III. STELLAR PLASMAS

MAGNETIC BRAKING

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ABSTRACT. The theory of angular momentum transport by magneto-thermo-centrifugal winds is reviewed, with special reference to recent observations of rapidly-rotating young solar-type stars.

1. INTRODUCTION

Although this is a Symposium on *Basic Plasma Processes on the Sun*, I need make no apology for extending my mandate to include other solar-type stars. The growing wealth of observations in the optical, radio and X-ray bands has firmly established the "Solar-Stellar Connection", in particular demonstrating the close correlation between the strength of solar-type magnetic activity and the star's rotation α. Single stars in young clusters show both stronger chromospheric activity and more rapid rotation (Noyes *et al.* 1984), supporting strongly the theoretical picture of a two-way interaction: magnetic coupling transfers angular momentum from the star, and the strength of the dynamo-built field declines with α. Further evidence comes from close binary systems, where the near synchronization (tidal or magnetic) of spin and orbital motion keeps the rotation rate high, so that older stars such as the RS CVn stars - some evolved off the main sequence - nevertheless remain magnetically active. Detailed comparison of braking models with observation may shed light on some of the most difficult problems in the area, such as the dependence on stellar rotation of the strength (and even the structure) of the dynamo-built fields, and the efficiency of coronal heating.

2. BRAKING BY WINDS: MATHEMATICAL THEORY

The basic idea of magnetic braking is now very familiar (Schatzman 1962; Mestel 1967a,b; 1968; Weber & Davis 1967; Goldreich & Julian 1970, and many other papers since). Gas outflowing as a wind from a rotating star but not subject to any torques conserves its angular momentum and so develops a rotational shear. In the simplest model the magnetic field permeating the gas is taken as symmetric about the rotation axis. Because of approximate field-freezing into the highly conducting gas, the shear acting on the meridional ("poloidal") component $\underset{\sim}{B}_p$ generates an azimuthal ("toroidal") component $\underset{\sim}{B}_t$, and the twisted field exerts torques which try to establish corotation of the gas with the star. The process is thus a competition between the speed v_p of the

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outflow that generates the shear and the Alfvén speed $v_{Alf} = B_p/(4\pi\rho)^{\frac{1}{2}}$ at which the field lines try to straighten. Intuitively one therefore expects that the gas will be kept in near corotation with the star as long $v_p << v_{Alf}$, so that a modest rate of mass loss can yield a disproportionately large rate of loss of angular momentum.

The essence of the mathematical theory describing the quasi-steady braking of an axisymmetric star is as follows.

The *steady-state magneto-kinematic equation* $\nabla \times (\underset{\sim}{v} \times \underset{\sim}{B}) = 0$ yields

$$\underset{\sim}{v} = \kappa \underset{\sim}{B} + \tilde{\omega}\alpha \underset{\sim}{t}, \tag{1}$$

equivalent to

$$\underset{\sim}{v}_p = \kappa \underset{\sim}{B}_p, \tag{2}$$

$$\Omega \equiv v_t/\tilde{\omega} = \kappa B_t/\tilde{\omega} + \alpha, \tag{3}$$

where $\underset{\sim}{t}$ is the unit toroidal vector, $\tilde{\omega}$ the axial distance, κ a scalar function, and α a constant on individual field lines. Equation (2) identifies poloidal field lines and the wind streamlines; equation (3) generalizes Ferraro's law of isorotation.

