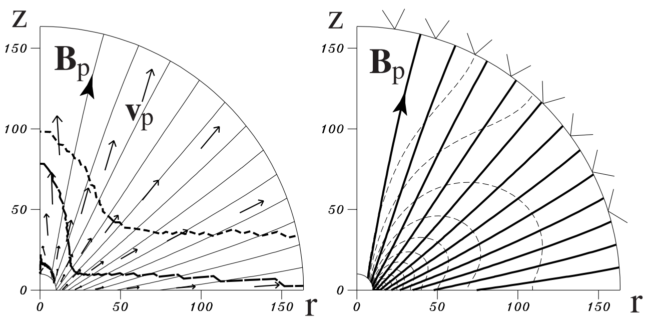

FIGURE 1. Simulation results for stationary MHD outflow obtained using a spherical coordinate system (Ustyugova *et al.* 1999). The solid lines represent the poloidal magnetic field, and the arrows the velocity vectors. The dashed lines in the left-hand plot represent the slow magnetosonic surface (the lowest dashed line), the Alfvén surface (the middle line), and the fast magnetosonic surface (the top line). In the right-hand panel, the dashed lines are surfaces of constant toroidal current density, while the lines on the outer boundary are the projections of the fast magnetosonic Mach cones.

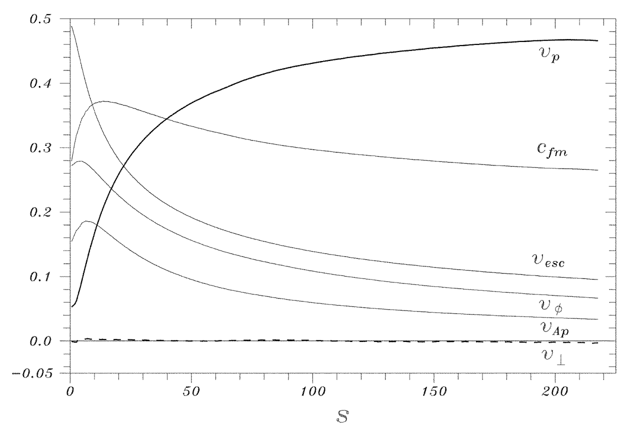

FIGURE 2. Dependences of different velocities on distance s measured in units of r_i (the inner radius of the disk) from the disk along the second magnetic field line away from the axis in Figure 1 (Ustyugova *et al.* 1999). This field line "starts" from the disk at $r \approx 6r_i$ where it has an angle $\theta \approx 28^o$ relative to the $z-$axis. The velocities are measured in units of $\sqrt{GM/r_i}$. Here, v_p is the poloidal velocity along the field line and $v_\perp$ is the poloidal velocity perpendicular to the field line. Also, v_{Ap} is the poloidal Alfvén velocity, c_{fm} is the fast magnetosonic velocity, and v_{esc} is the local escape velocity.

Li, & Blandford 1999; and Meier *et al.* 2001). Recently, laboratory experiments have succeeded in producing narrow high speed jets relevant to astrophysical jets (Hsu & Bellan 2002; Ciardi, *et al.* 2002).

A strong case for hydromagnetic jets as an explanation of jets in protostellar systems emerges because the temperature of the inner regions of these systems is insufficient to permit driving by thermal or radiation pressure (Königl & Ruden 1993). Observations of optical, proto-stellar jets (Bührke *et al.* 1988) reveal jet velocities $\sim 200-400$ km/s, which are comparable to the Keplerian disk velocity close to the star's surface.

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Magnetohydrodynamic Origin of Jets from Accretion Disks

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Abstract. The jets observed to emanate from many compact accreting objects may arise from the twisting of a magnetic field threading a differentially rotating accretion disk which acts to magnetically extract angular momentum and energy from the disk. Two main regimes have been discussed, *hydromagnetic outflows*, which have a significant mass flux and have energy and angular momentum carried by both matter and electromagnetic field and, *Poynting outflows*, where the mass flux is negligible and energy and angular momentum are carried predominantly by the electromagnetic field. Here, we give an overview of theoretical and simulation studies of non-relativistic hydromagnetic and Poynting jets and new theoretical work on relativistic Poynting jets from accretion disks.

