect to see a different resultant radial distribution of temperature in the disk and a different emergent spectrum (see Sec. VII). In particular, for the transport to occur *globally* (as indicated by the observations of flat-spectrum sources), from the star-disk interface to the outer edge of the disk $\sim 10^4$ times larger, we can restrict our search of unstable modes in nearly Keplerian disks to one-sided disturbances, in which circular streamlines distort to elliptical ones with foci at the star (Adams et al. 1989).

Unless such instabilities run away catastrophically to form a companion star (a possible origin for binary systems), we may speculate that they self-regulate the amount of mass in the disk during the infall phases (Fig. 1*b*, and *c*) so that it never consists of a large fraction of the total. Thus, we expect that some appreciable fraction f_1 of the total infall rate onto the system (mostly onto the disk), $\dot{M}_{\text{infall}}$, must make its way through the disk at an accretion rate $\dot{M}_{\text{acc}}$, eventually to be deposited onto the star,

$$\dot{M}_{\text{acc}} = f_1 \dot{M}_{\text{infall}}. \tag{8}$$

As the detailed calculations of Adams et al. (1989) demonstrate, disks with temperature profiles consistent with the observational requirements suffer a strong one-armed spiral instability (see Fig. 6) if they have masses equal, say, to one-half or one-fourth of the total; thus a reasonable estimate for f_1 ranges from 1/2 to 3/4. We emphasize that this estimate for f_1 represents a *time-averaged* value; the instantaneous rate at which mass from the disk actually empties onto the star may suffer large fluctuations about an average value. In particular, the arguments of Sec. VII suggest that disk accretion induced by global gravitational instabilities may have intrinsic tendencies to occur in a nonsteady manner.

VI. STELLAR WINDS AND BIPOLAR FLOWS: THEORY

The matter entering the protostar from the disk at the rate $\dot{M}_{\text{acc}}$ carries a relatively large specific angular momentum. For example, compared to the specific angular momentum of disk matter circling in orbit just outside of a star's equator, the angular momentum per unit mass of a uniformly rotating polytrope of index 1.5 (a good model for a fully convective low-mass protostar) is a small number $b = 0.136$, even if the star were to rotate at the brink of rupture (James 1964). Thus, the disk feeds material to the star that contains, per gram, about $b^{-1} \sim 7$ times more angular momentum than can be absorbed dynamically by the star. Hence, once the star begins to accrete matter from a centrifugally supported disk, it can increase its mass M_*, at most, only by an additional small fraction b before it reaches breakup. In practice, the star may be able to accept even less matter from the disk because even the matter accumulated by the star through direct infall carries nonzero values of specific angular momentum (see, e.g., Durisen et al. 1989*b*).

As the mass that can be accumulated by a star through direct infall from a realistically rotating molecular cloud core amounts to only several percent

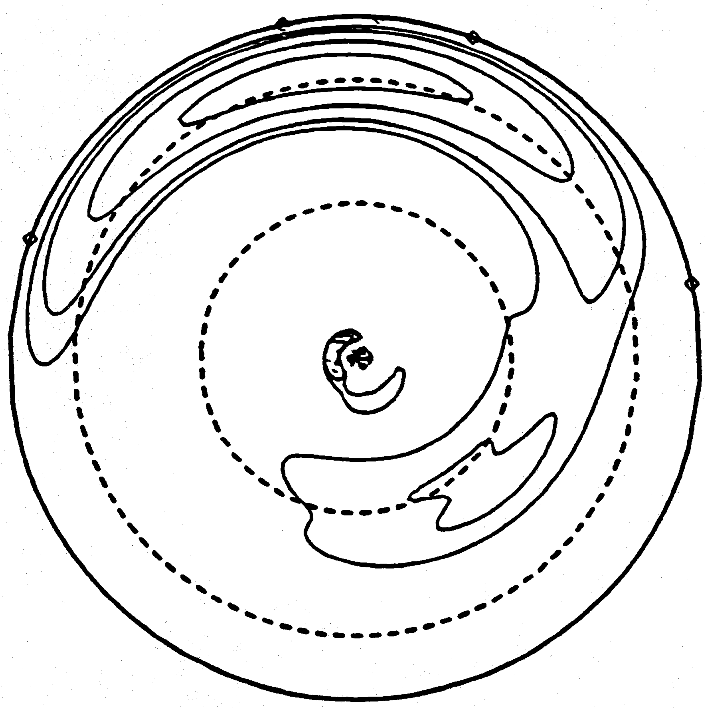

Figure 6. Contour plot of the lowest-order, growing, $m = 1$, mode in a star/disk system where the disk's mass equals the star's mass, but the disk's radius is 10^4 times larger than the star's radius (figure from Adams et al. 1989).

of a solar mass, and as the additional amount that can be accreted from a disk amounts to a small fraction of this tiny value, how do stars of sunlike masses and larger ever form?

Paczynski (1991) and Popham and Narayan (1991) calculated star plus disk models where the star, spun to (slightly faster than) breakup, can viscously transport angular momentum to an adjoining disk, and thereby continue to accrete an indefinite amount of mass. This model presumes the existence of a sufficiently efficient viscous coupling between star and disk as to transport away the excess specific angular momentum (above the fraction b) brought in by disk accretion, an assumption that may not hold if the mechanism of disk accretion is gravitational rather than viscous. In any case, their mechanism of continuous outward angular-momentum transport cannot work for T Tauri stars, which are observed to rotate at speeds significantly less than breakup. Furthermore, the almost ubiquitous presence of energetic bipolar flows among deeply embedded sources suggests that quiescent viscous torques do not provide the primary source of relief for the angular momentum difficulty of protostars in their main accretion phase.

Such objects evidently find a more spectacular way of shedding the excess angular momentum brought in by disk accretion. Shu et al. (1988) proposed that the protostar could fling off the excess in a powerful wind once the protostar gets spun to breakup by the disk, if the star possesses sufficiently strong surface magnetic fields (see also Hartmann and MacGregor 1982).

Schematically, the process works as follows. Like other rotating stars with outer convection zones, a protostar probably has a network of open and closed magnetic field lines that protrude from its surface (Fig. 7). An ordinary (ionized) stellar wind blows out along the open field lines; this O-wind may have a higher intensity than that characteristic of a normal star of the same spectral type, luminosity class, and rotation rate, because of the enhancement of stellar dynamo action associated with circulation currents induced by the disk through Ekman pumping (see Sec. VII).

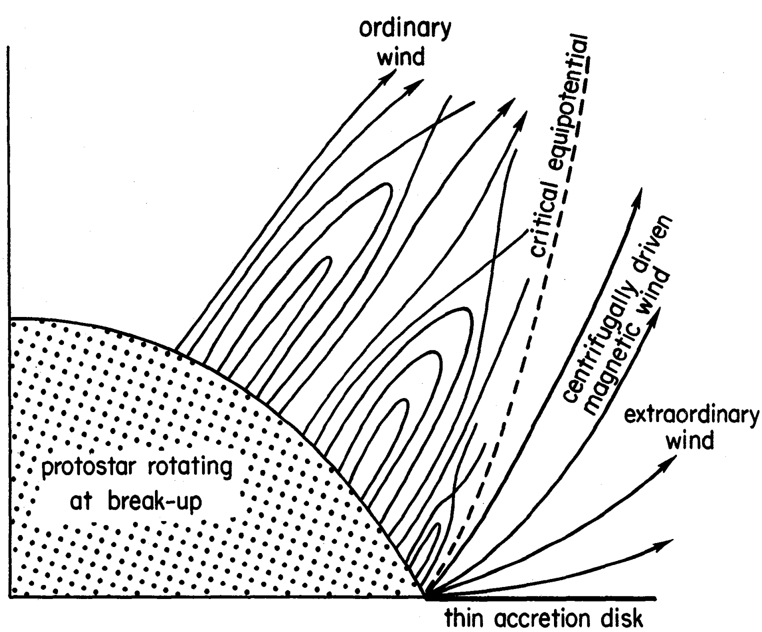

Figure 7. Schematic model for heavy mass-loss driven magnetocentrifugally from the equator of a protostar which is forced to rotate at breakup because of accretion through a Keplerian disk. All streamlines to the right of the critical equipotential (marked as a dashed curve) correspond to possible "downhill" paths for electrically conducting gas attached to magnetic lines of force that rotate rigidly with the star. The large slip velocities present in the interfaces between the X-wind and the O-wind and between the X-wind and the disk may lead to turbulent mixing layers that have observable radiative signatures (figure adapted from Shu et al. 1988).