The *continuity equation* $\nabla.(\rho \underset{\sim}{v}) = 0$ constrains κ:

$$\rho\kappa \equiv \rho v_p/B_p = \eta = \text{constant on poloidal field-streamlines.} \tag{4}$$

The *toroidal component of the non-viscous equation of motion*

$$\rho \underset{\sim}{v}.\nabla(\Omega\tilde{\omega}^2) = \tilde{\omega}\,[(\nabla \times \underset{\sim}{B}_t) \times \underset{\sim}{B}_p/4\pi].\underset{\sim}{t} \tag{5}$$

combines with (2) and (4) to yield

$$-\tilde{\omega}B_t/4\pi + \eta\Omega\tilde{\omega}^2 = -\beta/4\pi \tag{6}$$

where β is again constant on field-streamlines. Multiplication by the strength $d\Phi = B_p dA$ of an infinitesimal poloidal flux tube of local area dA converts (6) into

$$(-\tilde{\omega}B_t B_p/4\pi)dA + (\rho v_p \Omega\tilde{\omega}^2)dA = (-\beta/4\pi)d\Phi, \tag{7}$$

interpretable as showing that the transport across dA of angular momentum by the gas outflowing along $\underset{\sim}{B}_p$, plus the similar transport by the moment of the Maxwell stresses (Lüst & Schlüter 1955), yields the total constant transport $-\beta/4\pi$ per unit poloidal magnetic flux.

Equations (3), (4) and (6) combine into

$$\Omega(1 - 4\pi\eta^2/\rho) = (\alpha + \eta\beta/\rho\tilde{\omega}^2), \tag{8}$$

where

$$4\pi\eta^2/\rho = 4\pi\rho v_p{}^2/B_p{}^2 = v_p{}^2/v_{Alf}{}^2. \tag{9}$$

Near the star the wind speed $v_p << v_{Alf}$, but with ρ decreasing outwards $4\pi\eta^2/\rho$ increases, becoming unity at the Alfvénic point P_A where

$$v_p = v_{Alf} = B_p/(4\pi\rho)^{\frac{1}{2}} \quad \text{and} \quad \rho = \rho_A = 4\pi\eta^2. \tag{10}$$

Thus by (8), Ω (and likewise B_t) will become singular at P_A unless

$$\alpha = -\eta\beta/\rho_A\tilde{\omega}_A{}^2 = (-\beta/4\pi)/\eta\tilde{\omega}_A{}^2 \tag{11}$$

by (10), or

$$(-\beta/4\pi)d\Phi = (\rho v_p/B_p)\alpha\tilde{\omega}_A{}^2(B_p dA) = (\rho v_p dA)\alpha\tilde{\omega}_A{}^2: \tag{12}$$

the angular momentum transported per second along the flux tube is equivalent to that carried by the steady matter flux $(\rho v_p dA)$ if it were kept rotating with angular velocity α out to the Alfvénic point – "effective corotation" out to P_A. At any point the actual values of Ω and B_t are given by

$$\frac{\Omega}{\alpha} = \frac{(1 - \rho_A\tilde{\omega}_A{}^2/\rho\tilde{\omega}^2)}{(1 - \rho_A/\rho)} = \frac{[1 - (v_p/v_A)(B_A\tilde{\omega}_A{}^2/B_p\tilde{\omega}^2)]}{(1 - \rho_A/\rho)} \tag{13}$$

on use of (4) (with $(B_p)_A, (v_p)_A$ shortened to B_A, v_A), and

$$\tilde{\omega}B_t = -4\pi\eta\alpha\tilde{\omega}_A{}^2(1 - \tilde{\omega}^2/\tilde{\omega}_A{}^2)/(1 - \rho_A/\rho). \tag{14}$$

$B_p\tilde{\omega}^2$ is approximately constant along a field–streamline, while ρ_A/ρ, v_p/v_A both decrease strongly towards the coronal base; hence close to the star $\Omega \simeq \alpha$, and by (14) the magnetic transport term in (6) dominates by the factor $\simeq (\tilde{\omega}_A/\tilde{\omega})^2$. Sufficiently far beyond P_A, $\Omega\tilde{\omega}^2 \simeq \alpha\tilde{\omega}_A{}^2$, and the material transport term in (6) ultimately dominates. Over a wide domain encompassing P_A both terms are comparable, and even well within P_A, Ω already lags markedly behind α. It must be emphasized that this does not contradict the basic result (12): "effective corotation" is an exact prescription for computing the angular momentum *transport* along a field–streamline, carried *jointly* by the magnetic and material stresses (Mestel 1967a,b; 1968).