INTRODUCTION

Powerful, highly-collimated, oppositely directed jets are observed in active galaxies and quasars (see for example Bridle & Eilek 1984), and in old compact stars in binaries - the "microquasars" (Mirabel & Rodriguez 1994; Eikenberry *et al.* 1998). Further, highly collimated emission line jets are seen in young stellar objects (Bührke, Mundt, & Ray 1988). Different models have been put forward to explain astrophysical jets (Bisnovatyi-Kogan & Lovelace 2000). Recent observational and theoretical work favors models where twisting of an ordered magnetic field threading an accretion disk acts to magnetically accelerate the jets. Two main regimes have been considered in theoretical models, the *hydromagnetic regime* where the energy and angular momentum is carried by both the electromagnetic field and the kinetic flux of matter, and the *Poynting flux regime* where the energy and angular momentum outflow from the disk is carried predominantly by the electromagnetic field.

HYDROMAGNETIC JETS AND OUTFLOWS

The theory of the origin of hydromagnetic outflows has been developed by many authors (Blandford & Payne 1982; Pudritz & Norman 1986; Sakurai 1987; Koupelis & Van Horn 1989; Lovelace, Berk & Contopoulos 1991; Pelletier & Pudritz 1992; Königl & Ruden 1993; Cao & Spruit 1994; Contopoulos & Lovelace 1994; Ostriker 1997; Beskin & Malyshkin 2000). Physical understanding of the hydromagnetic outflows from disks has developed from MHD simulations (Uchida & Shibata 1985; Bell 1994; Ustyugova *et al.* 1995, 1999; Romanova *et al.* 1997, 1998; Ouyed & Pudritz 1997; Krasnopolsky,

Figure 1 shows the nature of the stationary MHD outflows found by Ustyugova *et al.* (1999). The outflows are approximately spherical, with only small collimation within the simulation region. The *collimation distance* over which the flow becomes collimated (with divergence less than, say, 10°) is much larger than the size of the simulation region. Because the flow is super fast-magnetosonic (see below), the flow is not expected to magnetically collimate at larger distances. Notice for example that the solar wind is also super fast-magnetosonic and essentially uncollimated. The collimation observed in protostellar jets at small distances, for example, in HH 30 observed with the space telescope, may result from the pressure of the external medium (Lovelace *et al.* 1991).

The conclusion by Heyvaerts and Norman (1989) that MHD jets always collimate is based on the premise that the entire half-space above the disk is described by a stationary solution $\Psi(r,z)$ of the Grad-Shafranov equation. In reality, MHD jets result from initial conditions at the disk (as in the simulation studies) which result in the outward propagation of the jet into a highly conducting external (interstellar or intergalactic) medium. At a given time, the region inside the surface $r = R\,(z,t)$ is filled by the jet flow and the region outside it is filled by external plasma. The net axial current carried by the jet $2\pi \int_0^R rdrJ_z$ (for say $z > 0$) is necessarily zero for any (z,t). Therefore, the toroidal magnetic field is zero on the surface R , and there is *no* magnetic confinement of the flow. Of course, it is possible that a fraction of the flux Ψ is collimated along the $z-$axis while the remaining flux is "anti-collimated" (that is, pushed away) from the $z-$axis. Exactly this situation is found to occur for the Poynting jets discussed in the next section. This case has the physical condition that the overall outflow carries no net magnetic flux, $2\pi \int_0^R rdrB_z = 0$.

Figure 2 shows the variation of the flow speed and other characteristic speeds along a sample streamline. The outflow accelerates from thermal velocity ($\sim c_s$) to a much larger asymptotic poloidal flow velocity of the order of one-half $v_K = \sqrt{GM/r_i}$ where M is the mass of the central object and r_i is the inner radius of the disk. This asymptotic velocity is much larger than the local escape speed and is larger than fast magnetosonic speed by a factor of ~ 1.75. The *acceleration distance* for the outflow, over which the flow accelerates from ~ 0 to, say, 90% of the asymptotic speed, occurs at a flow distance of about $80r_i$. Close to the disk the outflow is driven by the centrifugal force while at all larger distances the flow is driven by the magnetic force which is proportional to $-\nabla(rB_\phi)^2$, where B_ϕ is the toroidal field.

