The inner magnetized regions of circumstellar accretion discs

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Abstract

In this lecture, I will briefly address several phenomena expected when magnetic fields are present in the innermost regions of circumstellar accretion discs: (i) the magneto-rotational instability and related "dead zones"; (ii) the formation of magnetically-driven jets and the observational constraints derived from Classical T Tauri stars; (iii) the magnetic star-disc interactions and their expected role in the stellar spin down.

It should be noted that the magnetic fields invoked here are organized large scale magnetic fields, not turbulent small scale ones. I will therefore first argue why one can safely expect these fields to be present in circumstellar accretion discs. Objects devoid of such large scale fields would not be able to drive jets. A global picture is thus gradually emerging where the magnetic flux is an important control parameter of the star formation process as a whole. High angular resolution technics, by probing the innermost circumstellar disc regions should provide valuable constraints.

Key words:
Accretion, accretion discs, Magnetohydrodynamics (MHD), Stars: pre-main sequence, Stars: magnetic fields, Stars: rotation, ISM: jets and outflows

1 Introduction

Actively accreting classical T Tauri stars (CTTS) often display supersonic collimated jets on scales of a few 10-100 AU in low excitation optical forbidden lines. Molecular outflows observed in younger Class 0 and I sources may be powered by an inner unobserved optical jet. These jet signatures are correlated

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with the infrared excess and accretion rate of the circumstellar disc (Cabrit et al., 1990; Hartigan et al., 1995). It is therefore widely believed that the accretion process is essential to the production of jets.

Most of observed jets are extremely well collimated, with an opening angle of only some degrees. On the other hand, the derived physical conditions show that jets are highly supersonic. Indeed, emission lines require a temperature of order $T \sim 10^4$ K, hence a sound speed $C_s \sim 10$ km/s while the typical jet velocity is $v_j \sim 300$ km/s. The opening angle θ of a ballistic hydrodynamic flow being simply $\tan\theta = C_s/v_j$, this provides $\theta \sim 5^o$, nicely compatible with observations. Thus, jets could well be ballistic, with an inertial confinement farther up. But the fundamental question is *how does a physical system produce an unidirectional supersonic flow ?* This simply implies that confinement must be closely related to the acceleration process. To date, the only process that proved to be able to do this is the action of a large scale magnetic field anchored on the driving engine. This is the reason why current sophisticated jet models are computed using Magneto-Hydrodynamics (hereafter MHD).

For quite a while however, the precise physical connection remained a matter of debate: do the jets emanate from the star, the circumstellar disc or the magnetospheric star-disc interaction? This issue seems now to be almost settled: while all these wind components are probably present, magnetized disc winds would be responsible for most of the mass loss (see Ferreira et al. 2006a) [1] .

The basic and universal accretion-ejection mechanism would then be the following. An accretion disc around a central object can – under certain conditions and whatever the nature of this object (star or compact object) – drive jets through the action of large scale magnetic fields, threading the disc. These fields would tap the mechanical energy released by mass accretion within the disc and transfer it to an ejected fraction. The smaller the fraction, the larger the final jet velocity. One thing that must be understood is how the presence of these jets modifies the nature of the underlying accretion flow. Many papers in the literature actually assume that the accretion disc resembles a standard accretion disc (hereafter SAD), as first described by Shakura and Sunyaev (1973). Thus, although it was soon recognized that ejection and accretion were tightly related (Blandford, 1976; Lovelace, 1976; Blandford and Payne, 1982; Pudritz and Norman, 1983, 1986; Lovelace et al., 1987; Konigl, 1989), a truly self-consistent model appeared only lately (Ferreira and Pelletier, 1995; Ferreira, 1997; Casse and Ferreira, 2000a; Ferreira and Casse, 2004). To date, this is the only published MHD model that describes in a self-consistent way the physics of an accretion disc threaded by

[1] Throughout this paper the term "disc wind" designates a magnetically driven jet from an accretion disc. Note that in the literature it is sometimes used to refer to a thermally driven, uncollimated outflow.

a large scale magnetic field and giving rise to self-collimated jets. We term such a disc a Jet Emitting Disc, hereafter a JED. This model is unique in the sense that it provides both the physical conditions within the disc required to steadily launch jets and the distributions of all quantities in space (although the self-similar assumption used introduces some unavoidable biases). Several attempts to tackle the accretion-ejection connection were made in the past, e.g. Wardle and Königl (1993); Ferreira and Pelletier (1993a); Li (1995, 1996); Ogilvie and Livio (1998, 2001) and some are still being done (e.g. Campbell 2005 and references thererein), but these are not self-consistent models. For instance, while many of these works used also the self-similar ansatz, allowing a priori to take into account all dynamical terms, the authors made strong approximations in order to simplify the problem. The most common one was to assume a static vertical equilibrium, which leads to an underestimate of the magnetic compression and thereby to an overestimate of the allowed parameter space. Another example is the seminal work of Wardle and Königl (1993), where the authors replaced the mass conservation equation for the neutrals by the relation $\rho u_z = Cst$, which naturally always leads to a positive vertical velocity. It turns out, and this is the nasty thing about the accretion-ejection connection, that absolutely all dynamical terms are important and none can be dropped out in the equations. A self-consistent model must therefore take them all into account.

In this lecture, I first expose why the presence of a large scale vertical magnetic field should be expected in circumstellar accretion discs. Section 3 is then devoted to the major effect of such a field on standard accretion discs, namely the triggering of the magneto-rotational instability and the possible dead zone. Section 4 summarizes the current understanding of accretion-ejection systems with an emphasis on the underlying disc properties. Section 5 addresses new results on the star-disc interaction obtained with two MHD codes, VAC and PLUTO. It will be argued that such an interaction probably leads to a systematic spin up of the protostar. I will finally expose the basic concepts of the only model so far that allows a magnetic brake down of a low-mass protostar during its embedded stage.

2 Large scale magnetic fields in discs

2.1 Magnetic fields around YSOs

Where does this magnetic field come from? Let's face it: we don't know. There are two extreme possibilities. The first one considers that the field has been advected by the infalling material, leading to a flux concentration in the inner disc regions. The second one relies on a local dynamo action in the disc. Most probably, the answer lies between these two extreme cases, although I prefer the first possibility for the following reasons.

The necessary condition for launching a self-collimated jet from a Keplerian accretion disc is the presence of a large scale vertical magnetic field close to equipartition (Ferreira and Pelletier, 1995), namely

$$B_z \simeq 0.2 \left(\frac{M}{M_\odot}\right)^{1/4} \left(\frac{\dot{M}_a}{10^{-7} M_\odot/yr}\right)^{1/2} \left(\frac{r}{1\ \mathrm{AU}}\right)^{-5/4+\xi/2} \mathrm{G}, \tag{1}$$

where ξ is the disc ejection efficiency as measured by a varying disc accretion rate, namely $\dot{M}_a \propto r^\xi$. The value of this magnetic field is actually far smaller than the one estimated from the interstellar magnetic field assuming either ideal MHD $B \propto n$ or $B \propto n^{1/2}$ (Heiles et al., 1993; Basu and Mouschovias, 1994). Indeed, if we take the fiducial values $n \sim 1$ cm^{-3} and $B \sim 4$ μG observed within dense clouds and use the law $B \propto n^{1/2}$ (Crutcher, 1999), we get a magnetic field at 1 UA ranging from 10 to 10^3 G (depending on the density)! Thus, the main problem is to get rid off the magnetic field during the infalling stage. This issue is still under debate. However, it seems that building up accretion discs threaded by a large scale magnetic field seems to be rather straightforward (see for instance the 3D collapse simulations of Banerjee and Pudritz 2006).

Another indirect argument in favor of advection is provided by a statistical analysis. Using a sample of CTTS, Ménard and Duchêne (2004) found that CTTS are oriented randomly with respect to the local interstellar field. This sort of implies that magnetic fields play no role in enforcing the direction of the final angular momentum. However, sources with strong outflows display discs mostly perpendicular to the field (i.e. jets are aligned to it as first found by Strom et al. 1986), whereas sources with no jet detection are parallel. That could be a hint that, only in the former case, field dragging leads to the presence of inner JEDs. Finally, one might object that fields of the strength shown in Eq. (1) are impossible in accreting systems. Not only this would be devoid of any firm physical ground but stronger fields were actually already detected! Indeed, using the spectro-polarimeter ESPadOnS, Donati et al. (2005) found a ~kG field at 0.05 AU around FU Ori (a field that is actually larger than equipartition!).

2.2 Magnetic field advection in SADs

We thus assume that the outer parts of accretion discs are threaded by a large scale vertical (B_z) magnetic field of some unknown strength. The presence of such a field is the outcome of the infalling stage. But the actual field distribution in the disc, namely the function $B_z(r)$, depends on the interplay between field advection due to the accretion motion and diffusion: one needs an accretion disc theory.

A Standard Accretion Disc (SAD) is a disc where a turbulent viscosity ν_v allows an outward transport of angular momentum which then drives an inward accretion motion. By construction, the effective Reynolds number is $\mathcal{R}_e = ru_r/\nu_v \sim 1$ (Pringle, 1981). This "viscosity" must be of turbulent origin as collisions between particles provide a totally negligible normal viscosity. Now, in turbulent media, all transport coefficients are usually comparable leading also to heat (conductibility) and magnetic fields (diffusivity) turbulent transport. One consequence is an effective magnetic Reynolds number $\mathcal{R}_m = ru_r/\nu_m \simeq \mathcal{R}_e$, with $\nu_m \simeq \nu_v$ the turbulent magnetic diffusivity. But having $\mathcal{R}_m \sim 1$ in a disc implies that the poloidal field is straight, almost purely vertical (Heyvaerts et al., 1996). Under these circumstances, no magnetized jets can be launched from a SAD as the field is not bent enough[2]. This is consistent indeed with a SAD, as found by Lubow et al. (1994a). They investigated the advection of a large scale magnetic field by a SAD and, indeed, always obtained straight magnetic field lines. They concluded that, unless the magnetic Prandtl number $\mathcal{P}_m = \nu_v/\nu_m$ is unrealistically high, no magneto-centrifugal winds could be launched from a SAD. Again, the only torque taken into account was the viscous torque: the torque due to the magnetic field was simply neglected. Thus, according to the above arguments, they could only obtain $\mathcal{R}_m \sim \mathcal{R}_e \sim 1$, namely straight field lines. But note that $\mathcal{R}_m \sim 1$ does not mean that the magnetic field lags behind while the disc material is accreted. This is what one gets when rigid boundary conditions are applied on potential fields, as was done by these authors. In fact, as a result of the interplay between advection (due to accretion) and turbulent diffusion, the large scale magnetic field scales as $B_z \propto r^{-\mathcal{R}_m}$: the field strength in a SAD is therefore increasing towards the center.

Let us define the disc magnetization $\mu = B_z^2/\mu_o P$ where P is the gas pressure, as a measure of the dynamical importance of the field. In order to launch jets, one must have equipartition fields, namely $\mu \sim 1$. Now, hydrostatic vertical equilibrium in a non self-gravitating disc gives

$$P = \frac{\dot{M}_a \Omega_K^2 h}{6\pi\nu_v} \propto r^{-3/2-\delta} \tag{2}$$

where $\dot{M}_a$ is the disc accretion rate, Ω_K the Keplerian rotation law and $h(r) \propto r^{\delta}$. Since δ is always close to unity in circumstellar discs (and in most discs

[2] Blandford and Payne (1982) showed that magneto-centrigugally jets require a magnetic field bent by more than 30^o with respect to the vertical. This requires a magnetic Reynolds number $\mathcal{R}_m \sim r/h$, much larger than unity in a keplerian disc.

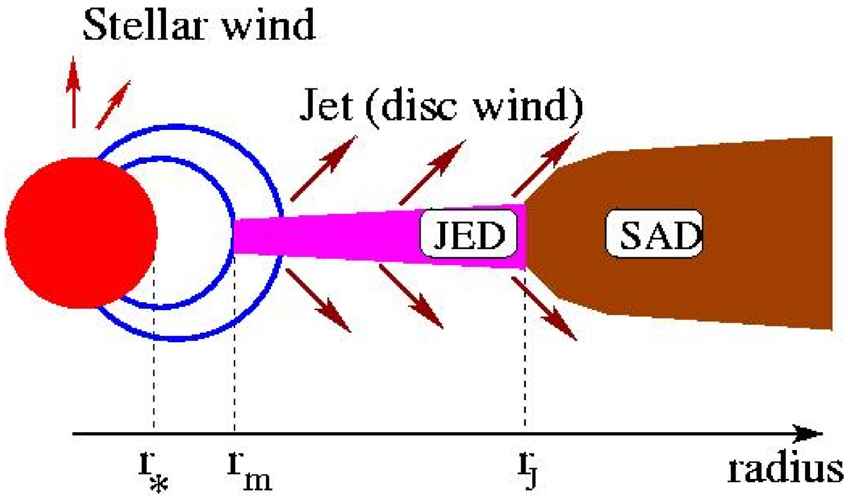

Fig. 1. Proposed paradigm for interpreting observations of the disc innermost regions around accreting protostars. A large scale magnetic field B_z is threading the disc an allows a transition from a standard, non ejecting disc (SAD) to a jet emitting disc (JED). The transition radius r_J is unknown and is probably varying in time for a single object and from one object to another. The truncation radius r_m, where the accretion flow is diverted and columns are formed, is discussed in section 5.

around compact objects). One therefore gets a disc magnetization μ which naturally increases towards the inner regions (Ferreira et al., 2006b).