If open field lines circulate into the equatorial regions of the protostar, which spins at breakup by assumption, the O-wind may intensify even more via the X-celerator mechanism (Shu et al. 1988) and become an extraordinary wind. The magneto-centrifugal acceleration in the X-wind takes place much as described in Sec. III on the Blandford-Payne mechanism for disk winds, except that the geometry here no longer selects against field lines (equivalent to streamlines for a conducting gas) that emerge in the "downhill" directions of the effective (corotating) potential. To carry away the requisite mass-loss rate $\dot{M}_w$, the sonic transition of the X-wind must occur typically in the photosphere of the star, where the material, being cool, has a relatively low ionization fraction. Thus, the X-wind is predominantly a *neutral* or *lightly ionized* wind (Natta et al. 1988; Ruden et al. 1990).

Consider the quasi-steady state wherein spinup by disk accretion at a rate $\dot{M}_{acc}$ is balanced by spin-down via a magnetized stellar wind at a rate $\dot{M}_w$ in such a way that at each stage of the process the star continues to rotate exactly at breakup. Suppose the specific angular momentum carried away in the X-wind, measured as an average over all mass-carrying streamlines, equals some multiple $\bar{J}$ of the specific angular momentum of the material in circular orbit at the equator of the star. If we further assume the equatorial radius to be proportional to the stellar mass, as roughly true for a star actively burning deuterium at its center, we find that the requisite mass-loss rate $\dot{M}_w$ equals some fraction f_2 of the disk accretion rate $\dot{M}_{acc}$ (cf. Eq. 1 of Shu et al. 1988):

$$\dot{M}_w = f_2 \dot{M}_{acc} \qquad \text{with} \qquad f_2 = \frac{1-2b}{\bar{J}-2b} \tag{9}$$

where $b = 0.136$ is the pure number defined earlier. (If only the equator of the protostar rotates near breakup, the effective value of b that appears in Eq. [9] could be [much] smaller.) To derive the above expression for f_2, we have assumed that all of the excess angular momentum brought in by accretion above that needed to keep the star exactly at breakup is carried away by the wind, i.e., that none of it is viscously transported to the disk.

The extent by which $\bar{J}$ exceeds unity depends on the strength of the stellar magnetic field in the neighborhood of the X-point of the equipotential. For escaping streamlines to have a finite terminal velocity v_w at infinity, the (rotating) magnetic field lines must exert sufficient torque as to make $\bar{J}$ exceed 3/2. For example, in order for the terminal velocity of the X-wind to have a measured value equal to the breakup speed at the equator of the star (typically $\sim$150 km s^{-1} for low-mass protostars), $\bar{J}$ must have a value at least equal to 2. In fact, because some of the angular momentum is asymptotically carried away by magnetic stresses rather than all by the fluid, detailed numerical calculations suggest that the actual needed value of $\bar{J}\sim 3.6$, implying a fiducial value for $f_2 \sim 0.2$. Combining Eqs. (8) and (9), we now obtain a relationship between inflow and outflow rates,

$$\dot{M}_w = f\, \dot{M}_{infall} \tag{10}$$

with $f \equiv f_1 f_2 \sim 0.1$ to 0.2 typically, in very rough agreement with the observations.

On the other hand, if the magnetic field emerging from the neighborhood of the X-point of the effective potential had infinite strength, X-wind gas would be flung to infinity in the entire sector from the equator to the critical equipotential surface (the dotted curve in Fig. 7). Because this equipotential surface (the last rigid "wire" to which an attached bead could still slide "downhill") bends up to turn asymptotically parallel to the rotation axis of the star (at a radius = $\sqrt{3}R_e$), the strong X-wind could act to focus and collimate an otherwise isotropic and more mild O-wind into an (ionized) optical jet (one on each side of the equator).

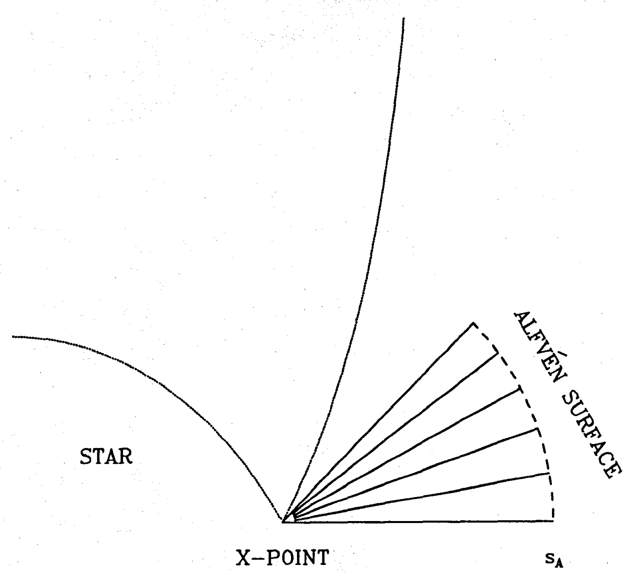

Figure 8. Spreading of streamlines for an X-celerator driven wind in the meridional plane, computed in a frame that corotates with the star. The bottom-most streamline has been constrained to flow horizontally across the surface of a perfectly flat disk; the uppermost streamline has also been taken to be straight for sake of simplicity. In a more realistic model, the slip interfaces between the X-wind and the O-wind and between the X-wind and the disk surface are likely to be turbulent. A dashed curve marks the Alfvén surface, where the magnetic stresses no longer dominate over the gas inertia. The dotted curve marks the location of the critical equipotential surface that, together with its counterpart in the lower half of the diagram (not drawn), crosses the stellar surface at an X-point. (figure from Shu et al., in preparation).

For stellar magnetic fields of finite strength, the situation becomes more complicated. The method of matched asymptotic expansions provides a tractable approach to the resultant magnetohydrodynamic problem if we adopt the assumption of axial symmetry (Shu et al. 1991). Figure 8 gives a sample solution for a case with a dimensionless stellar field strength that produces Alfvén crossing at a distance equal to one additional stellar radius from the equator of the star. In this calculation, we have used upper and lower boundary conditions that artificially constrain the limiting streamlines, $\psi = 0$ and $\psi = 1$, to be perfectly straight (see below). Interior to the Alfvén surface defined by $\mathcal{A} = 0$, where $\mathcal{A}$ is an Alfvén discriminant, the stream function ψ satisfies an elliptic partial differential equation:

$$\nabla \cdot (\mathcal{A} \nabla \psi) = \mathcal{Q}. \tag{11}$$

This equation has the formal structure of steady-state heat conduction; here, however, we are concerned with the spreading of streamlines, and not the spreading of heat. The source term $\mathcal{Q}$ responsible for streamline spreading in the meridional plane equals a sum of three terms that are each proportional to the derivatives (across streamlines) of quantities that are conserved along streamlines: Bernoulli's constant $H(\psi)$, the ratio of magnetic and mass fluxes $\beta(\psi)$, and the specific angular momentum carried in both matter and field $J(\psi)$. Smooth passage through the sonic transition near the X-point of the effective potential yields $H(\psi)$; for the part of the problem much beyond the immediate neighborhood of the X-point of the critical equipotential, $H(\psi)$ goes to 0 in the limit of a cold flow. The requirement that the flow accelerates smoothly through the Alfvén surface, $\mathcal{A} = 0$, places a constraint on the normal derivative of ψ, $\nabla\mathcal{A} \cdot \nabla\psi = \mathcal{Q}$, that determines the spatial location and shape of the Alfvén surface and the functional form of $J(\psi)$ if we are given $\beta(\psi)$. In other words, the geometry of the problem is entirely determined if we know the strength of the stellar magnetic field and the way that mass is loaded onto the open field lines at the equatorial belt of the protostar.

For the dimensional parameters that approximately apply to HH 7–11 or L1551, we find that a poloidal speed of $\sim$150 km s^{-1} (corresponding to $\bar{J} \sim 3.6$) can be achieved at the Alfvén surface when the equator of the protostar has open photospheric magnetic fields of the strength $\sim$4 kilogauss. Such fields appear reasonable if we extrapolate from the product of field strength and filling factor, 1 to 2 kilogauss, inferred for weak-line T Tauri stars by Basri and Marcy (1991) from the Zeeman broadening of spectral lines with high Landé-g factors.

It may be informative to elucidate why we believe the X-celerator mechanism to work for a star, but not for a disk (Sec. III). To carry the observed mass-loss rate, the density of the wind at the sonic surface must be large, close to photospheric values. Thus, the sound speed is low, and the sonic transition can occur only near X-points where the effective gravity vanishes, because in a frame which corotates with the footpoint of the magnetic field, the magnetic stresses cannot help the gas make a sonic transition (see Eq. [3] of Shu et al. 1988). On the other hand, if the gas is to continue to accelerate to higher speeds by magneto-centrifugal flinging after having made the sonic transition, the Lorentz force must be able to overcome both inertia and gravity. Because the densities are very high at the sonic surface, a smooth transition from a gas-pressure driven flow to a magnetohydrodynamically driven flow requires either the magnetic field **B**, or its curl, $\nabla \times \mathbf{B}$, to be large, or both (cf. Eq. [2]). The geometry of a thin disk requires $\nabla \times \mathbf{B}$ to be large, but to keep the gas sub-Alfvénic in the post-sonic-transition region, **B** also has to be large. This combination proves fatal, we believe, to disk-wind models in the protostellar context.

