1. INTRODUCTION

Over 30 years ago jets entered the lexicon of astrophysical phenomena when it was realized that the double radio sources observed in some galaxies are the product of tightly collimated bidirectional outflows. Since then there has been substantial progress in both jet theory and jet observation. We now know that in addition to active galaxies, jets are found in mass transfer binaries, Galactic microquasars, protostellar systems, and pulsars such as that in the Crab Nebula. Collimated outflows are clearly a natural occurrence, not solely the result of some rare and unusual circumstance in the core of a few active galaxies.

Much of the theoretical work on jets has concentrated on broad phenomena when it was realized that the double radio sources observed in some galaxies are the product of tightly collimated bidirectional outflows. Since then there has been substantial progress in both jet theory and jet observation. We now know that in addition to active galaxies, jets are found in mass transfer binaries, Galactic microquasars, protostellar systems, and pulsars such as that in the Crab Nebula. Collimated outflows are clearly a natural occurrence, not solely the result of some rare and unusual circumstance in the core of a few active galaxies.

The rotation of the central object, however, is also a source of energy, even if the central object is a black hole. The Blandford & Znajek (1977) mechanism, for example, envisions field lines connected directly to a rotating black hole; the rotation drives outward-going Alfvén waves. Punsly & Coroniti (1990) suggested an indirect MHD mechanism for accomplishing much the same end by having the fields anchored in plasma orbiting near the rotating black hole. In either case, some of the black hole's spin energy is extracted and transported outward as a Poynting flux.

Numerical simulations have long played an important role in examining and validating many aspects of jet physics. The work of Shibata & Uchida (1985) provides one of the earliest examples. They simulated an accreting disk of rotating gas transfixed by a vertical magnetic field. In this simulation the radial field formed in the disk by infall is twisted into a strong toroidal field, which provides vertical magnetic pressure forces to drive an outflow that is collimated along the vertical field. Since this pioneering effort, many simulations have been performed featuring improved resolution, additional physical mechanisms, longer time evolution, larger spatial domains, and even full three dimensions. However, most of these jet simulations have, like that of Shibata & Uchida, featured a large-scale poloidal magnetic field as part of the initial conditions.

In some cases a further simplification has been made: the accretion disk itself is collimated along the vertical field. Since this pioneering effort, many simulations have been performed featuring improved resolution, additional physical mechanisms, longer time evolution, larger spatial domains, and even full three dimensions. However, most of these jet simulations have, like that of Shibata & Uchida, featured a large-scale poloidal magnetic field as part of the initial conditions.
by treating the underlying accretion disk as a boundary condition. In such studies the focus is on the ability of this large-scale field to accelerate and collimate gas supplied by the disk boundary (e.g., Meier et al. 1997; Romanova et al. 1997; Ouyed & Pudritz 1997; Krastopolsky et al. 1999). Details differ from simulation to simulation in such things as the assumed magnetic field and the properties of the disk boundary conditions; but, taken as a whole, the results provide strong overall support for the magnetic launch and collimation scenario.

There remain, however, several significant questions that cannot be answered by simulations of this sort. First, do jets always require a net large-scale poloidal field? Under what circumstances do such fields develop? Can they be generated by a dynamo operating within the disk, or must a net field be brought in from a large radius and concentrated near the center? What is the rate at which mass is injected from the disk into the jets? These questions must be addressed with fully global accretion simulations that do not assume the presence of a large-scale field as an initial condition.

In contrast to simulations that have focused exclusively on jets, the simulations developed for this series of papers have been directed toward the dynamical properties of the accretion disk itself. Using a general relativistic MHD (GRMHD) code, we have investigated accretion disk structures that form from an initial condition consisting of an isolated torus of gas fully enclosing a weak initial magnetic field. An advantage of such an initial condition is that it is independent of the boundary conditions. Rather than assuming a preexisting large-scale field configuration, we can study the circumstances under which such large-scale fields might develop naturally in the accretion flow.

