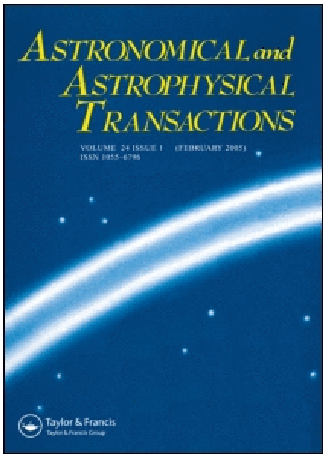

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C. Sauty [a]; E. Trussoni [b]; K. Tsinganos [c]
[a] Observatoire de Paris, DAEC, Meudon, cedex
[b] Osservatorio Astronomico di Torino, Pino Torinese
[c] Physics Dept, University of Crete, Greece

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CRITERIA FOR THE COLLIMATION OF WINDS INTO JETS

C. SAUTY[1], E. TRUSSONI[2], and K. TSINGANOS[3]

[1] *Observatoire de Paris, DAEC, F-92195 Meudon cedex*
[2] *Osservatorio Astronomico di Torino, I-10025 Pino Torinese*
[3] *University of Crete, Physics Dept, GR-71003 Greece*

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Continuous mass loss in the form of winds or jets is a widespread phenomenon in astrophysics, since it is observed in association with almost all stellar objects. By using a simple model for magnetized and rotating stationary MHD outflows from a central gravitating body, it is shown that magnetocentrifugally driven winds can collimate into thin jets for appropriate distributions of the various energetic terms across their source. This quantitative criterion suggests various scenarios for the evolution of collimated jets into ordinary winds, as the central star loses angular momentum and depending on whether the jet is under- or over-pressured. Finally, the equally important roles of gas pressure and magnetic fields in the confinement of jets are emphasized.

KEY WORDS Hydromagnetics, winds, jets, Sun, stars

1 INTRODUCTION

Young stellar objects are observed to be associated with accretion discs and jets (O'Dell and Wen, 1994; Ray *et al.*, 1996). A consequence of the continuous feeding of the central object by the accreting mass is to give rise to an outflowing counterpart. Such well-collimated outflows remove efficiently the strong excesses of angular momentum and magnetic flux from the forming star. Furthermore, as the central object evolves, its jet (Figure 1(a)) may undergo a transition to become an ordinary wind (Figure 1(b)). Then, during its main sequence life the star maintains a more spherical or bipolar outflow which still carries away mass and angular momentum at lower rates, although frequently with higher speeds. Despite the fact that the presence of accretion discs certainly seems crucial for the formation of jets, their presence does not necessarily imply that the outflowing plasma will eventually form a rather well collimated jet propagating at large distances.

The aim of the present contribution is to give some clues on how one might explain the various observed morphologies. From the technical point of view, one approach is to investigate solutions of the familiar MHD equations for a heated

