A CRITIQUE OF CURRENT MAGNETIC-ACCRETION MODELS FOR CLASSICAL T TAURI STARS

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ABSTRACT

Current magnetic-accretion models for classical T Tauri stars rely on a *strong, dipolar* magnetic field of stellar origin to funnel the disk material onto the star and assume a steady-state. In this paper, I critically examine the physical basis of these models in light of the observational evidence and our knowledge of magnetic fields in low-mass stars and find it lacking. I also argue that magnetic accretion onto these stars is inherently a time-dependent problem and that a steady state is not warranted.

Finally, directions for future work toward fully consistent models are pointed out.

Subject headings: accretion, accretion disks — MHD — stars: formation — stars: magnetic fields — stars: pre–main-sequence

1. INTRODUCTION

It is now widely accepted that classical T Tauri stars are accreting material from their circumstellar disks. So far, this is the most successful explanation for the large fluxes in the blue and UV regions of the spectrum that characterize these stars (Haro & Herbig 1954; Walker 1956) and the filling in or "veiling" of the absorption spectrum longward of 3800 Å (Joy 1949; Haro & Herbig 1954). The strongest evidence for accretion is the detection of redshifted absorption in upper Balmer and permitted metallic lines, at velocities of up to several hundred km s^{-1}, in T Tauri stars with large optical and UV excesses (Walker 1972; Edwards et al. 1994, and references therein).

Notwithstanding the observational evidence, our understanding of the accretion process is sketchy at best. Early attempts (Bertout, Basri, & Bouvier 1988; Basri & Bertout 1989), following the original suggestion by Lynden-Bell & Pringle (1974), modeled the accretion region as a shear boundary layer at the star-disk boundary. These empirical, equatorial boundary layer models remained popular until the early 1990s, even though they are inconsistent with the detection of redshifted *absorption* lines. Currently the most popular picture is that pioneered by Königl (1991), wherein the disk material is funneled onto the star by a dipolar stellar magnetic field. A variation of this scenario has been put forward by Shu and collaborators (Shu et al. 1994a, 1994b; Ostriker & Shu 1995). A major appeal of these models is their supposed ability to account for the fact that classical T Tauri stars are slower rotators than their naked siblings (see, e.g., Edwards et al. 1993, and references therein).

The purpose of this paper is to point out that the magnetic topology assumed in these models is inconsistent with all the available data on T Tauri stars (TTS) and our knowledge of magnetic fields in low-mass stars, a state of affairs analogous to that in the late 1980s, when boundary layer models were widely popular, even though they were inconsistent with the spectroscopic evidence for non-equatorial mass infall.

The paper is organized as follows: in § 2 I critically review the observational evidence that is inconsistent with the current assumption of a strong stellar dipolar field; the large-scale structure of a stellar field consistent with the data, and its implications for the accretion process, are examined in § 3 and § 4, respectively; and § 5 contains my conclusions.

2. THE OBSERVATIONAL EVIDENCE

Over the last two decades a wealth of data has accumulated that indicates that the magnetic activity of T Tauri stars (both classical and their naked siblings) is a scaled-up version of that in low-mass, main-sequence stars, wherein the stellar magnetic field is spatially intermittent and its topology varies with time (see, e.g., Rosner, Golub, & Vaiana 1985, and references therein).

The most dramatic evidence for dynamo-driven magnetic activity in T Tauri stars, akin to that found in late-type dwarfs, is the unexpected discovery that T Tauri stars— both classical (CTTS) and naked (NTTS) T Tauri stars—are strong emitters of X-rays (with typical X-ray luminosities a factor 10^2–10^4 above solar levels), and that this emission is time variable (for a review, see Feigelson, Giampapa, & Vrba 1991; and Montmerle et al. 1993). Recent observations of the Orion region with *ASCA* (Yamauchi et al. 1996) confirmed previous findings that this emission can be attributed to a multitemperature plasma with at least two components: one at $k_B T = 0.7$–1 keV and a second at $k_B T \gtrsim 2$ keV.

