THE EFFECT OF A STELLAR MAGNETIC VARIATION ON THE JET VELOCITY

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ABSTRACT

Stellar jets are normally constituted by chains of knots with some periodicity in their spatial distribution, corresponding to a variability of order of several years in the ejection from the protostar/disk system. A widely accepted theory for the presence of knots is related to the generation of internal working surfaces due to variations in the jet ejection velocity. In this paper we study the effect of variations in the inner disk-wind radius on the jet ejection velocity. We show that a small variation in the inner disk-wind radius produces a variation in the jet velocity large enough to generate the observed knots. We also show that the variation in the inner radius may be related to a variation of the stellar magnetic field.

Subject headings: accretion, accretion disks — ISM: Herbig-Haro objects — ISM: jets and outflows — stars: magnetic fields — stars: pre–main-sequence — stars: winds, outflows

1. INTRODUCTION

Stellar jets are observed in the form of chains of emitting nebulae called Herbig-Haro (HH) objects. HH objects are generally thought to be shock-heated density condensations traveling along the outflows formed by star-disk systems during the process of star formation (e.g., Reipurth & Bally 2001).

Since their discovery, several scenarios have been suggested for the origin of steady outflows from young stellar objects. In the stellar wind model, material is accelerated by thermal pressure gradients (e.g., Cantó 1980). Magnetohydrodynamics (MHD) models rely on the magnetocentrifugal launching mechanism (Blandford & Payne 1982). For the \( \dot{X} \)-wind scenario, the jet is magnetically driven from the so-called “\( \dot{X} \)-annulus,” where the young star’s magnetosphere interacts with the disk (Shu et al. 2000). In the disk-wind scenario, the jet is launched from an extended region of the disk surface (e.g., Ferreira 1997). The analytical models mentioned above have focused mainly on the steady state aspect of the ejection phenomena.

Unsteady periodic ejections with timescales of the order of several rotation periods of the inner disk radius have been obtained by numerical simulations (e.g., Ouyed & Pudritz 1997; Goodson & Winglee 1999; Matt et al. 2002). Observations indicate a significantly longer timescale is associated with the appearance of knots in stellar outflows. Nearly all observed jets present small-scale knots up to 0.1 pc from the central source, with a spacing between the knots corresponding to a timescale of \( \approx 1–20 \) yr (e.g., HH30, Burrows et al. 1996; HH111, Reipurth et al. 1992; RW-Aur, López-Martín et al. 2003). More fragmented knots are observed on a typical timescale of \( \approx 10^2–10^3 \) yr at larger distances from the source. While the long-term variation of jets can be explained by variations in the accretion rates (e.g., during FU-Orionis phases), the possible origin of small-scale knots is still unclear.

A first possibility, suggested by similarities with extragalactic jets, is that the knots may be formed by hydrodynamic Kelvin-Helmholtz instabilities (Micono et al. 1998), MHD Kelvin-Helmholtz reflective pinch mode instabilities (Cerqueira & de Gouveia Dal Pino 1999), or current-driven instabilities (Frank et al. 2000). Numerical simulations by these authors have shown that the shocks generated by plasma instabilities are weaker than those seen in observations once radiative cooling is taken into account. Moreover, some jets (e.g., HH212, Zinnecker et al. 1998) show a remarkable symmetry on both sides of the central star-disk system. This may be difficult to explain assuming that the instabilities are triggered by small-scale perturbations. In addition, and more importantly, optical images in many cases show that the compact knots have a bow-shock structure (e.g., HH111, Reipurth et al. 1992) that is difficult to obtain supposing that the knots are formed by instabilities.

A widely accepted theory for the presence of knots in stellar jets is related to the generation of internal working surfaces due to supersonic variations of the ejection velocity at the base of the jet (Raga et al. 1990). Indeed, numerical simulations, using velocity variations of about 10%–20% of the average velocity, have been able to reproduce the morphology and the emission properties of the knots in HH objects (e.g., Esquivel et al. 2007).

In this paper we show that velocity variations which lead to the creation of knots similar to the observed ones may be generated by variations of the inner disk-wind radius. Furthermore, we suggest that these variations may be related to changes in the stellar magnetic field.

