Simulations of Jet Production in Magnetized Accretion Disk Coronae

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Abstract. We describe the results of over 40 MHD simulations, performed on Caltech/JPL parallel supercomputers, of the coronae of magnetized accretion disks around compact objects. All produce some type of outflow from the disk. Our parameter study investigated the character of the outflow as a function of the strength of the initial poloidal magnetic field and its angle with respect to the disk rotation axis. When the radial component of the field is significant, this outflow takes the form of a collimated jet ejected from the center of the accretion disk. The jet velocity is a strong function of the strength of the initial magnetic field: for Alfvén velocities *(VA)* below the escape speed (V_{esc}), the jet velocity is of order *VA*, but for *VA* only slightly above V_{esc}, the jet velocity is an order of magnitude or more greater.

This "magnetic switch" behaves similarly for a broad range of magnetic field polar angle. However, when the initial coronal magnetic field is nearly completely dominated by an axial component, the import ante of the central jet diminishes and the outflow becomes dominated by a poorly-collimated wind from a broader region of the accretion disk.

The magnetic switch may have applications to galactic and extragalactic radio sources and other objects.

1. Introduction

Jets are observed to be ejected from a variety of astrophysical objects: the centers of radio galaxies and quasars, very old stars (planetary nebulae, accret ing neutron star and black hole binary systems), and very young stars (e. g., primitive solar system nebulae).

One of the most promising models for producing jets in accreting systems is the Bland ford-Payne magnetohydrodynamic (MHD) process in the coronae of thin accretion disks (Blandford & Payne, 1982). In this model a magnetocentrifugally-driven wind develops if the angle θ between the poloidal magnetic field component $B_P (\equiv [B_R^2 + B_Z^2]^{1/2})$ and the disk rotation axis is greater than 30°. Because of differential rotation, the magnetic field lines wrap around the Z axis, providing a Lorentz hoop stress that collimates the flow into a jet.

A variety of approaches have been used to study this mechanism in more detail. Most analytical or semi-analytical approaches assume open magnetic field lines anchored in an infinitely thin and dense Keplerian disk and then look for steady-state, collimated wind solutions in the disk corona (Kudoh & Shibata, 1996; Li *et al.*, 1992, hereafter LCB; Lovelace *et al.*, 1986). These studies have been successful in obtaining jet flow speeds with Lorentz factors $\Gamma \equiv (1 - v^2/c^2)^{-1/2} > 10$ (LCB), typical of the fastest quasar jets observed. Most numerical approaches, on the other hand, assume a disk of finite thickness, initially t breaded by open magnetic field lines, and use angular moment urn loss and accretion in the disk to produce a θ > 30° condition which then ejects some of the disk material itself in a jet (Uchida & Shibata, 1985; Stone & Norman, 1994; Bell & Lucek, 1995; Matsumoto, 1997; Koide *et al.*, 1997). However, the jets produced in these simulations appear to have a low Γ factor. A few investigators have studied the anchored field line/ejected disk corona case numerically (Ustyugova *et* al., 1995; Ouyed & Pudritz, these proceedings)} However, none of these have obtained jet flow speeds approaching the highly relativistic ones predicted by the semi-analytic studies either.

2. "Weakly" Magnetized Accretion Disks and Coronal Flow

Assuming that the magnetic field is anchored in a dense disk and that the disk corona is the source of the jet, is more than a convenient mathematical simplification. A good parameter for measuring the importance of magnetic forces is the ratio of the Alfvén velocity $[V_A \equiv B/(4\pi\rho)^{1/2}]$ to the escape velocity. (Here B is the magnetic field strength in Gauss and ρ is the mass density of the plasma.) Standard accretion disk models like the a-model (Shakura & Sunyaev, 1973) naturally have a low Alfvén velocity in the disk:

$$V_{\rm A,\,disk} \equiv \frac{B}{(4\pi\rho_{\rm disk})^{1/2}} \sim \alpha^{1/2}\frac{H}{R}V_{\rm esc} \ll V_{\rm esc} \qquad (1)$$

where $V_{\rm A,\,disk}$ and $V_{\rm esc}$ are the local Alfvén and escape velocities respectively in the disk, $H << R$ is the disk half-thickness at some radius R, and $\alpha < 1$ is the usual viscosity parameter. While even weak disk fields can have significant consequences for the internal structure of the disk and probably play an important role in the viscosity itself, the effective α from local shear instabilities is expected to be small ($\alpha \sim 0.01$; Brandenburg *et al.*, 1996; Stone *et al.*, 1996). Since the inward accretion drift time scale is several α^{-1} orbital times, the assumption of a coronal magnetic field anchored at a fixed radius is a good approximation for dynamical calculations.

Furthermore, although the Alfvén velocity is low in the disk, it can still be high in the corona because of the low density there. (Such conditions are not unlike those in the solar atmosphere and corona.) In fact, there exists a critical coronal density $\rho_{\rm crit} \equiv \alpha (H/R)^2 \rho_{\rm disk}$ such that if $\rho_{\rm corona} \lesssim \rho_{\rm crit}$, then $V_{\rm A,\,corona} \gtrsim V_{\rm esc}$. For coronae of sufficiently low density then, magnetic forces can dominate gravitational forces and significantly affect the dynamics of the corona, even though they are not strong enough to disrupt the disk itself.