The first paper of this series, De Villiers et al. (2003; hereafter Paper I), presented an overview of a series of three-dimensional GRMHD simulations of Keplerian accretion disks orbiting Kerr black holes (the “KD” simulations). The second paper of the series (Hirose et al. 2004; hereafter Paper II) discussed the overall magnetic field configurations in the accretion flow. Among the results noted in Paper I was the appearance of funnel-wall jets in all simulations. Such jets were also observed in a previous simulation using a pseudo-Newtonian potential rather than full GR (Hawley & Balbus 2002; hereafter HB02). Kato et al. (2004) have also carried out similar pseudo-Newtonian simulations and analyzed in greater detail the jets that were produced. More recently, jets were noted in two-dimensional axisymmetric GR simulations by McKinney & Gammie (2004). What is noteworthy in all these simulations is the natural emergence of jets from accretion disks that did not have large-scale fields included as an initial condition. These results have the potential to greatly improve our understanding of the link between accretion and large-scale outflows, since these jets arise self-consistently from the accretion process. The success—or failure—of any self-consistent simulation to produce jets will help to define the range of circumstances that determines jet production in astrophysical systems.

In this paper we focus on the properties of the outflows seen in the KD simulations. Relevant background material from Paper I is briefly outlined in § 2. In § 3 we discuss the internal structure of the outflows and their integrated fluxes of mass, energy, and angular momentum. In § 4 we contrast our results with those of other simulations. Lastly, we provide a summary of our conclusions in § 5.

2. OVERVIEW OF SIMULATIONS

We solve the equations of ideal MHD in the metric of a rotating black hole. The specific form of the equations we solve and the numerical algorithm incorporated into the GRMHD code are described in detail in De Villiers & Hawley (2003a). For reference we reiterate here the key terms and the definitions of the primary code variables.

We work in the Kerr metric, expressed in Boyer-Lindquist coordinates $(t, r, \theta, \phi)$, for which the line element has the form
\[ ds^2 = g_{tt} dt^2 + 2g_{t\phi} dt d\phi + g_{rr} dr^2 + g_{\theta\theta} d\theta^2 + g_{\phi\phi} d\phi^2. \]
We use the metric signature $(-+, +, +, +)$. The determinant of the metric is $g$, and $g^{1/2} = \alpha \sqrt{\gamma}$, where $\alpha$ is the lapse function, $\alpha = 1/(-g^{\mu\nu})^{1/2}$, and $\gamma$ is the determinant of the spatial 3-metric. We follow the usual convention of using Greek characters to denote full spacetime indices and Roman characters for purely spatial indices. We use geometrodynamical units where $G = c = 1$; time and distance are in units of the black hole mass, $M$.

The state of the relativistic test fluid at each point in the spacetime is described by its density, $\rho$; specific internal energy, $\epsilon$; 4-velocity, $U^\mu$; and isotropic pressure, $P$. The relativistic enthalpy is $h = 1 + \epsilon + P/\rho$. The pressure is related to $\rho$ and $\epsilon$ through the equation of state of an ideal gas, $P = \rho \Gamma (\Gamma - 1)$, where $\Gamma$ is the adiabatic exponent. For these simulations we take $\Gamma = 5/3$, unless otherwise indicated. The magnetic field of the fluid is described by two sets of variables: the constrained transport magnetic field, $F_{\mu} = \langle J|B|^2$, and the magnetic field 4-vector, $(4\pi)^{1/2} b^\mu = F^\mu W$. The ideal MHD condition requires $U^\mu F_{\mu\nu} = 0$. The magnetic field $b^\mu$ is included in the definition of the total 4-momentum, $S^\mu = (\rho + |b|^2)|W| U^\mu$, where $W$ is the Lorentz factor. We define auxiliary density and energy functions $D = \rho W$ and $E = Dw$ and transport velocity $V^\mu = U^\mu/|U|$.