This paper is organized as follows. In \( \S \) 2 we determine the effect of a variation in the inner disk radius on the jet velocity. In \( \S \) 3 we suggest that the variation in the inner radius could be connected to a change of the stellar magnetic field. In \( \S \) 4 we discuss the approximation used and the limitations of our model. Conclusions are given in \( \S \) 5.

2. THE EFFECT OF A VARIABLE INNER EJECTION RADIUS ON THE JET VELOCITY

We first discuss some asymptotic properties of a single field line anchored to an accretion disk as a function of the conserved invariants and the field line’s footpoint radius. Then we extend the discussion to an ensemble of field lines anchored at a range of footpoint radii and derive mean values for the asymptotic velocity of a disk wind extending from an inner to an outer finite...
radius. Lastly, we discuss the consequences of a time-dependent inner disk-wind radius on the average velocity of the outflow.

2.1. Asymptotics along a Single Field Line
In steady state, axisymmetric, ideal MHD, a number of quantities are conserved along a field line, i.e., a surface of constant magnetic flux \( \Psi \). Among these invariants are the mass-to-magnetic flux ratio \( k(\Psi) \), the total angular momentum \( L(\Psi) \), and the corotation frequency or angular velocity \( \Omega(\Psi) \) (e.g., Blandford & Payne 1982; Pelletier & Pudritz 1992). These invariants are defined as

\[
k(\Psi) = 4\pi \rho V_p B_p, \tag{1}
\]

\[
L(\Psi) = R \left[ V_\phi - \frac{B_\phi}{k(\Psi)} \right], \tag{2}
\]

\[
\Omega(\Psi) = \frac{1}{R} \left[ V_\phi - \frac{k(\Psi)}{4\pi \rho} B_\phi \right], \tag{3}
\]

where \( \rho \) is the mass density, \( V_p, V_\phi, B_p, \) and \( B_\phi \) are the poloidal and toroidal components of the velocity and magnetic fields, and \( R \) is the distance from the star. The invariants \( k(\Psi), \Omega(\Psi), \) and \( L(\Psi) \) and the footpoint \( R_0(\Psi) \) are assumed to be known or given functions.

The invariant \( L(\Psi) \) includes contributions from the twisted magnetic field. The relation

\[
L(\Psi) = R_A^2(\Psi) \Omega(\Psi) \tag{4}
\]

holds, where \( R_A(\Psi) \) is the Alfvén radius. Equation (4) determines \( R_A(\Psi) \) and subsequently \( \lambda(\Psi) = R_A(\Psi) / R(\Psi) \) (the magnetic lever arm) completely. Here, and in the following, all quantities with an explicitly indicated functional dependency on \( \Psi \), e.g., \( L(\Psi) \), are taken to depend on the magnetic flux \( \Psi \) only. Similarly, all quantities with a subscript \( 0 \) are taken to be defined at the base of the field line, i.e., the equatorial plane. The corotation frequency \( \Omega(\Psi) \) can be easily evaluated at the equator—where the poloidal velocity component \( V_p \) vanishes—and equals the orbital frequency at the equatorial plane \( \Omega_0 \),

\[
\Omega(\Psi) = \Omega_0(\Psi). \tag{5}
\]

Another conserved quantity is given by the total energy \( E(\Psi) \) per unit mass, with contributions from kinetic, thermal, and gravitational energy \( \Phi \) and the energy of the electromagnetic field (Poynting flux),

\[
E(\Psi) = \frac{1}{2} V^2 + \frac{\gamma}{\gamma - 1} \frac{P}{\rho} + \Phi - \frac{\Omega(\Psi) R B_\phi}{k(\Psi)}, \tag{6}
\]

where \( P \) is the thermal pressure and \( \gamma \) is the ratio of specific heats. The asymptotic jet velocity along a field line \( \Psi \) anchored at the footpoint radius \( R_0(\Psi) \) can be estimated from the Bernoulli equation for large distances and negligible enthalpy as (Michel 1969)

\[
v_\infty(\Psi) \sim \sqrt{2\Omega(\Psi) R_0(\Psi) \lambda(\Psi)}. \tag{7}
\]

The asymptotic jet velocity varies from field line to field line. However, we are ultimately interested in a typical value representative of the jet as a whole. Therefore, we weigh the asymptotic velocity with the mass flux carried by the field line. The asymptotic mass flux along a field line per unit magnetic flux (for one hemisphere) can be estimated by rewriting equation (1) as

\[
\frac{\partial M}{\partial \Psi} = \frac{k(\Psi)}{2}, \tag{8}
\]

where \( M \) is the wind mass-loss rate.