An X-celerator driven wind from a protostar (or star-disk interface) fares better, because we require only **B** to be large, and we can relatively easily believe that magnetic fields *rooted in the deep convection zones of a rapidly* *rotating star* could achieve the requisite values. High-mass stars that lack outer convection zones have greater theoretical difficulty generating surface magnetic fields, although the instability mechanism discussed by Balbus and Hawley (1991; see also Fricke [1969] for specific application to the case of stars) offers a promising mechanism for such objects to generate moderate-to-small magnetic fields provided only that they suffer differential rotation of dynamical significance for the structure of the star (with Ω decreasing radially outward). Indeed, X-celerator driven winds in such a situation may offer an explanation why the terminal velocities of winds from high-mass protostars have such low values (100 to 200 km s^{-1}) in comparison with the high values (1000 to 2000 km s^{-1}) that characterize (probably radiation-driven) winds from main-sequence and evolved stars of early spectral type (see, e.g., Chiosi and Maeder 1986).

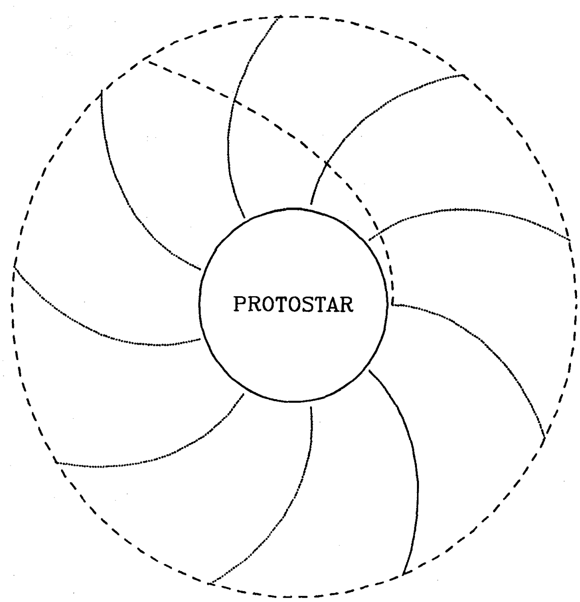

Figure 9. Field lines (dotted curves) and trajectory of a fluid element (dashed curve) for an X-celerator driven wind in the equatorial plane of an inertial frame of reference. The protostar rotates in a counter-clockwise sense, so that open field lines form spirals that trail in the sense of rotation. The dashed circle marks the location of the Alfvén surface in a calculation where we have explicitly assumed axial symmetry (figure from Najita et al., in preparation).

In Fig. 9, we plot magnetic field patterns and fluid trajectories of the flow in the equatorial plane of our X-celerator model. In a frame that rotates with the same angular speed as the star, streamlines and field lines coincide (dotted curves). In a stationary frame, however, the trajectory of an element of electrically conducting gas takes the form indicated by the dashed curve. The fluid element starts by making a sonic crossing close to the photosphere of the star. The acceleration to sonic speeds (7 or 8 km s^{-1} typically) can occur purely by the outward push of photospheric gas pressure in a small equatorial belt where the effective gravity nearly vanishes because we have assumed that the protostar rotates at breakup speeds. As the open magnetic line of force to which the gas is tied rotates in a counter-clockwise direction with the star, the excess rotation (above local Keplerian values) enforced by the field accelerates the gas along the field line (with negligible slip for the typical degrees of ionization computed by Ruden et al. [1990]), until the gas reaches the Alfvén velocity at the position of the dashed circle. At this point, if the gas speed exceeds the local escape velocity by a significant margin, the Alfvén speed will also typically equal a healthy fraction (say, ~90%) of the final terminal speed. Because the magnetic field does not have infinite strength, it cannot keep the gas corotating indefinitely with the star, and the pattern of field lines (in either the rotating or inertial frame) form trailing spirals. The field lines and streamlines form a similar lagging pattern at other latitudes, except that they climb in the vertical direction (toward the rotational poles) as the gas flows away from the protostar.

As already mentioned, the boundary conditions in Fig. 8 on the uppermost and bottom-most streamlines have been taken to be perfectly straight so as not to bias the results for the interior. When we do this, the intermediate streamlines have a tendency also to remain straight, with the gas at higher latitudes having larger densities and slower terminal speeds (for a magnetic field distribution that starts off approximately uniform at the sonic surface). We are currently trying to compute the amount of bending which takes place when the shape of the upper streamline is left more realistically as a free boundary (pressure balance between an X-wind and an O-wind). When such a calculation becomes available, we will be able to estimate the degree of collimation of the X-wind; in particular, we will be able to compute the angle dependence $P(\theta)$ of the rate of momentum injection (needed for any *a priori* calculation of the shapes of the swept-up shells of molecular gas):

$$\frac{\dot{M}_w v_w}{4\pi} P(\theta). \tag{12}$$

Even without detailed knowledge of $P(\theta)$, a simple order of magnitude estimate using the theory developed so far yields the velocity of the swept-up bipolar lobes of molecular gas roughly as the geometrical mean between the terminal speed of the wind v_w and the effective sound speed of the ambient molecular cloud core $a_{\rm eff}$ (Shu et al. 1991; see also Chapter by Königl and Ruden). For v_w measuring hundreds of km s^{-1}, and $a_{\rm eff}$ ranging from a fraction of a km s^{-1} (low-mass cores) to greater than 1 km s^{-1} (high-mass cores), we would then predict typical lobe speeds $\sim (v_w a_{\rm eff})^{1/2} \sim 10$ km s^{-1}, in rough agreement with the observations.

For well-collimated flows, $P(\theta)$ will have values larger than unity for a small range of angles θ near the two poles, 0 and π. Our previous discussion leads us to suspect that younger sources, which may have stronger magnetic fields, possess more highly collimated outflows. We speculate that as the sources age, their rotation speeds progressively fall below the critical rate (see Sec. VII), their magnetic fields weaken, and their outflows fan out more in polar angle and sweep out an increasingly greater solid angle of the overlying envelope of gas and dust. In this fashion may the protostars in stage *c* of Fig. 1 become optically revealed as the T Tauri stars of stage d.

VII. REVEALED T TAURI STARS

The ability to study T Tauri stars at optical and near-infrared wavelengths has yielded tremendous observational dividends. Foremost in importance among the discoveries of the past decade has been the realization that many of the photometric and spectroscopic peculiarities (variability, strong emission lines, infrared and ultraviolet excesses, strong outflows) long known to characterize these systems (cf. the reviews of Herbig 1962*b*; Kuhi 1978; Cohen 1984; Imhoff and Appenzeller 1987; Bertout 1986,1989; Appenzeller and Mundt 1991) may owe their explanation to the presence of circumstellar disks. Several independent lines of evidence lead to the conclusion that circumstellar disks surround many young stellar objects, e.g., the large polarizations often seen in these sources (Elsasser and Staude 1978; Bastien and Menard 1988; Bastien et al. 1989), or the asymmetry of forbidden O I line profiles (Appenzeller et al. 1984; Edwards et al. 1987), or the motions of SiO masers (Plambeck et al. 1990). Here, however, we shall concentrate on the nature and implications of the infrared and ultraviolet excesses in T Tauri stars.

We begin with the ultraviolet excess (Kuhi 1974; Herbig and Goodrich 1986), which has been plausibly linked by Hartmann and Kenyon (1987*a*), and Bertout et al. (1988) to the action of a boundary layer between the star and the disk (see also Chapter by Basri and Bertout). For a conventional boundary layer to arise (gas heated to high temperatures by the frictional rubbing of a rapidly rotating disk against a slowly rotating star), a disk must abut the star.