A clue to the origin of the X-ray emission in T Tauri stars is that these objects extend the range of the correlation between X-ray luminosity, L_X, and stellar equatorial velocity $v \sin i$ (Bouvier 1990; Damiani & Micela 1995) which was found for low-mass dwarfs and other magnetically active objects (Rosner et al. 1985). This so-called rotation-activity connection supports the view that TTS stars have coronae similar to those in late-type dwarfs, wherein spatially *intermittent* magnetic fields with a time-variable *topology* are responsible for the heating of the outer atmosphere (Rosner et al. 1985).

The strong variability of the X-ray emission in T Tauri stars (see Montmerle et al. 1983) lends further support to the idea that the magnetic field in TTS is spatially intermittent and its topology varies with time. X-ray flares in these objects are typically 10^2–10^3 (and up to 10^4) times stronger than the strongest solar flares (Preibisch, Zinneker, & Schmitt 1993, and references therein). Note that there is

no free energy in a global dipolar field, and therefore such a field cannot drive a flaring coronae (see, e.g., Priest 1982; Haisch, Strong, & Rodonò 1991).

Additional evidence for the dynamo-origin of the stellar magnetic field in TTS is the detection at microwave wavelengths of nonthermal radio emission from some of these objects (see, e.g., Chaing, Phillips, & Lonsdale 1996, and references therein; White, Pallavicini, & Kundu 1992). The luminosity of the nonthermal radio emission in TTS, L_R, extends to larger luminosities the correlation between L_X and L_R found for low-mass dwarfs and other magnetically active stars (Güdel & Benz 1993). This correlation, which extends over 6 orders of magnitude in L_X and L_R, suggests that the heating of active coronae and particle acceleration are closely linked, as is the case in solar flares. Güdel & Benz note that, according to the relation between L_X and L_R, the quiet Sun stands out as a radio-underluminous star. Therefore, stars that follow the L_X-L_R relation, and in particular T Tauri stars, must have a much higher level of magnetic activity in their *quiescent* state than the quiet Sun.

Irregular variability at all wavelengths is a distinguishing characteristic of TTS (Joy 1945). However, it is now well established that for many objects the optical light curves have a periodic modulation (see, e.g., Herbst et al. 1994; Safier 1995, and references therein) due to stellar spots that are colder than the surrounding photosphere. These cold spots are regions of enhanced magnetic flux in the photosphere (see, e.g., Gray 1989), and they provide additional evidence that the magnetic field in TTS is spatially intermittent.

Recent Doppler imaging of the NTTS V410 Tau (Hatzes 1995, and references therein) and HDE 283572 (Joncour, Bertout, & Bouvier 1994) confirms the presence of cold stellar spots but also shows that their surface distribution on these stars is fundamentally different from that on the Sun. In these NTTS, the spots are concentrated toward the polar caps, without any symmetry with respect to the stellar equator, and the fractional coverage of the stellar surface is very large. This is consistent with the results of Schüssler et al. (1996) and DeLuca, Fan, & Saar (1997), who found that with increasing angular velocity and depth of the convection zone, the Coriolis force dominates over magnetic bouyancy, and the mean latitude of flux-emergence shifts towards the poles. The magnetically active regions in the Sun, where X-ray flares originate, are located above sunspots, and outside these regions the magnetic field is much weaker. If the Sun is any guidance, then these findings for NTTS may be interpreted as suggesting that their magnetic fields near the stellar equator are significantly weaker than those in the cold spots.

Finally, the failure to measure the surface fields of TTS by Zeeman polarization experiments (Johnstone & Penston 1986, 1987), and their detection by means of Zeeman broadening observations (Basri, Marcy, & Valenti 1992), are consistent with a complex magnetic topology. Recall that the Zeeman polarization depends on the orientation of the magnetic field with respect to the observer; if the field orientation changes along the line of sight, then the contributions to the local polarization from different magnetic field elements will nearly cancel out. On the other hand, the Zeeman-broadening technique (Robinson 1980) relies on the measurement of Zeeman-induced changes in the *unpolarized* line profiles.