For simplicity we assume the wind zone to be isothermal with sound speed a_w. The wind speed then satisfies the generalized *Bernoulli integral*

$$H \equiv \tfrac{1}{2}(v_p{}^2 + \Omega^2\tilde{\omega}^2) + a_w{}^2\log\rho - GM/r - \alpha\Omega\tilde{\omega}^2 = E$$

$$= \text{constant on field-streamlines.} \tag{15}$$

The new terms are the local rotational energy $\Omega^2\tilde{\omega}^2/2$ per gram and the term $\alpha\Omega\tilde{\omega}^2$ coming from $(\underset{\sim}{j} \times \underset{\sim}{B}/c).\underset{\sim}{v} = \tilde{\omega}\alpha(\underset{\sim}{j} \times \underset{\sim}{B}/c).\underset{\sim}{t}$, the rate of working per gram of the same magnetic torque that gives the outflowing gas angular momentum. If the structure of the field $\underset{\sim}{B}_p$ is supposed known, then substitution of equations (4) and (13) converts equation (15) into a relation $H(\rho,r) = E$ between ρ and the convenient monotonic coordinate r. Once the choice (11) has ensured that Ω and B_t are non–singular, then all non–singular solutions for ρ pass automatically through the Alfvénic point P_A. The critical points of the curve $H = E$ are at the intersections (ρ_s,r_s), (ρ_f,r_f) of $\partial H/\partial r = 0$ and $\partial H/\partial\rho = 0$, corresponding respectively to the slow and fast magnetosonic wave speeds v_s and v_f. Smoothness of flow through these points yields the two conditions

$$H(\rho_s, r_s) = H(\rho_f, r_f) = E \tag{16}$$

which suffice to fix the solution along each field–streamline in terms of the non–dimensional parameters

$$\ell_w = GM/Ra_w{}^2, \quad k = \alpha^2R^3/GM, \quad \zeta_w = B_0{}^2/8\pi(\rho_0)_w a_w{}^2, \tag{17}$$

where B_0 and $(\rho_0)_w$ are standard values, eg at the point where the field-line considered emerges from the coronal base (Weber & Davis 1967; Goldreich & Julian 1970; Sakurai 1985).

In a slow rotator such as the Sun, the terms in Ω in (15) are small compared with the thermal pressure term, and so merely perturb slightly Parker's (1963) thermally-driven wind: the gas undergoes strong acceleration through the slow point where the speed is essentially the sound speed, followed by a much slower subsequent acceleration. By contrast, in a moderately rapid or very rapid rotator, the terms in Ω take over before the Alfvénic point is reached, yielding a net extra driving term $\alpha^2\tilde{\omega}^2(\Omega/\alpha)(1 - \Omega/2\alpha)$. Well within P_A where $\Omega \simeq \alpha$ this term $\simeq \frac{1}{2}\alpha^2\tilde{\omega}^2$, and even at P_A it is typically of the same order. This "centrifugal wind" effect must be included in estimates of the dependence of the rate of braking on the star's rotation.

From (4) and (10), at a general point

$$\tfrac{1}{2}\rho v_p{}^2/(B_p{}^2/8\pi) = 4\pi(\rho v_p/B_p)^2/\rho = \rho_A/\rho = (v/v_A)(B_A/B_p). \tag{18}$$