POYNTING JETS

The powerful jets observed from active galaxies and quasars are probably *not* hydromagnetic outflows but rather Poynting flux dominated jets. The motion of these jets measured by very long baseline interferometry correspond to bulk Lorentz factors of $\Gamma = O(10)$ which is much larger than the Keplerian disk velocity predicted for hydromagnetic outflows. Furthermore, the low Faraday rotation measures observed for these jets at distances <kpc from the central object implies a very low plasma densities. Similar arguments indicate that the jets of microquasars are *not* hydromagnetic outflows but rather Poynting jets. Poynting jets have been proposed to be the driving mechanism for

gamma ray burst sources (Katz 1997) Theoretical studies have developed models for Poynting jets from accretion disks (Lovelace, Wang, & Sulkanen 1987; Lynden-Bell 1996, 2003; Romanova & Lovelace 1997; Levinson 1998; Li *et al.* 2001; and Lovelace *et al.* 2002; and Lovelace & Romanova 2003). Stationary Poynting flux dominated outflows were found by Romanova *et al.* (1998) and Ustyugova *et al.* (2000) in axisymmetric MHD simulations of the opening of magnetic loops threading a Keplerian disk.

We first summarize the theory of *non-relativistic* Poynting jets which is based on the Grad-Shafranov equation. We show results of *non-relativistic*, axisymmetric MHD simulations which support this theory. Finally we give the results of solving the *relativistic* Grad-Shafranov equation.

Consider the coronal magnetic field - such as that shown in the lower part of Figure 3 - of a differentially rotating Keplerian accretion disk. That is, the disk is perfectly conducting, high-density, and has a small accretion speed ($\ll v_K$). Further, consider "coronal or force-free" magnetic fields in the non-relativistic limit. We use cylindrical (r,ϕ,z) coordinates and consider axisymmetric field configurations. Thus the magnetic field has the form $\mathbf{B} = \mathbf{B}_p + B_\phi\hat{\phi}$, with $\mathbf{B}_p = B_r\hat{\mathbf{r}} + B_z\hat{\mathbf{z}}$. Because $\nabla\cdot\mathbf{B} = 0$, $\mathbf{B} = \nabla\times\mathbf{A}$ with $\mathbf{A}$ the vector potential. Consequently, $B_r = -(1/r)(\partial\Psi/\partial z)$,$B_z = (1/r)(\partial\Psi/\partial r)$, where $\Psi(r,z) \equiv rA_\phi(r,z)$. The $\Psi(r,z) =$ const lines label the poloidal field lines; that is, $(\mathbf{B}_p\cdot\nabla)\Psi = 0 = \mathbf{B}\cdot\nabla\Psi = 0$. Note that $2\pi\Psi(r,z)$ is the magnetic flux through a horizontal, coaxial circular disk of radius r. The magnetic field threading the disk at $z = 0$ is assumed to evolve slowly so that it can be considered approximately time-independent, $\Psi(r,z=0) = \Psi_0(r)$. However, the magnetic field above the disk will in general be time-dependent, $\Psi = \Psi(r,z,t)$, due to the differential rotation of the disk.

The *non-relativistic* equation of plasma motion in the corona of an accretion disk is $\rho d\mathbf{v}/dt = -\nabla p + \rho\mathbf{g} + \mathbf{J}\times\mathbf{B}/c$, where $\mathbf{v}$ is the flow velocity, p is the pressure, and $\mathbf{g}$ is the gravitational acceleration. The equation for the $\mathbf{B}$ field is $\nabla\times\mathbf{B} = 4\pi\mathbf{J}/c$, because the displacement current is negligible in the non-relativistic limit. In the "coronal or force-free plasma limit," the magnetic energy density $\mathbf{B}^2/(8\pi)$ is much larger than the kinetic or thermal energy densities; that is, sub-Alfvénic flow speeds $\mathbf{v}$, $\mathbf{v}^2 \ll v_A^2 = \mathbf{B}^2/4\pi\rho$, where v_A is the Alfvén velocity. The force equation then simplifies to $0 \approx \mathbf{J}\times\mathbf{B}$; therefore, $\mathbf{J} = \lambda\mathbf{B}$ (Gold & Hoyle 1960). Because $\nabla\cdot\mathbf{J} = 0$, $(\mathbf{B}\cdot\nabla)\lambda = 0$ and consequently $\lambda = \lambda(\Psi)$, as well-known. Thus Ampère's law becomes $\nabla\times\mathbf{B} = 4\pi\lambda(\Psi)\mathbf{B}/c$. The r and z components of this equation imply $rB_\phi = H(\Psi)$, and $dH(\Psi)/d\Psi = 4\pi\lambda(\Psi)/c$, where $H(\Psi)$ is another function of Ψ. Thus, $H(\Psi) =$ const are lines of constant poloidal current density; $\mathbf{J}_p = (c/4\pi)(dH/d\Psi)\mathbf{B}_p$ so that $(\mathbf{J}_p\cdot\nabla)H = 0$. The toroidal component of Ampère's law gives the *non-relativistic* Grad-Shafranov (GS) equation,