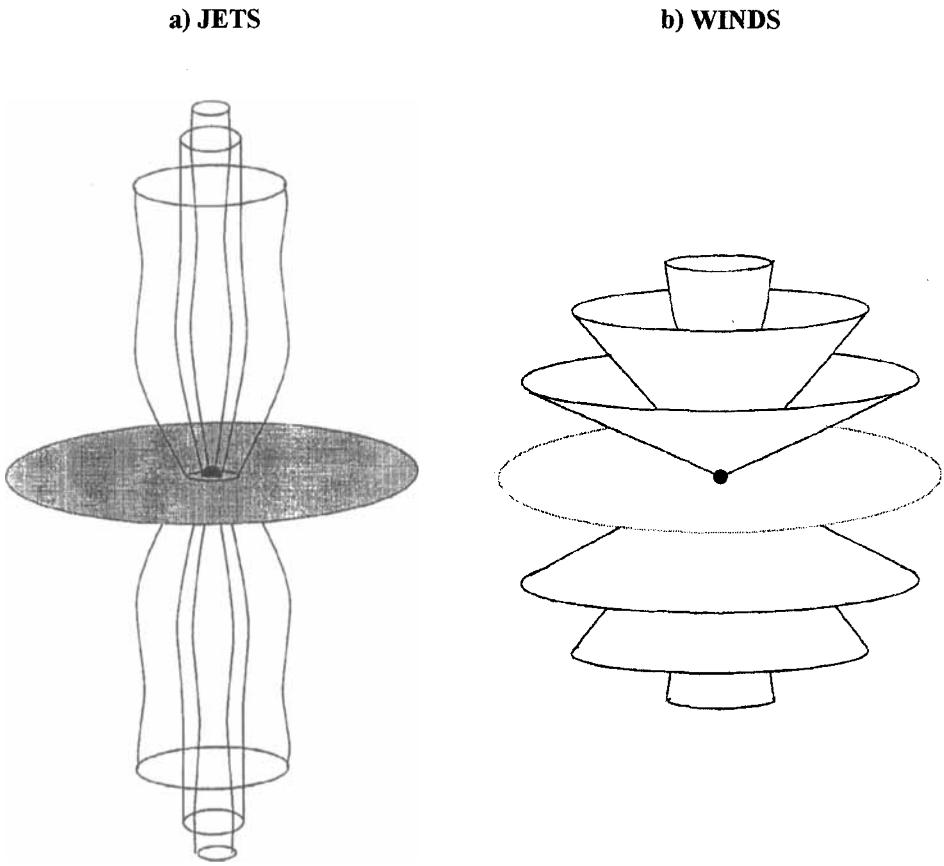

Figure 1 Two sketches showing two extreme morphologies of stellar outflows, as suggested by observations. Young stars are often associated with accretion discs and very well collimated jets (a). Main sequence stars like our Sun have a wind expanding roughly spherically (b).

plasma which is initially at rest and rotates with the stellar magnetosphere. The goal is to show how this plasma eventually escapes out of the gravitational potential well of the central body with rather high terminal speeds and some degree of collimation.

Some of the knots and shocks seen in jets are probably due to episodic ejections of mass, similar to solar coronal mass ejections and as such they may require time-dependent simulations (Kudoh and Shibata, 1997; Ouyed and Pudritz, 1997). However, the timescales for the variation of the outflow as inferred from the knot's proper motions and their spacing, can be months or years (Zinnecker, 1998). Some analogy can be drawn with the solar wind where a global steady outflow is observed during the quiet phase of the solar cycle and lasts for months. Thus the knotty structure of jets may also be the result of a rather slowly pulsating flow (Smith, 1998). In that respect, the study of steady flows is equally important. However the steady axisymmetric MHD equations are highly non-linear and cumbersome to solve. Nevertheless, part of their difficulty may be overcome without losing their fundamental non-linear character by using the so-called self-similar approach (Tsinganos *et al.*, 1996; Lima *et al.*, 1997; Vlahakis and Tsinganos, 1998). This amounts to splitting the variables and scaling all functions with one of the spatial coordinates. Radially self-similar solutions are usually applied to model disc winds ejected from the external part of an accreting system (Blandford and Payne, 1982; Contopoulos, 1995; Ferreira, 1997; Ostriker, 1997). Conversely, meridionally self-similar models are aimed at producing solutions describing the part of the ejection directly emanating from the star and its close vicinity (Sauty and Tsinganos, 1994: hereafter ST94; Trussoni *et al.*, 1997: hereafter TTS97).

The *combination* of gas pressure and magnetic fields may be essential in accelerating and confining jets. However, for the sake of simplicity these two ingredients are seldom considered simultaneously. For instance, most disc-wind models assume a cold plasma such that the magnetic field both accelerates and confines the flow (e.g., Blandford and Payne, 1982). Conversely, models similar to those using a De Laval nozzle mechanism (e.g., Kim and Raga, 1991) employ only gas pressure confinement. An MHD approach combining the magnetic and thermal effects has been proposed recently in a series of papers (Trussoni *et al.*, 1996, TTS97). In ST94 a criterion for the collimation of winds into jets was proposed by assuming that pressure does not play a direct role in confining the jet. We shall discuss some extensions of these results by including both ingredients and show how this suggests a scenario for the evolution of stellar jets into stellar winds.