The Alfvénic surface S_A (defined by all the points P_A) thus separates the domains nearer the star where the magnetic energy dominates from the more distant regions where the kinetic energy dominates. This suggests that near the star the wind will not distort the field but will rather be channelled parallel to $\underset{\sim}{B}_p$, whereas well beyond S_A the field will be passive, being pulled out to follow the wind. The detailed structure of the magnetic field should be determined by the remaining component of the equation of motion:

$$B_p{}^2(\nabla \times \underset{\sim}{B}_p) - [(\nabla \times \underset{\sim}{B}_t).\underset{\sim}{B}_p]\underset{\sim}{B}_t = -4\pi\rho\underset{\sim}{B}_p \times \underset{\sim}{F}_p , \tag{19}$$

this being the balance of the trans-$\underset{\sim}{B}_p$ component of the total magnetic force density and the corresponding component of the density $\rho\underset{\sim}{F}_p$ of non-magnetic (including inertial) forces. Well within S_A the term in $\underset{\sim}{B}_t$ causes a deviation from the vacuum condition $\nabla \times \underset{\sim}{B}_p = 0$ of order $(B_t/B_p)^2 \simeq (\tilde{\omega}_A B_A/\tilde{\omega} B_p)^2(\alpha\tilde{\omega}_A/v_A)^2 << 1$ if $\tilde{\omega} << \tilde{\omega}_A$; and likewise the effect of the thermal, inertial and gravitational terms in $\underset{\sim}{F}_p$ is small for $\tilde{\omega} << \tilde{\omega}_A$. Thus in the simplest model, with $\underset{\sim}{B}_p$ having a dipolar angular distribution over the stellar surface, the field near the star is well approximated by $\underset{\sim}{B}_p = (B_0R^3/r^3)(\cos\theta, \frac{1}{2}\sin\theta, 0)$ in spherical polars. Curl-free dipolar field lines leaving the star at colatitudes $\theta_s \simeq \pi/2$ will reach the equator at $r = R/\sin^2\theta_s << r_A$, so that the approximation $\nabla \times \underset{\sim}{B}_p = 0$ is still valid, whereas those leaving near to the poles will extend to regions close enough to S_A for the forces driving the wind to be able to distort the field into a more rearly radial structure. We are thus led to a multi-component corona, comprising a "dead zone" with hot gas trapped by closed field lines, and a "wind zone" with outflowing, somewhat cooler gas (Mestel 1968; Pneuman & Kopp 1971). This picture is vindicated qualitatively by X-ray observations of bright domains contiguous with the coronal holes associated with the wind.

3. BRAKING LAWS

The direct loss of angular momentum J from the star comes from the outflowing gas in the wind zone. With the radial structure adopted for the field at the Alfvénic surface S_A, approximated as a sphere, the "effective corotation" result (12) yields

$$-dJ/dt \simeq (8\pi/3)(\rho_A v_A r_A{}^2)\alpha r_A{}^2 \tag{20}$$

where α is now identified with the assumed uniform rotation of the star. Along each field-streamline

$$4\pi\rho_0 v_0/B_0 = 4\pi\rho_A v_A/B_A = B_A/v_A \tag{21}$$

from (4) and (10), so that

$$-dJ/dt \simeq (2/3)(B_A r_A{}^2)^2\alpha/v_A. \tag{22}$$

The simplest field models (Weber & Davis 1967; Sakurai 1985) replace Fig. 1 by a field that is radial all the way back to the star, so that $B_A = B_0(R/r_A)^2$ and (22) becomes

$$-dJ/dt \simeq (2/3)(B_0R^2)^2\alpha/v_A. \tag{23}$$

In a slow rotator the mainly thermal driving yields $v_A \simeq (2 - 3)a_w$, virtually independent of α. From (23), $-\dot{J}$ is independent of ρ_0: with the radial field structure, a higher ρ_0 at fixed a_w yields a greater mass loss rate, but the associated reduction in r_A compensates exactly. If further one assumes a linear relation $B_0 \propto \alpha$ for the dynamo-built field, then with $J = kMR^2\alpha$, (23) yields $d\alpha/dt \propto -\alpha^3$, which integrates to $\alpha \propto t^{-\frac{1}{2}}$, identical with the Skumanich law (1972), inferred from comparison of the solar rotation with observations of late-type stars in the younger Pleiades and Hyades clusters.

