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Magnetic Effects in Clouds*

L. Mestel

Research Centre for Theoretical Astrophysics,
University of Sydney, Sydney, N.S.W. 2006, Australia.
Permanent address: Physics and Astronomy Division,
University of Sussex, Falmer,
Brighton BN1 9QH, England.

Abstract

A realistic study of the structure and evolution of an interstellar gas cloud must take cognisance of the flux from the galactic magnetic field threading the cloud. If the non-dimensional mass-to-flux ratio is below a critical value, the forces exerted by the locally distorted field can balance gravity in the two trans-field dimensions, while Alfvénic turbulent motions yield support along the field. A super-critical cloud, collapsing with its flux virtually frozen in, may fragment into sub-condensations following spontaneous flattening along the field. Within a sub-critical molecular cloud, the very low degree of ionisation allows the magnetic forces to redistribute flux through the cloud, so that locally denser regions may become super-critical and condense out of the cloud. The Maxwell stresses also transport angular momentum efficiently from a slowly contracting condensation to the surroundings. If flux leakage remains slow throughout all the pre-opaque phases, the magnetic forces and the associated turbulent motions may shift the ultimate mass spectrum towards the high mass end. Most of the remnant flux may be lost by magnetic buoyancy during the pre-main sequence epoch, so possibly supplying a power source for the T Tauri phenomenon.

1. Magnetism and Gravitation

The galactic magnetic field alters radically the classical problems of the equilibrium, collapse and fragmentation of self-gravitating gas clouds. The maximum mass that can be supported by thermal pressure alone is essentially the Jeans mass $M_J \simeq 2{\cdot}5\, a^3/G^{\frac{3}{2}}\rho^{\frac{1}{2}}$, where ρ is the density and $a = (\mathcal{R}T/\mu)^{\frac{1}{2}}$ is the sound speed in a gas at temperature T and of mean molecular weight μ. In a typical molecular cloud with $\rho = n\, m_H$, $n = 10^{10}$–$10^{11}\,\mathrm{m}^{-3}$, $T = 10$–20 K and $\mu = 2{\cdot}34$, then $M_J \simeq (1$–$10)M_\odot$, to be compared with cloud masses that are $10^4 M_\odot$ or more. It had in fact long been recognised that the forces exerted by flux from the galactic magnetic field trapped within the cloud could be the reason why the clouds do not collapse in a free-fall time and fragment spontaneously into sub-condensations of stellar mass. In fact, when the suggestion was first made by Fermi and others that there should exist a large-scale galactic field, strong enough to mediate the conversion of a fraction of the kinetic energy of clouds into cosmic ray energy, and to align the anisotropic dust grains thought responsible

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for the polarisation of starlight, sceptics objected that such a field would prevent altogether the formation of stars. The implicit reasoning was as follows. Let the idealised local background have uniform density $\rho_0 = n_0 m_H$ and be permeated by a field $\mathbf{B}_0$ which can be taken as uniform if the currents maintaining it flow in circuits large compared with the radius R_0 of the cloud considered. After contracting to radius R, a spherical cloud of mass $M = 4\pi\rho_0 R_0^3/3$ exerts a mean gravitational force density $\simeq(GM/R^2)(\rho_0 R_0^3/R^3) = GM^2/(4\pi/3)R^5$. The same contraction increases the strength of the frozen-in field to $B = B_0 R_0^2/R^2$, and once $R \ll R_0$, the radius of curvature R of the locally distorted field lines implies a current density $j = B/\mu_0 R$ and so an outwardly acting Lorentz force density $\simeq B^2/\mu_0 R = B_0^2 R_0^4/\mu_0 R^5 = (B_0^2/\mu_0 R^5)(3M/4\pi\rho_0)^{\frac{4}{3}}$. Thus in order that self-gravitation can enforce indefinite contraction against the opposing magnetic forces, we apparently require $M > M_B$, where the 'magnetic Jeans mass' for a cold medium is given by