2.3 A picture for the innermost disc regions

A picture, that can be applied to accretion discs around both young stars and compact objects, is now gradually emerging (fig. 1). A large scale magnetic field is thought to be dragged in by the accretion flow and concentrated in the innermost disc regions. Although of very small strength in these outer regions, such a field triggers the magneto-rotational instability (see Section 3 and Balbus 2003), producing thereby a standard accretion disc with no ejection (note that a thermally driven or photo-evaporated disc wind is of course clearly possible). When the field reaches equipartition at some transition radius r_J, the accretion flow switches from a SAD to a JED, giving rise to self-confined jets. The physics of this inner disc is then no longer governed by the radial transport of angular momentum but vertical instead.

Note that this simple picture is consistent with all theoretical aspects known to date. It is just a natural consequence of assuming the presence of a large scale vertical field. The real unknown is then the magnetic flux available in the disc. Because each object has its particular history, due to unique initial conditions, one might expect to have also different values of the transition radius r_J, from one object to another. Note also that this picture, initially designed for low-mass young stars (CTTS), should also hold for more massive objects, like Ae-Be protostars.

3 Accretion in the outer Standard Accretion Disc

3.1 The magneto-rotational instability

The magneto-rotational instability or MRI has been first described within an astrophysical context by Balbus and Hawley (1991) and has since been actively worked out by various groups. Accretion requires removal of angular momentum throughout the disc and in particular, also in the inner regions where the disc is not self-gravitating. The theory of SAD requires therefore a "viscous" stress tensor component such that $\sigma_{r\phi} = \rho\nu_v r\partial\Omega/\partial r$ where $\nu_v = \alpha_v C_s h$ is the turbulent viscosity, with C_s the local sound speed and α_v a parameter measuring the efficiency of that turbulent transport (Shakura and Sunyaev, 1973). One can express this slightly differently: we need to find a mechanism providing a stress $\sigma_{r\phi}$ such that $\alpha_v = \sigma_{r\phi}/P$ (where P is the thermal pressure measured at the disc midplane) is large enough, namely compatible with observational constraints. These are in fact poorly known in young stars. But statistical arguments based on the lifetime of accretion discs provides values of α_v ranging from 0.1 to 0.01.

How can this be done? If the flow is unstable and if this instability is self-sustained, then one might expect large fluctuations inside the medium with some correlations (somehow, like interactions between eddies mimicking collisions between particles). From the momentum equation, one can derive the following expression

$$\sigma_{r\phi} = \rho < u_r u_\phi > - < \frac{B_r B_\phi}{\mu_o} > \qquad (3)$$

where the symbol $<>$ describes some spatio-temporal average over the fluctuating velocity and magnetic field components. The first term in the rhs is called the Reynolds stress whereas the second the Maxwell stress.

In an un-magnetized ($B = 0$) inviscid rotating flow, the specific angular momentum is conserved ($\Omega r^2 = Cst$). The stability of such a flow is then subjected to the Rayleigh criterion. In a keplerian flow, any small deviation from the radial equilibrium gives rise to a restoring centrifugal force and the flow is stable. It is said that Hydrodynamic (HD) keplerian flows are linearly stable. But what about a large deviation? The search for a parametric HD instability has recently finally gone into an end. To make a long story short, it seems that, yes, under some circumstances, HD sheared flows can indeed be unstable but, no, the amount of angular momentum transport is in practice negligible. More precisely, the value that can be expected for α_v (due to correlations between u_r and u_ϕ only) cannot be larger than 10^{-6} (Lesur and Longaretti, 2005). Only magnetic fields could thus do the job in astrophysical systems.

The current globally accepted paradigm is that this is done through an MHD turbulence that is triggered and sustained within the disc. So far, MRI seems to be the best candidate for providing such a self-sustained turbulence in keplerian discs (see Balbus 2003; Lesur and Longaretti 2007 and references therein). But note that there are many more MHD instabilities possible in discs and that this issue is not yet totally settled (see for instance Blokland et al. 2007).

The physical idea being MRI is the following. Take an ideal MHD disc in radial keplerian equilibrium and make a small, say negative, radial displacement ξ_r (so that $u_r = d\xi_r/dt$). This displacement bends the field lines which then react by a magnetic torque. This torque acts to decrease the angular velocity, thereby enhancing the displacement. Thus, the presence of a magnetic field (the argument holds for both B_z or B_ϕ) introduces a destabilizing agent and a keplerian flow is no more stable! To some extent though. Indeed, the magnetic tension does introduce a stabilizing effect. To put the things into more light, we can write the radial displacement equation

$$\frac{d^2\xi_r}{dt^2} \simeq (\Omega^2 - \Omega_K^2)r + \frac{F_r}{\rho} \tag{4}$$

where F_r is mainly the magnetic tension effect. If this last term is dominant, then the mechanism described above, which relies on angular momentum transfer between matter and the field, cannot work anymore. Let us examine this into more details.

Take the full set of ideal MHD equations, simplify them as much as you can and then linearize them by looking for modes with no radial propagation, namely $e^{i(\omega t - kz)}$ (since the instability is local). We therefore take an axisymmetric, incompressible[3], isothermal disc threaded by a homogenous vertical field B_z. After some lengthy algebra, we obtain the dispersion relation of our modified Alfvén waves

$$\omega^4 - \omega^2\left(2k^2v_A^2 + \frac{d\Omega^2}{d\ln r} + 4\Omega^2\right) + k^2v_A^2\left(k^2V_A^2 + \frac{d\Omega^2}{d\ln r}\right) = 0 \tag{5}$$

[3] The mechanism described above relies on pure magnetic effects: we should therefore be able to catch them with modified Alfvén waves.

where $v_A = B/\sqrt{\mu_o \rho}$ is the Alfvén speed and $\Omega = \Omega_K$ is the disc material angular velocity. An instability occurs whenever the pulsation ω is imaginary [4], namely when

$$k^2 v_A^2 < -\frac{d\Omega^2}{d \ln r} \quad (6)$$

In a keplerian disc there is thus a minimum wavelength λ_{min} above which all wavelengths are unstable. The existence of λ_{min} (or k_{max}) comes from the stabilizing effect of the magnetic tension. It can then be easily shown that the most unstable mode (thus the most probable) has a dynamical growth rate such that $kv_A = (\sqrt{15}/4)\Omega$: MRI is therefore a highly dynamical instability. However, for this instability to be of any relevancy, this most unstable mode must have a wavelength smaller than the disc scale height. Indeed, if $\lambda > h$, the disc will be barely affected. As a consequence, MRI is expected to set in only if $\mu = V_A^2/\Omega^2 h^2$ is smaller than a number close to unity (or a plasma beta greater than unity, Balbus and Hawley 1991).

This analytical prediction has been verified with MHD numerical simulations by several authors (see Balbus 2003 and Lesur and Longaretti 2007 for a review). These simulations were able to follow the linear stage of the instability and compare the growth rates with analytical calculations. The mechanism of the instability has thus been confirmed. More importantly, MRI leads indeed, in its non linear stage, to a self-sustained turbulent situation with anomalous transport of angular momentum. The main transport is due to the Maxwell stresses, the Reynolds stresses being typically 10 times smaller. However, the value of the turbulent parameter α_v derived from these simulations is a complicated matter as it depends on the numerical resolution used, the boundary conditions and both the initial value and geometry of the magnetic field (see scalings found in Lesur and Longaretti 2007 and Pessah et al. 2007). It is known for instance that higher values of α are obtained with vertical fields but there is no simulation to date able to run with a large scale magnetic B_z field (ie non zero flux) as initial condition. A very interesting scaling has been however found, with possible direct implications for our picture of the inner regions of accretion discs. Indeed, it seems that

$$\alpha_v \simeq \mu^{1/2} \quad (7)$$

holds for small values of the magnetic field ($\mu \ll 1$). Taking this expression at face value (risky) and the radial profile of the disc magnetization $\mu(r)$ discussed in the previous section, we obtain that at the transition radius r_J one has $\alpha_v \sim \mu \sim 1$. These are actually the best conditions for magnetically

[4] namely $\omega^2 = i\gamma^2$ where γ is the instability growth rate.

driving self-confined jets from the JED (see for instance the required values of the magnetic diffusivity in JEDs as found in Ferreira 1997).

3.2 The dead zone

The above calculation of the MRI assumed ideal MHD. But this is valid only if there is a good coupling between the disc material (mostly neutrals) and the magnetic field. This actually requires a minimum level of ionization in the disc. Taking this into account requires to deal with non ideal effects. Without going into lengthy calculations, we can instead estimate their importance. Indeed, MRI will set in only if the Alfvénic time scale is shorter than the diffusion time scale due to electron-ions collisions (providing an Ohmic resistivity η). This criterion translates into a magnetic Reynolds number $\mathcal{R}_m = hV_A/\eta$ that must be greater than unity (using the fact that the maximum wavelength is the disc scale height h). Note that this line of arguments remains valid for any MHD instability triggered inside the disc, not only MRI. While a standard accretion disc relies on turbulent "viscosity" (Shakura and Sunyaev, 1973), a jet emitting disc (hereafter JED) requires a turbulent magnetic diffusivity (Ferreira and Pelletier, 1993a). Both accretion flows must then verify $\mathcal{R}_m > 1$ (or even larger, see Fleming et al. 2000).

Using Spitzer's value for η one gets [5]

$$\mathcal{R}_m \simeq 10^{13}\zeta^{1/2}x_e \left(\frac{\varepsilon}{0.01}\right)\left(\frac{r_o}{1\,AU}\right) \tag{8}$$

where $\varepsilon = h/r$ is the disc aspect ratio, x_e is the ionization fraction and ζ a dimensionless quantity that depends on the nature of the accretion flow. For a SAD $\zeta = \alpha_v$ is the viscosity parameter, much smaller than unity (Gammie, 1996), whereas $\zeta = \mu$ for a JED, where $\mu = B_z^2/(\mu_o P) \sim 1$ is the disc magnetization (Ferreira and Pelletier, 1995). Thus, around 1 AU, an ionization fraction x_e greater than 10^{-13} for a JED ($\alpha_v^{-1/2}$ times larger in a SAD) is necessary in order to allow MHD turbulence, hence sustain accretion. Gammie (1996) showed that a zone within the accretion disc could have a too low ionization fraction for this coupling to occur. The reason is the following. In the outer parts of the disc, the column density is low enough to allow ionization by cosmic rays (with an ionization rate $\xi_{CR} \simeq 10^{-17}$ s^{-1}, Spitzer and Tomasko 1968), whereas in the innermost parts where $T_o > 10^3$ K, collisional ionization is enough to maintain $x_e \simeq 10^{-11}$. But there is an intermediate zone where the disc would be both too dense (column density $\Sigma > 10^2$ g cm^{-2}, Umebayashi

[5] Gammie (1996) gives an equivalent expression but using Hayashi (1981)'s resistivity and the central disc temperature T_o instead of the disc aspect ratio $\varepsilon = h/r$.

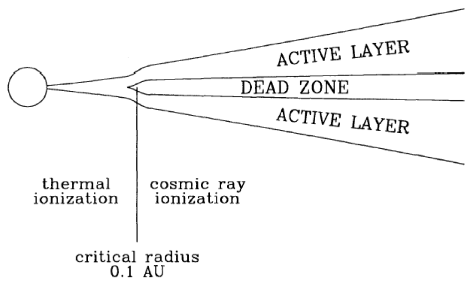

Fig. 2. Sketch of the radial and vertical stratification of a SAD, according to Gammie (1996). In the outer dense parts of the disc, accretion can only proceed at the surface, where cosmic ray ionization is enough to maintain material well coupled to the magnetic field, thereby allowing MRI to set in. From the point of view of accretion, the zone located at the disc midplane is thus expected to be "dead".

and Nakano 1981) and too cold ($T_o < 10^3$ K). This defines the "dead zone" where MRI cannot set in and, thereby, no accretion is believed to be possible (fig. 2). In fact, accretion remains possible at the disc upper layers, ionized by cosmic rays, but material at the disc midplane would not accrete.

According to Gammie (1996), the outcome of these dead zones would be unsteady accretion events due to the slow mass accumulation at their outer edge. This is easy to understand. The disc accretion rate is $\dot{M}_a = 2\pi r \Sigma u_r$ and thus scales like Σ_a in the active layer. By construction $\Sigma_a \simeq 10^2$ g cm^{-2} is a constant and is therefore not following the usual radial scaling that is required to maintain $\dot{M}_a$ constant through the disc. Thereby, one has $\dot{M}_a(r)$ and mass conservation implies unsteady events. Despite the fact that this has been said more than 10 years ago, there is (to my knowledge of course) no model nor numerical simulation addressing this fundamental question. Note also that dead zones would be very interesting for planet formation as they provide a region where dense material has time (since there is no accretion) to create planet cores (see eg. Glassgold et al. 1997; Matsumura and Pudritz 2003, 2006).

Although everybody agrees on these issues, the localization and radial extent of the dead zone is still a matter of debate. The reason is twofold: (1) the difficulty raised by using an underlying accretion disc model consistent with the calculation of the ionization structure; (2) the uncertainty on the composition of the circumstellar material, in particular dust and metallic grains.