The field structure of Fig. 1 ensures that in fact only part of the flux emanating from the star is coupled with the wind, so that $B_A r_A{}^2 < B_0R^2$ and by (22) the braking rate is reduced. Also, in a rapid rotator the centrifugal driving term in (15) forces the gas to reach the local Alfvénic speed at a smaller value of r_A, again reducing the braking; this is shown by v_A (a monotonically decreasing function of r) appearing in the denominator in (22). The model of Mestel and Spruit (1987) takes some account of these effects. The approximately dipolar field has $B \simeq B_0(R/r)^3$ out to the radius $\bar{r} < r_A$, and $B \simeq B_0(R/\bar{r}^3)(\bar{r}/r)^2$ beyond $\bar{r}$. (Since along dipolar field lines $\tilde{\omega}^2 \propto r^3$, $B\tilde{\omega}^2 \simeq$ constant along a field line, as quoted above). Thus (22) becomes

$$-dJ/dt \simeq (2/3)\alpha(B_0R^2)^2/(\bar{r}/R)^2 v_A \tag{24}$$

$$= (2/3)(B_0{}^2\alpha^3{}_\odot)R^4(\alpha/\alpha_\odot)^3/(\bar{r}/R)^2 v_A \tag{25}$$

if again $B_0 \propto \alpha$. The details of the wind speed, and in particular its value v_A at P_A, depend on the value of $a_w{}^2/\alpha^2\tilde{\omega}_A{}^2$. In the limit with this parameter zero, $\Omega/\alpha = (1 - \surd(19/27))$ at P_A, yielding the value $4\alpha^2\tilde{\omega}_A{}^2/27$ for the centrifugal driving term in (15). With a_w finite, even in a rapid rotator there will always be a cone about the axis in which thermal driving dominates. A tolerable approximation for v_A is given by

$$v_A/a_w = \{[v_{th}(r_A)/a_w]^2 + \alpha^2\tilde{\omega}_A{}^2/3a_w{}^2\}^{\frac{1}{2}} \tag{26}$$

where $v_{th}(r_A)$ (given by (15) with the terms in Ω dropped) is $(2 - 3)a_w$ typically. Substitution of (26) into (24) yields an equation for r_A and so also for v_A/a_w, which now depends not only on a_w and on α but also on ρ_0. As a convenient alternative to (24) or (25), we can write

$$-dJ/dt \simeq [(8\pi/3)\alpha_\odot(\rho_0)_\odot R^4(v_0/a_w)(a_w)_\odot]K(\alpha/\alpha_\odot), \tag{27}$$

where the bracket is the angular momentum that would be carried off from the Sun by the solar wind if there were no solar magnetic field, and the function $K(\alpha/\alpha_\odot)$ takes care of the different effects studied above. In Mestel & Spruit (1987) the coronal sound speed a_w was assumed independent of α, but the relation $\rho_0 \propto B_0 \propto \alpha$

was adopted as the simplest way of accounting for the optically thin X-ray emission rates $L_X \propto \rho_0^2$ which are observed to be consistent with $L_X \propto \alpha^2$ (Pallavicini *et al.* 1981; Micela *et al.* 1984); whence $K = (R/\bar{r})(r_A/R)^2(\alpha/\alpha_\odot)^2$.

The remaining unknown is the extent $\bar{r}$ of the dead zone. On the equator beyond $\bar{r}$ the oppositely directed field lines are held apart by a thermal pressure equal to the magnetic pressure $[B_r(r,\pi/2)]^2/8\pi$. At $\bar{r}$ these field lines start to diverge from the equator to form the separatrix S between the wind and dead zones, with the analogous boundary condition

$$p + B^2/8\pi = \text{continuous across S.} \qquad (28)$$

The difference in p across S becomes especially marked when the rotation is high. Just outside S, in the wind zone, we have seen that the centrifugal force of near corotation assists in the outward acceleration of the gas, so reducing ρ and p below that in a thermally driven wind. By contrast, in the dead zone the same centrifugal force must be included in the condition of hydrostatic balance that holds along each field line. Again we take the dead zone isothermal, with the sound speed a_d. At the point (r,θ) along the field line leaving the star's surface at (R,θ_0)

$$\frac{\rho}{\rho_0} = \exp\left[-\left[\frac{GM}{Ra_d^2}\right]\left(1-\frac{R}{r}\right) + \frac{1}{2}\left[\frac{\alpha^2R^2}{a_d^2}\right]\left[\frac{r^2\sin^2\theta}{R^2} - \sin^2\theta_0\right]\right]. \qquad (29)$$