2 STEADY AXISYMMETRIC FLOWS

Winds and jets can be considered as outflows of a fully ionized plasma flowing at a bulk speed $\boldsymbol{V}$ and carrying a magnetic field $\boldsymbol{B}$ in the gravitational field of a central body of mass $\mathcal{M}$. The ideal gas has a mean density ρ, pressure P and temperature $T \propto P/\rho$. The plasma dynamics is assumed to be governed by the ideal MHD equations. In particular, under steady and axisymmetric conditions, free integrals exist that remain constant on the magnetic surfaces generated by the revolution around the magnetic/flow symmetry axis of the system of a magnetic fieldline (Tsinganos, 1982): if A denotes the magnetic flux, they are $\Psi_A(A)$, the ratio of the mass and magnetic fluxes, $L(A)$, the *total* specific angular momentum carried by the flow *and* magnetic field, and $\Omega(A)$, the corotation frequency or angular velocity of each streamline at the base of the flow.

Furthermore, it is well known that the poloidal (p) and azimuthal (φ) components of the magnetic field and the velocity can be expressed in terms of these free integrals and the poloidal Alfvén Mach number $M = \sqrt{4\pi\rho(V_p^2/B_p^2)}$, using spherical (r, θ, φ) or cylindrical (ϖ, φ, z) coordinates (for details see Sauty and Tsinganos, 1994). The two integrals $L(A)$ and $\Omega(A)$ are not independent if the flow is trans-Alfvénic. In such a case, at the cylindrical distance ϖ_a of the Alfvén point ($M = 1$) from the field/flow axis of a flux tube labelled by A, they are related by $\varpi_a^2 = L/\Omega$.

Moreover the product of the two: $L\Omega = \varpi_a^2\Omega^2$ is known as the energy of the magnetic rotator (Belcher and MacGregor, 1976). In a cold plasma, ΩL is the potential energy which provides the energy to magnetically launch and confine the plasma.

3 A MERIDIONAL SELF-SIMILAR MODEL

By assuming self-similarity in the meridional (θ) direction, we can obtain solutions of the MHD equations in order to study the physical properties of the outflow close to its rotational axis. As in this region the contribution of the magnetocentrifugal forces to acceleration is small, the effect of a thermal driving force is important. This implies that the structure of the gas pressure in the flow is essential. To construct an exact analytical model (see ST94 for more details), we assume that the Alfvén surfaces are *spherical*, $M = M(R)$ and that the magnetic flux function A is expressed as the product of a function of the spherical radius, r, times some known function of the colatitude θ.

For convenience, all the variables are normalized to their respective values at the Alfvén surface along the rotation axis, $r = r_*$. In particular, we define the dimensionless radial distance $R = r/r_*$ and Alfvén speed $V_*^2 = B_*^2/4\pi\rho_*$, where B_*, V_* and ρ_* are the poloidal magnetic field, poloidal velocity and density along the polar axis at the characteristic radius r_*. For the magnetic flux function A we define its dimensionless form by

$$\alpha(R,\theta) = \frac{A(r,\theta)}{2r_*^2 B_*}. \tag{1}$$

Altogether, the four main assumptions of this meridionally self-similar model can be summarized as follows

$$\rho(R,\alpha) = \frac{\rho_*}{M^2(R)}(1+\delta\alpha), \quad \Psi_A^2 = 4\pi\rho_*(1+\delta\alpha), \tag{2}$$

$$\varpi^2(R,\alpha) = r_*^2 G^2(R)\alpha, \quad \varpi_a^2(\alpha) = r_*^2\alpha, \tag{3}$$

$$L\Psi_A = \lambda r_* B_* \alpha, \quad L\Omega\Psi_A^2 = \lambda^2 B_*^2 \alpha, \tag{4}$$