2.2. Asymptotic Mean Jet Velocity
The jet is assumed to originate in a disk wind. Mass loading onto the disk wind shall be efficient only for magnetic flux surfaces in the range \( \Psi_m \leq \Psi \leq \Psi_{ex} \), where the flux surfaces \( \Psi_m \) and \( \Psi_{ex} \) are anchored at \( R_m = R_0(\Psi_m) \) and \( R_{ex} = R_0(\Psi_{ex}) \), respectively. The mean or typical velocity of the jet given by the superposition of the asymptotic velocities along the field lines anchored between \( R_m \) and \( R_{ex} \) can be determined by weighting the jet velocity by the mass flux carried along every field line,

\[
\langle v \rangle = \frac{\int_{R_m}^{R_{ex}} v_\infty dM}{\int_{R_m}^{R_{ex}} dM}. \tag{9}
\]

So far, we have assumed steady state and axisymmetry. In the following we add two assumptions, namely, Keplerian rotation and self-similarity (see, e.g., Vlahakis & Tsinganos 1998).

Assuming Keplerian rotation for the plasma in the equatorial plane fixes the corotation frequency to

\[
\Omega(\Psi) = \Omega_0[R_0(\Psi)] = \sqrt{\frac{GM}{R_0^3(\Psi)}} \tag{10}
\]

The magnetic flux is a power law in \( \Omega \), i.e., \( \Psi \sim \Omega^{-\alpha} \). Self-similarity fixes the power-law index to \( \alpha = 1/2 \) (Blandford & Payne 1982). In terms of the footpoint radius \( R_0 \), \( \Psi \) is given by

\[
\Psi = \Psi_{ex}(R_0/R_{ex})^{3/4}. \tag{11}
\]

Self-similarity also fixes the mass-to-magnetic flux function \( k(\Psi) \) to

\[
k(\Psi) = k_{ex}(R_0/R_{ex})^{-3/4} \tag{12}
\]

and renders the lever arm independent of the magnetic flux surface [i.e., \( \lambda(\Psi) = \lambda \)]. Keplerian rotation and constant \( \lambda \) is sufficient to completely determine the asymptotic velocity along a field line as

\[
v_\infty(\Psi) = \lambda \sqrt{\frac{2GM}{R_0(\Psi)}}, \tag{13}
\]

while the wind mass-loss rate is determined by equation (8) and yields

\[
dM = \frac{3}{8} k_{ex} \Psi_{ex} \frac{dR_0}{R_0}. \tag{14}
\]

Finally, defining \( \chi = R_{ex}/R_m \), the average velocity (eq. [9]) can be integrated, giving

\[
\langle v \rangle = 2 \frac{\chi^{1/2} - 1}{\ln \chi} v_\infty(R_{ex}), \tag{15}
\]

where \( R_{ex} \) is assumed fixed. The velocity variation as a function of the inner radius, for different values of \( \chi \), is shown in Figure 1.
As this figure clearly shows, an increase in the inner radius causes a decrease in the average velocity. Moreover, the velocity variation is nearly independent of the ratio $\chi_0 = R_{ex}/R_{m,0}$ (where $R_{m,0}$ is a position of the inner radius chosen as a reference). To produce a variation of 20% in velocity it is necessary to have a variation of order of 30% in the inner radius. Figure 1 shows also the difference between the velocity of the disk wind at the inner radius (Keplerian, see eq. [13]) and the average velocity (defined by eq. [9]).