3. The MHD Simulations

We performed formally non-relativistic, two-dimensional axisymmetric magnetohydrodynamic simulations of magnetized accretion disk corona, which include toroidal, as well as poloidal, magnetic field and velocity components and gravity.

3.1. Initial and boundary conditions

For initial conditions we assume that open poloidal magnetic field lines (B_{ϕ_0} = O) are anchored in a dense, cold, Keplerian rotating accretion disk and thread a warm, also rot ating, corona whose temperature is a significant fraction of virial but is still bound in a relatively thin structure. Permeating the entire region is a tenuous, hot halo ($\rho_{\rm halo}/\rho_{\rm corona}$ = 0.1) in pressure equilibrium with the corona. The strength of all parameters varies with cylindrical radius as a Shakura & Sunyaev (1973) accretion disk, such that the ratio of local velocities in the corona (V_A, $V_{\rm sound}$, and $V_{\rm wind}$) to the local escape velocity $V_{\rm esc}$ *is* a constant. The parameter *VA*, $_{\rm corona}/V_{\rm esc}$ is allowed to vary between 0.2 and 4,0 and the initial poloidal field angle θ_0 at the base of the corona (Z = O) is allowed to vary between 8° and 83° [log(tan 0) = –0.861 to 0.889 in steps of 0.25].

For boundary conditions, we assume "flow through" conditions on the Z = $Z_{\rm max}$ and R = $R_{\rm max}$ outer boundaries and cylindrical symmetry at $R = O$. The Z = O boundary condition perpetuates the initial conditions by continually injecting magnetized coronal material in a low velocity wind ($V_{\rm wind}/V_{\rm esc} \sim$ 0.05). The simulations were dimensionless (G = M = 1.0). To re-int reduce scales, one must specify a measure of the compact object radius $R_{\rm g}$, and the Keplerian velocity at $R_{\rm g}$ [$V_{\rm g} = (GM/R_{\rm g})^{1/2}$], with the unit of time being their ratio.

3.2. The code

The simulations were performed using the code FLOW (Lind *et al.*, 1989) on the Caltech Intel Touchstone Delta parallel supercomputer using 150 radial by 300 axial zones. Tests of HD and MHD against analytic solutions indicate that the code maintains second-order accuracy in space and time. The main limitations are boundary reflection effects, which limit the total evolution time, and the cylindrical $R = O$ condition, which generates a small error in flow near the axis. Both affect very low $V_{\rm A}/V_{\rm esc}$ *flows,* Nevertheless, in this regime our simulations obtain results similar to others who computed low to moderate Alfvén velocity flows (Ustyugova *et al.*, 1995; Ouyed & Pudritz, these proceedings).

3.3. Relativistic flow

As long as $V_{\rm A}^2 < C^2$ and $V_{\rm sound}^2 < C^2$, even a non-relativistic code, can be used to investigate relativistic flow. If, in the relativistic fluid equations, the inertia of the magnetic and thermal energy densities are ignored and the equations are written in terms of the three spatial components of the four-velocity (u $\equiv \Gamma \mathbf{V}$), then their form is *identical* to that of the non-relativistic equations, save for one ternl—the electric force normal to the magnetic field lines—but this also does not become important until $V_A^2/c^2 \sim 1$. Note that, even when this term is important, it does not change the velocity or angular momentum along a field line, but does determine the exact angle and shape of the field lines (LCB). In addition. even when $V_A^2/c^2 \sim 1$, including the inertia terms would make the fluid only twice as heavy as in our simulations and the electric force would modify the tangent of the magnetic field by, at most, about a factor of 2. Since we investigate a wide range of magnetic field angle, and find similar effects throughout, we conclude that our non-relativistic MHD code, with suitable transformations and care in analyzing the results, can be used to investigate flow for $V_A/C \sim 1$ and $\Gamma_{jet} \rightarrow$ cm.

4. Results

Our results are reported and discussed in more detail elsewhere (Meier *et al.*, 1997). Here we give a brief account of them.

4.1. Typical model

Most thoroughly studied was the θ_0 = 54° branch, with V_A/V_{esc} ranging from 0.2 to 1.5. In addition, we computed a V_A/V_{esc} = O case as a control experiment, which produced no outflow of any kind. The V_A/V_{esc} = *1.2* case, which we studied most, produced a strong, well-collimated jet after only two inner disk rotation periods. The internal flow velocity of the jet was $\sim 19\ V_{esc_0}$, where V_{esc_0} is the escape velocity measured at the point in the disk where the coronal field has a maximum value (R. = 7.2). Because of the high velocity, with the mass flux limited to that injected at the base of the corona, the resulting jet density was very light ($\rho_{jet}/\rho_{ext} \equiv \eta \sim 10^{-3}$), so its actual propagation velocity across the grid was only $\sim 1\ V_{esc_0}$. The calculation was continued to late times (10 inner disk rotation periods), well after the jet propagated off the Z = Z_{max} boundary and reached a steady state.