$$P(R,\alpha) = \frac{1}{2}\rho_* V_*^2 \Pi(R)[1+\kappa\alpha]. \tag{5}$$

In the frame work of this model, the asymptotic equilibrium is governed by the following five parameters: $\delta = \Delta\rho/\rho$ which controls the increase ($\delta > 0$) of the density with colatitude; $\kappa = \Delta P/P$ which controls the increase ($\kappa > 0$) or decrease ($\kappa < 0$) of the pressure with colatitude; $\nu = \sqrt{2\mathcal{G}\mathcal{M}/r_*}/V_*$, the ratio of the escape speed to the Alfvén speed along the polar axis, at the Alfvén distance, where $\mathcal{G}$ is the gravitational constant; λ, which controls the rotational rate, and finally $Q = \Pi(r=\infty)$, in dimensionless form the asymptotic value of the pressure. Note that here and in the rest of the text Δ is the difference of a quantity between a non-polar line and the polar axis. Note that λ^2 is a crucial parameter as it is proportional to the magnetic rotator energy, ΩL.

The efficiency of the magnetic rotator in confining the jet independently of the thermal gas pressure term depends on how much its energy is used to accelerate the jet. This is measured by the following parameter evaluated at the stellar surface r_o and normalized for convenience to the magnetic rotator energy,

$$\frac{\epsilon}{2\lambda^2} = \frac{E_{\mathrm{MR}} - E_{R,o} + \Delta E_G^*}{E_{\mathrm{MR}}}, \tag{6}$$

where $E_{\mathrm{MR}} = L\Omega$ is the energy of the magnetic rotator, $E_{R,o} = V_{\varphi,o}^2/2$ is the centrifugal energy at the footpoints and

$$\Delta E_G^* = -\frac{\mathcal{GM}}{r_o}\left[1 - \frac{T_o(\alpha)}{T_o(\mathrm{pole})}\right] \propto -\frac{(\delta-\kappa)\nu^2\alpha}{R_o(1+\alpha)}. \tag{7}$$

Along the polar axis the acceleration of the wind is done by thermal processes alone. On a non-polar line, if the relative increase of the density is not compensated by an identical relative increase of the pressure ($\delta - \kappa > 0$), then thermal driving is inefficient along this streamline and the assistance of magnetocentrifugal driving is needed. In other words, part of E_{MR} must be used to accelerate the non-polar flow. In that sense, $\epsilon/2\lambda^2$ is the relative amount of energy of the magnetic rotator that is not used to accelerate the flow out of the gravitational well or to make the flow rotating. In other words $\epsilon/2\lambda^2 < 0$ corresponds to *inefficient magnetic rotators* while $\epsilon/2\lambda^2 > 0$ to *efficient magnetic rotators* as far as magnetic confinement is concerned.

4 A CRITERION FOR THE COLLIMATION OF JETS

A fourth constant of the motion expresses conservation of energy along a given streamline. In fact, by projecting the momentum equation along this streamline, we obtain the generalized classical Bernoulli integral (ST94). It turns out that a more crucial constant is a slightly modified version of this integral if we introduce the *converted enthalpy*. This new quantity is defined at a given r along a given flux tube A by

$$\tilde{h}(r,A) = h(r,A) + \int_r^{\infty} \frac{g(r',A)}{\rho(r',A)V_r(r',A)}\,\mathrm{d}r' - h(\infty,A), \tag{8}$$

where q is the net local volumetric heating/cooling rate, and h the usual enthalpy of a perfect monatomic gas ($\gamma = 5/3$). In a thermally driven wind, all thermal input (internal enthalpy plus external heating provided along the flow) are not necessarily fully converted into other forms of energy, unless the terminal temperature is exactly zero. There always remains some asymptotic thermal content in the form of enthalpy. Thus the heat content that is really used by the flow is the *converted enthalpy* $\tilde{h}$.