It may seem paradoxical that T Tauri stars, which are fully convective (Stahler 1988, and references therein), exhibit all the characteristics of dynamo-generated magnetic fields, when a variety of theoretical arguments suggest that the dynamo operative in the Sun, and late-type dwarfs with radiative cores, requires a strong toroidal field anchored at the base of the convection zone (Parker 1975; Golub et al. 1981).

However, Rosner (1980) has discussed how a magnetic dynamo could operate in a fully convective star, and Durney, De Young, & Roxburgh (1993) first showed in detail that such dynamo is possible. Recent observations by Hodgkin, Jameson, & Steele (1995) show that fully convective, young main-sequence stars have coronae and suggest that the nature of the stellar dynamo in these objects is different from that in stars with radiative cores. In other words, although the nature of the dynamo in fully convective stars is not fully understood, the observations show that fully convective stars *do* have active coronae. Therefore, there is no reason why, in principle, the magnetic fields in T Tauri stars could not be generated by a stellar dynamo.

To summarize, the currrent observations of T Tauri stars are consistent with the notion that their magnetic field is generated by a stellar dynamo. Dynamo fields cannot be axisymmetric (Cowling 1934), and, therefore, the magnetic field in these stars is most likely spatially intermittent, with a variety of structures on different spatial scales, and highly variable with time.

3. THE LARGE-SCALE STRUCTURE OF THE STELLAR MAGNETIC FIELD

In current models of magnetic accretion, the region of interaction between the stellar field and the circumstellar disk is located at distances of order several stellar radii. It is usually argued that the complexity of the stellar field is irrelevant because, at these distances, the dipolar component of the field will dominate over larger multipoles. This is correct, but one may not conclude that the field has a closed geometry at large distances from the star. The reason is that the stellar magnetic field is embedded in hot, coronal gas, the presence of which is overlooked in current magnetic-accretion models, and the thermal pressure eventually dominates over magnetic forces. Once this happens, if the thermal pressure is large enough to overcome the gravitational force, a stellar wind will flow (Parker 1963). This stellar wind will drag and stretch the magnetic field lines, changing the field topology into a configuration wherein the field lines are mostly radial away from the star. Note that, close to the star, regions with both open and closed field lines will exist, but the large-scale structure of the field will be predominantly open; Figure 1 illustrates this situation.

Does thermal pressure ever dominate over magnetic stresses in a typical TTS corona? To answer this question, one has to find the ratio of the gas to magnetic pressure, $\beta \equiv p_{gas}/p_{mag}$, as a function of distance from the coronal base. To quantify matters, consider the following example given by Rosner et al. (1995). Assume, in light of the observed magnetic activity of these stars, that the surface field can be regarded as a random and sparse spatial superposition of surface dipoles of fixed strength, with a mean separation d between dipoles. At a height $z \equiv r - R_*$ above the stellar surface, the mean field-strength, B, from this flux distribution is a function of z and d only, and $B(z) \propto d^3(z^2 + d^2)^{-3/2}$. Hence, in the far field $(z \gg d)$, $B(z) \propto (d/z)^3$ to first order in d/z. This means that the mag-
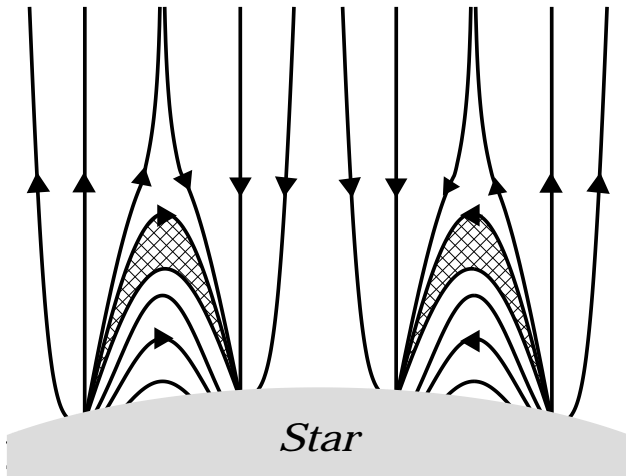
FIG. 1.—Sketch of the magnetic topology in the corona of a T Tauri star with a thermally driven wind. Two magnetic elements of opposite polarity are shown (not to scale); hatched regions correspond to magnetic flux tubes with an enhanced density, which would appear as X-ray-bright coronal loops. For clarity, the magnetic topology in the photosphere, chromosphere, and transition region are not shown. Arrowheads indicate the polarity of the field lines.