In Paper I we presented results of a series of high- and low-resolution simulations, the KD (Keplerian disk) set of disk models. The emphasis in this paper is on the high-resolution models, which use $192 \times 192 \times 64 (r, \theta, \phi)$ grid zones. These are designated KD0, KDI, KDP, and KDE, and they differ in the spin of the black hole around which they orbit, with $a/M = 0$, 0.5, 0.9, and 0.998, respectively. The initial condition consists of an isolated gas torus orbiting near the black hole, with a pressure maximum at $r \approx 25M$ and a slightly sub-Keplerian initial distribution of angular momentum throughout. The initial magnetic field consists of loops of weak poloidal field lying along isodensity surfaces within the torus. We choose parameters for the initial tori that keep the inner edge of the disk and the location of the initial pressure maximum constant as the black hole rotation parameter is varied. Table 1 lists the general properties of the models, where $a/M$ is the spin of the black hole, $\beta$ is the initial (volume-integrated) ratio of gas to magnetic pressure, $r_m$ is the inner edge of the disk (in the equatorial plane), $r_{\text{max}}$ is the location of the pressure maximum (also in the equatorial plane), and $T_{\text{orb}}$ is the orbital period at the pressure maximum in units of $M$. In the following sections $T_{\text{orb}}$ is used as the unit of evolution time of the simulations. For reference we also list $r_{\text{m}}$, the location of the marginally stable orbit, and $T_{\text{orb(3d)}}$, the orbital period of a test particle at $r_{\text{m}}$ in units of $M$. The inner and outer boundaries of the radial grid, $r_{\text{min}}$ and $r_{\text{max}}$, are also given. All radii in the table are in units of $M$. The $\theta$ grid ranges over 0.045$\pi \leq \theta \leq 0.955\pi$, and the $\phi$ grid spans the quarter plane, $0 \leq \phi \leq \pi/2$.

Although we focus our analysis on the KD models, we also draw upon the results of a number of other models to establish the ubiquity of jets under a variety of initial conditions. Most of these other models were run at lower resolution, using $128 \times 128 \times 32$ grid zones. The effects of numerical resolution are gauged in part by a direct comparison between KDP and the lower resolution model KDPI (Paper I). We examine the series of models that began with constant angular momentum tori
around holes with \( a/M = -0.998, 0.0, \) and 0.9, the “SF” models described in De Villiers & Hawley (2003b; hereafter DH03); these initial tori are hotter (thicker) and have an initial pressure maximum closer to the black hole. We also consider the presence (or absence) of jets in several pseudo-Newtonian simulations. Finally we present some results from two new simulations, a high-resolution simulation with \( \Gamma = 4/3 \) and a disk with an initial toroidal field.

2.1. Simulation Diagnostics

Three-dimensional numerical simulations generate an enormous amount of data, only a representative sample of which can be examined. Our analysis is based on a specific set of volume- and shell-averaged history data taken every 1M in time and complete data snapshots taken every 80M in time. In addition, the complete density field is output every 2M for use in making highly time-resolved animations. Although the details of the history calculations are given in Paper I, we provide here a brief summary to clarify the calculations of mass, energy, and angular momentum transported by the unbound outflows. The flux \( \mathcal{F} \) of a given quantity through a shell at radius \( r \) is computed using

\[
\langle \mathcal{F} \rangle(r) = \int \int \mathcal{F} \sqrt{-\mathbf{g}} \, d\theta \, d\phi,
\]

where the bounds of integration range over the full \( \theta \) and \( \phi \) computational domains. We evaluate the rest-mass flux \( \langle \rho U^\mu \rangle \), energy flux \( \langle T^\mu_\nu \rangle \), and angular momentum flux \( \langle T^\mu_\phi \rangle \). Since we are interested in material that leaves the computational domain through the outer boundary, we refer to these fluxes as ejection rates when they are evaluated at the outer boundary, \( r_{\text{out}} \). Gas that leaves the grid at the inner radial boundary is considered to be accreted into the black hole; we refer to this as the black hole accretion rate. We account for changes in the total mass, energy, and angular momentum in terms of the corresponding net fluxes that leave the grid, either into the hole or out of the grid, by integrating over time,

\[
\{ \mathcal{F} \}_\text{out} = \int dt \langle \mathcal{F} \rangle(r_{\text{out}}),
\]

for example.