Beyond the point where the components along $\underset{\sim}{B}$ of gravity and centrifugal force balance, ρ and p exponentiate outwards, requiring by (28) a corresponding increase in the relative jump in B. To make a first estimate of $\bar{r}$ the field lines are given the vacuum dipolar form $\sin^2\theta/r$ = constant, and condition (28) is applied just at the point $(\bar{r},\pi/2)$ in the form $(B^2/8\pi)_W = (\rho a^2)_d$ with ρ given by (29). As for the wind zone, the relations $B_0 \propto \alpha$, $(\rho_0)_d \propto B_0 \propto \alpha$ are adopted. As expected, the growth in B_0 with α yields an increasing dead zone, but later the growth in $\bar{r}$ is halted and reversed by the strong centrifugal effect on the dead zone pressure.

The angular momentum loss, fixed by the function K in (27) can be parametrized by $\dot{\alpha} \propto -\alpha^p$, with the index p increasing from $\simeq 2.3$ to $\simeq 2.9$ as $\alpha/\alpha_\odot$ increases from unity to the maximum possible value $\simeq 80$. Over the same range K increases by 8 x 10^4, as compared with the value 5 x 10^5 predicted by the Skumanich law p = 3. If instead of increasing like B_0, ρ_0 is assumed independent of α, equation (27) yields smaller values for p, reducing the maximum braking rate by another factor $\simeq 6$. The reason is partly an increase in $\bar{r}$ but primarily the increasing importance of the centrifugal driving term, which prevents the increase in r_A with α from being large enough to compensate for the constancy of ρ_0 (*cf.* the remark following equation (23)).

4. COMPARISON WITH OBSERVATION

Too much attention should not in fact be paid to the precise values of the index p. As pointed out by Vilhu & Moss (1986), allowing a_W rather than ρ_0 to vary suitably with α can yield again the Skumanich index p = 3. Because of the many uncertainties – eg in dynamo theory and in the theory of coronal heating – one is well advised to be guided by observation. Support for the dead zone model comes from observations of non-thermal radio emission from stellar coronae (*cf.* Mutel in Havnes *et al.* 1988). Further support comes from combined optical and X-ray studies of the rapidly rotating (P = .5d) G8-KO dwarf star AB Doradus (Collier Cameron *et al.* 1988). The EXOSAT data show that the part of the spectrum 1–10 keV is well fitted by thermal bremsstrahlung from gas at T $\simeq$ 1.7 x 10^7K, with an emission measure that yields an emitting volume greater than 10 times the volume of the star.

This high energy emission is not subject to eclipse, consistent with its origin in a dead zone extending out to 2–3 R (rather than in the cooler, more compact regions associated with star-spots and like them subject to eclipse, which are responsible for the .1–1 keV part of the spectrum). Observations in Hα (Collier Cameron & Robinson 1989a, b) show transient absorption features, which are consistent with clouds of HI transiting the stellar disk and scattering chromospheric Hα photons. The crucial observation is that the line-of-sight velocity component of the absorbing material always matches that of the obscured stellar limb as the cloud enters and emerges from transit, showing clearly that it is in *near corotation* rather than in a Keplerian orbit. From the rate at which they cross the profile the clouds appear to form at 3–4R, outside the Keplerian corotation radius $r_K = (GM/\alpha^2)^{1/3} \simeq 2.6R$.

It is gratifying to a theorist to have further confirmation of the picture of magnetic control of rotation. There is also evidence of slow outward motion of the Hα clouds, suggesting that the "dead" zone in fact makes some contribution to mass and angular momentum loss: this is discussed further below.