netic field varies very rapidly with height as soon as one reaches heights comparable to the typical correlation length of magnetic flux-bundles at the stellar surface.

Next, one has to find the radial dependence of the gas pressure. In light of the strong dependence of the magnetic field strength with distance, it suffices to use an approximation for the gas pressure that is appropriate for distances of order a few stellar radii. Therefore, assume that the run of gas pressure with distance is that for a static, isothermal atmosphere, where $p_{gas}(r) \propto \exp\left[R_*(R_* - r)/rH\right]$. Here $H \equiv 2k_B T_*/m_H g_*$ is the gravitational scale height, T_* is the (constant) temperature, and g_* is the gravitational acceleration at the stellar surface. The isothermal assumption is obviously an oversimplification, but it should be noted that, starting at the coronal base, the temperature will first increase with distance before it starts to drop.

In the Sun, the gas temperature at the base of the corona is not uniform. The highest temperatures, $T_* = 2.5 \times 10^6$ K, occur in active regions, wherein the magnetic field topology is mostly closed. These active regions appear in X-ray images as bright, closed loops. On the other hand, the lowest value of T_*, $T_* = 1.4–1.8 \times 10^6$ K, is found in the so-called coronal holes, which are regions of open magnetic field lines where the bulk of the solar wind originates.[1] Bame et al. (1975) deduced the gas temperature above coronal holes out to distances $\lesssim 4 R_\odot$ and found it to be roughly constant. Therefore, and given the observational evidence for large amounts of hot coronal gas up to $T \sim 10^7$ K in TTS (see § 2), the isothermal assumption with $T_* = 10^6$ K should give an appropriate *lower* limit to the gas pressure out to $r \sim$ few R_*.

The distance r_T where the gas pressure comes to dominate the magnetic pressure is found by solving the equation

$p_{gas}(r_T) \sim B^2(r_T)/8\pi$. In this example, the gas pressure will always overcome the magnetic pressure far enough from the star because $B^2 \propto r^{-6}$, whereas p_{gas} approaches asymptotically a finite value. To be definite, I have plotted in Figure 2 the solution for r_T as a function of d for different values of B_* and n_* for a typical T Tauri star with $M_* = 0.5$ $M_\odot$ and $R_* = 2.5 R_\odot$, and using $T_* = 10^6$ K.

The values of n_* and B_* that I have chosen should be typical. At the base of the solar corona $n_* = 10^8–10^9$ cm^{-3}, and because the transition regions between the chromospheres and coronae of TTS are denser than in the Sun (see, e.g., Brown, de Ferraz, & Jordan 1984) these values of n_* are probably lower limits. As for the value of B_*, Basri et al. (1992) detected *photospheric* fields ~ 1 kG, a typical value in sunspots (see, e.g., Priest 1982). Because the field strength inferred in active regions above sunspots is a factor ~ 10 smaller than the photospheric field (Priest 1982), and, to repeat, because the *photospheric* field strengths measured in TTS are similar to those in the Sun, a value $B_* = 100$ G is probably typical at the base of the TTS's coronae. However, and for completeness, I have plotted in Figure 2 also the solution for $r_T(d)$ for the unlikely case $B_* = 10^3$ G.