For instance, Glassgold et al. (1997) use for the disc a minimum solar nebula approximation, namely no accretion and rather ad-hoc prescriptions for the radial profiles $\Sigma(r)$ and $T(r)$ but take into account X-ray illumination. They obtain a dead zone that could range between 1 and 10 or even 30 AU. Fromang et al. (2002) take a SAD illuminated by X-rays as well but obtain results that are highly dependent on the value of the turbulent parameter α_v. For instance, for $\alpha_v = 10^{-3}$ all the disc is dead whereas there is no dead zone at all for $\alpha_v > 0.1$. Matsumura and Pudritz (2003) argued that using a passive disc model with no accretion, as developed by Chiang and Goldreich (1997), is more correct for describing a dead zone. Considering then X-rays, cosmic rays and radioactivity as ionization sources, these authors found a dead zone between 0.2 and 3 AU. However, it is probably X-rays rather than cosmic rays that are responsible for most of disc ionization, a point made early on by Glassgold et al. (1997).

But, again, one should be cautious as the underlying disc model is also assuming a value for the disc column density Σ and this is the most important control parameter. Moreover, it is not clear at all that the vertical stratification envisioned (actually assumed) in the dead zone could indeed be maintained. For instance, Fleming and Stone (2003) performed a 3D simulation of the MRI in a stratified disc using the shearing sheet approximation and with an Ohmic resistivity $\eta(z)$. The profile used was such that only the disc upper layers where MRI unstable while the disc midplane was stable. But the non linear stage of the simulation showed that the disc midplane was also providing accretion (although at a smaller pace than the surface). This is due to the strong Maxwell stresses developed at the surface that are coupled to Reynolds stresses which, in turn, generate a turbulent mixing and mass exchange between the two regions: the dead zone is not quite so dead.

4 Jets and Jet Emitting Discs

4.1 The three basic steady state jet models

There are 3 classes of magnetized jet models (see fig. 3), depending on the origin of the ejected material and on which energy reservoir is tapped:

Stellar winds assume that field lines are anchored into the rotating star. Both mass and energy are extracted from the star itself (Weber and Davis, 1967; Hartmann et al., 1990; Sauty et al., 2002).

disc-winds assume that field lines are anchored into the accretion disc alone. Both mass and energy is therefore extracted from the underlying disc. If jets are launched from a large radial extension in the disc, then one obtains the situation calculated by Blandford and Payne (1982) and extended by Vlahakis et al. (2000). If, on the contrary, the large scale magnetic field is assumed to thread the disc only in a tiny region at the disc inner boundary, then one gets the "X-wind" picture (Shu et al., 1994; Shang et al., 2002)

Magnetospheric winds are winds that are produced at the interaction between the accretion disc and the protostellar magnetic field. They can carry away mass coming from the disc and tap rotational energy from the star

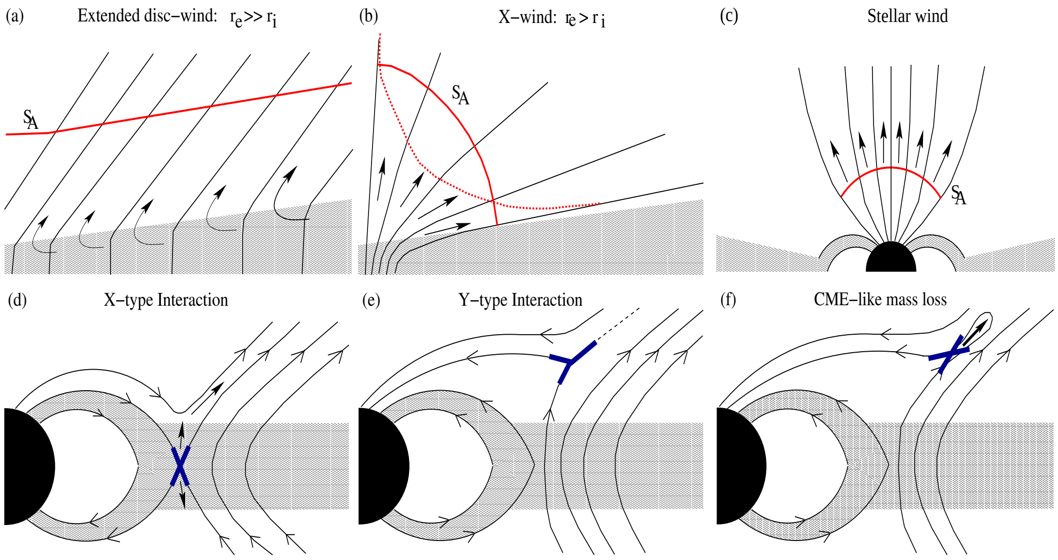

Fig. 3. **Top:** Classes of published stationary MHD jets for YSOs. When the magnetic field is threading the disc on a large radial extension (a: extended disc wind) or a small disc annulus (b: X-wind), jets are accretion-powered. They are mostly pressure-driven when the field lines are anchored onto a slowly rotating star (c: stellar wind). The corresponding Alfvén surfaces S_A have been schematically drawn (thick lines). In the X-wind case, two extreme shapes have been drawn: convex (solid line) and concave (dashed). **Bottom:** Sketch of the two possible axisymmetric magnetospheric configurations: (d) X-type neutral line driving unsteady Reconnection X-winds, when the stellar magnetic moment is parallel to the disc field; (e) Y-type neutral line (akin the terrestrial magnetospheric current sheet) when the stellar magnetic moment is anti-parallel (or when the disc field is negligible). (f) A CME-like ejection is produced whenever the magnetic shear becomes too strong in a magnetically dominated plasma. Such a violently relaxing event may occur with any kind of anti-parallel magnetospheric interaction (even with an inclined dipole). The thick lines mark the zones where reconnections occur. Taken from Ferreira et al. (2006a).

(Ferreira et al., 2000; Matt et al., 2002).

All these models suffer from simplifying assumptions but disc-wind models have been more developed: the capability of T Tauri stars to drive massive jets still needs to be proven (see arguments developed in Ferreira et al. 2006a), while magnetospheric winds require numerical simulations that are difficult to control. But all these jet models share the same physics and are therefore described by the same set of MHD equations.

Magnetized jets are assumed to be steady, axisymmetric, non-relativistic and described as a single fluid within the ideal MHD framework. They can be viewed as made of nested magnetic surfaces (defined by a constant magnetic flux $a(r, z) = Cst$) that are anchored on a rotating object (star, disc or some star/disc interface). For convenience, I hereafter focus on the disc (fig. 4). The

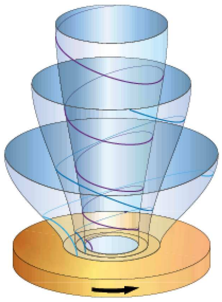

Fig. 4. Axisymmetric jets are made of magnetic surfaces of constant magnetic flux nested around each other and anchored in the disc. Each surface behaves like a funnel whose shape depends on the transfield equilibrium. Solving the jet equations requires to specify several quantities (see text).

axisymmetric poloidal magnetic field writes

$$\vec{B}_p = \frac{1}{r}\nabla a \times \vec{e_\phi} \tag{9}$$

and is usually assumed to be of bipolar topology. The toroidal magnetic field is tremendously important: it produces the magnetic braking on the underlying disc, feeds the jets with energy (MHD Poynting flux) and provides the so-called “hoop-stress” that confines them. This hoop-stress can be understood as the Laplace force due to presence of a vertical current I, flowing inside a magnetic surface, and the toroidal field. Note however that, because currents must be closed, not all field lines can be self-collimated (Okamoto, 2003). Now, jet acceleration depends also on this current. Indeed, since ejected mass is accelerated by MHD forces (see Eq. 15 in Ferreira 1997), there is a transfert from the MHD Poynting flux to the kinetic energy flux: $I = 2\pi r B_\phi/\mu_o$ decreases as we go downstream, allowing thereby a current closure. So, from this simple argument, two important issues of MHD jets can be understood: (1) the degree of asymptotic collimation (value of I) is related to the acceleration zone, which is located upstream; (2) any approximation made on the poloidal current has dramatic consequences on jet dynamics.

The source of such an electric current is the unipolar induction effect (also known as the Barlow wheel experiment). Any rotating conductor embedded in a magnetic field will produce an electromotive force, which drives a current. So, the current that feeds magnetized jets is driven by the underlying *resistive* accretion disc. This is a very strong constraint, neglected in studies where the disc is a mere boundary condition. A very important byproduct is the possibility to launch *asymmetric jets*. Indeed, a symmetric (bipolar) magnetic configuration threading the disc produces one electromotive force. But there are two independent electric circuits, one for each jet. A slight mismatch between the two ”resistances” (ie, the ambient media) can produce two different jets with different mass loads and power (Ferreira and Pelletier, 1995).

A magnetized jet is described by 8 unknown variables: density ρ, velocity $\vec{u}$, magnetic field $\vec{B}$ (flux function a and toroidal field B_ϕ), pressure P and temperature T (the energy equation is usually replaced by a polytropic state equation). There are 8 partial differential equations (PDE) allowing to solve the full 2D problem: mass conservation, momentum conservation, induction equation, equation of state (perfect gas) and a polytropic law. Because jets are non-dissipative structures in ideal MHD, there are 5 invariants along each magnetic surface (Jacobi integrals) fixed by boundary conditions. These invariants are

(1) Mass flux to magnetic flux ratio $\eta(a)$

$$\vec{u}_p = \frac{\eta(a)}{\mu_o \rho} \vec{B}_p \tag{10}$$

(2) Magnetic surface rotation $\Omega_*(a)$

$$\Omega_*(a) = \Omega - \eta \frac{B_\phi}{\mu_o \rho r} \tag{11}$$

(3) Total specific angular momentum $L(a) = \Omega_* r_A^2$

$$L(a) = \Omega r^2 - \frac{r B_\phi}{\eta} \tag{12}$$

(4) Total specific energy $E(a)$

$$E(a) = \frac{u^2}{2} + H + \phi_G - \Omega_* \frac{r B_\phi}{\eta} - \mathcal{H} \tag{13}$$

(5) Specific entropy $K(a)$

$$P = K(a) \rho^\gamma \tag{14}$$

where r_A is the cylindrical radius where the poloidal velocity reaches the poloidal Alfvén velocity, H the enthalpy, $\mathcal{H}$ an heating term (zero if jets are adiabatic) and γ the polytropic index.

Solving the set of MHD equations requires to specify 8 boundary conditions. One of them is obtained by assuming $\Omega_* = \Omega_K$, ie. magnetic surfaces rotating at Keplerian speeds (sub-Keplerian speeds for the ejected mass). The second fixes the entropy K: for instance $K = 0$ in Blandford and Payne (1982) "cold" jets and $K \neq 0$ in Contopoulos and Lovelace (1994); Vlahakis et al. (2000) and all numerical simulations of jets. Three other constraints arise from MHD regularity conditions. Stationarity requires that, as ejected material is accelerated along each magnetic surface (fig. 5), it becomes super slow-magnetosonic (SM), then super-Alfvénic (A) and finally super fast-magnetosonic (FM). This

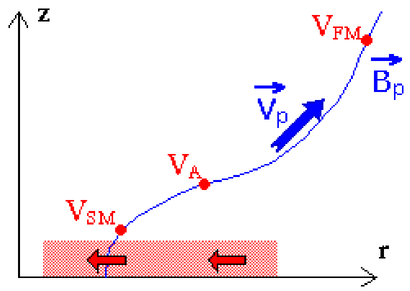

Fig. 5. Once material has left the resistive MHD zone, it is frozen in a particular field line and encounters the three MHD critical points. Note however that the existence of a smooth transition between resistive and ideal MHD regimes already selects a region in the parameter space.

leaves us with *3 free and independent* boundary conditions that must be specified at each magnetic surface (anchored at r_o). Note the crucial fact that, because jet models are in ideal MHD, they do not address the mass loading by the underlying disc: the mass flux is simply parametrized by η.

Instead of solving 8 coupled equations, the problem can be formulated in a much more compact way using the invariants. This leads us to the following Grad-Shafranov equation for an adiabatic jet

$$\nabla \cdot (m^2 - 1)\frac{\nabla a}{\mu_o r^2} = \rho \left\{ \frac{dE}{da} - \Omega \frac{d\Omega_* r_A^2}{da} + (\Omega r^2 - \Omega_* r_A^2)\frac{d\Omega_*}{da} - \frac{C_s^2}{\gamma(\gamma-1)}\frac{d\ln K}{da} \right\} + \frac{B_\phi^2 + m^2 B_p^2}{\mu_o}\frac{d\ln\eta}{da} \quad (15)$$

where $m^2 \equiv u_p^2/V_{Ap}^2$ is the Alfvénic Mach number and $C_s^2 = \gamma k_B T/\mu m_p$ is the jet sound speed. This equation provides the transverse equilibrium (ie. the degree of collimation) of a magnetic surface. Mathematically, this is a mixed-type PDE providing $a(r, z)$ for a given set of invariants (η, Ω_*, L, E and K). Indeed, the flow is hyperbolic between the cusp (hopefully located at the disc surface) and the SM surface, elliptic between the SM and the FM surface, and again hyperbolic downstream the FM surface. Now, while elliptic zones are fully determined by boundary conditions (like eg the Laplace equation), hyperbolic domains require to be computed as initial value problems. We have therefore to face two overwhelming difficulties: (1) both shape and localization of the parabolic surfaces (SM and FM) are unknown; (2) the choice of boundary conditions is terribly large.