Then we can define a generalized energy integral in the form,

$$\tilde{E}(A) = \frac{1}{2}V_p^2 + \frac{1}{2}V_\varphi^2 - \frac{\mathcal{GM}}{r} - \frac{\Omega}{\Psi_A}\varpi B_\varphi + \tilde{h}. \tag{9}$$

Thus, at a given radial distance r along the streamline labelled by A, the conserved energy $\tilde{E}(A)$ represents the sum of the kinetic, gravitational, Poynting and converted thermal energy flux densities per unit of mass flux density. Note that the Poynting flux is always positive because ΩB_φ is negative by construction (Tsinganos, 1982).

The variation of the conserved energy $\tilde{E}(A)$ across the streamlines normalized to the energy of the magnetic rotator gives an extra parameter of our model ϵ' that can be evaluated at the base r_o of the flow

$$\begin{aligned} \frac{\epsilon'}{2\lambda^2} &= \frac{\rho(R,\alpha)\tilde{E}(\alpha) - \rho(R,\mathrm{pole})\tilde{E}(\mathrm{pole})}{\rho(R,\alpha)L(\alpha)\Omega(\alpha)} \\ &= \frac{\Delta\left[\rho_o(E_{\mathrm{MR}} - E_{R,o} + \tilde{h}_o + E_{G,o})\right]}{\rho_o(\alpha)E_{\mathrm{MR}}}, \end{aligned} \tag{10}$$

where E_{MR}, $E_{R,o}$ and $\tilde{h}_o$ have been defined above and $E_{G,0} = -\mathcal{GM}/r_o$ is the gravitational energy at the stellar surface. The subscript "o" means that the quantity is evaluated at r_o. Assuming that the flow is cylindrically collimated, the same constant $\epsilon'/2\lambda^2$ evaluated at infinity contains only the variation across the lines of the Poynting flux plus the rotational energy. The sum of these two terms is positive on a non-polar line and zero along the axis (ST94) so the difference is positive. Thus cylindrical collimation can be achieved only if $\epsilon'/2\lambda^2$ is positive. This extends the criterion obtained in ST94 to the $\kappa \neq 0$ case. Note that in ST94, κ was zero and therefore ϵ' was identical to ϵ. This is obvious in Figure 2 where we see that the curve $\kappa = 0$ is quantitatively the same for G_∞ versus $\epsilon/2\lambda^2$ (left) and versus $\epsilon'/2\lambda^2$ (right).

We conclude that cylindrical collimation occurs only if there is an excess of the total amount of energy $\tilde{E}$ available on a non-polar line as compared to the polar axis. It turns out with our assumptions that ϵ' contains two terms,

$$\epsilon' \equiv \epsilon + \kappa\frac{V_\infty^2}{V_*^2}. \tag{11}$$

The first term is ϵ which simply accounts for the efficiency of the magnetic rotator as we have already seen. The second term is the variation across the lines of the heat content that is finally converted into kinetic energy. This term is linked with the variation across the lines of the pressure and thus is the efficiency of the pressure gradient in collimating (if $\kappa > 0$) or decollimating (if $\kappa < 0$) the flow. Hence *under-pressured* jets (i.e. jets with a pressure along the polar axis lower than the pressure at their outer edge, $\kappa > 0$) can be confined either because of the pressure gradients or because the magnetic rotator is efficient ($\epsilon > 0$). Conversely *over-pressured* jets (i.e. jets with a pressure along the axis higher than the pressure at their outer edge, $\kappa < 0$) are confined only if the magnetic rotator is efficient enough ($\epsilon > |\kappa|V_\infty^2/V_*^2$).

5 APPLICATION TO THE COLLIMATION OF JETS

For an asymptotically cylindrical jet, the force balance across the flux tube becomes $F_C + F_B + F_P = 0$ where F_C is the centrifugal force, F_B is the total magnetic transverse force (magnetic tension + magnetic pressure gradient) and F_P is the transverse pressure gradient. We deduce from this the asymptotic values of M_∞, G_∞ and V_∞/V_*, once we have fixed $\epsilon/2\lambda^2$, for given values of $\kappa/2\lambda^2$ and Q. In Figures 2(*a*), (*b*), (*c*), (*d*), the asymptotic cylindrical radius of a flux tube, G_∞, is plotted versus $\epsilon/2\lambda^2$. The corresponding G_∞ is also plotted versus $\epsilon'/2\lambda^2$, Figures 2(*a*'), (*b*'), (*c*'), (*d*').