$$\begin{aligned} M_B &\simeq (B_0/\rho_0^{\frac{2}{3}})^3(1/\mu_0 G)^{\frac{3}{2}}(3/4\pi)^{\frac{1}{2}} \\ &= [(B_0/B_{\text{fid}})/(n_0/n_{\text{fid}})^{\frac{2}{3}}]^3 3\times 10^6\, M_\odot\,, \end{aligned} \tag{1}$$

where normalisation is in terms of currently accepted fiducial values $B_{\text{fid}} = 3\times10^{-10}$ T $= 3\times10^{-6}$ G and $n_{\text{fid}} = 10^6$ m^{-3} $= 1$ cm^{-3}. Here M_B can be written $\simeq(4\pi/3)\rho_0\lambda_B^3$ with the 'magnetic Jeans length' $\lambda_B = (B_0/\rho_0)(1/\mu_0 G)^{\frac{1}{2}}(3/4\pi)^{\frac{1}{2}}$.

This argument, leading to a minimum mass $\simeq 3\times10^6\,M_\odot$, is in fact overstated, and with its assumption of isotropic accumulation of gas is rather misleading, since one may reasonably expect the field will tend to channel flow parallel to itself. It is more enlightening to begin by accepting that a roughly spheroidal cloud has formed, oblate about the direction of $\mathbf{B}_0$, of mean density ρ, with semi-major and semi-minor axes (R, z) and so of mass $M = 4\pi\rho R^2 z/3$, and containing trapped flux $F = \pi B R^2 = \pi B_0 R_0^2$. Here B is the mean large-scale field strength within the cloud, and R_0 is now defined as the value of the semi-major axis R at which the frozen-in field is undistorted and so has the background value B_0. Once $R \ll R_0$, the magnetic curvature yields a force density $\simeq B^2/\mu_0 z = F^2/\mu_0\pi^2 R^4 z$, opposing a gravitational force density $\simeq GM\rho/R^2 \simeq GM^2/(4\pi/3)R^4 z$. The ratio of magnetic to gravitational force density then yields the non-dimensional flux/mass parameter, conveniently normalised to

$$f = 2F/\pi^{\frac{3}{2}}(\mu_0 G)^{\frac{1}{2}}\,M\,. \tag{2}$$

For indefinite gravitational collapse to be possible, f must be less than a critical value f_c of order unity, corresponding to a critical mass M_c for a prescribed flux:

$$M > M_c(F) = 2F/\pi^{\frac{3}{2}}(\mu_0 G)^{\frac{1}{2}} f_c\,. \tag{3}$$

This is equivalent to

$$(\mu_0 G)^{\frac{1}{2}}\sigma/B > 2/\pi^{\frac{3}{2}} f_c\,, \tag{4}$$

where $\sigma \doteq M/\pi R^2$ is the projected area density. If clouds form from the galactic background $\mathbf{B}_0$, ρ_0 under strict flux-freezing, one can define an *accumulation length* L_0 along $\mathbf{B}_0$ by

$$M/F \simeq \sigma/B \simeq \rho_0 L_0/B_0 \,, \tag{5}$$

and the condition for the indefinite collapse of the spheroidal cloud in the trans-field directions is

$$\begin{aligned} L_0 > (L_0)_c \;&=\; [B_0/(\mu_0 G)^{\frac{1}{2}}\rho_0](2/\pi^{\frac{3}{2}} f_c) \\ &\approx 700(B_0/B_{\rm fid})/(n_0/n_{\rm fid})\,\text{parsec}\,. \end{aligned} \tag{6}$$

This is again of the order of the magnetic Jeans length λ_B, but the contrast with the argument that assumes isotropic contraction and so leads to (1) is clear, for no length-scale *across* the field enters the criterion (6). A thinner flux tube contains a correspondingly smaller mass, but once the mass in a cylinder of length $>(L_0)_c$ has accumulated down the field into a spherical or oblate body, its self-gravitation will always exceed the mean magnetic force generated during subsequent contraction perpendicular to $\mathbf{B}_0$.