3. THE EFFECT OF A MAGNETIC FIELD CYCLE ON THE INNER DISK-WIND RADIUS

So far, we have dealt exclusively with the magnetic field of the disk wind. These magnetic flux surfaces carrying the disk wind are necessarily open to infinity. A nonopen field line cannot contribute to the global outflow. In the following we consider the superposition of a dipolar magnetosphere carried by the central star, contributing magnetic flux $\Psi_{dp}$ and the previously discussed magnetic field $\Psi$ which threads the disk and carries the outflow. Then the total magnetic flux is given by $\Psi + \Psi_{dp}$.

This configuration leads to three different kinds of magnetic field lines as is illustrated in Figure 2. First, there are field lines anchored in the polar region of the stellar surface. While these field lines are open to infinity, they are assumed to carry only negligible mass flux and do not contribute significantly to the outflow. Second, closed dipolar field lines anchored at lower latitudes of the stellar surface do not contribute to the outflow in our model. Lastly, there is a global open magnetic field threading the disk, which carries the disk wind.

The transition from the closed dipolar field lines to the open disk wind, i.e., the location of the innermost open flux surface $\Psi_{dp}$ in the equatorial plane, is given by the location of the saddle point of the magnetic flux distribution

$$\left. \frac{d}{dR_0} (\Psi + \Psi_{dp}) \right|_{R_0 = R_{m,0}} = 0,$$

where the dipolar field on the equatorial plane is given by a time-varying dipole moment $m(t)$ as

$$\Psi_{dp} \propto m(t) \frac{1}{R_0},$$

and $\Psi$ is given by equation (11). Since both magnetic flux distributions are known, the inner radius of the disk wind is given in terms of the time-varying stellar magnetic field (defined on the stellar surface) from equation (16) as

$$R_{in} \propto |m(t)|^{4/7}.$$

For simplicity, the stellar magnetic dipole moment is assumed to vary in time as

$$m(t) = m_0 [1 + \delta \sin (2\pi t/\tau_e)],$$

where $\tau_e$ and $\delta$ are the period and the amplitude of the stellar magnetic field variation, and $m_0$ is the stellar magnetic dipole moment at $t = 0$.

The inner radius of the ejecting region depends on the variation in the magnetic field intensity and changes on timescales of $\tau_e$. The average velocity will therefore change periodically following the magnetic field cyclic variation.

A plot of the velocity variations, that is, the difference between the maximum and minimum velocities, normalized to the average velocity during one cycle of the stellar magnetic field, is shown in Figure 3 as a function of $\delta$. The different curves are calculated corresponding to different values of $\chi_0 (=R_{ex}/R_{m,0})$. A
50% variation in the stellar magnetic field produces a ~20% jet
velocity variation, nearly independent of the value of \( \chi_0 \).

We have assumed that the only effect of a stellar magnetic field
variation on the ejected jet is a change in \( R_{in} \). Actually, some of
the open stellar flux will thread the disk. Therefore, the disk mag-
netic field (including also the contribution from the open stellar
magnetic flux) will not decrease as \( \sim R_0^{5/4} \) and will change with
time following the change in the stellar dipole.

A change in the stellar magnetic field produces changes in
the mass flux and velocity. In fact, the lever arm \( \lambda \) is a function
of the magnetic field and of the mass flux (that is also a function
of the magnetic field). Ferreira (1997) and Casse & Ferreira (2000)
studied models of self-similar MHD accretion disks driving jets.
These authors showed that \( \lambda^2 \sim 1 + 1/(2\xi) \), where \( \xi \) is the ejec-
tion parameter (that relates the mass flux to the radius by \( M \sim r^4 \)).
In addition, \( \xi \sim 0.1\mu \), where \( \mu \) is the magnetization parameter
[given by \( \mu = B^2/(8\pi p) \), where \( p \) is the thermal pressure].

In this case, equations (13) and (14) may be written as

\[
e_v(\Psi) = J(\Psi) \left[ \frac{2GM}{R_0(\Psi)} \right],
\]

where \( J(\Psi) \) is the mean jet velocity, \( M \) is the mass flux, \( \xi \) is the
magnetization parameter, and \( R_0 \) is the stellar radius.