The work of Section 3 assumes the simple dynamo law $B_0 \propto \alpha$, which is in fact responsible for most of the α-dependence of the Skumanich law and its modifications. In a recent survey, Stauffer and Hartmann (1987) (see also Gray 1982) have produced evidence for a rapid early spin-down on the main sequence, with a decline from $35\alpha_\odot$ to $5\alpha_\odot$ occurring in a few x 10^7yr, and at a rate weakly dependent on α, with subsequent braking given by the Skumanich law $\dot{\alpha} \propto -\alpha^3$. These results *prima facie* set problems for standard wind theory. One can get a weak dependence on α by supposing that the dynamo dipole component saturates at $\alpha = \tilde{\alpha}$; an essentially thermal wind would then yield $\dot{\alpha} \propto -\alpha^p$ with $p \leq 1$ for $\alpha > \tilde{\alpha}$, while when α is large enough for centrifugal driving to dominate, the very weak dependence p = 1/3 results. But all the deviations from the conditions leading to the Skumanich law – the dead zone, centrifugal driving and now dynamo saturation – also reduce $\dot{\alpha}$, and it is not easy to produce an early braking time of a few x 10^7yr. For example, if the present-day Sun were losing half its angular momentum in 3 x 10^9yr, and if dynamo saturation occurs at $\tilde{\alpha} = 20\alpha_\odot$, then spin-down of the whole Sun from $70\alpha_\odot$ to $20\alpha_\odot$ would require 7.5 x 10^7yr; while if $\tilde{\alpha} = 10\alpha_\odot$, spin-down to $10\alpha_\odot$ requires 4.8 x 10^8yr. The problem is made still harder by revised estimates (Pizzo *et al.* 1983) that put the solar braking time at $\simeq$ 2 x 10^{10}yr (implying $r_A \simeq 12R$ and a surface field strength $\simeq$ 1.4 Gauss in the wind zone). Stauffer and Hartmann suggest that their observations imply an initial spin-down of just the convective envelope, which has only $\simeq$ 1/20 of the total moment of inertia of the Sun. The consequent constraints on core-envelope magnetic coupling necessary for such a relative shear to persist for a few x 10^7yr are severe though perhaps not completely implausible (*cf.* Mestel & Weiss 1987, where the same problem is discussed but with a "fossil" differential rotation supposed maintained for 2 x 10^9yr). Even so, one is prompted to look for a more efficient braking process, operative at high rotations.

5. EXPLORATIONS

There are many possible variations on the simple wind models discussed above. The high coronal temperature inferred for AB Doradus could imply also a coronal base density for rapid rotators higher than the $\rho_0 \propto B_0$ relation adopted by Mestel & Spruit. As noted above, for a given surface field strength an increase in ρ_0 has primarily the effect of moving the Alfvénic surface inwards without changing much the rate of braking. If the dipolar component of the dynamo-built field saturates at $\alpha = \tilde{\alpha}$, but quadrupolar and higher components increase, one can imagine models in which the "dead" zones extend further than the Alfvénic surface in the wind zone, i.e. $\bar{r} > r_A$.

For a given total magnetic flux F_A crossing S_A, the rate of braking given by expression (22) increases as v_A decreases. This result must be treated with caution: a hypothetical cool corona ($T < 10^5$K) with the same base density as in the Sun will not expand under its thermal pressure, for the density exponentiates down to such low values that the pressure far from the star would be balanced by that of the interstellar medium (Parker 1963). However, once a hot dense wind that has passed through the slow magnetosonic point has ceased to receive heat from the star, it will tend to cool adiabatically so that its Mach number increases still further. If the gas is then shocked, the compressed subsonically moving gas behind the shock cools radiatively and so becomes still denser. The shock may then define a new "coronal base" from which the gas expands into a new standard wind flow (Li Jianke & Mestel, in preparation), but again in a rapid rotator centrifugal driving will force v_A up, and there is no obvious reason why the net braking rate should increase.