The value $(L_0)_c$ in (6) is of the same order as the minimum Rayleigh–Taylor–Parker instability wavelength (Parker 1969; Mouschovias 1974), so it may be argued that an accumulation length of the order of a kiloparsec is not necessarily an embarrassment. Nevertheless, one is led to wonder whether strict flux freezing always holds during the condensation of clouds from the background galactic medium, or whether sometimes clouds form following a violent input of energy, e.g. from a supernova, with an effective turbulent resistivity yielding a net diffusion of gas across the field (cf. Sweet 1950), so that $(L_0)_c$ becomes a notional length. It can reasonably be argued that the problem of cloud *formation* should be mentally separated from that of the equilibrium, collapse and fragmentation of star-forming clouds.

The discussion has a natural bifurcation at this point:

(1). *Super-critical clouds* have $f < f_c$ or $M > M_c(F)$, and so cannot be prevented from gravitational collapse by magnetic forces alone. Further, as again no length across $\mathbf{B}$ enters the collapse condition (4), a thin flux-tube within a super-critical cloud contains enough mass for a cold, *super-critical oblate fragment* to form following flow of the gas down the field lines. Fragmentation of an already super-critical cloud is thus not forbidden even under strict flux-freezing, but the bodies so forming within a cloud that is moderately super-critical ($f \lesssim f_c$) will themselves be 'strongly magnetic', with magnetic energy close to the gravitational. We return to these questions later.

(2). *Sub-critical clouds* have $f > f_c$. Even if the cloud is cool, weakly turbulent and slowly rotating, indefinite collapse in the two trans-field dimensions is prevented by the magnetic forces that develop as the frozen-in field is distorted. After dissipation of macroscopic kinetic energy, the cloud should settle into a state of equilibrium with magnetic forces providing the bulk of the support against gravity in the two dimensions perpendicular to $\mathbf{B}_0$, and with the length scale along $\mathbf{B}_0$ fixed by the thermal and turbulent energies. The formation of gravitationally bound sub-condensations within a sub-critical cloud depends on flux leakage (see below).

For approximate treatments of the equilibrium that use the Chandrasekhar–Fermi (1953) virial theorems, see Mestel (1965); Strittmatter (1966); Mestel and Paris (1984). For detailed model construction, see Parker (1973, 1974), Mouschovias (1976*a*, 1976*b*); Mestel and Ray (1985); Tomisaka *et al.* (1988); Barker and Mestel (1990). When $M \lesssim M_c(F)$, the virial treatment yields a trans-field radius

$$R \simeq R_0[1 - (M/M_c)^2], \quad R_0 = (F/\pi B_0)^{\frac{1}{2}}, \tag{7}$$

with an associated mean field strength within the cloud

$$B = B_0(R_0/R)^2 \approx (\mu_0 G)^{\frac{1}{2}} M/R^2 \approx (\mu_0 G)^{\frac{1}{2}} \rho z, \tag{8}$$

where factors of order unity are ignored. When $M \ll M_c(F)$, then $R \simeq R_0$ and $B \simeq B_0$: a slight distortion of the background field suffices to generate the forces able to balance gravity.

The prediction of bodies held essentially in magneto-gravitational equilibrium received its first observational support in 1968 when Verschuur succeeded in picking up the Zeeman effect on the 21 cm line. It was gratifying to learn that the kinematic, flux-freezing condition—high B correlated with high n—and the dynamical condition—M close to the virial estimate (8)—were both apparently vindicated. For a resumé of the present-day observational evidence, see the papers by Myers and Goodman and by Troland in Beck *et al.* (1990).