Cabrit et al. (1990; see also Chapter by Edwards et al.) find an interesting correlation between the strength of the near-infrared excess (a measure of the amount of disk just beyond the boundary layer) and the strength of Hα emission (a measure of T Tauri wind power). Systems inferred to be missing inner disks (radii equal to a few stellar radii) exhibit relatively little wind power (weak-line T Tauri stars), whereas systems with appreciable inner disks have relatively strong winds (classical T Tauri stars). As millimeter and submillimeter investigations reveal no strong biases with respect to the question of whether weak-line and classical T Tauri stars possess outer disks (radii equal to thousands of stellar radii), Edwards and her colleagues interpret their finding to imply that pre-main-sequence and protostellar winds are driven by disk accretion (one that extends virtually right up to the surface of the star). At first sight, this interpretation seems to bode well for the X-celerator theory, which pinpoints the interface between star and disk as the source of the wind. In particular, theories that use self-similar solutions for disk winds (Blandford and Payne 1982; Konigl 1989) have a difficult time explaining why missing just a small portion (the innermost part) of a disk should completely change the nature of the solution. However, a deeper probing of the T Tauri results shows that they also pose a severe problem, potentially fatal, for X-celerator models.

For the X-celerator model of Sec. VI to work, we need to posit that the equator of the star, to which the open magnetic field lines are tied, rotates at breakup. The very fact that T Tauri stars possess boundary layers (if this is the correct interpretation for the ultraviolet excess) demonstrates, however, that their equatorial regions are not rotating at breakup. Even more damaging, direct spectroscopic investigation shows that most T Tauri stars rotate at speeds about an order of magnitude *slower* than breakup (Vogel and Kuhi 1981; Bouvier et al. 1986*a*; Hartmann et al. 1987; see the review of Bouvier 1991). Yet, Edwards et al. (see also Chapter by Hartmann et al.) deduce that T Tauri stars also satisfy an inflow-outflow relationship of the form of Eq. (9), $\dot{M}_w = f_2 \dot{M}_{acc}$, with f_2 numerically not very different from the value that we have deduced from X-celerator theory ($f_2 \sim 0.2$).

How can we resolve the discrepancy? Galli and Shu (in preparation) propose that a partial solution may be found by asking how an accretion disk actually tries to spin up a slowly rotating star. Their answer involves, not the inward diffusion of vorticity as in the models of Paczynski (1991) and Popham and Narayan (1991), but *Ekman pumping*. The mathematics for Ekman pumping is fairly involved; however, the basic physics is simple, and easily explained by analogy to spinup in a cup of tea.

Suppose we place a cup of tea on a turntable spinning at angular speed Ω, and we ask how long it takes for the liquid, initially at rest, to spin up to the same angular speed as the boundaries of the cup. If spinup occurred by the diffusion of vorticity from the boundary layer to the interior, the theoretical answer, in order of magnitude, would be given by the familiar formula:

$$t_{\rm diff} \sim L^2/\nu \qquad (13)$$

where L represents the typical dimension of the cup, and ν gives the kinematic viscosity of tea. Putting in characteristic numbers, we would deduce $t_{\rm diff} \sim$ tens of minutes. In fact, the actual spinup time empirically measures more in the neighborhood of tens of seconds, and is theoretically given as the geometric mean between the viscous diffusive time scale Eq. (13) and the overturn time scale Ω^{-1}:

$$t_{\rm spinup} \sim (t_{\rm diff}/\Omega)^{1/2}. \qquad (14)$$

The formula (14) represents essentially the time that it takes tea to circulate from the interior of the cup to the boundaries, where it can quickly match its rotational speed to that of the cup (Fig. 10a). In other words, Ekman pumping sets up a secondary circulation that brings the tea to the boundary layer rather than waiting for the effects of the accelerating agent to diffuse to the interior. Because the bottom of the cup speeds up the tea in contact with it relative to the tea in the interior, the pressure gradient needed to balance

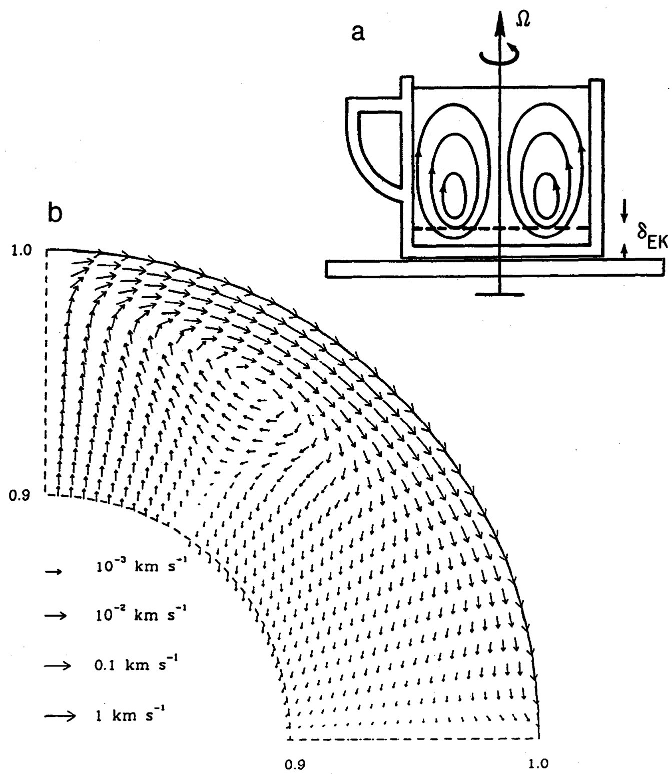

Figure 10. (a) Ekman pumping in a teacup (schematic). A thin Ekman layer δ_{EK} forms above the bounding surface perpendicular to the axis of rotation. (b) Ekman pumping in a T Tauri star modeled via the spinup of a (uniformly rotating) polytrope of index 1.5 with turbulent viscosity estimated by convective mixing-length theory. A secondary circulation pattern sets up in the meridional plane when the star is subjected to a shear stress at its surface that is highly concentrated toward the equatorial plane. The velocity scale corresponds to a model in which $M_* = 0.8$ $M_\odot$, $R_* = 3$ $R_\odot$, $\Omega_* R_* = 20$ km s^{-1}, and $\nu = 10^{14}$ cm^2 s^{-1}. The radial scale from 0.9 R_* to 1.0 R_* has been expanded for clarity of display (figure from Galli 1990).

the centrifugal force of rotation differs in the two regions. The difference in pressure distribution outside and inside the boundary layer then pumps the fluid from the former to the latter in such a way as to set up the secondary circulation depicted in Fig. 10a.

We believe an analogous situation to hold for the case of T Tauri stars, except that gravity, in addition to pressure and inertial terms, enters the equation for force balance. Schematically, the presence of an abutting disk tries to spin up the equatorial regions of a star faster than its poles. The excess of centrifugal support in the equatorial regions necessitates less vertical pressure support against gravity there than elsewhere. As a consequence, the unbalanced horizontal pressure gradients will push gas from the surface layers of the rest of the star toward the equator. Continuity requires those streamlines converging onto the equator of the star that do not blow outwards in a wind to resubmerge into the interior of the star, establishing a circulation pattern that looks as depicted in Fig. 10b.

Figure 10b represents a detailed calculation from Galli's (1990) doctoral thesis of Ekman pumping in a polytrope of index 1.5 (to represent a fully convective star). The calculation, however, contains only one of two possible effects that can lead to a secondary circulation: the *barotropic* response to the spinup induced by the frictional torque of an adjacent disk highly concentrated toward the equatorial regions. The rotational evolution, with $d\Omega/dt > 0$, is assumed to proceed keeping the zeroth-order rotation spatially uniform, which leads to perturbations in the angular speed of the stellar interior that are stratified across cylinders. If the true surfaces of constant Ω in T Tauri stars correspond—as they do in the case of the convection zone of the Sun (Dziembowski et al. 1989)—not to cylinders, but to radial cones, a second effect can arise because the surfaces of constant pressure do not then correspond to surfaces of constant density. In this situation, the *baroclinic* generation and maintenance of meridional mass flow may compete with or even dominate over the barotropic part. In other words, the meridional flows induced by a quasi-steady balance between spinup through a disk and spin-down via a wind may make the application of Fig. 10b to actual T Tauri stars somewhat oversimplified.

Nevertheless, Fig. 10b shows a robust qualitative feature of great importance to our current discussion: the circulation of matter in the photospheric and subphotospheric layers toward the equator. If such layers carry stellar magnetic field lines along with them, then we can easily visualize how open field lines of kilogauss strength might be brought from the rest of the star that rotates slowly (as required to satisfy the spectroscopic observations of T Tauri stars) to an equatorial band that rotates quite rapidly (by definition, at "breakup" when we enter the disk proper). Conceivably, the Doppler imaging (Vogt 1981) of classical T Tauri stars could test whether circulation currents of this geometry cause starspots to migrate toward the equator on time scales (radius of star divided by photospheric circulation speed) of weeks to months. If so, the resulting ejection of matter in a powerful magnetocentrifugal wind along open field lines might then qualitatively proceed, more or less, as we described in Sec. VI for an X-celerator-driven flow. As it takes on the order of a day for circulation-driven, inhomogeneous, magnetic structures to migrate through the sonic surface of a wind originating from an equatorial belt equal to a few percent of the radius of the T Tauri star, we can now understand why optical spectral lines believed to be wind diagnostics (in contrast to general photospheric diagnostics) should show profiles that vary wildly on that sort of time scale (G. Basri, personal communication).