In this paper we are concerned mainly with the mass and energy fluxes in the jet. The jet is defined as that portion of the flow that is \textit{outbound} and \textit{unbound}. Specifically, it is where the radial momentum is positive and the specific energy at infinity is greater than unity, i.e., \( -hU_t < 1 \). These criteria are met only in the axial funnel and the funnel-wall regions of the simulation volume. There is considerable internal substructure within this unbound outflow. An overview of the unbound outflow is presented in Figure 1, which shows an azimuthally averaged composite taken from late-time data for model KDP. The color contours depict gas density on a logarithmic scale. The boundary of the jet is indicated by a thick white contour; the contour nearest the corona marks the boundary between unbound funnel gas and bound coronal gas, while the contour nearest the funnel indicates where the unbound positive radial mass flux has dropped below the limit used to define the funnel-wall jet. The boundary of the jet is indicated by a thick black contour; the contour nearest the hole marks the boundary between bound funnel gas and bound coronal gas, while the contour nearest the funnel indicates where the unbound positive radial mass flux has dropped below the limit used to define the funnel-wall jet. The boundary of the jet is indicated by a thick black contour; the contour nearest the hole marks the boundary between bound funnel gas and bound coronal gas, while the contour nearest the funnel indicates where the unbound positive radial mass flux has dropped below the limit used to define the funnel-wall jet.

3. UNBOUND OUTFLOWS IN KD MODELS

3.1. Overview of Jet Structure

The jet is defined as that portion of the flow that is \textit{outbound} and \textit{unbound}. Specifically, it is where the radial momentum is positive and the specific energy at infinity is greater than unity, i.e., \( -hU_t > 1 \). These criteria are met only in the axial funnel and the funnel-wall regions of the simulation volume. There is considerable internal substructure within this unbound outflow. An overview of the unbound outflow is presented in Figure 1, which shows an azimuthally averaged composite taken from late-time data for model KDP. The color contours depict gas density on a logarithmic scale. The boundary of the jet is indicated by a thick white contour; the contour nearest the corona marks the boundary between unbound funnel gas and bound coronal gas, while the contour nearest the funnel indicates where the unbound positive radial mass flux has dropped below the limit used to define the funnel-wall jet. The boundary of the jet is indicated by a thick black contour; the contour nearest the hole marks the boundary between bound funnel gas and bound coronal gas, while the contour nearest the funnel indicates where the unbound positive radial mass flux has dropped below the limit used to define the funnel-wall jet.
origin of the continuous region of unbound, outgoing radial mass flux. The base of the jet is located where the inner torus, coronal envelope, and axial funnel intersect. In the high-spin model KDP the base is relatively stable and is marked by a ring of denser gas. In low-spin models the base is not so sharply defined, instead showing both spatial and temporal variability.

Although the continuous outflow begins at this point, mass is loaded into the jet along an extended region that we call the injection region. This begins at a radius close to that of the marginally stable orbit, \( r_{\text{ms}} \), hence moving inward with increasing black hole spin. Figure 2 shows the time-integrated mass flux in the funnel-wall jet as a function of radius, \( M_{\text{jet}}(r) = (\rho U^r)_{\text{jet}} \), for all four KD models. The curves show two features: a region of relatively steep slope near the black hole indicating rapid matter injection near the base, and a more extended region of shallow slope at larger radii indicating that mass entrainment takes place along the jet-corona boundary. In each case, 20%–30% of the ultimate mass flux is injected within a radius \( \sim 3r_{\text{ms}} \), with the remainder of the mass accumulated more gradually at a considerably larger radius. Because \( r_{\text{ms}} \) decreases with increasing spin, the mass injection region shrinks in absolute terms when the black hole spins more rapidly. One way the mass-loading of the jet can be enhanced by effects due to black hole spin is by the increased density and pressure found in the accretion disk inner torus. As discussed in Paper I, this inner torus moves closer to the black hole and becomes denser and thicker with increasing black hole spin.