A crucial feature that was not predicted is the typical Doppler width of the molecular cloud lines, which is a clear indication of strong supersonic 'turbulent' velocities v_t providing support along the field lines. In a markedly flattened cloud, the Spitzer–Ledoux (e.g. Mestel 1965) one-dimensional equilibrium condition yields

$$v_t^2 \simeq 2\pi G \rho z^2 \simeq 3GMz/2R^2, \tag{9}$$

whence from (8) and (9)

$$B/(\mu_0 \rho)^{\frac{1}{2}} \simeq v_t. \tag{10}$$

If the turbulence is strong enough to keep the cloud nearly spherical, there are just changes of order unity in the numerical factors. Thus in a cloud with $M \lesssim M_c$, so that B is given by (8), the equilibrium achieved by the lateral adjustment of B and ρ and the longitudinal adjustment of ρ inevitably yields by (10) an Alfvén speed v_A that is close to the turbulent speed v_t: the term 'Alfvénic turbulence' is almost tautological. If v_t is constant ('isothermality'), then this simple discussion suggests that clouds with different (near-critical) parameters f should yield the relation $B \propto \rho^{\frac{1}{2}}$ (independent of the cloud mass). In general, (10) can be rewritten

$$B \propto (\mu_0 G)^{\frac{1}{2}} \rho^{\frac{2}{3}} M^{\frac{1}{3}} [v_t^2/(GM/R)]^{\frac{2}{3}}, \tag{11}$$

so if instead v_t is always a constant fraction of the free-fall speed, (11) predicts the '$\gamma = \frac{4}{3}$' relation $B \propto \rho^{\frac{2}{3}}$ for clouds of nearly the same mass, essentially because by (9) the turbulence then maintains a nearly isotropic relation $z \propto R$ between the different clouds.

A plausible model for the turbulent motions is of a field of torsional Alfvén waves propagating outwards along $\mathbf{B}$ with decaying amplitude (Arons and Max 1975; Zweibel and Josafatsson 1983; and others). The simplest illustrative example (Dewar 1970; Shu *et al.* 1987) has a mean field $\bar{\mathbf{B}} = \bar{B}\hat{\mathbf{z}}$ along which travel transverse waves with amplitudes b, δv, originating near $z = 0$. For $z > 0$

$$\delta\mathbf{B} = \mathbf{b}\sin(kz - \omega t)\exp(-k_1 z), \quad b = (\delta v/v_A)\bar{B}, \tag{12}$$

yielding the time-average of the fluctuating Lorentz force density

$$\mathbf{f} = |\mathbf{b}\exp(-k_1 z)|^2 (k_1/2\mu_0)\hat{\mathbf{z}},$$

acting along $\bar{\mathbf{B}}$. A non-zero damping constant k_1 is essential for a net input of momentum into the cloud, but the associated energy input need not be as heat but may be as macroscopic kinetic energy (e.g. rotational energy). Less idealised models will have the lines of the mean field $\bar{\mathbf{B}}$ curving so as to contribute to the Lorentz force opposing gravity in the lateral directions. Waves travelling along a field of decreasing mean strength yield a net force along $\mathbf{B}$ even without any intrinsic damping, a property exploited in models of Alfvén-wave-driven stellar winds (Belcher 1971; Hollweg 1973). The source of the motions could be waves from the magnetic braking of sub-condensations (cf. below), or winds from proto-stars, or local expanding HII regions from hot stars newly arrived on the main sequence. It is also possible that cool disc-like clouds in magneto-gravitational equilibrium may be spontaneously unstable against the conversion of gravitational into kinetic energy. Although there is evidence for flattened structures in some clouds (e.g. Sargent *et al.* 1988), it appears that usually the turbulent motions are strong enough to yield a longitudinal force of the same order as the lateral, so keeping the clouds moderately oblate. The advantage of Alfvén waves is that in the linear approximation there are no density variations, so reducing the energy dissipation consequent on compression and shock formation. Damping via ion–neutral collisions (cf. Section 2) will occur, but there is a constraint that heat input must not raise the cloud temperature above the observed maximum of 20 K (Zweibel and Josafatsson 1983). The theoretical picture of combined magnetic and turbulent support of molecular clouds is eminently plausible, but proper understanding of the structure, origin, maintenance and decay of the turbulence remains a major problem. The rate of input of wave energy into the clouds is an extra parameter to be included in both micro- and macro-studies of star formation in different types of galaxy.