The above speculation, that T Tauri winds originate basically in a highly magnetized and rapidly rotating boundary layer, amounts essentially to a compromise solution between a stellar-driven wind and a disk-driven wind. In this compromise, we make use of the known ability of a star to hold onto very strong magnetic fields, and we take advantage of the natural tendency for disks to possess very rapid rotation. More work, however, needs to be done on this problem before we can claim that we can safely import to this complex situation the main features of the solutions for more simple models.

We should note, however, that *closed* field lines may play a competing role to open ones. If stellar circulation currents force closed field lines of sufficient strength to thread through the disk (a process made easier at a late stage when the disk accretion rate onto the star drops to small values), and if the star becomes magnetically coupled largely to those parts of the disk that rotate fairly slowly, then we might have another explanation as to why T Tauri stars have projected rotational speeds $v \sin i \lesssim 15$ km s^{-1} (see Konigl [1991]). For this braking mechanism to work, however, the same magnetic torques associated with fields threading the gaseous disk need to empty out much of the rapidly rotating inner portions of the disk, which would otherwise act to spin up the star, rather than spin it down. Weak-line T Tauri stars may represent (evolved) systems missing such inner disks, and it is interesting to note that they do not possess the powerful (X-celerator-driven) winds that classical T Tauri stars do.

We turn now to the dynamics of the disk proper, and the origin of the infrared excesses of T Tauri stars. Since the pioneering work of Mendoza (1966,1968), astronomers have known that T Tauri stars emit significantly more infrared radiation than other stars of their spectral type (typically K subgiants). From the start, thermal emission from circumstellar dust has been suspected as the culprit; however, the geometry of the dust distribution remained obscure until Lynden-Bell and Pringle (1974) suggested that it might lie within a viscous accretion disk. In particular, Lynden-Bell and Pringle pointed out that an optically thick disk, vertically thin but spatially extended in distance ϖ from the rotation axis over several orders of magnitude, and possessing a power-law distribution of temperature,

$$T_{\rm D} \propto \varpi^{-p} \tag{15}$$

would exhibit a power-law infrared energy distribution,

$$\nu F_\nu \propto \nu^n \qquad \text{where} \qquad n \equiv 4 - \frac{2}{p}. \tag{16}$$

Viscous accretion-disk theory for a system in quasi-steady state predicts $p = 3/4$, i.e., $n = 4/3$.

Rucinski (1985) examined the observational evidence and he concluded that, while the infrared excesses of T Tauri stars did well approximate a power-law distribution, the mean value of n is appreciably smaller than 4/3 (see also Rydgren and Zak 1987). Adams et al. (1987) noted that the presence of considerable infrared emission without a correspondingly large visual extinction of the central star implies that the circumstellar distribution of dust must exist in orbit in virtually a single plane about the star. In such a situation, a *non-accreting*, optically thick, flat, and radially extended disk would intercept 25% of the starlight and reprocess it to near- to far-infrared wavelengths, with disk-temperature and spectral-energy distributions also very nearly satisfying Eqs. (15) and (16), except that the coefficient of proportionality would be exactly known. Unshielded large dust particles at a distance r from a star would, of course, have a temperature distribution that satisfies $T_D \propto r^{-1/2}$. However, astronomers believe that young nebular disks are completely thick at optical wavelengths even vertically through the disk, so that starlight cannot shine directly on dust particles by propagating through the midplane, but must come at oblique angles from near the limbs of a star of size R_*. This effect introduces an extra geometrical factor of essentially R_*/r into the relation between absorbed starlight ($\propto [R_*/r]r^{-2}$) and emitted infrared radiation ($\propto T_D^4$), which makes the reprocessing temperature law $T_D \propto r^{-3/4}$ rather than $r^{-1/2}$. To remind ourselves of this important difference between the reprocessing in an opaque flat disk with faces perpendicular to the z direction and in an airless planetary surface with projected area perpendicular to r, we write the former law as $T_D \propto \varpi^{-3/4}$.

Adams et al. (1987) also demonstrated that the few cases of T Tauri systems that did show the canonical value $n = 4/3$ could indeed be explained in terms of *passive* or *reprocessing* disks; whereas, those disks that have the most extreme infrared excesses and, therefore, exhibit the most quantitative evidence for a nonstellar contribution to the intrinsic luminosity, correspond, not to $n = 4/3$, but to $n = 0$ (Adams et al. 1988; see Fig. 11a,b). Flat-spectrum sources with $n = 0$ require $p = 1/2$, i.e., $T_D \propto \varpi^{-1/2}$. A comprehensive survey by Beckwith et al. (1990; see also Chapter by Beckwith and Sargent) fitted temperature power-laws to the infrared excesses of classical T Tauri stars, and found that the derived values for the exponent p range from the extremes $p = 3/4$ (steep-spectrum sources) to $p = 1/2$ (flat-spectrum sources), with more sources having the latter value than the former. Strom claims, however (see his chapter in this volume), that the "flatness" of the spectra of many of the sources disappears if one uses long-wavelength measurements with smaller beams than the IRAS measurements.

Kenyon and Hartmann (1987) proposed that *geometric flaring* in purely reprocessing disks could flatten out the spectral energy distributions of the spatially flat models that predict $n = 4/3$, and thereby remove much of the discrepancy between theory and observation. A greater interception and scattering of stellar photons would also better account for the relatively large fractional polarization detected for some T Tauri stars (Bastien et al. 1989). In contrast, Hartmann and Kenyon (1988) note that the source FU Orionis, which may be undergoing enhanced accretion through a disk (for a dissenting view, see Herbig 1989*a*), *does* show the canonical value $n = 4/3$ (see also Adams et al. 1987), and they conclude that nebular disks may transfer much, if not most, of their mass onto the central stars through such recurrent episodic outbursts. Finally, Kenyon and Hartmann (1991) conclude that, as a class, many of the FU Orionis variables may still be surrounded, not by flared disks,

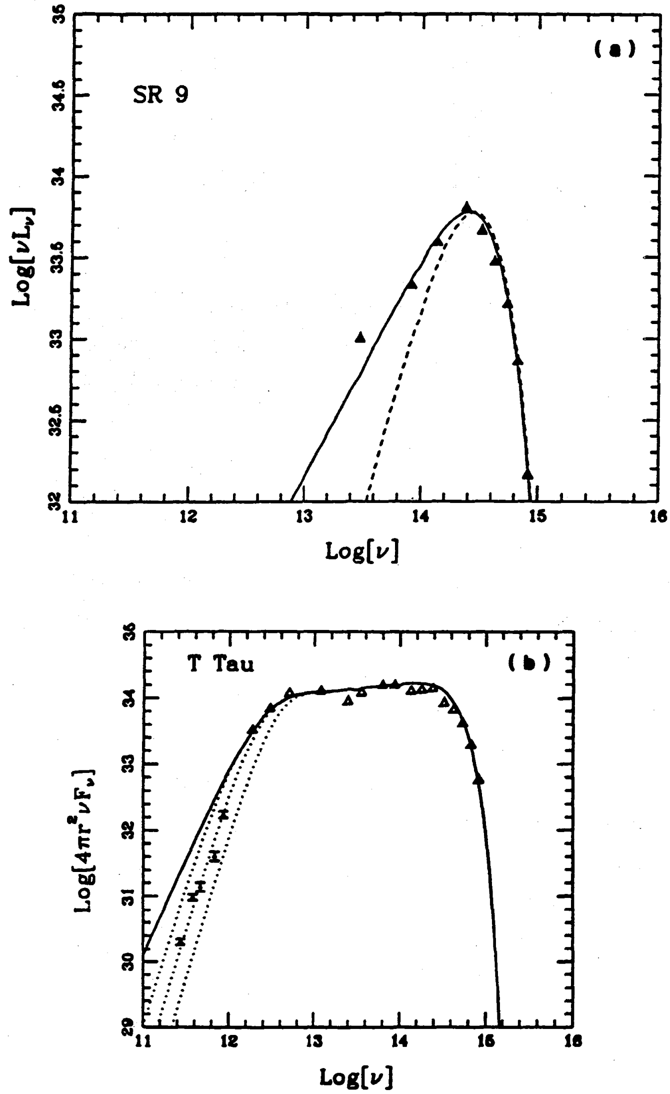

Figure 11. (a) A steep-spectrum T Tauri star, SR 9. The dashed curve shows the expected spectrum for a star alone; the solid, for a star plus a spatially flat, but radially extended, disk that locally reprocesses the starlight that falls on it (figure from Adams et al. 1987) . (b) A flat-spectrum T Tauri star, T Tau itself. The curves show models (from Adams et al. 1988) for a star plus a spatially flat disk, with the latter heated by stellar photons as well as possessing an intrinsic source of luminosity with a temperature distribution, $T_D \propto \varpi^{-p}$, where $p = 0.515$. The disk extends from the stellar surface to a radius $R_D = 120$ AU, the minimum value needed to fit the extent of the flat portion of the infrared data points (triangles). The actual size also cannot be much larger before suffering truncation by the companion. (The system corresponds to a binary, but the model is insensitive to which star has the flat-spectrum disk, or even if both stars have disks.) The disk model in the solid curve has infinite mass and is optically thick at all frequencies; the dotted curves that turn over progressively toward higher frequencies because of declining optical depth have disk masses, respectively, of 1, 0.1, and 0.01 $M_\odot$. Measurements at submillimeter and millimeter frequencies (data points with error bars) can establish the disk mass, provided the dust opacities in the disk correspond with those of the model (figure from Adams et al. 1990).