The results in Figure 2 show that, for $d \lesssim 0.1R_*$ and $B_* \leq 10^3$ G, $r_T \lesssim 3R_*$. Therefore, for small mean separations of the magnetic elements on the stellar surface, the thermal pressure dominates over the magnetic stresses beyond $\lesssim 3R_*$ *right above* the closed structures. Between these loops, if the thermal pressure is large enough to drive a stellar wind, the radial field lines will extend all the way to the coronal base (see Fig. 1). Even if $d \sim R_*$, with $B_* \leq 10^3$ G, thermal pressure will dominate most of the corona beyond $r \ll 3R_*$, and the argument is as follows: because the number of magnetic elements that can be accommodated on the stellar surface with a mean separation d scales as d^{-2}, if $d \sim R_*$, only a *few* such elements will be present. Right above these large loops, magnetic stresses may be dominant (depending on the value of B_*) beyond $\sim 3R_*$; however, between these loops, again, the field lines will be radial all the way to the coronal base if a thermally driven
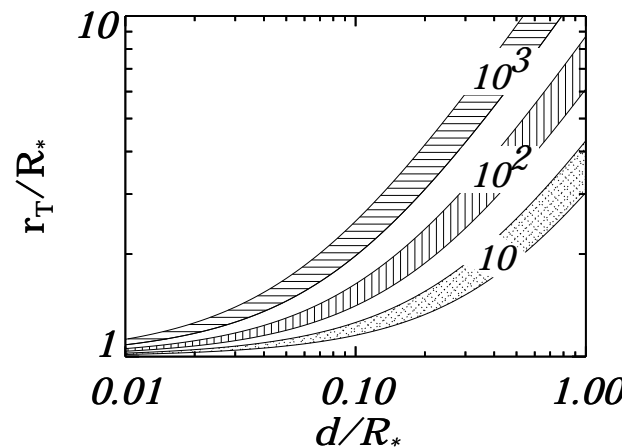

FIG. 2.—Radius where the thermal pressure dominates over magnetic stresses, r_T (in units of the stellar radius R_*), as a function of the mean dipole separation at the stellar surface, d (also in units of R_*), for a T Tauri star of mass $M_* = 0.5 M_\odot$, radius $R_* = 2.5 R_\odot$, and coronal temperature $T_* = 10^6$ K. Each hatched region corresponds to one value of the magnetic field strength at the base of the corona, B_*, and to coronal-base densities, n_*, in the range $10^8–10^9$ cm^{-3}. Upper and lower curves bounding each region correspond to $r_T(d)$ for $n_* = 10^8$ cm^{-3} and $n_* = 10^9$ cm^{-3}, respectively; the labels indicate the value of B_* in gauss.

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[1] One of the most striking features of X-ray images of the solar corona is the presence of bright loops against a background of diffuse emission and large, X-ray-dark regions—the so-called coronal holes. However, one may not conclude that the hot gas ($T \gtrsim 10^6$ K) is confined to the coronal loops. The X-ray contrast between the loops and the coronal holes is due to the enhanced densities in the former (see, e.g., Munro & Jackson 1977; Priest 1982).

wind is present (Fig. 1), and a large fraction $\sim (d/R_*)^2 \sim 1$ of the corona will be threaded by open field lines. Moreover, in light of the evidence for a predominance of high-latitude cold spots, as a result of the Coriolis force dominating over magnetic buoyancy (see § 2), most of these large, closed magnetic structures are likely to be away from the stellar equator and will not intersect the disk's midplane.[2]

In summary, the value $r_T \lesssim 3R_*$ is a likely *upper* limit to the distance from a typical TTS where thermal pressure starts to dominate over magnetic stresses, and the critical question now is whether a thermally driven stellar wind is possible.