Models which assume clouds to form from the intercloud medium under strict flux freezing will necessarily have all the cloud field lines 'infinite'. Equilibria which depend on an outwardly acting lateral magnetic force at all points likewise require that no cloud field lines have detached to form O- and X-type neutral points (e.g. Mestel and Strittmatter 1967). As noted by McKee and colleagues (personal communication), it appears difficult for such models to yield mean fields

$\bar{B}$ within the clouds that are much above $4B_0$, a factor 2 or more below the observed mean field in some cases. The reason appears to be that the magnetic curvature force density exerted by undetached field lines would become too large for the gravitational force density in the low density outer regions of clouds (Mestel 1966; Mouschovias 1976*a*, 1976*b*). Relaxing the conditions on the field line topology does allow $\bar{B}/B_0$ to increase (Barker and Mestel, in preparation), but at the cost of requiring that support in some of the outer regions of the cloud must be non-magnetic (turbulent, centrifugal).

2. Flux Leakage

All main sequence stars are magnetically 'weak': a plausible inward extrapolation of observed fields—even Babcock's record of $3 \cdot 4 \times 10^4$ G—yields a ratio of magnetic to gravitational energies $\simeq (F/(\mu_0 G)^{\frac{1}{2}} M)^2 \ll 1$. It is quite impossible to account for this by extending the picture of anisotropic accumulation under flux freezing—the length L_0 would have to be 40–1000 kpc. It is clear that *at some epoch*, the bulk of the primeval flux within the mass that will ultimately reach the main sequence must leak out, so that **B** ceases to be important for the overall dynamics, but the present evidence is that this does not occur until the later stages of star formation (cf. Section 4). In the early, molecular cloud phase, flux leakage appears to be 'slow', occurring in a time scale τ_d that, while much shorter than the Hubble time, is still longer than the instantaneous free-fall time τ_f. The magnetic forces remain dynamically important, but relaxation of the strict flux-freezing constraint allows the cloud to evolve.

'Ambipolar diffusion' or 'plasma drift' in a lightly ionised cloud is simply pictured as the drift of the ionised component ('plasma') and the inductively coupled **B** field relative to the neutral bulk of the gas. The magnetic force acts directly only on the plasma, while the gravitational force is felt essentially by the neutrals. The magnetic force causes the plasma to drift at a rate $\mathbf{v}_d$ limited by plasma–neutral friction, yielding

$$\mathbf{v}_d = (\mathbf{j} \times \mathbf{B})/\alpha(n_i/n)\rho^2 \,, \tag{13}$$

where n_i and n are the respective number densities of ions and neutrals. The frictional coefficient $\alpha = \langle \sigma_{in}(v_T)_n \rangle / m_n$, where m_n is a mean neutral particle mass, σ_{in} the momentum transfer collision cross section, and $\langle \sigma_{in}(v_T)_n \rangle$ the average over the thermal velocity distribution of the neutral particles; whence the number of collisions per second felt by an atomic or molecular ion (more massive than a neutral hydrogen or helium atom or molecule) is $1/\tau_{in} = n\langle \sigma_{in}(v_T)_n \rangle$.

Consider now an oblate cloud in approximate magneto-gravitational equilibrium in two dimensions, as discussed above. The same magnetic forces that hold it up are responsible for evolution of the cloud through flux leakage. A typical leakage time averaged over the cloud is

$$\begin{aligned} \tau_d \simeq \frac{R}{v_d} &\simeq \alpha(n_i/n)(M^2/F^2)(9\mu_0/16)(R/z) \\ &\simeq 10^{14}(n_i/n)(R/z)\ \mathrm{yr} \,, \end{aligned} \tag{14}$$