but by dusty (infalling?) envelopes that contribute significant amounts of far-infrared emission through the reprocessing of starlight.

There exists a difficulty with explaining the most extreme flat-spectrum T Tauri stars (which includes some of the most famous examples in the class: T Tauri itself, HL Tau, DG Tau, etc.) as passive, geometrically flared disks. To reprocess the starlight into the observed amounts of infrared excess would require unrealistic ratios of the dust photospheric height H to disk radius ϖ. Appendix B of Ruden and Pollack (1991; see also Kusaka et al. 1970) gives the temperature of a flared disk illuminated by a central star of radius R_* and effective temperature T_* (at a radius $\varpi >> R_*$) as

$$T_{\rm D}(\varpi) = T_* \left[\frac{2}{3\pi} \left(\frac{R_*}{\varpi} \right)^3 + \frac{1}{2} \left(\frac{R_*}{\varpi} \right)^2 \left(\frac{H}{\varpi} \right) \left(\frac{{\rm d}\ln H}{{\rm d}\ln \varpi} - 1 \right) \right]^{1/4}. \quad (17)$$

Notice that the canonical law $T_{\rm D} \propto \varpi^{-3/4}$ dominates at small radii ϖ (where $H << R_*$), while shallower temperature gradients can arise at larger radii if H varies as a power of ϖ with an exponent s greater than unity (a geometrically flared disk). If the dust does not settle to the midplane but remains suspended in the gas (an inconsistent assumption if reprocessing disks are not turbulent), vertical hydrostatic equilibrium requires $H \propto \Omega^{-1} T_{\rm D}^{1/2}$, so that $s = 9/7$ and $p = 3/7$ in the outer parts of a self-consistent, Keplerian, reprocessing disk. Such a disk typically requires $H/\varpi \gtrsim 1/3$ at large ϖ to produce an appreciable departure from the canonical case (Kenyon and Hartmann 1987). A circumstellar layer of dust with such an aspect ratio would conflict, however, with the observed fact that unobscured T Tauri stars are an order of magnitude more numerous than heavily extincted sources (in the Taurus molecular cloud). Moreover, for sources not seen *through* the disk, since $n = +4/3$ for $p = 3/4$ but $n = -2/3$ for $p = 3/7$, the spectrum would *fall* from the near- to mid-infrared but then *rise* toward the far infrared. Reprocessing of starlight in either flared disks or warped disks cannot generally reproduce the *single* power-law index, $n = 0$ from the near infrared to the far infrared, observed for the most extreme flat-spectrum sources.

Binarity is known to add complications in the case of T Tauri, and other sources of confusion may contribute to the observed departure from the expected spectral slope of $n = 4/3$. Taking the point of view, however, that physical significance could be assigned to the frequent occurrence of the single exponent $n = 0$, which neither viscous accretion nor passive reprocessing adequately explains, Adams et al. (1988) speculated that the flat-spectrum T Tauri stars must have *active* disks where the processes of transport of mass (inward) and angular momentum (outward) involve *nonlocal* mechanisms. Adams et al. (1989; see also Chapter by Adams and Lin) worked out one such mechanism: the generation of one-armed spiral density waves ($m = 1$ disturbances) by eccentric gravitational instabilities that involve the displacement of both star and disk from the center of mass of the system. Shu et al. (1990) gave an analytic demonstration, under a restrictive set of circumstances, that such modal eccentric instabilities would arise whenever the disk mass as a fraction of the whole (star plus disk) exceeds $3/4\pi \sim 0.24$. They also pointed out that the temperature law $T_D \propto \varpi^{-1/2}$ plays a special role in the linear theory of $m = 1$ disturbances in that this distribution keeps the fractional amplitude of the wave (ratio of disturbance to unperturbed surface densities) constant throughout much of the disk. As the fractional amplitude in a full theory provides a measure of the nonlinearity of the wave, Shu et al. (1990) speculated that nonlinear propagation and dissipation of excited $m = 1$ waves would provide the feedback that accounts for the temperature law $T_D \propto \varpi^{-1/2}$ in the first place. Ostriker and Shu (work in progress) are pursuing the correct nonlinear derivation of this result, and they are also investigating the possibility that stellar companions in binary systems may provide an alternative source for the excitation of $m = 1$ density waves.

We find it informative to comment on the basic reason why either low-freqeuncy unstable normal modes or fairly distant companions can provide, in principle, a redistribution of mass and angular momentum of a different qualitative nature than viscous accretion. A disturbance which rotates at a uniform angular speed Ω_p (see Sec. VIII) contains excess energy E_{wave} and angular momentum J_{wave} above the unperturbed values. These are related to each other by the formula (cf. Lynden-Bell and Kalnajs 1972),

$$E_{wave} = \Omega_p J_{wave}. \tag{18}$$

The same proportionality applies *locally* to the wave energy densities $\mathcal{E}_{wave}$ and angular-momentum densities $\mathcal{J}_{wave}$ (per unit area),

$$\mathcal{E}_{wave} = \Omega_p \mathcal{J}_{wave} \tag{19}$$

if the spiral waves are tightly wrapped (Lin and Shu 1964; Toomre 1969; Shu 1970), as they inevitably are for disturbances in most of a radially extended nebular disk (cf. Adams et al. 1989; Shu et al. 1990). Moreover, $\mathcal{E}_{wave}$ and $\mathcal{J}_{wave}$ are both positive or both negative depending on whether the wave rotates faster or slower than the matter; i.e., $\mathcal{E}_{wave}$ and $\mathcal{J}_{wave} > 0$ outside corotation, $\Omega_p > \Omega$, whereas $\mathcal{E}_{wave}$ and $\mathcal{J}_{wave} < 0$ inside corotation, $\Omega_p < \Omega$, where $\Omega(\varpi)$ represents the local angular speed of the disk.

On the other hand, in order for a gas element to remain in *circular* orbits when some external torque changes its angular momentum at a rate $\dot{J}$, its energy must change at a rate

$$\dot{E}_{circ} = \Omega \dot{J}. \tag{20}$$

When spiral density waves inside (outside) corotation dissipate they give a negative (positive) amount of angular momentum to the local gas, causing that gas to move to a radially smaller (larger) orbit. However, the energy lost (gained) by the gas element (a factor of Ω_p times the angular momentum lost [gained]) in the process of damping the wave will generally be inappropriate for the gas to remain in a circular orbit with the modified angular momentum. The excess orbital energy will be converted to heat at a rate

$$\dot{H} = (\Omega_p - \Omega)\dot{J} \tag{21}$$

when the orbit of the gas element circularizes by hydrodynamic dissipative mechanisms.

Note that as the gas loses (gains) angular momentum inside (outside) corotation, the rate of conversion of energy into heat, $\dot{H}$, is everywhere positive. Note also that deep inside the disk where $\Omega >> \Omega_p$, the rate of heat gain is very large per unit rate of angular momentum transferred. In other words, although a slowly rotating disturbance in the outer part of the disk (either a normal mode or a companion star) cannot directly supply much energy to the inner disk (only a factor $\Omega_p \times \dot{J}$), it potentially comprises a huge reservoir for dumping the angular momentum of the system. Removing angular momentum from the part of the disk inside of corotation (by direct near-resonant coupling in $m = 1$ density waves) results in accretion. When the basic orbits cannot become very noncircular (because tightly wrapped waves cannot yield highly eccentric orbits without producing either radiative damping or streamline crossing and dissipative shocks), the heat released by this nonviscous form of accretion will provide, we believe, the luminosity excess that one sees in the flat-spectrum T Tauri stars.