To answer this question, I have computed a Weber & Davis (1967) wind model using the one-dimensional generalization by Sakurai (1985), with M_*, R_*, T_*, β_*, and Ω_* as input parameters, where $\beta_* = 8\pi n_* k_B T/B_*^2$ and Ω_* is the stellar angular velocity. The Weber & Davis model (originally applied to the Sun) assumes a steady state, polytropic flow confined to the equatorial plane of a rotating star; the generalization by Sakurai allows for arbitrary rotation rates. Although this wind model strictly applies to the stellar equatorial plane, it should be accurate enough for my purposes because of the small geometrical thickness of accretion disks around TTS (see, e.g., Bertout et al. 1988).

For the fiducial values $M_* = 0.5\ M_\odot$, $R_* = 2.5\ R_\odot$, $T_* = 10^6$ K, $\beta_* = 3.5 \times 10^{-4}$ (corresponding to $B_* = 100$ G and $n_* = 10^9$ cm^{-3} at $T_* = 10^6$ K), and $\Omega_* = 10^{-5}$ s^{-1}, which gives a stellar rotation period of 7.3 days, a typical value for CTTS (Bouvier et al. 1993), I was able to find a wind solution that becomes trans-sonic at $r_s = 1.07R_*$, trans-Alfvénic at $r_A = 1.59R_*$, and trans-fast-magnetosonic at $r_F = 1.62R_*$.

If one varies the values of B_* and n_* while holding the other parameters constant, the values of r_s, r_A, and r_F change by only a few percent, with $r_F - r_A \ll 1$. This shows that these stellar winds are thermally driven, even though TTS are relatively fast rotators, and therefore the stretching of the field lines is very effective.

If one chooses $M_* = 1\ M_\odot$ and $R_* = 3\ R_\odot$ (this value of R_* corresponds to an age ~ 0.5 Myr according to the evolutionary tracks of D'Antona & Mazzitelli 1994), then $r_s = 2.4R_*$, $r_A = 2.8R_*$, and $r_F = 3.6R_*$; while for a 0.3 $M_\odot$ TTS with $R_* = 2.2\ R_\odot$ (D'Antona & Mazzitelli 1994) one finds that $r_s = 1.01R_*$, $r_A = 1.48R_*$, and $r_F = 1.49R_*$.

In summary, for any reasonable choice of TTS stellar and coronal parameters, one finds that a thermally driven stellar wind must be present,[3] and an *upper* limit to the size of closed magnetic structures that may interact with the disk is

$r \lesssim 3R_*$. Therefore, there is no basis for assuming a closed field geometry further out.

4. CONSEQUENCES FOR MAGNETIC ACCRETION

Current models of magnetic accretion rely on a closed stellar magnetic field to mediate the accretion of disk matter and to spin down the star, and a steady state is assumed. In one way or another, the corotation radius r_{co}, the radius where the Keplerian angular velocity equals the stellar angular velocity Ω_*, plays a fundamental role in these models, with $r_{co} = 5R_*(M_*/0.5\ M_\odot)^{1/3}(\Omega_*/10^{-5}\ s^{-1})^{-2/3}$ $(R_*/2.5\ R_\odot)^{-1}$. The results of the preceding section show that there is no basis for assuming a closed field geometry beyond $r \lesssim 3R_* \sim 0.5r_{co}$.

Does this matter?

Ostriker & Shu (1995) argue that the stellar field will truncate the disk just inside of r_{co}, and the field lines will be open just beyond r_{co}. In this ad hoc way, they bypass the issue of field shearing by the Keplerian flow. However, if the stellar field lines are radial beyond $r \lesssim 3R_*$, disk disruption may occur only inside of $\lesssim 3R_*$, if at all. Hence, the Ostriker & Shu model is missing an important piece of physics. This casts serious doubts on the validity of their model (see also Wang 1997).

The more physically realistic model of Ghosh & Lamb (1978) invoked by Königl (1991) relies on the balancing of the spin-up and spin-down torques from inside and outside r_{co}, respectively, and includes field shearing due to the disk flow. In this model, the growth of the toroidal field, B_φ, is limited by reconnection, and it is assumed that a steady state is reached. However, note that the wind flow described in § 3 becomes trans-fast-magnetosonic beyond $r_F = 1.5$–$3R_* = 0.3$–$0.7r_{co}$, and, therefore, no torques can be exerted on the star from beyond $r = 0.3$–$0.7r_{co}$ as long as the field lines are filled with a stellar wind.