if the cloud is near-critical. It is seen that the value of the ratio n_i/n is crucial. In the original discussion (Mestel and Spitzer 1956), the process was applied to HI clouds with dust obscuring the galactic UV ionising radiation. It was estimated that attachment to dust grains and dissociative recombination would reduce n_i/n rapidly to such a low value that τ_d would in fact be less than the free-fall time, so that before even a super-critical cloud could begin its collapse, the bulk of the flux threading it would leak out to join the local galactic field, with the excess energy in the cloud field being thermalised through the ion–neutral friction as the field lines straighten. Subsequent estimates using a revised value for σ_{in} (Osterbrock 1961) and taking account of ionisation by the galactic cosmic ray flux showed that in typical HI clouds τ_d is certainly much longer than τ_f and may exceed the Hubble time. In molecular clouds n_i/n is significantly lower, but recent detailed discussions (Norman and Heyvaerts 1985; Nakano and Umebayashi 1986*a*, 1986*b*; Nakano 1988) suggest that τ_d is still longer than τ_f by a factor 10 or more during the whole of the molecular cloud phase. Further, as pointed out by Nakano (1976, 1983) (see also Mouschovias and Morton 1991), diffusion will tend to be non-uniform within a cloud, because of the dependence on n_i/n, which decreases with increasing n roughly like $1/n^{\frac{1}{2}}$. A cloud core of radius R' and density ρ' held in magneto-gravitational equilibrium has magnetic force density $|\mathbf{j}\times\mathbf{B}| \simeq G\rho'^2 R'$, yielding by (13) a local diffusion time $R'/v_d \simeq (\alpha/G)(n_i/n)' \propto 1/(n')^{\frac{1}{2}}$. Thus a core initially denser than the mean over the cloud will lose flux to the rest of the cloud, and will steadily adjust to a new state of magneto-gravitational equilibrium at still higher density, leading to a further reduction in n_i/n and so to further differential flux leakage within the cloud. The picture is essentially as originally discussed but on a smaller scale, with the cloud core steadily losing flux to the rest of the cloud instead of to the intercloud medium. Thus Nakano's argument shows how an initially sub-critical core can become super-critical and go over into collapse, with its mass determined essentially by ambipolar diffusion, while the bulk of the cloud remains magnetically supported. This fragmentation process is thus closer in spirit to the old Jeans picture than to the Hoyle (1953) picture, in which sub-condensations have to separate out of a collapsing cloud (cf. Section 4). However, once the first generation of stars has formed within the cloud, the local ionisation level may be raised to a level which severely limits further diffusion, so that fragmentation monitored by flux leakage effectively ceases.

3. Angular Momentum

A great advantage of flux leakage being slow even in a lightly ionised molecular cloud is that the strong field can transport away efficiently most of the embarrassingly high angular momentum that a cloud or a sub-condensation would inherit from both the galactic rotation and the interstellar turbulence. The transport can be from the cloud to the inter-cloud medium, or from a cloud core to the outer parts of the cloud. A contracting central mass initially in corotation with its surroundings will tend to spin up, so generating an Alfvén wave which tries to restore corotation. An estimate for the characteristic braking time τ_b is given by the time of travel of the waves through a surrounding region of density ρ_0 with the same moment of inertia as the central mass. The results are to some extent model-dependent (Mouschovias and Paleologou 1979, 1980; Mestel and Paris 1984), but for a near-critical mass with roughly *radial* external field lines,

$$\tau_b = (M/F)(\mu_0/\rho_0)^{\frac{1}{2}}(\rho_0/\rho)^{\frac{2}{5}} . \tag{15}$$