Beyond the injection region lies the body of the jet, where there is significant unbound outward mass flux. The unbound mass flux is not a transient feature in these models; it is established early in the simulation along with the quasi-steady accretion flow into the black hole and endures for the length of the simulation. Animations of unbound density reveal that the funnel-wall jet is highly dynamical, and its shape fluctuates on a timescale comparable to the orbital period of the main disk body. This variability is due in part to a changing dynamical balance between forces in the magnetically dominated funnel and the coronal envelope, as well as variations in the rate at which matter is injected into the jet.

While the bulk of the unbound mass flux is confined to the funnel wall, there is also a tenuous, fast, unbound funnel outflow. A defining feature of the funnel is that hydrostatic equilibrium is not possible there; material with significant specific angular momentum cannot get into the funnel, and any low-l material must either accrete into the hole or leave as an outflow. As discussed in Paper II, in these simulations the axial funnel contains a large-scale organized poloidal magnetic field similar to a split monopole. This field is created during the initial accretion, as long radial lengths of field are drawn out in the plunging inflow. Some of these field lines contained within the dense accretion flow near the equator slip outward into the region of very low mass density as the mass attached to them drains into the black hole. As they do so, they expand outward, creating the large-scale poloidal field that occupies the axial funnel. If the black hole is rotating, the funnel field is wound up and the toroidal field is amplified by frame dragging. Once the field in the funnel becomes comparable in magnitude to the adjacent field in the corona, its strength saturates and little further flux is drawn into the funnel.

Finally, asymptotic properties of the jet are determined by its properties in the outflow region, which is the distant region well away from the disk and the black hole. In these simulations the outer grid boundary is located at \( r_{\text{max}} = 120M \). Hence we compute various outflow quantities at \( r = 100M \) for comparison among the different simulations.

Figure 1 also includes a vector field showing the direction, but not the magnitude, of net pressure gradient and Lorentz forces (eq. [5]). In the inner torus and adjacent corona, the orientation of the net poloidal force is largely perpendicular to isodensity contours, consistent with approximate hydrostatic equilibrium. The magnitudes of the gas pressure gradient and Lorentz force are comparable, but since the inner torus and the thin corona surrounding it are highly dynamical, one force may dominate...
the other at different times. In the funnel outflow and jet regions, the net poloidal force is radially outward, and the Lorentz force is dominant. The toroidal component of the Lorentz force (not shown) is prominent through the base of the jet, the injection region, and the jet body, while the toroidal component of the gas pressure gradient is negligible. The toroidal component of the Lorentz force corresponds to the familiar $B_t \nabla \times B_t$, Maxwell stress in nonrelativistic MHD. The strength of the toroidal Lorentz force is correlated with the spin of the black hole; it is weakest in model KD0 and strongest in model KDE. The picture that emerges is that the funnel-wall jet is launched when pressure gradients and Lorentz forces push gas up through the inner torus into the injection region, where the magnetocentrifugal acceleration occurs. In the funnel and along the funnel wall, a strong radially directed Lorentz force accelerates the unbound material outward.

3.2. Cross Section through the Jets

Figure 3 presents a cut through a representative region of the jet body for all four KD models. It depicts four sets of variables as functions of the polar angle, $\theta$, from the north polar axis to the equator; the variables are azimuthally averaged and taken at a radial distance of $10r_{\text{ms}}$.