This situation will last until a flux tube is filled with disk matter all the way to the star. Because of the higher density of disk matter, the Alfvén speed along the flux tube will be lower than that when wind matter is present, and this *may* result in $r_F < R_*$; if this is the case, then disk torques from beyond $\sim 0.5r_{co}$ will act on the star. To within 1 order of magnitude, disk matter will fill a flux tube all the way to the star in a free-fall time, t_{ff}. However, a strong toroidal field $B_\varphi \sim B_r$ will develop on a timescale $t_\varphi \sim |\Omega_K - \Omega_*|^{-1}$ ($\sim t_{ff}$ outside the neighborhood of r_{co}) where Ω_K is the local Keplerian angular velocity in the disk. The presence of this toroidal field has important consequences.

First, a realistic stellar wind will not be confined to the equatorial plane, but it will fill a large fraction of 4π steradians. Because the field lines carried by this wind are mostly radial, only the field near the stellar equatorial plane will be strongly sheared by the disk matter. Thus, a strong gradient in B_φ^2 will exist parallel to the axis of rotation; and the flow, and the field, will be deflected *away* from the disk on a timescale $t_z \sim h\rho_d^{1/2}/B_\varphi$, where h is the disk scale height at a distance r from the star and ρ_d is a typical density in the disk *outer* atmosphere at r. The disk matter will start shearing the field as soon as the field penetrates to a disk depth where $\rho_d \sim B^2/v_K^2$, were v_K is the local Keplerian speed. Therefore, $t_z \sim (h/r)t_{ff} < t_{ff}$, and *before* disk matter fills a stellar flux tube all the way to the star, this flux tube will be disconnected from the disk outer layers and the loading of disk

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[2] In light of the observed correlation between L_R and L_X, it is unlikely that most of the closed field lines are in a few, very large loops. Although the emission measures derived from the hot component of the X-ray emission can be explained in terms of a few, dense, very large magnetic loops with material at temperatures $\sim 10^7$ K, such isolated loops cannot account for the nonthermal electrons required to explain the observed nonthermal radio luminosities. On the other hand, a multitude of smaller loops undergoing frequent flaring can explain the observed X-ray luminosities at $T \sim 2$ keV, the nonthermal radio emission, *and* the correlation between L_R and L_X (Güdel 1997).

[3] I do not propose that the observed outflows from young stellar objects are thermally driven stellar winds. These winds are much too weak to account for the inferred mass outflow rates and mechanical luminosities. I have just shown that TTS must have thermally driven stellar winds, which probably coexist with the more powerful, observed outflows, the origin of which is still unknown.

matter onto the field will stop. Eventually, as new matter flows from the stellar surface dragging along magnetic field lines, the original geometry may be reestablished, but this shows that the assumption of steady state accretion on timescales $\sim t_{ff}$ is unwarranted and that the transmission of torques to the star from beyond $r \sim 0.5 r_{\rm co}$ is very ineffective.

One may argue that, by averaging over several rotation periods, a net torque will act on the star from beyond ~ 0.5 corotation radii. However, the above analysis shows that disk-torques from a point at a distance r from the star will last for a fraction h/r of the free-fall time, where h is the disk's local scale height, before the magnetic field detaches from the disk; therefore, the duty cycle of the disk torques from beyond $\sim 0.5 r_{\rm co}$ is likely to be very small.

Second, strongly sheared fields are subject to MHD instabilities (both ideal and resistive), which relieve the field shear on a timescale $t_{\rm ins}$, typically $t_{\rm ins} \ll t_{\varphi}$, with a dramatic change in topology.