Compared with the instantaneous free-fall time τ_f,

$$\tau_b/\tau_f \simeq (M/M_c)(\rho/\rho_0)^{\frac{1}{10}} \simeq (M/M_c) , \tag{16}$$

since the dependence on the density ratio (ρ/ρ_0) is so weak. Thus the ratio M/M_c not only discriminates between sub-critical and super-critical clouds, but it determines whether the braking time is less than or greater than the free-fall time. This suggests strongly that the rotational history of a mass may be very sensitive to modest variations in the value of the parameter $f = 2F/\pi^{\frac{3}{2}}(\mu_0 G)^{\frac{1}{2}}M$. A cloud forming as a super-critical body $[f < f_c,\ M > M_c(F)]$ with the expected high angular momentum will contract in approximate magneto-centrifugo-gravitational equilibrium, at the rate fixed by the braking time τ_b (as long as the diffusion time $\tau_d > \tau_b$). An initially sub-critical cloud $[M < M_c(F)]$ will achieve rapid corotation with the surroundings, but the degree of contraction is limited; however, under slow flux leakage M_c decreases and the cloud continues to contract quasi-statically in the time $\tau_d \gg \tau_f > \tau_b$, with corotation being maintained and the ratio of centrifugal force to gravity steadily declining. The contraction becomes dynamic when M_c has become close to M, and once $M > M_c$ braking becomes ineffective, and the cloud again collapses in free-fall and spins up. However, because so much angular momentum is removed during the earlier near-corotation phases, magneto-centrifugo-gravitational equilibrium is now not achieved until much higher densities than for a cloud which begins its life as super-critical. Detailed analytical and numerical work (Gillis *et al.* 1974, 1979; Mestel and Paris 1979, 1984; Nakano 1989) confirms and refines these qualitative conclusions. There is also some tentative observational support, e.g. a correlation of higher Ω values (40–50 times the usual values) in cloud cores with masses that are higher than the norm and presumed to be super-critical (see references in Shu *et al.* 1987).

It should be noted that even with the most efficient magnetic braking, as in an initially sub-critical mass, for which the bulk of the 'angular momentum problem' is resolved before the body becomes super-critical, further collapse under angular momentum conservation is likely to be halted by the growth of centrifugal force at densities within the opaque domain but well below main sequence densities. Enough angular momentum is left to allow the formation of dense disc-like bodies, the precursors of binary systems and solar systems. A complete theory, able to predict the frequency distribution of binary systems, would need to incorporate subtle details such as field line detachment, which may sometimes lower the efficiency of braking.

The apparent ability of the magnetic field to cope with the angular momentum problem is striking. However, until far more is known about angular momentum transport in non-magnetic systems, e.g. via spiral shocks (Spruit 1987), it is premature to claim that star formation cannot occur without a large-scale magnetic field being present. This is an important question, since if (as seems likely) galactic magnetic fields are not essentially cosmological but are built up by dynamo action over times exceeding typical galactic rotation times, then it may be necessary to account for the formation of the first proto-stars by processes not involving magnetic fields.

4. Masses of the Proto-stars

The presence of the magnetic field in molecular clouds and its likely persistence up to and perhaps beyond the onset of high opacity inevitably affects the discussion of the expected mass spectrum of stars reaching the main sequence. The prediction of contrasting rotational histories for bodies that are initially sub-critical or super-critical in M/F must also be incorporated, perhaps leading to a picture with 'bimodal star formation' (Shu *et al.* 1987; Lizano and Shu 1989; cf. Herbig 1962). Consider first a gravitationally contracting, slowly rotating mass which for convenience is called a 'cloud', even if it has emerged by the diffusion of flux out of the core of a massive molecular cloud (cf. Section 2). As noted earlier, further fragmentation of a super-critical cloud ($f < f_c$) is not forbidden by the magnetic forces even if the flux is now frozen in, provided flow down the field lines can occur. Application of the virial theorem (Campbell and Mestel 1987) shows that within such a cloud, supposed flattened along the field to a thickness fixed by the turbulent velocity v_t (treated as just a greatly enhanced sound speed), the minimum possible gravitationally bound fragment has mass