4.2. Variations with magnetic field strength: The magnetic switch

Additional calculations were performed with higher and lower magnetic field strengths, Increasing VA to 1.5 V_{esc} increased the jet speed by about 70%, but *decreasing* V_A to ().7 V_{esc} decreased the internal jet speed by a factor of ~ 40 to 0.5 V_{esc_0} ! Further decreases in the VA/V_{esc} parameter resulted in jets with speeds of a few tenths of V_{esc_0}. There apparently is a boundary at $V_A/V_{esc} \sim 1$, below which jets have velocities $\sim VA$, but above which the jets attain high speeds. For classification purposes, we shall call the low-speed jets Type 1 flows and the high-speed jets Type 2 flows.

A simple physical explanation of this "magnetic switch" is as follows. When $V_A/V_{esc} < 1$, the flow is dominated by gravitational, not magnetic, forces. The jet is accelerated out of the potential well by the recoil of the toroidal magnetic pressure, similar to a Parker wind. When $V_A > V_{esc}$, magnetic forces dominate and gravitational effects are unimportant. The jet accelerates via the magneto-centrifugal effect, as occurs in pulsar wind solutions.

4.3. Variations with initial **poloidal** field angle

We investigated the magnetic switch effect further for a wide range of poloidal field angles. We found that the sharp transition from Type 1 to Type 2 jets occurs for all magnetic field angles investigated. Even the $\theta < 30°$ cases produced Type 2 jets if the magnetic field strength was great enough, and Type 1 jets otherwise. Therefore, while the magnetic switch is a strong function of field strength, it is not a strong function of the angle of the field at the base of the corona.

Why is this the case? Steady-state solutions predict no outflow for $\theta < 30°$. The answer is in the non-steady development of the jets at early times. Initially the field is twisted by differential rotation, yielding a strong toroidal component, The hoop stress from this component drives an inward flow just above the disk, pinching the field lines inward, enhancing a $\theta > 30°$ condition for those with $\theta_0 > 30°$ *and producing such a condition even when* $\theta_0 < 30°$ *initially.* Having achieved this critical Blandford-Payne condition dynamically, the system proceeds to drive an outflow which eventually collimates. Following the flow along a field line, we find it first bends inward toward the axis just above the disk and then turns outward to produce the outflow and jet. The importance of the angle θ_0 in non-steady simulations, then, is not that it is above or below 30°, but that it determines how much of the magnetic field is radial and hence can be wound into a toroidal field to initiate this process.

The very small angle cases ($\theta_0 \sim$ 8°-140) show an additional effect. As θ_0 decreases, keeping the magnetic field strength constant, the strength of the jet diminishes and an uncollimated or poorly-collimated wind appears from much of the disk. The wind velocity also seems to exhibit the magnetic switch effect as well, with its velocity low when $V_A/V_{esc} < 1$ and high otherwise.

5. Discussion

From our simulations we find that jets are ubiquitous, being produced by magnetized accretion disks in a wide range of parameter space, and that the magnetic switch is a robust effect as well. Scaling our results to the black hole and protostellar cases, we find the following. For the black hole case, making the transformation $V \rightarrow \Gamma V$, we find jet velocities of $V_{jet} \sim$ 0.2-0.5 c when the switch is off, and $\Gamma_{jet} \sim$ 5–20 when the switch is turned on. For the protostar case, $V_{jet} \sim 80\text{–}200$ km s^{-1} when the switch is off and $V_{jet} \sim$ 2,000-8,000 km s^{-1} when turned on. Whether or not such high velocity protostellar jets occur will depend on their accretion disks developing hot and readily -replenishable coronae.

Our results are consistent with both the semi-analytic results of LCB (when suitable relativistic transformations are taken into account) and of Ustyugova *et al.* (1995), and Ouyed & Pudritz (these proceedings) in their respective regions of parameter space.

Applications of the magnetic switch may be considerable. The galactic superluminal source GRO J1655-40 displayed such a bimodal behavior. remaining a weak radio source until $\sim$12 days after the X-ray outburst (Harmon *et* al.. 1995; Tingay *et* al., 1995), and only then producing a strong radio jet. The ejection of the jet appeared to coincide with the formation of a hot corona and a drop in the X-ray flux, troth indicators that the switch may have turned on because of processes occurring when accreting near the Eddington limit (Meier, 1997).

The strength of jets in some active galactic nuclei may also be explained by this mechanism. The properties of FR I and FR II radio sources closely resemble those of Type 1 and Type 2 jets, respectively (Meier *et* al., 1997). We are working on a more detailed analysis to see if the mechanism can be applied to the Ledlow & Owen (1996) relation, for example. Finally, such a mechanism also may provide an alternative model for the broad absorption line (BAL) quasars–another class of radio weak objects with strong winds (Stocke *et al.*, 1992). The small incidence of BALs in the quasar population would not be caused by covering factor but by the low incidence of super-Eddington accretion in quasars. The anomalous radio weakness of the BAL class may be because of the presence of a super-Eddington wind, which would shut off the magnetic switch.

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