Probably the most hopeful prospect for enhanced braking is by a development of the "dead zone" picture. As already noted, beyond $r \simeq r_K$ equation (29) predicts that ρ increases outwards along a field line. Cool HI condensations as observed in AB Doradus may form by radiative recombination and the associated energy loss; a reduction in a_d would lead to further compression, to be made up by further outflow of heated gas from the star into the lower corona. The centrifugal pull $\rho\Omega^2 r$ on a high density condensation in turn causes further distortion of the closed field loops, leading ultimately to field line reconnection and detachment – an effective "macro-resistivity", allowing flow across the field, driven by centrifugal force against the restraining Lorentz force. Although such a trans-field flow cannot be described by the perfect conductivity theory of Section 1, nevertheless one expects again that the slower the magnetically-controlled outflow across $\underset{\sim}{B}$, the greater the distance to which the field can enforce near corotation. And if also $\bar{r} > r_A$, this transport of angular momentum by HI clouds forming in and breaking out of the "dead zone" will dominate over the normal braking from the wind zone.

Further calculations (Barker & Mestel, in preparation) should clarify the details of the expected field structure and the formation of condensations. The very high temperature 1.7×10^7K inferred from the X-ray emission for AB Doradus would severely limit the density increase predicted by (29), but the above thermal instability could set in if the coronal dead zone is inhomogeneous, with parts of it having a temperature closer to the solar value of $\simeq 2 \times 10^6$K. The likely sensitivity of the results to the assumed magnetic field structure and to the coronal base density may yield clues as to dynamo action and to coronal heat input at high rotations.

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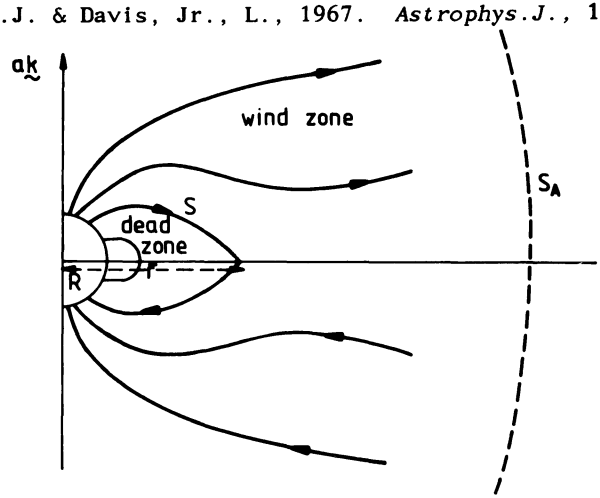

Figure 1. Schematic magnetic field model

DISCUSSION

WEISS: Some models of young stars on the main sequence assume that only the envelope is spun down, so that strong differential rotation persists for 10^8 years or longer. Do you think that this is likely?

MESTEL: One reason why I am trying to increase the efficiency of braking in the earliest phases (*i.e.*when the star is rapidly rotating) is because I am unhappy about the interpretation of observations that requires spin-down of just the envelope. You recall that in our paper we noted what severe constraints are imposed on internal magnetic fields coupling core and envelope, if such a strong shear is supposed to persist for so long.

CHITRE: The helioseismological data seem to suggest that the differential rotation observed at the solar surface is impressed even in the interior. Is this indicating the influence of the magnetic field linking the core with the envelope on the rotation profile?

MESTEL: My impression is that the interpretation of the helioseismological data is still ambiguous. Maybe Dr Christensen-Dalsgaard will comment.

DALSGAARD: I agree with your impression that the question of the core rotation has not been settled. I think we can definitely say that the region outside $0.2R_{\odot}$ cannot be rotating at substantially more than the surface rate. Some observations indicate splitting for $\ell = 1$ modes corresponding to almost twice the surface angular velocity. If true, these observations would indicate a very small core rotating quite rapidly.

MESTEL: That would still restrict severely the strength of the field penetrating into this very small core.