The top graph in each panel shows mass flux, $\rho U^r$, which defines the jet boundaries. The second row of graphs shows density (solid lines) as well as gas (dashed lines) and magnetic pressure (dashed lines). The third row of graphs shows the fluid (solid lines) and electromagnetic (dotted lines) components of outward energy flux (gaps in the curves correspond to radially inward flux). The bottom graphs show the absolute value of the magnetic stress.

![Figure 3](image-url)
since the low-density gas has been heated by shocks driven into the funnel from denser regions of the flow. The \( \theta \) profile of the magnetic pressure is much more shallow than that of the gas pressure in all the models. In the \( a/M = 0 \) model (KD0), magnetic pressure is roughly constant through the funnel, dips slightly through the jet and corona, then rises slightly through the body of the disk. In models KDI, KDP, and KDE, magnetic pressure is roughly constant in the funnel and jet, dipping slightly in the corona near the jet boundary before rising again as one moves toward the equator. As pointed out in Paper I (see Fig. 8 in that paper), the total pressure is relatively smooth through the corona into the funnel. The funnel is magnetically dominated, the main disk is gas-pressure-dominated, and the corona has a ratio of pressures near unity. The jet is found just inside the region where \( \beta = P_{\text{gas}}/P_{\text{mag}} \) drops below unity.

The third graph from the top shows the fluid and electromagnetic contributions to outward energy flux (\( T_{\text{r}}^{(\text{FL})} \) and \( T_{\text{r}}^{(\text{EM})} \)) normalized to the maximum of the absolute total energy flux, which is found near the equator in all cases. The fluid component varies in sign in the main disk and corona, as would be expected for turbulent motions. Its absolute value shows a steady drop from the main disk body through the corona. In the unbound portions of the outflow, energy flux is everywhere outward. Despite the fact that the mass flux is largely confined to the funnel wall, the fluid energy flux varies much more slowly with polar angle within the funnel; this is because the low mass density is compensated by a high specific enthalpy and higher outflow velocity. The electromagnetic component of the energy flux is extremely weak when \( a/M = 0 \) (KD0) and becomes more prominent in the funnel with increasing black hole spin, but it never dominates the fluid component.

In the lower graph we plot the absolute value of the magnetic stress, \( \| b' b_0 \| \), again normalized to its maximum value on the slice of constant radius. This quantity exhibits a strong dependence on black hole spin in the jet and the funnel. In model KD0, the stress drops sharply in crossing from the main disk body to the corona. Although stress is variable through the corona, the peaks are roughly level. Peak stress in the jet is slightly weaker than in the corona and is extremely weak in the funnel. The models with a rotating hole also show a drop in stress in the corona, relative to the main disk. However, stress in the jet and funnel becomes progressively more dominant with increasing spin, with the maximum stress found inside the funnel in models KDP and KDE. In these two models the stress is strongest where the density is low.

The fluid velocity (not shown) also makes a sharp transition through the jet, from predominantly orbital motion in the corona to radial or helical motion in the funnel. As noted in Paper II the magnetic field lines also undergo a corresponding transition through the funnel wall. The velocity is predominantly radial on the funnel side and toroidal on the corona side of the funnel-wall jet. The radial velocity increases sharply through the funnel wall. The mass-weighted mean of \( V' \) within the funnel-wall jet ranges from \( \approx 0.3c \) for model KD0 to \( 0.5c \) for model KDE. The values are comparable to the toroidal velocities found in the inner torus, which range from \( \approx 0.4c \) in KD0 to \( 0.6c \) in KDE. The outward speed at a fixed angle is established quite close in; the flow accelerates from near its base to 10–20M (the end of the acceleration zone moves outward with decreasing black hole spin) and then retains that speed all the way to the outer boundary at 120M.