There are three different domains where two of the three forces dominate the transverse momentum balance:

(1) $F_C + F_P \approx 0$, the thermally confined regime where the jet is centrifugally supported and pressure confined;

(2) $F_C + F_B \approx 0$, the magnetocentrifugal regime where the jet is magnetically confined and centrifugally supported;

(3) $F_B + F_P \approx 0$, the thermally supported regime where the jet is magnetically confined and pressure supported.

Note that in Figure 2(*a*), where $Q = 0$, besides the case of isopressured collimated jets ($\kappa = 0$) studied in ST94, cylindrical equilibrium can be found with negative ϵ if $k > 0$ and larger G_∞ if $\kappa < 0$. These solutions correspond to situations where the pressure becomes negligible asymptotically but it is still very important in the intermediate region between the source and infinity. Nevertheless they correspond to the same set of values for G_∞ and $\epsilon'/2\lambda^2$) as we see in Figure 2(*a*'). As the asymptotic pressure Q increases from Figure 2(*a*) to 2(*d*), the magnetocentrifugal domain shrinks around the curve $\kappa = 0$.

In Figure 2(*b*'), (*c*'), (*d*'), we see that the transition between the various asymptotic regimes of the forces correspond practically to the minimum of $\epsilon'/2\lambda^2$. For under-pressured jets ($\kappa > 0$) it almost corresponds also to the maximum in radius G_∞ (see Figure 2) while for over-pressured jets ($\kappa < 0$) it corresponds to a maximum in velocity (not shown here).

Finally, the case of over-pressured jets is very similar to the case of isopressured outflows ($\kappa = 0$) studied in ST94. If the magnetic rotator is efficient enough ($\epsilon/2\lambda^2 > |k|V_\infty^2/V_*^2$) cylindrical collimation is obtained and a jet can form. Otherwise the outflow remains radially expanding. Under-pressured outflows ($\kappa > 0$) behave quite differently as cylindrical collimation is always obtained due to the combination of pressure and magnetic confinement. As far as we know from observations of the solar wind it seems that faster winds and lower pressure flows come from the pole such that $\kappa > 0$. In the case of jets from young stars we may also infer that the central part of the winds seems to be less dense and with lower pressure than the surrounding molecular flows, suggesting again $\kappa > 0$. Although these conclusions may not be definitive we shall reduce our next discussion to the case of under-pressured outflows.