This is the reason why *there is no consistent solution for general 2D steady-state MHD jets* yet. This is an unsolved mathematical problem. Either one solves the time-dependent problem with an MHD code, or one uses some trick. The method using separation of variables leads to self-similar solutions that incorporate all dynamical terms (all the physics) but with biases due to the impossibility to treat radial boundary conditions. Other methods address this point but then usually crudely simplify some equations: for instance by assuming the shape of the magnetic surfaces (Lery et al., 1999) or by solving only in the elliptic domain (Fendt and Camenzind, 1996; Shang et al., 2002) [6].

This variety of approaches has probably contributed to the idea that the jet phenomenon was still not understood. While, indeed, there is no ultimate model addressing everything, it is however fair to say that most of the basic physical ingredients have been understood for steady state jets. We understand how jets can be accelerated, why collimation occurs and we are able to relate the asymptotic jet velocity to the initial available energy. This is enough to compare models with observations and derive strong constraints. But the importance of the jet internal stratification on its collimation, propagation and instability properties is still a very active research field.

4.2 Constraints from T Tauri jets

A review of the observational properties of outflows from young stars can be found in Cabrit (2002). A consensus seems to have slowly emerged these last years. It is now being gradually accepted that most of the ejected mass in jets comes from the accretion disc, even if other ejection events are coexisting. The major uncertainties lies in the estimate of the mass fluxes ratio, but an ejection to accretion rate ratio $\dot{M}_j/\dot{M}_a$ of about 10% is indeed hardly compatible with stellar winds (see Ferreira et al. 2006a for more details and the reviews made in Protostars & Planets V for a slightly different perspective).

Basically, explaining the jet phenomenon from low-mass young stars with only stellar winds faces the quite overwhelming task of finding a means to produce the observed huge mass loss rates. Because ejected material, in this case, is almost at the stellar surface hence deep down in the stellar potential well, only pressure forces can lift it up. This pressure cannot be thermal as it would lead to temperatures of order 10^6 K. Given the large mass loss rates, it would lead to emission losses that are not observed. Another possibility is pressure due to Alfvén waves, as already proposed in the 80's (Hartmann and MacGregor, 1980; DeCampli, 1981; Hartmann et al., 1982). If one assumes that these waves are present (triggered for instance by the accretion shock) then one can safely expect that some of their energy will be converted to the plasma, acting as some pressure but much less dissipative. However, the efficiency of this conversion cannot be 100% which hints to the fact that the total energy

[6] In the X-wind case, boundary conditions were specified at the Alfvén surface itself and not at the FM surface as it should. Some "connection" has been made afterwards to match the sub-Alfvénic solution to an asymptotic solution.

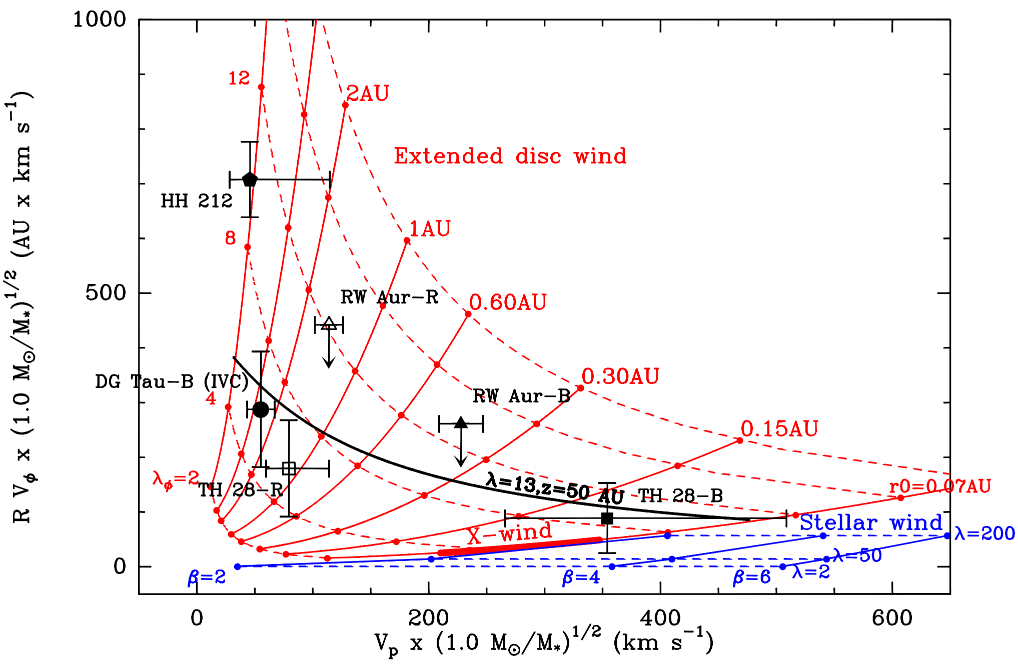

Fig. 6. Comparison of predicted specific angular momentum vs. poloidal velocities with observations of T Tauri microjets. Full and dashed curves show expected theoretical relations for MHD disc and stellar winds. Plotted in symbols are jet kinematics measured at distance $z \simeq 50$ AU in the DG Tau, RW Aur, and Th 28 jets. The infrared HH 212 jet is also shown for comparison (taken from Ferreira et al. 2006a). Note that the observed values were corrected by the inclination of the system.

available in the waves must be much larger than that provided to the flow. This can estimated quite easily.

The asymptotic poloidal velocity of stellar winds almost entirely depends on the asymptotic value of a parameter $\beta(> 2)$ alone that describes all pressure effects (Ferreira et al., 2006a), namely

$$v_p \simeq 250\, \sqrt{\beta - 2} \left(\frac{M}{M_\odot}\right)^{1/2} \left(\frac{R_*}{3R_\odot}\right)^{-1/2} km/s \qquad (16)$$

where R_* is the stellar radius. Observations tell that the range of jet poloidal speeds require $\beta \simeq 2-4$. The total power transferred to the two jets via either thermal or wave pressure gradients is then simply

$$L_\beta \simeq \bar{\beta}\frac{GM\dot{M}_j}{R_*} \simeq \bar{\beta}\left(\frac{\dot{M}_j}{\dot{M}_a}\right) L_{acc} \tag{17}$$

where $\bar{\beta}$ is the average value of β through the jet and $L_{acc} = GM\dot{M}_a/R_*$ the accretion luminosity onto the star. Taking $\bar{\beta} = 3$ and a one-sided ejection to accretion mass ratio ranging from 1% to 10% gives a total power that must be as high as 3% to 30% of the accretion luminosity. But this transferred power is itself a fraction of the total power that must be available to the ejected material. Since this power is presumably stored in some accretion-related turbulence, we can write $L_\beta = \eta L_{turb}$ where η is the efficiency of energy conversion. For instance DeCampli (1981) obtained an efficiency of only roughly 20% with a prescribed radial field and assuming undamped waves. This is a conservative value considering that the ejected plasma will also loose energy though radiation. Thus, the total net power L_{turb} that must be available for stellar winds must be as high as 15% to 150% of the accretion luminosity! Clearly this is very uncomfortable. Note that this poses an energetic problem only if one insists on explaining all the jet mass loss with stellar winds. This probably implies that "accretion powered" stellar winds (Matt and Pudritz, 2005a) only carry a small fraction of the observed jet mass flux in T Tauri stars.

On the other hand, published wide-angle disc winds from the co-rotation ("X-winds", Shu et al. 1994) have kinematical properties which are inconsistent with observations. The range in poloidal speeds is narrower than observed in some objects (50-400 km/s), and the frequent steep decline in poloidal speed towards jet edges is not explained. Moreover, if recent high resolution kinematic signatures are indeed probing jet rotation (Bacciotti et al., 2002; Woitas et al., 2005), then both X-winds and stellar winds are ruled out as the main mass loss mechanism for CTTS. Note however that one point in their favor would be that they make use of stellar magnetic fields only (that are observed), while we know less about the magnetic fields and ionization state of the disc. Anyway, this observational difficulty of the X-wind may be seen in fig 6 which is a plot of the specific angular momentum carried by one magnetic surface (anchored at a disc radius r_o) as a function of the jet poloidal speed. In this plot, red solid lines define a constant anchoring radius r_o whereas dashed lines a constant magnetic lever arm $\lambda \simeq 1 + 1/2\xi$ (where $\dot{M}_a \propto r^\xi$). Solving these problems for the X-wind would require a strong modification in the Alfvén surface and/or in the collimation of outer streamlines (to allow entrainment of slow ambient gas within 20-30 AU of the jet axis).

Observations clearly favor self-confined jets launched from some radial extension in the disc (say from 0.1 to 0.5-3 AUs). However, cold models (with

$\xi \sim 0.01$, hence $\lambda \sim 50$) are excluded as their large magnetic lever arm ($\lambda \geq 50$) predicts excessive jet rotation on observed scales (and too large velocities, Garcia et al. 2001). Only "warm" solutions with $\xi \sim 0.1$ ($\lambda \sim 10$) are fully compatible with current observations (mass flux, velocities, collimation). Such models require heat input at the upper disc surface layers in order to allow more mass to be loaded onto the field lines

The origin of this heat deposition remains an open question. It must lead to an increase of the disc temperature at the surface that must be several times higher than the midplane temperature (which depends of course on the disc radius and accretion rate). But then, for a given radius and accretion rate, one needs to compute all cooling and illumination processes in order to precisely determine how much energy deposition is required. Preliminary results show that this energy deposition cannot be due solely to illumination by stellar UV and X-ray radiation (Garcia et al., in prep). Alternatively, the turbulent processes responsible for the required magnetic diffusivity inside the disc might also lead to a turbulent vertical heat flux leading to dissipation at the disc surface layers. It is interesting to note that in current MHD simulations of the magneto-rotational instability a magnetically active "corona" is quickly established (Stone et al., 1996; Miller and Stone, 2000). Although no 3D simulation has been done with open magnetic field lines, this result is rather promising. Indeed, it might be an intrinsic property of the MHD turbulence in accretion discs, regardless of the launching of jets (see also observational arguments developed by Kwan 1997 and Glassgold et al. 2004).

4.3 Accretion-Ejection systems

In a standard accretion disc (ie. steady-state, no mass loss), the accretion rate is constant with radius. In a Jet Emitting Disc (JED) it cannot be anymore constant and one defines the ejection index as

$$\xi = \frac{d \ln \dot{M}_a}{d \ln r} \tag{18}$$

This parameter (which can vary in the disc) measures the local ejection efficiency. If this efficiency is constant throughout the disc (of inner and outer radii r_{in} and r_{out}), then mass conservation provides the ratio of ejection to accretion rates

$$\frac{2\dot{M}_j}{\dot{M}_a} = 1 - \left(\frac{r_{in}}{r_{out}}\right)^{\xi} \simeq \xi \ln \frac{r_{out}}{r_{in}} \tag{19}$$

Observations show that this ratio varies between 0.01 and 0.1, which implies similar values for the ejection efficiency ξ. A complete theory must provide the allowed values of ξ as a function of the disc properties. One must take into account the full 2D problem and *not* treat the disc as infinitely thin as in a standard disc theory. This requires to treat the physics of both accretion and ejection, from a quasi-Keplerian disc thread by a large scale magnetic field of bipolar topology.

4.3.1 Accretion

Steady-state requires the presence of an anomalous magnetic diffusivity ν_m allowing the accreting (and rotating) material to cross the magnetic field lines. Ambipolar diffusion could do the job but fully ionized discs (ie. around compact objects) also display self-collimated jets. So, if one looks for an "universal" model, one should use another prescription. We assume that an MHD turbulence, triggered and maintained inside the disc, can be described by a local "effective" diffusivity ν_m. Since rotation is much faster than accretion, there must be a higher dissipation of toroidal field than poloidal one. *A priori*, this implies a possible anisotropy of the magnetic diffusivities associated with these two directions, poloidal ν_m and toroidal ν'_m. Besides, such a turbulence might also provide a radial transport of angular momentum, hence an anomalous viscosity ν_v. To summarize, at least three anomalous transport coefficients are necessary to describe a stationary structure. We will use the following dimensionless parameters defined at the disc midplane:

$$\begin{aligned} \alpha_m &= \frac{\nu_m}{v_A h} \quad \text{level of turbulence} \\ \chi_m &= \frac{\nu_m}{\nu'_m} \quad \text{degree of anisotropy} \\ \mathcal{P}_m &= \frac{\nu_v}{\nu_m} \quad \text{magnetic Prandtl number} \end{aligned} \tag{20}$$

where $v_A = B_o/\sqrt{\mu_o \rho_o}$ is the Alfvén speed. A conservative picture of 3D turbulence would translate into $\alpha_m < 1$, $\chi_m \sim 1$ and $\mathcal{P}_m \sim 1$. But as stated before, the amount of current dissipation may be much higher in the toroidal direction, leading to $\chi_m \ll 1$. Moreover, it is not obvious that α_m must necessarily be much smaller than unity. Indeed, stationarity requires that the time scale for a magnetic perturbation to propagate in the vertical direction, h/v_A, is longer than the dissipation time scale, h^2/ν_m. This roughly translates into $\alpha_m > 1$. Thus, we must be cautious and should freely scan the parameter space defined by these turbulence parameters. The magneto-rotational instability (Balbus and Hawley, 1991) may well be the main source of turbulence in discs. Unfortunately, there is to date no MHD simulation (3D, with a non-zero poloidal flux) providing expected values for the "viscosity": the shearing sheet approximation brakes down in a few orbital times (Miller and Stone, 2000). No value for the magnetic diffusivity has ever been computed.