$$\Delta^\star\Psi = -H(\Psi)\frac{dH(\Psi)}{d\Psi} \ , \qquad (1)$$

for Ψ. Here, $\Delta^\star \equiv \partial^2/\partial r^2 - (1/r)(\partial/\partial r) + \partial^2/\partial z^2$ is the adjoint Laplacian operator. Note that $\Delta^\star\Psi = r(\nabla^2 - 1/r^2)A_\phi$ and that $H(dH/d\Psi) = 4\pi rJ_\phi/c$. From Ampère's law, $\oint d\mathbf{l}\cdot\mathbf{B} = (4\pi/c)\int d\mathbf{S}\cdot\mathbf{J}$, so that $rB_\phi(r,z) = H(\Psi)$ is $(2/c)$ times the current flowing through a circular area of radius r (with normal $\hat{\mathbf{z}}$) labeled by $\Psi(r,z)$= const. Equivalently, $-H[\Psi(r,0)]$ is $(2/c)\times$ the current flowing into the area of the disk $\leq r$. For all cases studied here, $-H(\Psi)$ has a maximum so that the total current flowing into the disk for $r \leq r_m$ is $I_{tot} = (2/c)(-H)_{max}$, where r_m is such that $-H[\Psi(r_m,0)] = (-H)_{max}$ so that r_m is less than the radius of the $O-$point, r_0. The same total current I_{tot} flows out of the region of the disk $r = r_m$ to r_0.

The function $H(\Psi)$ must be determined before the GS equation can be solved. $H(\Psi)$ is determined by the differential rotation of the disk: The azimuthal *twist* of a given field line going from an inner footpoint at r_1 to an outer footpoint at r_2 is fixed by the differential rotation of the disk. The field line slippage speed through the disk due to the disk's finite magnetic diffusivity is estimated to be negligible compared with the Keplerian velocity v_K. For a given field line we have $rd\phi/B_\phi = ds_p/B_p$, where $ds_p \equiv \sqrt{dr^2+dz^2}$ is the poloidal arc length along the field line, and $B_p \equiv \sqrt{B_r^2+B_z^2}$. The total twist of a field line loop is

$$\Delta\phi(\Psi) = \int_1^2 ds_p \, \frac{-B_\phi}{rB_p} = -H(\Psi)\int_1^2 \frac{ds_p}{r^2 B_p} \,, \tag{2}$$

with the sign included to give $\Delta\phi > 0$. For a Keplerian disk around an object of mass M the angular rotation is $\Omega_K = \sqrt{GM/r^3}$ so that the field line twist after a time t is $\Delta\phi(\Psi) = \Omega_0\, t[(r_0/r_1)^{3/2} - (r_0/r_2)^{3/2}] = (\Omega_0\, t)\; F(\Psi/\Psi_0)$, where r_0 is the radius of the $O-$point, $\Omega_0 \equiv \sqrt{GM/r_0^3}$, and F is a dimensionless function.

Equations (1) and (2) have been solved numerically by Li *et al.* (2001) and Lovelace *et al.* (2002) for an initial poloidal magnetic field as shown in the lower part of Figure 3. As the "twist," as measured by $\omega_0 t$, increases a high twist field configuration appears with a different topology. A "plasmoid" consisting of toroidal flux detaches from the disk and propagates outward. The plasmoid is bounded by a poloidal field line which has an $X-$point above the $O-$point on the disk. The occurrence of the $X-$point requires that there be at least a small amount of dissipation in the evolution from the poloidal dipole field and the Poynting jet configuration. The high-twist configuration consists of a region near the axis which is *magnetically collimated* by the toroidal B_ϕ field and a region far from the axis which is *anti-collimated* in the sense that it is pushed away from the axis. The field lines returning to the disk at $r > r_0$ are anti-collimated by the pressure of the toroidal magnetic field. The *poloidal field* fills only a small part of the coronal space.