In the inner portions of the disk, where $\Omega >> \Omega_p$, but still much beyond the star-disk boundary layer, we note that Eq. (21) for the disk heat input differs only by a factor of 3 from the corresponding formula given by viscous accretion-disk theory. (A much bigger difference exists beyond the corotation circle, where $\Omega_p >> \Omega$, but the physical dimensions of realistic disks may not extend to regions where the wave energy much dominates over the accretion energy.) The ability of wave transport (perhaps) to produce flat-spectrum sources may then depend critically on the fact that $m = 1$ wave dissipation yields a nonuniform rate of local disk accretion $\dot{M}_{acc}$, with $\dot{M}_{acc} \propto \varpi$ needed to produce a flat-spectrum source.

It might be argued that viscous-accretion disk models could also accomodate flat-spectrum sources if we allow ourselves the luxury of assuming $\dot{M}_{acc} \propto \varpi$. Hartmann and Kenyon (1988) dismissed this possibility on the grounds that it would require mass accretion rates in the outermost parts of a disk three or more orders of magnitude higher than the innermost regions. Compare the required behavior with theoretical models of viscous accretion. Independent of the details for the specification of the viscosity, the diffusive nature of the viscous process (if stable) inevitably yields in the relevant portions of the disk, a long-term evolution that satisfies the approximate condition, $\dot{M}_{acc}$ = a spatial constant (cf. Ruden and Lin 1986; Ruden and Pollack 1991). The same result need not apply to accretion disks driven by global instabilities. Indeed, the dynamics of quasi-steady wave generation, propagation, and dissipation will generally *not* correspond to the condition of steady-state accretion.

The implication that the (self-gravitating) disks in flat-spectrum T Tauri stars have much larger accretion rates in their outermost parts (say, $\sim 10^{-5}$ $M_\odot$ yr^{-1}) than their innermost parts (say, 10^{-8} to 10^{-7} $M_\odot$ yr^{-1}) may have important observational consequences. In particular, the tendency to pile up matter at a bottle-neck at the smallest radii leads, for a temperature distribution fixed by wave dissipation independent of the surface density distribution (see Shu et al. 1990), to a local lowering of Toomre's (1964) Q parameter. If Q falls sufficiently, violent $m > 1$ gravitational instabilities may make an appearance (see, e.g., Papaloizou and Savonije 1991; Tohline and Hachisu 1990). Such modes act over a much more restricted radial range than $m = 1$ modes, but they may induce a sudden release of the piled-up material onto the central star. Two very different gravitational modes of mass transport may therefore act: $m = 1$ waves, driven by distant orbiting companions or mild instabilities characteristic of the outer disk; and $m > 1$ instabilities with high growth rates characteristic of the inner disk. The relative ineffectiveness of the first mechanism for inducing mass transport at small ϖ leads to large amounts of matter caught in the "throat"; the clearing of this material by the second mechanism may then yield the sporadic "coughs" that Hartmann et al. identify in their chapter as FU Orionis outbursts. Much more work would have to be performed on this scenario, however, before we could characterize it as more than an interesting speculation.

VIII. BINARY STARS AND PLANETARY SYSTEMS

In their studies, Adams et al. (1989) and Shu et al. (1990) suggested that the existence of binary systems among the T Tauri stars may itself represent a manifestation of the prior action of $m = 1$ gravitational instabilities (see Chapters by Bodenheimer et al., and by Adams and Lin). Recent numerical simulations using smooth-particle-hydrodynamics codes (Lattanzio and Monaghan 1991; Bastien et al. 1991; F. C. Adams and W. Benz, personal communication) indicate that the formation of companion objects by non-axisymmetric disturbances (in particular, $m = 1$) in massive nebular disks may indeed be a viable physical process. Nevertheless, because the actual accumulation of mass occurs, at least in some systems, via *gradual* infall of a single molecular cloud core (cf. Secs. IV and V), we should not directly apply the results of simulations of fully accumulated stars plus disks to the actual binary-formation process. We believe it much more likely that *both* stars start off small (perhaps as the result of a runaway $m = 1$ gravitational instability), that the two protostars separately acquire inner accretion disks, that they may also jointly acquire a circumbinary disk, and that buildup of the entire system occurs substantially while the system lies deeply embedded within a common infalling envelope. Haro 6–10 and IRAS 16293–2422 may represent examples of such systems (see Zinnecker 1989*a*; G. A. Blake and J. E. Carlstrom, personal communication; for a review of alternative formation scenarios, see Pringle 1991).

An interesting aspect of the realization that binary-star formation may occur in the presence of circumstellar, or even circumbinary, disks concerns the eccentricity of the resulting orbits. Duquennoy and Mayor (1991) and, especially, Mathieu et al. (1989) have emphasized the important observational point that most binaries are probably born with eccentric orbits, and that the shorter-period ones only later become circularized by stellar tidal processes. This state of affairs should be contrasted with that applicable to planetary systems, where current belief holds that most planets are born with nearly circular orbits. How do we reconcile this difference if we believe that companion stars and planets are both born from disks?

A qualitative answer to this question may come from the work of Goldreich and Tremaine (1980) on the interactions of satellites and disks (see also Ward 1986; Chapter by Lin and Papaloizou). For satellite orbits of small eccentricity e and inclination i, a Fourier expansion of the potential associated with the satellite's gravity introduces a series of disturbances that have wave frequency ω given by (cf., e.g., Shu 1984)

$$\omega = m\Omega_s \pm n\kappa_s \pm p\mu_s \tag{22}$$

where m, n, and p are positive integers, with κ_s and μ_s being the epicyclic and vertical oscillation frequencies, respectively, when we expand the satellite's eccentric and inclined orbit about the circular orbit with angular rotation speed Ω_s. The integer m yields the angular dependence $\propto e^{-im\varphi}$ of the disturbance potential, whereas n and p represent the number of powers of e and $\sin i$ that enter the coefficient of the specific term in the series expansion. We refer to the terms corresponding to $n = p = 0$, which can arise even if the satellite orbit is circular, as being of *lowest order*; higher-order terms depend on the orbit having nonzero eccentricity e or inclination i.

As the time dependence of the disturbance potential $\propto e^{i\omega t}$ and the angular dependence $\propto e^{-im\varphi}$, the pattern speed of the disturbance equals

$$\Omega_p = \omega/m. \tag{23}$$

In a steady state (including dissipation), the satellite potential excites a density response in the disk that remains stationary in a frame that rotates at the pattern speed Ω_p. This response achieves its highest values at various resonance locations. *Corotation resonances* occur at radii ϖ where

$$\Omega(\varpi) = \Omega_p \tag{24}$$

Lindblad resonances occur where

$$\Omega(\varpi) \pm \frac{1}{m}\kappa(\varpi) = \Omega_p \tag{25}$$

and *vertical resonances* occur where

$$\Omega(\varpi) \pm \frac{1}{m}\mu(\varpi) = \Omega_p \tag{26}$$

with $\Omega(\varpi)$, $\kappa(\varpi)$, and $\mu(\varpi)$ representing, respectively, the circular, epicyclic, and vertical oscillation frequencies of the local gas. *Inner* resonances occur where the minus sign is chosen in Eqs. (25) and (26); *outer* resonances, where the plus sign is chosen. For a nonaxisymmetric disturbance with given pattern speed Ω_p, we generally encounter in succession an inner resonance, a corotation circle, and an outer resonance, as we move radially outwards through the disk. For a Keplerian disk, all three frequencies are equal, $\Omega = \kappa = \mu$, and special significance attaches to $m = 1$ disturbances because the entire disk can be in near-inner Lindblad or vertical resonance with a distant companion (where $\Omega_p \sim 0$).

Consider now the gravitational back-reaction of the density response in the disk on the exciting satellite. The gravitational potential associated with the density response, being nonaxisymmetric and time dependent in an inertial frame (but time independent in a frame that rotates at Ω_p), conserves neither the orbital angular momentum J_s nor the orbital energy E_s of the satellite; however, the combination $E_s - \Omega_p J_s$ = Jacobi's constant *is* conserved. As a consequence, the back-reaction from the disk pumps energy and angular momentum into the orbit of the satellite in the ratio Ω_p, $\dot{E}_s = \Omega_p \dot{J}_s$, with $\dot{J}_s$ equal to minus the torque on the disk. For the strongest resonances (corotation or Lindblad), where $n = p = 0$, Eqs. (22) and (23) state that $\Omega_p = \Omega_s$; consequently to zeroth order in e (and $\sin i$), these strong resonances contribute energy and angular momentum to the satellite in a ratio as to keep circular orbits circular. To linear order in e^2, however, the strongest resonances change the satellite's eccentricity at rates comparable to the weaker resonances that have $n = 1$ and $p = 0$. The former always lead to eccentricity damping (growth) for inner (outer) Lindblad resonances; the latter contribute to eccentricity changes depending on the sign in $\Omega_p = \Omega_s \pm \kappa_s/m$ and on whether the Lindblad resonances are inner or outer ones. The torque exerted by corotation resonances (and therefore their contribution to eccentricity changes) depends on the sign of the gradient of the vorticity divided by the surface density.