One important thing to realize is that the torque due to the large scale field (in a JED, this translates into the torque due to the jets) is always much larger than the "viscous" (turbulent) torque. Their ratio at the disc midplane writes

$$\Lambda = \frac{\text{magnetic torque}}{\text{viscous torque}} \simeq \frac{B_\phi^+ B_z/\mu_0 h}{\alpha_v P/r} \simeq \frac{r}{h}\frac{B_\phi^+}{B_z}\frac{B_z^2/\mu_0}{\alpha_v P} \sim \frac{r}{h} \tag{21}$$

where $h(r)$ is the local disc vertical scale height (defined with the gas pressure P) and B_ϕ^+ is the toroidal field at the disc surface. Now, such an engine can only be steady with a field close to equipartition, namely $B_z^2/\mu_0 \sim P$. A much smaller field triggers the magneto-rotational instability, while a much larger field simply forbids ejection (huge vertical compression, Ferreira and Pelletier 1995). Also, unless very special conditions are met (strong tubulence anisotropy, see Casse and Ferreira 2000a for more details), the toroidal field at the disc surface is of the same order than the vertical field.

4.3.2 Ejection

The disc is accreting because (i) the large scale magnetic field is extracting its angular momentum and (ii) mass can diffuse through that field. Now, if jets are to exist, then mass that has been loaded onto the field lines (ie. material located at the disc surface in ideal MHD regime) must be azimuthally accelerated by the magnetic field. This means that the magnetic torque must go from negative at the disc midplane to positive at the disc surface. This necessary condition requires a decrease of the radial current density J_r on a disc scale height (Ferreira and Pelletier, 1995). Since the magnetic field is compressing the disc (downward magnetic pressure gradient due to both bending B_r and shearing B_ϕ), the only force capable of pushing material up is the gas pressure gradient. This has two important consequences: (1) all terms in the disc vertical equilibrium are dynamically important and neglecting one leads to a wrong estimate of the disc ejection capability [7] ; (2) the correct range of ξ highly depends on the temperature profile $T(z)$.

Solving the full 2D problem without any approximation (keeping all dynamical terms) can be done using a method of variable separation. Such a separation is made by looking for solutions with the same functional dependency in (r, z)

[7] This is the reason why Li (1995) obtained a huge jet parameter space. Wardle and Königl (1993) got a similar parameter space but they simplified the mass conservation equation by using instead $\rho u_z = Cst$: this *always* leads to the development of an outwardly directed vertical velocity, whatever the magnetic compression.

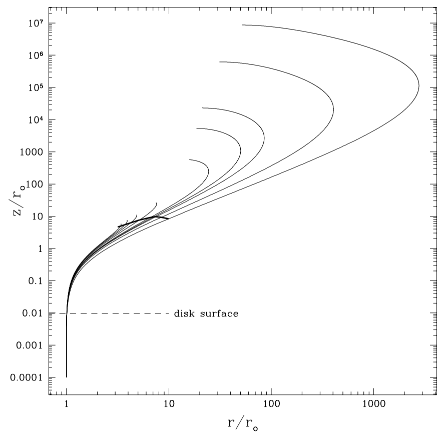

Fig. 7. Poloidal magnetic field lines for $\xi = 0.05, 0.04, 0.03, 0.02, 0.012, 0.01, 0.009, 0.007$ and 0.005 (the maximum radius increases with decreasing ejection efficiency). The thick line connects the position of the Alfvén point for each solution (here isothermal, Ferreira 1997). Note the logarithmic scales: small ξ jets recollimate with angles smaller than one degree.

than the dominant force, namely gravity (Ferreira and Pelletier, 1993a,b). This separation made, one obtains a system of coupled ODE that must be propagated along the self-similar variable $x = z/h(r)$. Starting at $x = 0$ in resistive MHD, the ideal MHD regime is reached above the disc ($x > 1$). Once in this regime, the 5 MHD invariants are fixed and 3 conditions remain thus to be imposed: this is done by the regularity conditions at the SM, Alfvén and FM points. In practice, this corresponds to fixing a parameter in order to smoothly cross each critical point met: μ is thus imposed by the slow point, ξ by the Alfvén one and the jet polytropic index γ by the fast point. General analytic links between the disc and usual jet parameters (such as λ, related to $L(a)$, and κ, related to $\eta(a)$) can be found in Casse and Ferreira (2000a).

4.3.3 Cold solutions

Because accretion discs are quasi-Keplerian, the enthalpy carried away by the outflowing mass at the disc surface is negligible (cold jets). The use of isothermal magnetic surfaces (Ferreira, 1997) or adiabatic ones (Casse and Ferreira, 2000a) introduces no significant change. All solutions display Alfvén surfaces located quite far from the disc (at $z_A \sim r_A$), with typical ejection efficiencies $\xi \sim 0.01$. They all recollimate towards the axis because of a dominant "hoop-stress" (Fig. 7) and terminate with a shock at the location of the FM point (which is uncrossable for these solutions). Note however that all solutions are super-FM in the conventional sense (ie. $u_p > V_{FM}$, see Ferreira 1997 for more details).

4.3.4 Warm solutions

Warm jets can be produced from quasi-Keplerian discs if the outflowing gas temperature undergoes a sudden rise above the disc. This can be done in astrophysical systems by two ways:

Illumination: if the central object has a hard surface (ie, is not a black hole), then the accretion shock produces hot spots that will illuminate (with UV or X-rays) and possibly photo-ionize the surface layers of the disc. The large X-ray activity of young stars may also contribute to this effect. Above a black hole, pair production and formation of an X-ray zone may also have the same effect on the disc upper atmosphere.

Local dissipation: this may be due to the dissipation of accretion power itself in the highly turbulent magnetized corona (or more correctly chromosphere) expected to be present above the turbulent disc. This is in fact suggested by both observations (Kwan, 1997) and numerical simulations (Miller and Stone, 2000).

Including the energy equation and using a prescription for additional heat deposition, Casse and Ferreira (2000b) showed that enhancing the temperature at the disc surface layers has dramatic effects (see Fig. 8). For example, the disc vertical equilibrium can be changed so that a balance can now be achieved with magnetic configurations much more bent. As a result, much smaller values of ξ, down to 0.001, can be obtained. On the other extreme, providing a large enthalpy allows more mass to be loaded ontothe field lines: these thermally and magnetically driven jets can accelerate up to $\xi \sim 0.1$ (jet parameter $\kappa \sim 1$). Jets 3 to 5 times slower but denser than in the "cold" case can be obtained.

The physics of such magnetized accretion-ejection systems (MAES), as understood through these semi-analytical works, has been confirmed by two independent groups, using two distinct numerical MHD codes (see figure 9). These are the only works where the mass load is computed in a consistent way with the jet acceleration. Other MHD simulations of jets driven by accretion discs usually do not compute the disc and simply assume this mass load.

There has been some claims in the literature that the magneto-centrifugal acceleration process was unstable (Lubow et al., 1994b; Cao and Spruit, 2002). The idea was the following. Start from a steady picture where the accretion velocity u_r at the disc midplane is due to the jet torque. It leads to a bending of the poloidal field lines described by an angle θ with the vertical. Now imagine a small perturbation δu_r enhancing the accretion velocity. Then, according to these authors, the field lines would be more bent (θ increases) which would lead to lower the altitude of the sonic point. Because the sonic point would be located deeper in the disc atmosphere, where the density is higher, more mass would be henceforth ejected which would then increase the total angular

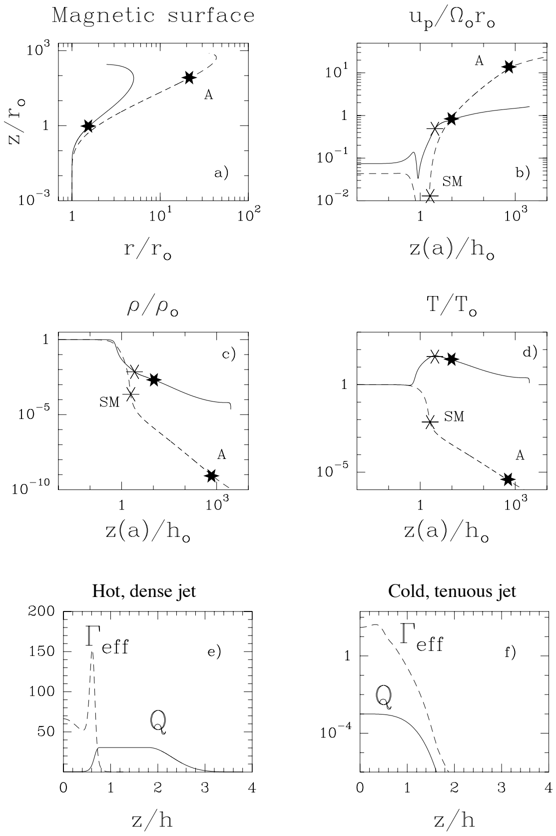

Fig. 8. Characteristic plots of two extreme solutions: **a)** poloidal cross-section of the magnetic surfaces anchored at a radius r_o; **b)** poloidal velocity of the plasma along a magnetic surface in units of the Keplerian speed at its footpoint; **c)** and **d)**, plasma density and temperature along a magnetic surface normalized to their values at the disc midplane. The warm denser jet is drawn in solid line while the cold and more tenuous jet is in dashed line. The cross symbolizes the locus of the slow-magnetosonic (SM) point and the star the locus of the Alfvén (A) point. Pannels **e)** and **f)** show the effective (viscous and Ohmic) heating term Γ_{eff} and the prescribed entropy source Q (taken from Casse & Ferreira 2000b).

momentum carried away by the jet. This means that the torque due to the jet is enhanced and will, in turn, act to increase the accretion velocity. Thus, according to these authors, the accretion-ejection process would be inherently unstable. This is wrong. In fact the whole idea of this instability is based upon a crude approximation of the disc vertical equilibrium. A magnetized disc is *not* in hydrostatic equilibrium. The magnetic field produces a strong vertical compression, comparable to the gravity. As a consequence, as θ is increased, *less* mass is being ejected, not more. This has been pointed out by Königl and Wardle (1996) and Königl (2004) and is indeed verified in full MAES calculations reported here.

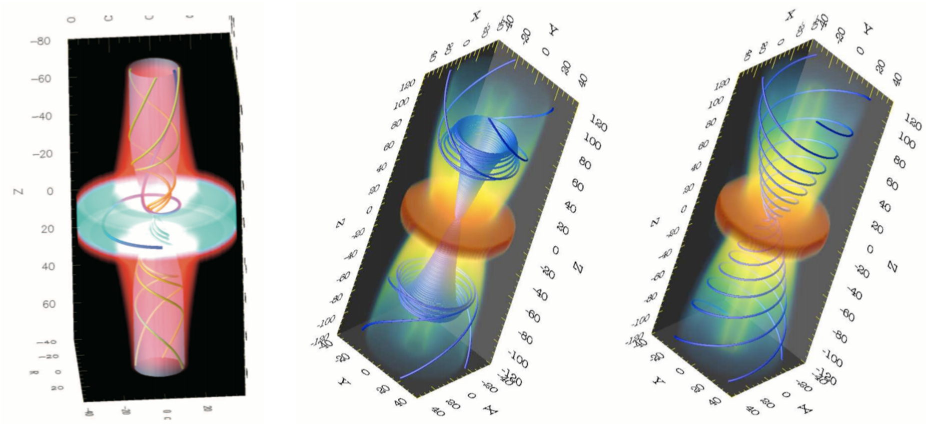

Fig. 9. **Left:** First MHD simulation of an accretion-ejection system done with the MHD code VAC (taken from Casse and Keppens 2002). The numerical experiment has confirmed that only an equipartition field can drive steady jets. **Right:** Two simulations done with the AMR MHD code FLASH taken from Zanni et al. (2007). Another important analytical result is confirmed: only a large magnetic diffusivity allows a steady state (rightmost image).

4.4 Comparison of JEDs and SADs

From the observational point of view, JED and SAD emission properties are quite different. While in a SAD all the released accretion power is radiated away at the disc surfaces, in a JED this power is feeding the two jets. As a consequence, only a small fraction (of order h/r) of the available power is put into the disc luminosity (Ferreira and Pelletier, 1995). Let us analyze this in more details.