Probably, the most relevant ideal-MHD instability in this problem is the Balbus-Hawley instability (Balbus & Hawley 1991). Although the growth time for this instability is of order the shearing timescale, t_{φ}, the nonlinear evolution of the instability results in a *turbulent* magnetic field for almost any initial field geometry (see Stone et al. 1996). Hence, on a timescale $\sim t_{\varphi}$ the disk will be endowed with a magnetized corona that is not simply connected to the star.

As for resistive instabilities, astrophysically relevant instabilities rely on resistivity of anomalous origin which becomes important after a finite stress threshold is crossed (see, e.g., Priest 1982). The outcome of these instabilities is usually a *sudden* change of topology—solar and stellar flares are prime examples—as opposed to a steady state configuration (see, e.g., Priest et al. 1986, and references therein; Forbes & Priest 1984). In the context of T Tauri stars, the simulations of Hayashi, Shibata, & Matsumoto (1996) show that when anomalous resistivity with a finite stress-threshold is included, the stellar field is sheared by the disk for a time $\sim t_{\varphi}$, and thereafter the stresses are *dynamically* relieved ($t_{\rm ins} \ll t_{\varphi}$) with a dramatic change in field topology. Therefore, there is no basis to assume that reconnection will result in a steady state topology that is known a priori.

In summary, the fact that the magnetic field will be mostly tangential to the disk surface beyond $\sim 0.5 r_{\rm co}$ means that there is no basis for the Ostriker & Shu (1995) claim that the disk will be truncated just inside $r_{\rm co}$ and that disk currents induced by a penetrating field result in an open field geometry beyond this point. Furthermore, due to the strong vertical gradients in B_{φ} that will develop if the stellar field is mostly radial, each interaction of the stellar field with the disk will last for less than about a local rotation period, and the disk torques on the star from beyond $r \sim 0.5 r_{\rm co}$ will be very ineffective because the field will be deflected away from the disk in much less than a free-fall time. This casts serious doubts on the applicability of the Ghosh & Lamb model to this problem. Finally, a realistic consideration of MHD instabilities shows that there is no basis to assume a steady state field topology that is known a priori.

5. CONCLUSIONS

I have critically examined the validity of current magnetic-accretion models for classical T Tauri stars in light of the available data and our knowledge of magnetic fields in active, low-mass dwarfs. I have shown that the hot coronae around TTS can sustain thermally driven stellar winds which change the field topology into a radial configuration beyond $\lesssim 3 R_*$. Therefore, there is no basis for assuming a closed field topology farther out.

Because the magnetic field will be mostly tangential to the disk surface beyond ~ 0.5 corotation radii, the disk will not be truncated beyond this point, and field shearing by the disk has to be included in any realistic model.

The strong vertical gradients in the toroidal field that will develop in the outer layers of the disk will limit each interaction between the field and the disk to about one local Keplerian period, while the disk torques from beyond $\sim 0.5 r_{\rm co}$ will be effective for only a small fraction of this time. Therefore, the assumption of steady state accretion is not warranted, and the duty cycle of disk torques on the star from beyond about half the corotation radius will be small.

Finally, a realistic consideration of MHD instabilities shows that there is no basis to assume a steady state field topology that is known a priori.

How, then, does nonequatorial accretion take place?

It is beyond the scope of this paper to present a fully consistent model. On the contrary, it is my purpose to point out that much work needs to be done before such model can be presented. In particular, we need a better understanding of the interaction of the stellar field with the disk atmosphere through MHD instabilities, and the inherent time-dependent nature of the problem cannot be ignored.

Magnetic accretion probably *does* take place but at distances smaller than those proposed in current models. I have shown that the stellar magnetic field is likely to be radial *beyond* $\lesssim 3 R_*$; inside this radius, magnetic accretion may be mediated by stellar magnetic loops. However, because these magnetic loops only exist well inside the corotation radius, the accreting material will spin up the star rather than spin it down (see Wang 1997). Therefore, the observational evidence for classical T Tauri stars being slower rotators than their naked siblings is still unexplained.

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