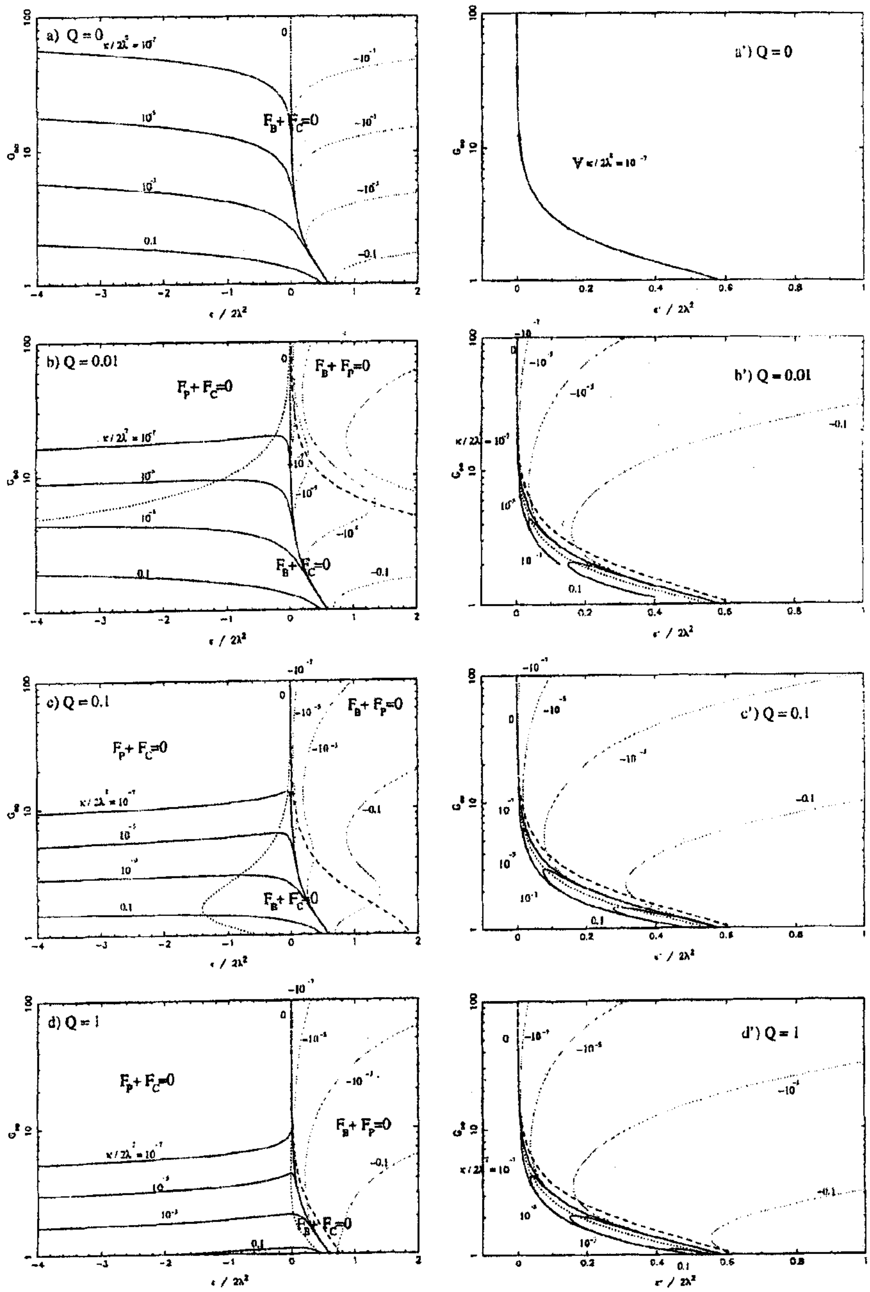

Figure 2 Plots of the asymptotic radius of cylindrically collimated outflows in units of the Alfvénic radius, $G_\infty = \varpi_\infty/\varpi_a$. The asymptotic radius is plotted versus $\epsilon/2\lambda^2$ in panels (*a*), (*b*), (*c*), (*d*) and versus $\epsilon'/2\lambda^2$ in panels (*a*'), (*b*'), (*c*'), (*d*'), for constant values of $\kappa/2\lambda^2$ and Q. The thermally confined region is separated from the magnetocentrifugal region by a dotted line while a dashed line separates the magnetocentrifugal domain from the thermally supported one.