The global energy budget in the JED is $P_{acc,JED} = 2P_{rad,JED} + 2P_{MHD}$ where P_{MHD} is the MHD Poynting flux feeding a jet, whereas the liberated accretion power writes

$$P_{acc,JED} \simeq \frac{GM\dot{M}_{a,J}}{2r_{in}} \left[\left(\frac{r_{in}}{r_J} \right)^{\xi} - \frac{r_{in}}{r_J} \right] \tag{22}$$

where r_{in} is the inner radius of the JED, probably the disc truncation radius (see fig. 1). The dynamical properties of a JED have been extensively studied in a series of papers (see Ferreira 2002 and references therein). It was shown earlier that the ratio at the disc midplane of the jet torque to the turbulent "viscous" torque is $\Lambda \sim r/h \gg 1$. This dynamical property has a tremendous implication on the JED emissivity. The JED luminosity comes from the accretion power dissipated within the disc by turbulence and transported away by photons, so that $2P_{rad,JED} = P_{diss}$. This dissipated power is very difficult to estimate with precision because it requires a thorough description of the turbulence itself. Thus, one usually uses crude estimates based on "anomalous" turbulent magnetic resistivity $\eta_m = \mu_o \nu_m$ (Joule heating) and viscosity $\eta_v = \rho \nu_v$ (viscous heating). This translates into $P_{diss} = P_{Joule} + P_{visc} = \int \eta_m J^2 dV + \int \eta_v (r\partial\Omega/\partial r)^2 dV$ where integration is made over the whole volume occupied by the JED. The importance of local "viscous" dissipation with respect to the MHD Poynting flux leaving the disc is approximately given by

$$\frac{P_{visc}}{2P_{MHD}} \sim \frac{1}{\Lambda} \tag{23}$$

which is much larger than unity: turbulent "viscosity" provides negligible dissipation in a JED. Joule heating arises from the dissipation of toroidal and radial current densities which are comparable [8]. One therefore gets $\eta_m J^2 \sim \nu_m B_z^2/\mu_o h^2 \sim \nu_v \rho \Omega^2 \sim \eta_v (r\partial\Omega/\partial r)^2$, for equipartition fields, isotropic magnetic resistivity $\eta_m = \mu_o \nu_m$ and a turbulent magnetic Prandtl number of order unity. This leads to

$$\frac{P_{Joule}}{2P_{MHD}} \sim \frac{1}{\Lambda} \tag{24}$$

namely a negligible effective Joule heating. Thus, the total luminosity $2P_{rad,JED}$ of the JED is only a fraction $1/(1+\Lambda)$ of the accretion disc liberated power $P_{acc,JED}$.

This translates right away into a lack of disc emission from the innermost ejecting parts: the spectral energy distribution would thus appear flatter than the usual $-4/3$ scaling. But more interestingly, the disc is much less dense than a corresponding SAD at the same radius (with same $\dot{M}_a$, Combet & Ferreira, submitted). This is a straightforward consequence of a much larger accretion velocity due to the dominant jet torque. While in a SAD, the turbulent "viscous" torque provides a sonic Mach number $m_s = u_r/C_s = \alpha_v h/r$, in a JED one gets $m_s \sim 1$. As a consequence, a JED fed with the same accretion rate than a SAD has a lower surface density $\Sigma(r)$, $\alpha_v h/r$ times smaller than that of the SAD [9]. This could lead to optically thin parts in the JED but, in any case, to a sharp decrease of the disc surface density at the SAD-JED transition

[8] Full computations of MAES show that the three magnetic field components are comparable at the disc surface, namely $B_\phi^+ \sim B_r^+ \sim B_z$ (Ferreira and Pelletier, 1995; Ferreira, 1997).

[9] Note that the deviation from the keplerian law is then also larger in a JED than in a SAD. One may indeed write the disc midplane angular velocity as $\Omega = \delta\Omega_K$, where $\delta^2 \simeq 1 - \frac{5}{2}\varepsilon^2 - \mu\varepsilon$ (Ferreira and Pelletier, 1995). Thus, while in a SAD the deviation is of order $(h/r)^2$ only because of the radial pressure gradient, a JED has a deviation of order $\sim (h/r)$ because of the magnetic tension. Nevertheless, to all practical means, a Keplerian rotation law remains a good approximation.

radius r_J. As already pointed out, this radius is unknown as it depends on the magnetic flux Φ available in the disc (and may vary from one object to another as discussed above), but may lead to observational investigations.

This last property is very interesting since Masset et al. (2006) have shown that a transition of this kind would act as a trap for low mass protoplanetary embryo ($M < 15\ M_\oplus$). Indeed, Type I inwards migration is due to the differential Lindblad torque arising from the planetesimal interaction with the viscous disc. But this negative torque is strongly reduced at the transition radius and balanced by the positive co-rotation torque (which is due to the exchange of angular momentum between the planetesimal and trapped disc material in its vicinity). Thus, these planetesimals would be halted at r_J which may be as large as 1 AU, long before the disc truncation radius due to the star-disc interaction.

Finally, JEDs seem to put into question the very existence of dead zones (only below the transition radius r_J of course). In steady state, the ionization fraction due to X-rays can be given by the generic expression $x_e = \sqrt{\xi_2/\beta n_H}$, where β is the dominant recombination rate and ξ_2 the secondary ionization rate [10] . For illustrative purposes, let us assume for simplicity that there is no metallic ion in the medium, so that β is the collisional recombination rate. In that case, we obtain

$$x_e = 7.5\,10^{-10}\, m_s^{1/2} \varepsilon^{3/2} \left(\frac{\xi_2}{10^{-17}\, s^{-1}}\right)^{1/2} \left(\frac{\dot{M}_a}{10^{-7} M_\odot/yr}\right)^{-1/2} \times \left(\frac{M_*}{M_\odot}\right)^{1/2} \left(\frac{r_o}{1AU}\right)^{1/2} \quad (25)$$

where $\varepsilon = h/r$, $m_s = u_r/C_s$ is the sonic Mach number within the accretion disc. This expression holds for both SADs and JEDs. In the case of a SAD, $m_s = \alpha_v \varepsilon$ whereas it is of order unity in a JED (Ferreira and Pelletier, 1995). This has two important consequences on X-rays capability to ionize the accretion disc. First, the column density through the disc itself is less, which allows for a deeper penetration of X-rays. Second, since $x_e \propto n_H^{-1/2}$, the recombination time scale is longer which enhances the ionization efficiency. Combining Eq. 25 and 8 gives an ionization rate in the JED

$$\xi_2 > 2.6\,10^{-19} \left(\frac{\varepsilon}{0.01}\right)^{-5} \left(\frac{\dot{M}_a}{10^{-7} M_\odot/yr}\right) \left(\frac{M_*}{M_\odot}\right)^{-1} \left(\frac{r_o}{1AU}\right)^{-3} \ \mathrm{s}^{-1} \quad (26)$$

[10] The dominant contribution to ionization by X-rays is due to K-shell electrons liberated with an energy $E \sim$ keV. These electrons will then collisionaly ionize and heat the gas. This is called the secondary ionization.

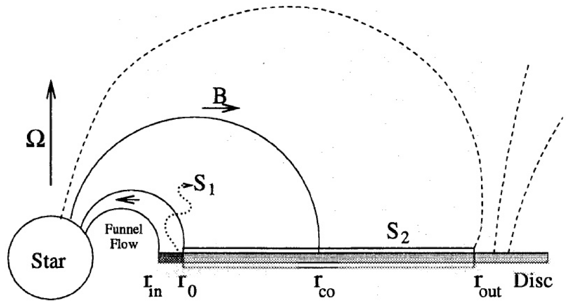

Fig. 10. Sketch of the Gosh & Lamb magnetic configuration adapted for young stars (Ghosh et al., 1977; Collier Cameron and Campbell, 1993; Armitage and Clarke, 1996). The stellar magnetic fields anchored below r_{co} spin up the star whereas those anchored beyond brake it down.

in order to have $\mathcal{R}_m$ larger than unity. This is much smaller than X-ray ionization rates computed with a "warm" solution and a typical X-ray spectrum (Ferreira, in prep). This is just an indication as Eq. (25) is probably an over simplification, but it indicates that if the inner regions of accretion discs are indeed driving magnetic jets, then it might be possible that a dead zone never settles in.

5 Star-disc interactions

5.1 The disc locking paradigm

Once they become visible in the optical, T Tauri stars exhibit rotational periods of the order of 10 days, which is much smaller than expected (Bouvier et al., 1997; Rebull et al., 2002). This implies a very efficient mechanism of angular momentum removal from the star during its embedded phase. Moreover, a T Tauri star seems to evolve with an almost constant rotational period although it undergoes some contraction and is still actively accreting disc material for roughly a million years. This is a major issue in star formation, unsolved yet, but one solution to this paradox is the star-disc interaction.

Angular resolution is not yet sufficient to directly image this interaction region (of size 0.1 AU or less: it would require optical interferometry) but there have been mounting spectroscopic and photometric evidences that the disc is truncated by a stellar magnetosphere and that accretion proceeds along magnetic funnels or curtains towards the magnetic poles (see Bouvier et al. 2007 and references therein). This observational puzzle gave rise to the so-called *disc locking paradigm*, where accretion discs have to, somehow, provide a means to brake down contracting and accreting protostars. Indeed, this seems almost paradoxical.

The first idea put forward came from the X-ray binary community (Ghosh et al., 1977; Collier Cameron and Campbell, 1993; Armitage and Clarke, 1996). The Ghosh & Lamb configuration assumes that the stellar magnetic field, a dipole, is threading the circumstellar accretion disc on a quite large radial extension (fig. 10). Let Ω_* be the angular velocity of the star. Its magnetosphere will try to make the disc material corotate with the protostar so that the sign of the torque depends directly on the relative angular velocity. Stellar magnetic field lines threading the disc beyond the rotation radius $r_{co} = (GM/\Omega_*^2)^{1/3}$ exert a positive torque, whereas they brake down the disc material below r_{co}. Let us also define the truncation radius r_t below which the stellar magnetic field is strong enough to "truncate" the disc by enforcing the material to flow along the field lines and no longer on the plane of the disc. A spin-down of the protostar can then arise only if the two torques balance each other, namely if the outer radius r_{out} where the magnetospheric field remains anchored in the disc is significantly larger than r_{co} (in order to enhance the braking torque).

Unfortunately, this idealized picture can probably not be maintained. The simple reason is that accretion onto the star and this "strict" disc locking mechanism are two contradictory requirements (see thorough discussion in Matt and Pudritz 2005b).

5.2 The formation of accretion curtains

One can safely realize that accretion onto the star can only proceed if $r_t < r_{co}$ [11] . In this situation, the stellar magnetic field can brake down both the disc and the material accreting in the funnel flows. This implies of course a stellar spin up by the disc material located below r_{co}. The disc locking paradigm assumes that stellar field lines remain anchored beyond r_{co}, giving hopefully rise to some angular momentum balance. But within this paradigm, the disc viscosity must be efficient enough so as to radially transport outwards both the disc and stellar angular momentum! This is unrealistic because the stellar angular momentum is far too large. Moreover, all numerical simulations done so far showed a fast opening of the field lines beyond r_{co} (through numerical reconnection), severing the causal link and thereby dramatically reducing this negative torque (Lovelace et al., 1995, 1999; Long et al., 2005). Although this effect is strongly dependent on the disc magnetic turbulent diffusivity, the main result is to spin up the star whenever $r_t < r_{co}$.

Therefore, an important question is what determines the disc truncation radius

[11] When $r_t > r_{co}$, this is called the "propeller" regime (Romanova et al., 2004, 2005), where the disc material is flung away from the central star. Although some episodic accretion events are reported in these simulations, it is not consistent with CTTS observations as one never gets a total disappearance of accretion signatures.

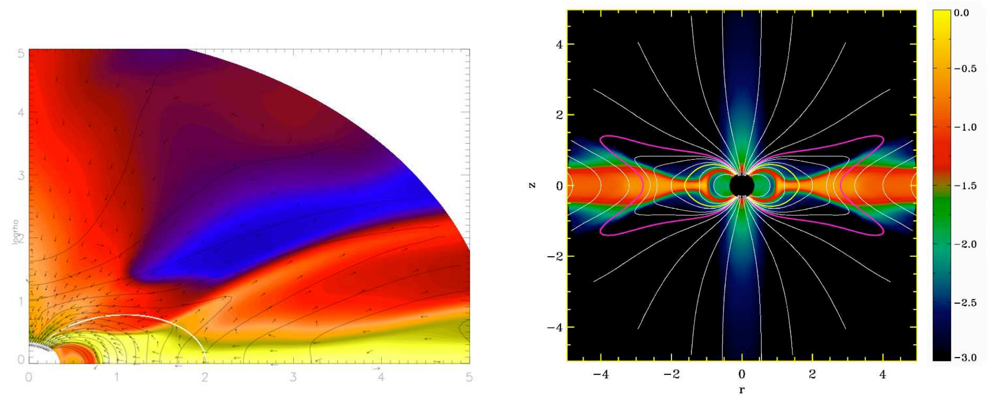

Fig. 11. Snapshots of MHD simulations of a 2D star-disc interaction. The stellar magnetic field is assumed to be a dipole field aligned with the rotation axis, whereas no field is included in the disc. Left: resistive MHD simulation done with the VAC code (Bessolaz et al, submitted). Right: resistive and viscous simulation done with PLUTO (Zanni et al, in prep). Starting from an initial condition out of equilibrium, both simulations converge and lead to the formation of accretion funnel flows at a radius consistent with Eq. (27). Although stellar field lines are quickly opened by the differential rotation, no X-winds are formed. As an outcome of this star-disc interaction, the star is being spun up by the accreting material. The co-rotation radius is marked by a white line in both simulations.

r_t? Let us put it that way. Can we, for a given accretion disc model and stellar dipole field strength, estimate where r_t will be located? In fact, two constraints must be simultaneously fulfilled for driving steady-state accretion funnel flows. First, the poloidal stellar magnetic field must be strong enough to halt the accretion motion, namely $B_z^2/\mu_o = \rho u_r^2$. Second, the disc thermal pressure must be able to lift material vertically in order to initiate the accretion funnel flow, $B_z^2/\mu_o \simeq P$. This last constraint is equivalent to an equipartition field, as in a JED and also proposed by Pringle and Rees (1972). Putting this two constraints together one derives a ratio of the disc truncation radius to the co-rotation radius (Bessolaz et al, submitted)

$$\frac{r_t}{r_{co}} = 0.66\ m_s^{4/7} B_*^{4/7} \dot{M}_a^{-2/7} M_*^{-10/21} R_*^{12/7} P_*^{-2/3} \tag{27}$$

where $m_s = u_r/C_s \sim 1$ is the disc midplane sonic Mach number and the disc accretion rate $\dot{M}_a$ has been normalized to 10^{-8} $\mathrm{M}_\odot$ yr^{-1}, stellar dipole field B_* to 150 G, mass to 0.5 $\mathrm{M}_\odot$, radius to 3$\mathrm{R}_\odot$ and period P_* to 8 days. These analytical constraints and estimates have been confirmed using MHD axisymmetric numerical simulations of a star-disc dipole interaction (fig. 11). This is therefore a robust result. It shows that truncating discs can be done with a *dipole field* of several hundreds of Gauss (not kG !), consistent with both observations of magnetic fields (Donati, private communication) and sizes of inner disc holes (Najita et al., 2007). However, it turns out that the disc accretion rate measured in these simulations is about 10^{-9} $M_\odot$ yr^{-1} only. In fact, it was implicitly assumed in Eq. 27 that B_* and $\dot{M}_a$ were independent variables, which is not true. Mass loading onto the stellar magnetosphere is actually a process not very different from accretion-ejection: the magnetic configuration determines a funnel whose shape controls the mass loading. In a steady-state approach, this arises from the requirement of a smooth crossing of the slow magnetosonic point. This is a very interesting point that deserves further investigations. To summarize, if observations confirm the existence of kG dipole fields then such a magnetic configuration is suitable for producing steady accretion columns. If, on the other hand, the mean value of the dipole component in accreting TTS is smaller than a kG, then accretion is proceeding along multipolar kG fields (as in Long et al. 2007).