In an untruncated Keplerian disk of uniform surface density, Goldreich and Tremaine (1980) demonstrate that the sum of the effects of all resonances has a slight dominance of eccentricity-damping resonances over eccentricity-exciting ones (see also Ward 1986 for estimates of the effects of surface-density gradients). This presumably explains why planets embedded in nebular disks acquire nearly circular orbits, although gap opening adds to uncertainties (see Lin and Papaloizou 1986*a*, *b*). The ability of strong resonances to open gaps is observed in the Saturnian ring-moonlet system (see the discussion of Cuzzi et al. 1984). For a massive companion such as another star, however, not only may the companion be born with an initial eccentricity, but wide gaps opened up on either side of the secondary star may lead to the disappearance of the strongest $m \neq 1$ resonances. The remaining resonances may amplify orbital eccenticities under some circumstances; for example, Artymowicz et al. (1991) find eccentricity growth in a numerical simulation of a binary surrounded by a circumbinary disk that is dominated by a $m = 2$ outer-Lindblad resonance with $\Omega_p = \Omega_s - \kappa_s/2$. On the other hand, the system GW Ori appears to be a binary with a circular orbit and with circumstellar and circumbinary disks (Mathieu et al. 1991). The inferred gap surrounding the companion eliminates the stronger $m \neq 1$ resonances; if true, the circumstellar disk may provide a circularizing influence on the binary orbit through the action of $m = 1$ inner-Lindblad near-resonances (Ostriker et al. 1991). In any case, if a companion star acquires large eccenticities, it may have a high efficiency for sweeping up the matter of the disk in which it is embedded.

In summary, the orbital characteristics of newly formed planets and companion stars might qualitatively differ from one another because the two types of objects originate by completely different processes (cf. Pringle [1991] for this point of view). On the other hand, even if they both descend from disks, we see that good mechanistic reasons exist why binary stars might acquire appreciably eccentric orbits, whereas planets usually do not. The difference in outcomes—binary stars or planets—might then depend on differences in initial conditions, namely, the total amount of angular momentum J_{tot} contained in the collapsing system. If J_{tot} is small, a substantial central mass could accumulate by direct infall before disk formation occurs. The stabilizing influence of a massive center of attraction may then prevent any runaway gravitational instabilties in the disk, leaving only the possibility of planetary condensation by chemical or physical means. If J_{tot} is large, the ratio of disk to star mass may reach order unity before the central object has grown sufficiently to prevent the gravitational fragmentation of the disk to form another body on an eccentric orbit. In this case, continued infall builds up a binary-star system.

IX. THE ORIGIN OF STELLAR AND PLANETARY MASSES

As the primary transformers of the simple elements that emerge from the big bang, stars constitute the fundamental agents of change in the visible universe. Of all its properties, a star's mass most affects its main-sequence appearance and later evolution. Given that giant molecular clouds, the raw material for forming stars, contain much more mass than required to make the final products, the most basic problem posed for theories of stellar origin concerns, then, how these objects acquire the masses that they do. At one time, astronomers sought for an answer to this important question in terms of a picture of hierarchical fragmentation, but severe doubt about this possibility has been cast by the realization that giant molecular clouds are not collapsing dynamically and have, in fact, generally a very low efficiency for stellar genesis (Zuckerman and Evans 1974).

The discovery that massive outflows ubiquitously accompany the birth pangs of stars has led to an alternative explanation, that forming stars might help to define their own masses. From this point of view, star formation represents basically an accretion process (a theoretical point of view with modern origins in the work of Larson [1969]), with the mass of the star continuing to build up until a powerful enough wind turns on to reverse the infall (first, over the poles, and later, even in the equatorial directions) and shut off the accretion. Indeed, given the basic angular-momentum difficulty faced by any condensing and rotating object which needs to contract by several orders of magnitude in linear scale, perhaps theorists should have anticipated all along that stars would not be able to form without simultaneously suffering tremendous mass loss. As it is, however, theorists from Laplace onwards managed to predict, in one form or another, only three of the four stages of star formation depicted schematically in Fig. 1.

A parallel situation holds for the issue of planetary masses, which we know today probably also results from a process of accumulation. At one time theorists sought to skirt the issue by postulating that the primitive solar nebula had only the amount of mass, when reconstituted to solar composition, needed to build the known planets. The planets then incorporated the masses that they did, essentially because that was all the material there was. (We exaggerate the degree to which "input equals output" in the so-called "minimum-mass model" because theorists still had the chore of explaining why the solar system formed with *nine* major planets, and not one, or a hundred.)

The modern view suggests that the birth of a star must indeed almost always proceed with the accompanying appearance of a disk. However, it also claims that the masses of the growing star and disk are unlikely to have a ratio of anything like 100 to 1. In the early stages, at least, the solar nebula may have had a comparable mass to the proto-Sun. Such a situation places sharp focus on the question how planets acquire the masses that they do if most disks have available more mass than needed to make the known planets.

A partial answer may have been supplied by the suggestion by Lin and Papaloizou (1985,1986*a*; see also their Chapter) that *tidal truncation* limits the maximum amount of mass acquirable by a *giant* planet. Tidal truncation works because a circularly orbiting body has a basic tendency to want to speed up more slowly rotating material exterior to its own orbit, and to slow down more rapidly rotating material interior to it. This tendency to give angular momentum to the matter on the outside and to remove it from matter on the inside (as already discussed for resonant interactions in Sec. VIII) secularly pushes gas on either side of a protoplanet away from its radial position, opening up a gap centered on the body's orbit. Viscous diffusion tends to fill in this gap, and for a true gap to form, the tidal torques have to dominate, leading, for a given level of turbulent viscosity, to a minimum protoplanetary mass that can open up a gap larger than its own physical size. (See Chapter by Lin and Papaloizou on the additional criterion that arises in a *gaseous* disk because the gap width must also exceed the vertical pressure scale-height, if gas pressure is not to close the gap.)

The observational knowledge that the tidal force of moonlets can clear analogous gaps in planetary rings (see, e.g., Showalter 1991) shows that this mechanism does have empirical validity; however, the simplest formulations of gap opening and shepherding theories do not yet yield good quantitative comparisons with observations of planetary rings and their moonlets (see, e.g., Goldreich and Porco 1987). Moreover, this mechanism applied to the solar nebula will not work to define the masses of the *terrestrial* planets, nor does it explain what happens to the excess disk material that does not go into planets. As already discussed in Sec. I, the fate of the solid debris may not pose a great problem given the demonstration that most of the space between the major planets yields secularly unstable orbits (see Chapter by Duncan and Quinn); however, the dispersal of the remnant gaseous component remains a serious problem, especially if massive protostellar and pre-main-sequence winds are driven by disk accretion.

A possible resolution of the problem may rest with Elmegreen's (1978*b*,-1979*c*) proposal that stellar winds erode disks, not by blowing the gas away (Cameron 1973), but by inducing turbulence in the wind-disk interface that enhances the *inward* viscous transport of matter. If, in turn, disk accretion drives powerful stellar winds, the process could continue to operate, in principle, until the disk has completely emptied onto the star. Indeed, the existence of a feedback loop—disk accretion drives wind → wind drives disk accretion, etc.—together with time delays, leads to the possibility of instabilities and limit cycles (see Bell et al. [1991] for an example of a negative feedback cycle leading to chaotic oscillations). Such feedback loops—modulated perhaps by magnetic cycles on the star itself—yield yet another possible mechanism for FU Orionis outbursts.

Finally, we note that disk accretion can be dammed by the formation of giant planets and the appearance of gaps across which matter does not flow. The removal of the gas from the outer regions of dammed disks may then require slow evaporation by ultraviolet radiation (Sekiya et al. 1980*a*; Hayashi et al. 1985; see review of Ohtsuki and Nakagawa 1988). The abundance of possibilities regarding the dispersal of the remnant gas in nebular disks yields a specific demonstration that we still possess ample problems to wrestle with before *Protostars and Planets IV* is written.

Acknowledgments. This work was supported in part by a NSF grant and in part under the auspices of a special NASA astrophysics theory program which funds a joint Center for Star Formation Studies at NASA-Ames Research Center, U. C. Berkeley, and U. C. Santa Cruz.