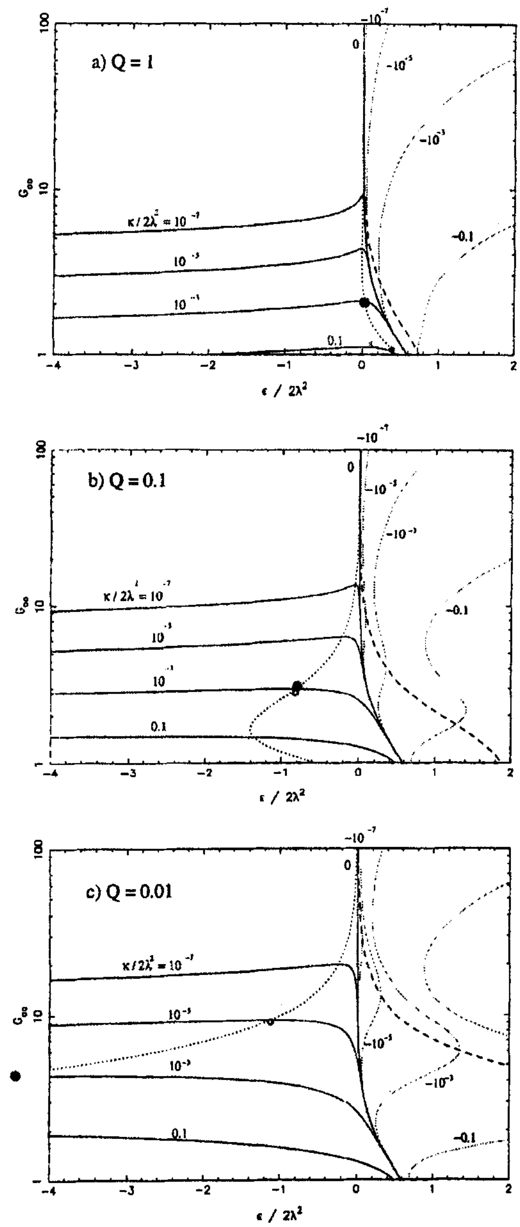

Figure 3 Two possible scenarios for the evolution of the jet radius under the condition of equal contribution to confinement by magnetic forces and gas pressure gradient in under-pressured flows. As the central star loses angular momentum, a first possible evolution indicated by a (•) is that the asymptotic pressure Q decreases successively from 1 to 0.1 to 0.01, as the magnetic efficiency of the rotator $\epsilon/2\lambda^2$ decreases, keeping $\kappa/2\lambda^2 = 0.001$ constant. A second possible evolution denoted by a ∘ is that both the asymptotic pressure Q and the latitudinal pressure variation $\kappa/2\lambda^2$ reduce as $\epsilon/2\lambda^2$ decreases from $Q = 1$ to 0.1 and then to 0.01 and $\kappa/2\lambda^2 = 0.1$ reduces to 10^{-3} and then to 10^{-5}.

6 A SCENARIO FOR THE EVOLUTION OF UNDER-PRESSURED JETS

In the case of under-pressured jets, as the flow extracts angular momentum from the central object, the initially efficient magnetic rotator of a newborn star becomes less and less efficient. Thus $\epsilon/2\lambda^2$ decreases from positive to negative values. An inspection of Figure 2 suggests that the originally magnetically confined jet now becomes pressure confined as the boundary dotted line is crossed. However, in a more realistic scenario, as the jet widens the asymptotic pressure Q may also decrease, as shown in Figure 3. Then, as $\epsilon/2\lambda^2$ decreases, the jet moves less quickly to the thermally confined regime. Among other possible scenarios an interesting one is to assume that the gas pressure and magnetic fields are equally important in the confinement of the jet (dotted line in the figures). As we have shown, this practically corresponds to minimum possible energy variation ($\epsilon'/2\lambda^2$) and maximum asymptotic radius. This might look like an optimum compromise between effective thermal acceleration and effective magnetic braking. Following this assumption we may go from Figure 3(a) to Figure 3(c) by reducing the asymptotic pressure Q but keeping $\kappa/2\lambda^2$ fixed. This is illustrated with a $\bullet$ in Figure 3. Thus, as the star evolves the jet radius widens but pressure and magnetic gradients continue to be equally important, although weaker in magnitude. Actually we may also expect that the latitudinal variation of the pressure κ also decreases, making the wind more homogeneous. Thus, another scenario is illustrated by a $\circ$ in Figure 3. In this case the outflow radius becomes even larger which implies that although the outflow is theoretically collimated, it may look like an ordinary wind from the observational point of view. Moreover, as the magnetic rotator reduces its strength, the thermal effects in the intermediate region become more and more important (the $\circ$ is farther and farther away from the $\kappa = 0$ curve) despite the fact that asymptotically we always have an equal contribution of the magnetic and pressure confinement, $F_B = F_P$.

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