Moreover, the "disc locking" picture *a la* Ghosh & Lamb is seriously put into danger as the inner disc radius r_t can be significantly smaller than the co-rotation radius r_{co}. That picture could still work if the field lines remain connected well beyond r_{co}, but this is not found in numerical simulations. Unless playing around with a non unity effective Prandtl number (i.e. turbulent magnetic diffusivity ν_m much larger than the viscosity ν_v), one obtains a star-disc interaction confined to a zone whose radial extent is small, located below r_{co}. This is a hint that one must probably find another means to evacuate the stellar angular momentum.

5.3 Protostellar magnetic braking

So, how to conciliate the formation of quasi-steady accretion columns (as observed) with the long term requirement of stellar spin down (observed as well)? Very simple (in 2D) magnetic configurations can then be designed (see fig. 3). The obvious way out is to propose that accreting stars are actually being spun down via *winds that would not exist without the presence of accretion.* This has lead Matt and Pudritz (2005a) to propose the name of "accretion powered stellar winds".

Accretion onto the star takes place along closed magnetospheric field lines, shocks the stellar surface and releases there most of its mechanical energy (through mostly UV emission). The idea is then that a fraction of this accretion-heated mass diffuses towards the magnetic pole until it reaches open field lines (see fig. 12). A warm stellar wind can then be initiated (forming a magnetic Y point point if a magnetic field is present in the disc, case (e) in fig. 3). The problem with T Tauri stars is that they are rotating at about 10% of their break-up speed. This translates into a totally negligible magnetic acceleration (the stellar material is far too deep in the gravitational potential well). One has therefore to rely on pressure-driven winds (see eg. discussion in Ferreira

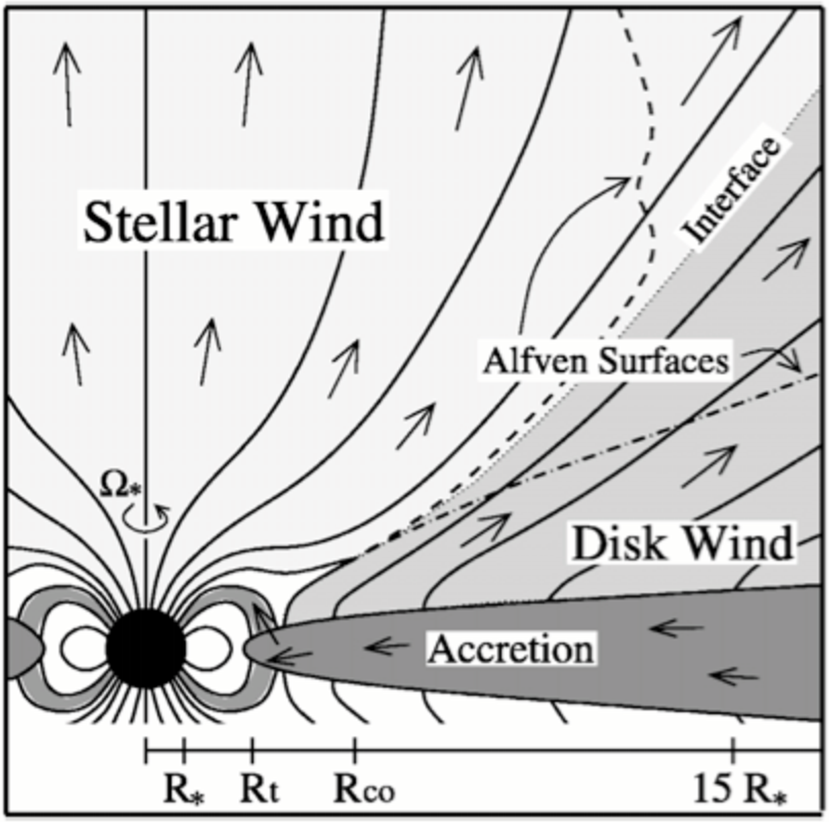

Fig. 12. Star-disc interaction in the case where the stellar magnetic moment is anti-parallel to the disc magnetic field. Here, the stellar spin down is done by a wide open stellar wind assuming no strong confinement by the outer disc wind (taken from Matt & Pudritz 2005b). Note that if the radial extent of the JED is small, then this is a field geometry alike the X-wind model (Shu et al., 1994).

et al. 2006a for "enhanced stellar winds"). Now, if that initial pressure is only thermal, then temperatures of several million degrees are required. This raises the critical issue of probably too strong emission losses due to this inner hot wind. The alternative is to rely on a turbulent Alfvén wave pressure that would be less dissipative (Hartmann and MacGregor, 1980; DeCampli, 1981). Note that the presence of turbulent MHD waves is indeed highly expected in this context.

It should be noted that current MHD numerical simulations of star-disc interaction (e.g. Long et al. 2005, Zanni et al.) do show a magnetic braking due to the opened stellar field lines. This has been interpreted as a "magnetic tower" since no real stellar wind was incorporated in the simulations. This is obviously a very promising issue. A thorough investigation should therefore be conducted in order to assess whether or not accretion-powered stellar winds of this kind can indeed (i) be dense enough and with a magnetic lever arm large enough to brake down the protostar, (ii) have radiative losses consistent with observations and (iii) do not pose any energetic problem like e.g. requiring to

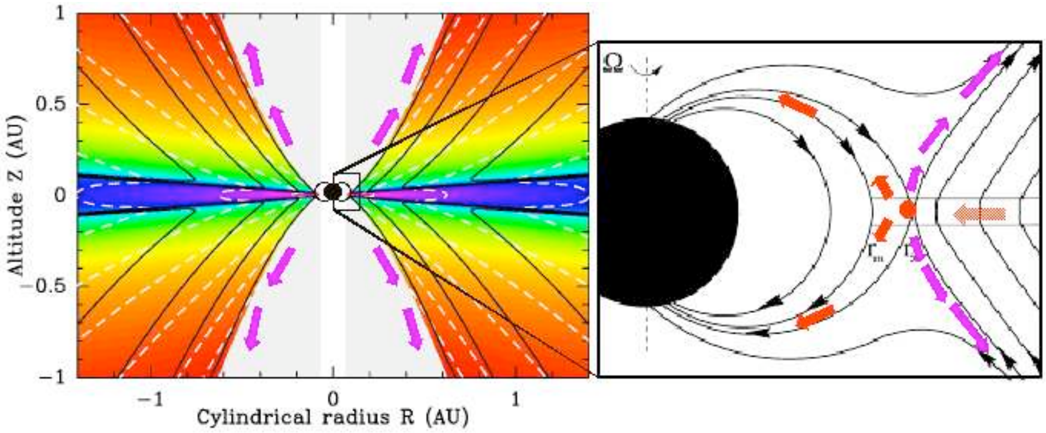

Fig. 13. Star-disc interaction where the stellar magnetic moment is parallel to the disc magnetic field. There are three distinct types of ejection: a stellar wind on the axis, a disc wind (MAES shown in colors) and a sporadic reconnection X-wind at the interface, braking down the protostar (Ferreira et al., 2000).

tap more than 50% of the accretion luminosity.

5.4 Reconnection X-winds

Accretion-powered stellar winds are somehow designed to explain the "disc locking" paradigm for T Tauri stars, namely to maintain their low rotation rate despite accretion. But how can we explain that T Tauri stars do *already* rotate at 10% of their break-up speed? Numerical simulations of the collapse of rotating magnetized clouds succeed nowadays in explaining the formation of protostellar cores at break-up speeds thanks to magnetic braking (Banerjee and Pudritz, 2006; Machida et al., 2006). However, it is doubtful that such a braking could provide much lower initial rotation rates. One must then rely on some interaction between the protostar and its disc during the embedded phase (Class 0 and possibly Class I).

To our knowledge, the only model that addresses this issue is the Reconnection X-wind model (Ferreira et al., 2000). In this model, it is assumed that the interstellar magnetic field is advected with the infalling material in such a way that a significant magnetic flux Φ is now threading the protostellar core and the inner disc regions (as simulations show). This self-gravitating core will develop a dynamo of some kind but whose outcome is assumed to be the generation of a dipole field with a magnetic moment parallel to the disc magnetic field (see fig. 13). This is clearly an assumption as there is no theory of such a constrained dynamo that takes into account both the presence of an initial strong fossil field and the outer disc (see however Moss 2004). The coexistence of this dipolar stellar field with the outer disc field generates an X-type magnetic neutral line (case (d) in fig. 3), where both fields cancel each other at a radius r_X. Note that such a magnetic configuration has been previously considered by Uchida and Low (1981) and Hirose et al. (1997), but without taking into account the stellar rotation.

Let us assume that at $t = 0$ the dipole is emerging from a protostar rotating at break-up speed so that $r_X = R_{*,0}$ with $R_{*,0} = R_*(t = 0)$, $M_{*,0} = M_*(t = 0)$, $\dot{M}_{a,0} = \dot{M}_a(t = 0)$, $\Omega_{*,0} = \sqrt{GM_{*,0}/R_{*,0}^3}$ and $B_{*,0} = B_*(t = 0)$. It is further assumed that the field threading the disc is strong enough to drive self-confined disc winds at these early stages. Then Eq. (1) applies and provides us the value of the required stellar field. What will be the consequences of this initial state?

From the point of view of the disc, nothing is changed beyond r_X: a disc wind is taking place in the JED and disc material accretes by loosing its angular momentum in the jets. At r_X however, magnetic reconnection converts closed stellar field lines and open disc field lines into open stellar field lines. Accreting material that was already at the disc surface at r_X is now loaded into these newly opened field lines (there is a strong upward Lorentz force above r_X). Since these lines are now rotating at the stellar rotation rate, they exert a strong azimuthal force that drives ejection. This new type of wind has been called "Reconnection X-winds". Although material is ejected along field lines anchored onto the star, this is not a stellar wind since material did not reach the stellar surface and thus did not loose its rotational energy: it is much easier to accelerate matter under these circumstances.

Reconnection X-winds are fed with disc material and powered by the stellar rotational energy. As a consequence, they exert a negative torque on the protostar which leads to a stellar spin down. On the other hand, an increase of the stellar angular velocity Ω_* is expected from both accretion and contraction, with a typical Kelvin-Helmoltz time scale of several 10^5 yrs. Because of the huge stellar inertia, the evolution of Ω_* with time must be followed on these long time scales. One assumption used to compute the angular momentum history of the protostar on those scales is that $r_X \simeq r_{co}$. Such an assumption relies on the possibility for the protostellar magnetosphere to evacuate angular momentum through violent ejection events (Reconnection X-winds) whenever $r_X > r_{co}$, while quasi steady accretion columns form when $r_X < r_{co}$. Consistently with $r_X \simeq r_{co}$, a constant fraction $f = \dot{M}_X/\dot{M}_a$ is assumed on these long time scales, where $\dot{M}_X$ is the ejected mass flux in Reconnection X-winds, as well as a constant magnetic lever arm parameter λ. *These winds are therefore best seen as violent outbursts carrying disc material (blobs?) and stellar angular momentum from the star-disc interaction, channeled and confined by the outer disc wind.* Note that a conventional stellar wind would of course take place and fill in the inner field lines with mass, but its effect on the stellar spin evolution has been neglected in this work.