An increase of the magnetic field in the disk (due to the increased
stellar contribution) produces an increase in the ejection parameter
\( \xi \) and of the mass flux and a drop in the asymptotic velocity.
Therefore, the jet velocity decreases further with respect to the
value calculated by equation (15).

The average velocity (eq. [9]) is given by

\[
\langle v \rangle = \frac{\int_{R_{in}}^{1} e_v(R_{in}) \lambda / \dot{J}(R_{in}) \xi^{-2} \, dR}{\int_{R_{in}}^{1} \dot{J}(R_{in}) \xi^{-2} \, dR},
\]

where \( R_{in} \) is the inner radius (assumed to be fixed) and \( e_v(R_{in}) \)
varies with the stellar magnetic field as a function of \( \dot{J}(R_{in}) \)
[\( \sim 1/(2\xi)^2 \) for \( \xi \ll 1 \)].

For \( \xi \rightarrow 0 \) (assuming that \( \dot{J} \) is converging to a large but finite
value) equation (22) reduces to equation (15).

The ratio between the external and inner radius \( \chi \) is difficult to
determine theoretically, but has been recently constrained by ob-
servations. In fact, interpreting the observed transverse velocity
shifts in T Tauri microjets as an indication of rotation, a range of
ejecting radii corresponding to \( \chi \approx 10 \) can be inferred (e.g.,
Ferreira et al. 2006).

Therefore, we may assume \( \xi \ll 1 \) and write the average out-
flow velocity normalized to \( \langle v \rangle \rightarrow \langle v \rangle_{R_{in}} \) as

\[
\langle v \rangle \approx \left( \frac{\langle v \rangle_{R_{in}}}{\langle v \rangle} \right)_{\xi=0} \frac{\xi_0}{\xi},
\]

where the scaling \( \langle v \rangle \sim 1/\sqrt{\xi} \) comes directly from \( e_v \sim \lambda \sim
1/\sqrt{\xi} \). From this equation, it becomes clear that the stellar mag-
netic field variation necessary to produce a 20% variation in the
average velocity may be much smaller than the 50% value deter-
mined from Figure 3. The 50% value represents an upper limit to the
necessary stellar magnetic field variation.

4. DISCUSSION

Is this section we discuss in some detail the approximations and
the limitations of the model. First, to use the results from the
standard disk-wind theory, we approximated the evolution of the
system as a sequence of steady state configurations.

We made two key assumptions on the behavior of the stellar
magnetic field. The first assumption of a dipolar stellar magnetic
field at the inner edge of the disk is standard in the magnetos-
hpheric accretion model for young stars (Königl 1991). On the
stellar surface, magnetic fields as large as several kG have been
observed (e.g., Johns-Krull 2007). The dipole component (~100–
500 G on the stellar surface), which decays more slowly than the
multipole components, dominates at large distances from the star.

The second hypothesis is that a dynamo mechanism operates
in protostars (Chabrier & Küker 2006; Dobler et al. 2006) with a
periodic or quasi-periodic variation of the magnetic field with a
typical timescale of a few to tens of years. A dynamo mechanism
is necessary to explain the observed magnetic field, because this
would be dissipated on scales of order of \( R_0/\xi \approx 100 \) yr, where
\( R_0 \) is the stellar radius and \( \xi \) is the turbulent magnetic diffusiv-
ity (e.g., Chabrier & Küker 2006). A dynamo mechanism, if
present, is not necessarily cyclic, as needed by our model. How-
ever, recently Sokoloff et al. (2008) proposed models of dyna-
mos in low-mass, fully convective stars with a cyclic magnetic
field. The photometric variability observed in T Tauri stars (e.g.,
Me’Nikov & Grankin 2005; Grankin et al. 2007) may be related to
a corresponding change in the stellar magnetic field (Armitage
1995).