$$M_{\rm min} \simeq (15/2)(2v_t^2 R/3GM)^2 M/(1 - f^2/f_c^2)\,, \tag{17}$$

where R is the instantaneous trans-field radius of the cloud of mass M. As anticipated in Section 1, the strength of the turbulence at each epoch is crucial. In a non-turbulent cloud, contracting isothermally at the low temperature of a typical molecular cloud, $M_{\rm min}$ from (17) steadily decreases, and unless f is very close to f_c, it will not differ much from the ordinary Jeans mass at the density and temperature considered (Mestel 1965). If as in Hoyle's (1953) original model, fragmentation occurs with maximum efficiency, ceasing only when the optical depth through the last fragment approaches unity, then the predicted typical masses are embarrassingly low (e.g. Mestel and Spitzer 1956; Low and Lynden-Bell 1976). At the other extreme, if input of energy keeps v_t^2 so high that the contraction of the cloud is nearly isotropic, then fragmentation is forbidden.

It should be emphasised that in the fragmentation problem the virial condition is only a censor. Even in the non-magnetic problem (cf. the contribution by Lattanzio and Monaghan 1992; present issue p. 559), it is a non-trivial dynamical problem to demonstrate that a mass well above the instantaneous virial limit can actually separate out of a gravitationally collapsing background, while the opacity limit may grossly underestimate the masses of the final fragments. It is however at least plausible that some fragmentation of a spontaneously flattening, collapsing super-critical magnetic cloud does occur. The important question is whether the presence of a strong **B** field so inhibits the dissipation of the turbulence that v_t in (17) is kept high, and the mass spectrum consequently shifted sharply towards higher masses. Such a result would be of great interest, not least for the galactic dynamo problem. In a rapidly rotating magnetic cloud, the relative orientation of the magnetic and angular momentum vectors is significant. If they are nearly parallel, flow down the field lines can again occur, and in fact the strong centrifugal forces of the orbital motion inhibit lateral collapse of the cloud and assist fragmentation. When the two vectors are more nearly orthogonal, flow of gas down the field lines will be inhibited, and fragmentation of a super-critical mass depends on the competition between flux leakage and angular momentum transport (Mestel 1965; Campbell and Mestel 1987).

The discussion has exploited the slowness of ambipolar diffusion, which relaxes the constraints imposed by strict flux freezing, while leaving the magnetic field dynamically important throughout the molecular cloud phase. As emphasised above, all stars have become magnetically 'weak' by the time they reach the main sequence. The particular physical process by which most of the primeval flux is lost may very well differ from one star-forming cloud to another. In a molecular cloud at densities beyond $\rho \approx (10^{16}\text{–}10^{17})\ m_H/\mathrm{m}^3$, the optical depth through a Jeans mass exceeds unity, and compressional heat generated during subsequent contraction is largely trapped, so that the temperature rises nearly adiabatically. Beyond $T = 10^3$ K, thermal ionisation increases n_i/n to the level at which flux freezing is normally a very good approximation. In the interval $10 < T < 10^3$ K, ρ increases by 10^4 or more. Nakano and Umebayashi (1986*b*) have argued that in this interval the field no longer moves with the plasma: the electrons are largely attached to the grains, and the currents maintaining the field are carried by ions which suffer a large Ohmic dissipation due to ion–neutral collisions, so destroying most of the primeval flux. However, the decline in the plasma density necessary for the Ohmic dissipation to be large may sometimes be offset by extra sources of ionisation, e.g. by X-rays from newly-formed stars in the same cloud complex (Silk and Norman 1983). A further possibility is that ionising sub-cosmic rays may be generated locally as runaway particles during magnetic reconnection processes as the field of a super-critical fragment detaches itself from the cloud field. Further, Population II (or III) stars may have formed from magnetised gas at 10^4 K (Hoyle 1953), so that collisional ionisation will have kept flux leakage by ambipolar diffusion to a negligible rate. One can hazard a guess that excess primeval flux would in any case be lost during the pre-main-sequence phase by magnetic buoyancy, in a thermal time-scale. Such a reservoir of magnetic energy is a tempting source to power the pre-main-sequence T Tauri phenomenon, and could also yield the extra angular momentum transport required to explain the normally slow rotations of the T Tauri stars.

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