The global picture is then the following. As the protostar is being spun down,

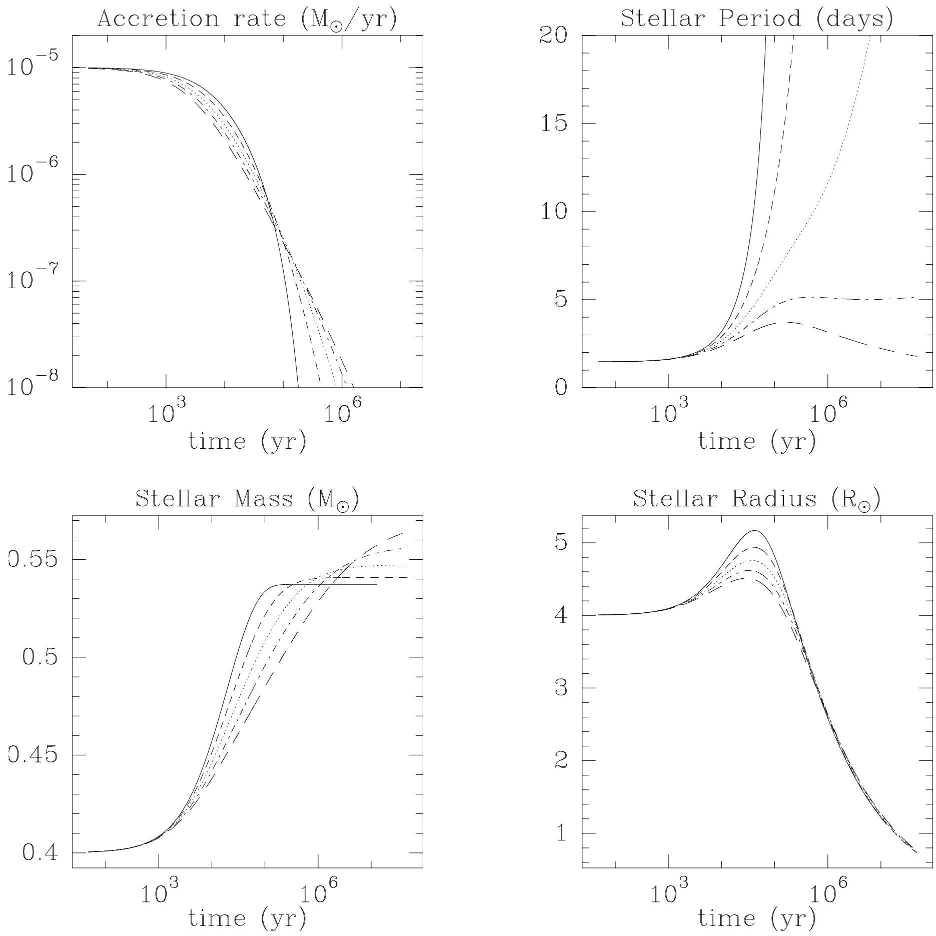

Fig. 14. Time evolution of disc accretion rate, protostellar period, mass and radius as function of n for $f = 0.1$ and $\lambda = 3$ with the following initial conditions: $R_{*,0} = 4R_\odot$, $M_{*,0} = 0.4M_\odot$, $\dot{M}_{a,0} = 10^{-5}\ M_\odot yr^{-1}$ and $T_* = 3000$ K. $n = 3$ (solid), $n = 3.41$ (dashed), $n = 3.87$ (dotted), $n = 4.4$ (dash-dotted) and $n = 5$ (long-dashed). For these reasonable values of the parameters, a very significant braking is obtained in only a few 10^5 yrs (taken from Ferreira et al. 2000).

the co-rotation radius r_{co} increases and so must r_X. The stellar dipole field is assumed to follow $B_{star} = B_*(r/R_*)^{-n}$ where the index n describes a deviation from a pure dipole in vacuum. Now, r_X is defined by the cancellation of the stellar and disc field, whose scaling is very different from the former (see Eq. 1). The only way to ensure $r_X \simeq r_{co}$ on these long time scales is then to decrease $\dot{M}_a$ in time as well. Note that this is not a surprise as the accretion rate onto the star is controlled by the star-disc interaction. Thus, while computing the stellar spin evolution in time $\Omega_*(t)$, starting from conditions prevailing in Class 0 objects and using f, λ and n as free parameters, one gets also $R_*(t)$, $M_*(t)$ and $\dot{M}_a(t)$. Note that this global process of angular momentum removal is intimately related to the magnetic history of the protostar-disc system. Two additional ingredients are thus necessary: the amount of magnetic flux

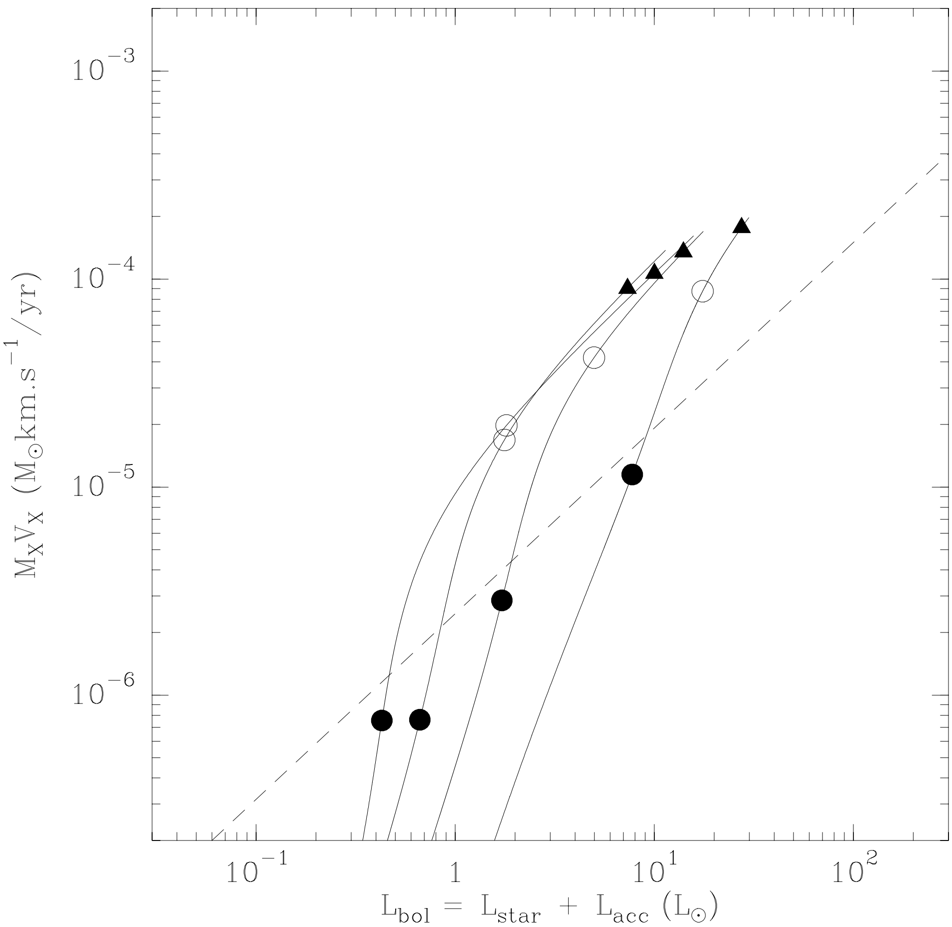

Fig. 15. Computed evolutionary tracks showing the relation between the reconnection X-wind momentum flux and the bolometric luminosity, for $n = 3.87$, $f = 0.1$ and $\lambda = 3$ for different initial conditions. The markers show selected times along each track: $t = 10^3$ yr (triangles), 10^4 yr (open circles) and 10^5 yr (filled circles, typical Class I age). Note the similarity with Fig. 5 in Bontemps et al. (1996). The dashed line represents their 'best-fit' correlation for Class I sources. Initial conditions, from right to left: (a) $M_{*,0} = 0.8 M_\odot$, $R_{*,0} = 6 R_\odot$, $T_* = 3900$ K; (b) $M_{*,0} = 0.4 M_\odot$, $R_{*,0} = 4 R_\odot$, $T_* = 3000$ K; (c) $M_{*,0} = 0.2 M_\odot$, $R_{*,0} = 3 R_\odot$, $T_* = 2800$ K; (d) $M_{*,0} = 0.2 M_\odot$, $R_{*,0} = 2 R_\odot$, $T_* = 2800$ K. The initial accretion rate $\dot{M}_{a,0} = 10^{-5} M_\odot yr^{-1}$ is the same for all cases (taken from Ferreira et al. 2000).

Φ threading the disc and how the stellar field B_* evolves with time (through dynamo). The calculations reported in Ferreira et al. (2000) were performed using simple assumptions about the dynamo and a more realistic modeling is needed. However, the results are already very promising (see Ferreira et al. 2000 for more details).

It was found that all low-mass Class 0 objects can indeed be spun down, from the break-up speed to about 10% of it, on a time scale consistent with the duration of the embedded phase for very reasonable values of the free parameters ($n = 3$ or 4, $f\lambda > 0.1$, see fig. 14). Stellar period, mass, radius and disc accretion rates were found consistent with values for T Tauri stars with a dipole field smaller than 1 kG. According to Bontemps et al. (1996), there is a decrease in time of all jet signatures (kinetic power P_X, momentum flux $F_X = \dot{M}_X V_X$ and ejection rate $\dot{M}_X$) that seems to follow a decrease of the accretion rate $\dot{M}_a$. Typical values are $F_X \simeq$ a few 10^{-4} $M_\odot kms^{-1}yr^{-1}$ for Class 0 sources and $F_X \simeq 10^{-6}$ to 10^{-5} $M_\odot kms^{-1}yr^{-1}$ for Class I sources; during this time, L_{bol} decreases from ~ 10 $L_\odot$ to $\sim L_\odot$, $\dot{M}_X$ from $\sim 10^{-6}$ $M_\odot yr^{-1}$ for the youngest Class 0 sources to $\sim 10^{-8}$ $M_\odot yr^{-1}$ for the most evolved Class I sources. This observation is naturally accounted by the Reconnection X-wind model (see fig. 15).

In a sense, the Reconnection X-wind does exactly what was demanded to the X-wind model in its earliest version (Shu et al., 1988). In this model, a protostar at breakup speed is launching a wind from its equator, hence providing angular momentum and energy to the wind. This model was discarded for TTS when it was shown that these stars were actually rotating much slower and a new version of the X-wind was then proposed (Shu et al., 1994). In the X-wind paradigm the only magnetic field is the stellar field. This is in strong contrast with the Reconnection X-wind where a disc magnetic field must be present. As shown above, this is the crucial ingredient that allows to go from a rapidly rotating, accreting and contracting Class 0 object to a slowly rotating (still accreting) Class II T Tauri star. Such an evolution was never demonstrated within the X-wind scenario. However, the final outcome of a Reconnection X-wind, obtained when the disc magnetic flux Φ is about to vanish, would resemble the X-wind (in terms of ejection).

This scenario offers therefore a natural and unique explanation for two important questions, why all T-Tauri stars rotate much slower than their break-up velocity and why outflows are more powerful during the early stages of star formation. Finally, note that the main difference between "accretion-powered stellar winds" and Reconnection X-winds relies on the stellar magnetic moment. In the former case, it is anti-parallel to the disc field while it is parallel in the latter. If the dynamo action explicitly assumed provides a magnetic field reversal, then recurrent transitions from one wind configuration to another can be expected.

6 Concluding remarks

The theory of *steady* jet production from Keplerian accretion discs has been completed in the framework of "alpha" discs. The physical conditions required to thermo-magnetically drive jets are constrained and all the relevant physical processes have been included. The role of large scale magnetic fields in discs has gradually emerged and it seems now (after more than 20 years) an unavoidable ingredient of star formation theory as a whole. The progresses in star-disc interaction reported here provide valuable insights but one should remain cautious as stellar magnetic fields are not aligned dipoles.

The amount of magnetic flux Φ in the disc is an unknown parameter but it is reasonable to assume that it scales with the total mass M. If this is verified then two important aspects could be naturally explained:

(1) Reconnection X-winds can brake down a protostar during the embedded phase, explaining that T Tauri stars rotate at about 10% of the break-up speed. Remarkably, the mystery of the low dispersion in angular velocities would be naturally accounted by a low dispersion in the ratio Φ/M (Ferreira et al., 2000). These winds are also a very promising means to drive time dependent massive bullets, channeled by the outer steady disc wind.

(2) The transition from Classes 0, I and maybe II (with both disc winds and reconnection X-winds) to Classes II and III (with stellar winds, "accretion powered" or not) would follow the evolution of the disc magnetic flux Φ, with a transition radius between the outer SAD and the inner JED decreasing in time.

Current up-to-date observational technics show that jets from classical T Tauri stars are arising from the innermost disc regions. For a long time, these regions remained out of reach of observational investigation. But the venue of interferometry and high resolution technics provides hope to directly probe these regions. In this perspective, one must be aware that jet emitting discs have very different emission properties than usual standard accretion discs. In these inner regions, magnetic fields may indeed have a very important role, more than just triggering the magneto-rotational instability.

I would like to stress that each individual star is probably affected by the amount of magnetic flux available, which depends on its history. At a given time, a pre-main sequence star would be accreting magnetically channeled material from its circumstellar accretion disc. The disc would be radially stratified, going from an inner JED driving jets to an outer SAD. The inner parts of the SAD could host a dead zone where planet cores are formed. Time dependent accretion events would be thus expected. Most of what was reported here designs thereby a rather general picture. But it was confronted only to T Tauri stars. What's going on in more massive stars? Is this picture also valid for embedded objects? These are open questions that require observational inputs.

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