Figure 4 shows results of *non-relativistic* axisymmetric simulations of the formation of a Poynting jet by Ustyugova *et al.* (2000). The flow near the $z-$axis is the Poynting jet and its physical properties agree with Grad-Shafranov solutions.

In the case of *relativistic* Poynting jets we hypothesize that the magnetic field configuration is similar that in the non-relativistic limit (Ustyugova *et al.* 2000; Lovelace *et al.* 2002). Thus, most of the twist $\Delta\phi$ of a field line of the relativistic Poynting jet occurs along the jet from $z = 0$ to $Z(t)$ as sketched in Figure 3, where $Z(t)$ is the axial location of the "head" of the jet. Along most of the distance $z = 0 - Z$ the radius of the jet is a constant and $\Psi = \Psi(r)$ for $Z \gg r_0$. Note that the function $\Psi(r)$ is different from $\Psi(r,0)$ which is the flux function profile on the disk surface. Hence, $r^2 d\phi/dz = rB_\phi(r,z)/B_z(r,z)$. We take for simplicity that $V_z = dZ/dt =$ const. We determine V_z in §3. In this case, $H(\Psi) = [r^2\Omega(\Psi)/V_z]B_z$, where the right-hand side can be written as a function of Ψ and $d\Psi/dr$. With $H(\Psi)$ known, the relativistic Grad-

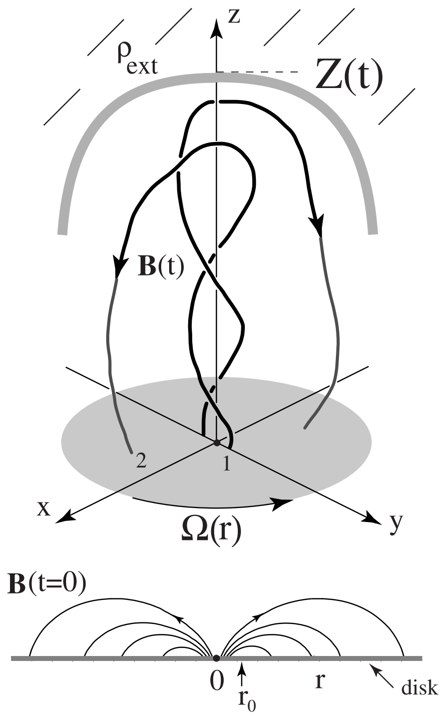

FIGURE 3. Sketch of the magnetic field configuration of a Poynting jet from Lovelace and Romanova (2003). The bottom part of the figure shows the initial dipole-like magnetic field threading the disk which rotates at the angular rate $\Omega(r)$. The top part of the figure shows the jet at some time later when the head of the jet is at a distance $Z(t)$. At the head of the jet there is force balance between electromagnetic stress of the jet and the ram pressure of the ambient medium of density ρ_{ext}.

Shafranov equation,

$$\left[1-\left(\frac{r\Omega}{c}\right)^2\right]\Delta^\star\Psi-\frac{\nabla\Psi}{2r^2}\cdot\nabla\left(\frac{r^4\Omega^2}{c^2}\right)=-\,H(\Psi)\frac{dH(\Psi)}{d\Psi}\;, \tag{3}$$

can be solved (Lovelace & Romanova 2003).

The quantity not determined by equation (3) is velocity V_z or Lorentz factor $\Gamma = 1/(1-V_z^2/c^2)^{1/2}$. This is determined by taking into account the balance of axial forces at the head of the jet: the electromagnetic pressure within the jet is balanced against the dynamic pressure of the external medium which is assumed uniform with density ρ_{ext}. This gives $(\Gamma^2-1)^3 = B_0^2/[8\pi R^{\,2}\rho_{ext}c^2]$, or for $\Gamma \gg 1$,

$$\Gamma \approx 8\left(\frac{10}{R}\right)^{1/3}\left(\frac{B_0}{10^3\mathrm{G}}\right)^{1/3}\left(\frac{1/\mathrm{cm}^3}{n_{ext}}\right)^{1/6}\;, \tag{4}$$

where $R\,\equiv r_0/r_g \gg 1$, with r_0 the $O-$point of the magnetic field, $r_g \equiv GM/c^2$, and B_0 the magnetic field strength at the center of the disk. This value of Γ is of the order of the Lorentz factors of the expansion of parsec-scale extragalactic radio jets observed