Finally, our results depend on the complex interaction between
the stellar magnetic field and the disk. This has been studied by a
number of authors both analytically and numerically (see, e.g.,
Uzdensky 2004 and references therein). It is well understood that
the differential rotation in the disk produces a winding-up of the
magnetic field lines on a timescale \( \tau = 2\pi/(\Omega - \Omega_s) \) (where \( \Omega_s \)
is the stellar rotation frequency) with a following opening of the
stellar magnetic field lines and a consequent outflow.

The later evolution of the system is unclear. Once the mag-
netic field lines are opened, they can stay in that configuration
indefinitely (Lovelace et al. 1995), allowing a disk wind to be
traditionally ejected from an inner radius \( R_{in} \) to an outer radius \( R_{ex} \).
In this case, the stellar magnetic field lines are connected to the disk
only in a small region around \( R_{in} \) (see Fig. 2). The inner radius \( R_{in} \)
can be determined as the distance from the star where the mag-
netic stress \( B_x B_y R_{in}^2 \) begins to dominate over the viscous stress
\( M d(l_{R_{in}})/dR_{in} \) (e.g., Wang 1996),

\[
R_{in} = B_p(t)^{9/7} \left[ \frac{2R_0^6}{M(GM_*)^{1/2}} \right]^{2/7},
\]

where \( R_* \) and \( M_* \) are the star radius and mass, \( \Omega = (GM/R_0^3)^{1/2} \),
and we assume \( B_x \approx B_y \) (with \( B_x \propto m \), defined by eq. [19]) and
a full penetration of the vertical magnetic field lines into the disk.
The same dependence of \( R_{in} \) on \( B_p \) is recovered, and the results
obtained in § 3 are unchanged.

Alternatively, the stellar-disk magnetic field lines might close
again via magnetic reconnection, producing quasi-periodic eje-
tions (e.g., Goodson & Winglee 1999; Matt et al. 2002). Our
results correspond to the hypothesis that the mass flux carried by
the quasi-periodic ejections is much less than the mass flux carried
by the disk wind. If on the other side the mass flux carried by
the coronal mass ejection is important, the stellar cycle would
also produce a variation in the jet velocity. Actually, in contrast
with the disk-wind case, for a coronal mass ejection an increase
We concentrated on a cyclic magnetic field as the origin of the variations in the inner launching radius for the disk wind. However, any physical process that leads to quasi-periodic variations of the launching region will result in similar variations of the mean asymptotic velocity—and therefore to the appearance of knots—as long as those variations are large enough.

It is unclear which process (other than a stellar magnetic field variation) might create variations on a ~1–20 yr timescale. Spruit & Taam (1993) showed that a relaxation oscillator may produce a change in the inner radius in X-ray binaries. The timescale of the oscillation depends on the details of the stellar magnetosphere and can vary by 2 orders of magnitude. A variation in the accretion rate may also produce a similar variation in the inner radius (see eq. [24], and the discussion by Blackman et al. 2001).

Hartigan et al. (2004) showed that short-period T Tauri binaries do not seem to have jet properties different from those of wider binaries or single stars, arguing in this way against binary pairs of the origin of jet knots. Hartigan et al. (2004) also argued that there is no correlation between the ejection of a new knot in CW Tau and an increase in source brightness. One might have expected this if the knot was caused by an accretion outburst. Finally, the suggestion from Ferreira et al. (2006) that the alternation between X- and Y-type interactions between stellar and disk magnetic fields is due to a magnetic field cycle could explain the observed periodicity.

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5. CONCLUSIONS

In this paper we explored the effect of a stellar magnetic field variation on the ejection velocity of protostellar jets. While large-scale knots may successfully be explained by a large increase in the accretion rate, we showed that small-scale knots may be produced by changes in the inner radius.

We showed that a stellar magnetic field variation, if present, represents a natural candidate to produce stellar knots similar to those observed. In fact, a stellar magnetic field variation produces a change in the inner radius and, therefore, in the jet velocity. We also estimated that a stellar magnetic field variation ≤50% produces a variation (~20%) in the average velocity large enough to produce knots similar to the observed. The stellar magnetic field periodic or quasi-periodic variation remains an hypothesis of our scenario, but future large-scale temporal observations of magnetic field in young stars, and progress in dynamo theory, will both help to provide an answer to this problem.

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