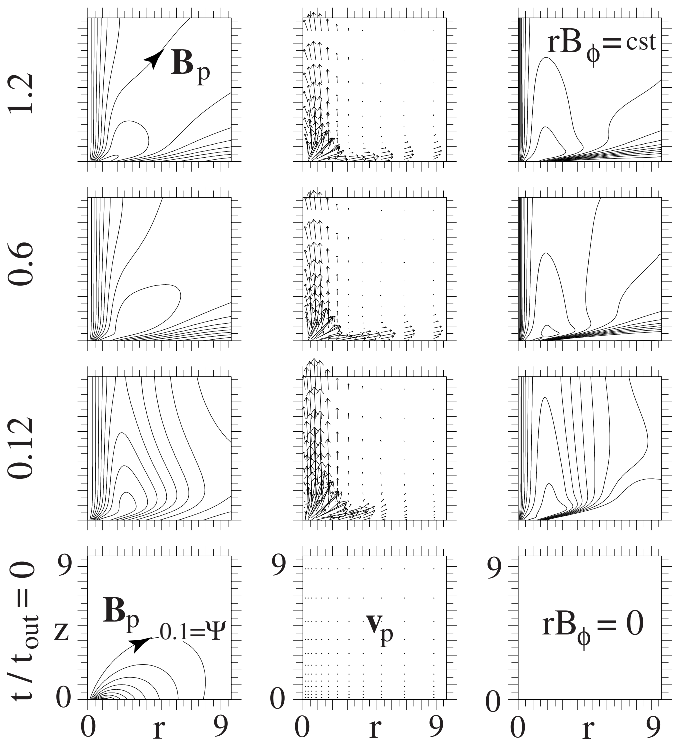

FIGURE 4. Non-relativistic time evolution of dipole-like field threading the disk from the initial configuration $t = 0$ (bottom panels) to the final quasi-stationary state $t = 1.2t_{out}$, where t_{out} is the rotation period of the disk at the outer radius R_{max} of the simulation region from Ustyugova *et al.* (2000). The left-hand panels show the poloidal field lines which are the same as $\Psi(r,z)$ =const lines; Ψ is normalized by Ψ_{max} and the spacing between lines is 0.1. The middle panels show the poloidal velocity vectors $\mathbf{v}_p$. The right-hand panels show the constant lines of $-rB_\phi(r,z) > 0$ in units of Ψ_{max}/r_0 and the spacing between lines is 0.1. For this calculation a 100×100 inhomogeneous grid was used with Δr_j and Δz_k growing with distance r and z geometrically as $\Delta r_j = \Delta r_1 q^j$ and $\Delta z_k = \Delta z_1 q^k$, with $q = 1.03$ and $\Delta r_1 = \Delta z_1 = 0.05 r_0$.

with very-long-baseline-interferometry (see, e.g., Zensus *et al.* 1998). This interpretation assumes that the radiating electrons (and/or positrons) are accelerated to high Lorentz factors ($\gamma \sim 10^3$) at the jet front and move with a bulk Lorentz factor Γ relative to the observer. The luminosity of the $+z$ Poynting jet is $\dot{E}_j = c \int_0^{r_2} rdrE_r B_\phi/2 = cB_0^2 \mathcal{R}\, r_g^2/3 \approx 2.2 \times 10^{45} (B_0/10^3\mathrm{G})^2 (\mathcal{R}/10)(M/10^9 M_\odot)^2$erg/s, where M is the mass of the black hole.

For long time-scales, the Poynting jet is of course time-dependent due to the angular momentum it extracts from the inner disk ($r < r_0$) which in turn causes r_0 to decrease with time (Lovelace *et al.* 2002). This loss of angular momentum leads to a "global magnetic instability" and collapse of the inner disk (Lovelace *et al.*1994, 1997, 2002) and a corresponding outburst of energy in the jets from the two sides of the disk. Such outbursts may explain the flares of active galactic nuclei blazar sources (Romanova & Lovelace 1997; Levinson 1998) and the one-time outbursts of gamma ray burst sources (Katz 1997).

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