We can measure the collimation of the jet by tracing the location of the boundary between the bound and unbound outflow as a function of \( r \). Generally speaking, the jets experience some collimation out to a radius between 20 and 40M and then maintain a more or less constant opening angle of \( \approx 0.6 \) rad. This suggests that the observed collimation is mainly due to the pressure in the surrounding corona near the hole.

3.3. Integrated Outflow Quantities

In Paper I we noted an increase in the prominence of the jet’s radial mass flux, in relation to typical values in the main disk body, with increased black hole spin (e.g., Fig. 10 of Paper I). To put this on a more quantitative footing, we examine the ejection rates of mass, energy, and angular momentum for unbound and bound material in the outflow region. Figure 4 shows a time history of these rates in the \( a/M = 0.9 \) model (KDP) over the full 8000M of simulation time; the time dependence of these quantities in the other KD models is similar. In the energy and angular momentum flux plot, the solid line is the fluid component of the flux, and the dashed line shows the contribution of the electromagnetic portion of the stress tensor.

There are two initial peaks in the jet mass flux around times 1400M and 2100M. The first peak is almost entirely made up of unbound material blown off the initial torus as the evolution begins. The second, larger flux peak is backlash from the initial plunging accretion into the black hole. By \( t \approx 3000M \) the accretion disk becomes more fully established, and the jet emerges as a stable structure with continuous ejection of unbound material at a mean rate of \( \lesssim 10^{-6} \) of the initial mass, energy, and angular momentum of the torus per unit time \( M \). The jet mass flux fluctuates about its mean on a timescale \( \sim 1000M \), comparable to the orbital period of the initial pressure maximum in the disk. We expect that the jet should persist until the reservoir of mass in the main disk is depleted, a process that would take about 10 times longer than the length of the simulation.

A quantitative depiction of the mass, energy, and angular momentum content of the jets is given in Table 2. For each quantity \( (Q) \), five columns are shown: the initial value, \( Q_0 \); the accretion rate of bound material through the inner boundary, \( \Delta Q_b \); the bound (coronal) ejection rate, \( \Delta Q_b \); the unbound (jet and funnel outflow) ejection rate, \( \Delta Q_a \); and the efficiency, \( \eta = \Delta Q_a/\Delta Q_b \), expressed as a percentage. These values are totals, expressed in code values, computed from time integration over the whole
dependent increase for mass, energy, and angular momentum, quantity accreted into the black hole. This ratio shows a spin-
simulation using equation (2). For comparison purposes we also

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|c|c|c|c|c|c|}
\hline
Model   & $M_0$ & $\Delta M_i$ & $\Delta E_i$ & $\Delta E_b$ & $\Delta \dot{M}$ & $\dot{E}_b$ & $f_\ell$ & $L_0$ & $\Delta L_i$ & $\Delta L_b$ & $\Delta L_u$ & $f_l$ \\
\hline
KDO...... & 156  & 24.9 & 1.14 & 0.17 & 0.007 & 154 & 22.7 & 1.14 & 0.22 & 0.01 & 902 & 77.1 & 6.85 & 0.44 & 0.006 \\
KDI...... & 258  & 36.1 & 2.70 & 0.60 & 0.017 & 255 & 32.8 & 2.69 & 0.94 & 0.029 & 1489 & 95.0 & 17.2 & 3.02 & 0.032 \\
KDP...... & 291  & 17.9 & 1.86 & 0.87 & 0.048 & 286 & 15.0 & 1.85 & 1.63 & 0.11 & 1652 & 33.5 & 9.12 & 4.10 & 0.12 \\
KDE...... & 392  & 14.4 & 8.35 & 4.78 & 0.33 & 386 & 9.34 & 8.33 & 9.67 & 1.04 & 2255 & 14.1 & 55.2 & 23.0 & 1.63 \\
KDPi...... & 151  & 5.73 & 0.48 & 0.32 & 0.016 & 148 & 4.79 & 0.47 & 0.54 & 0.11 & 835 & 10.8 & 2.34 & 1.52 & 0.14 \\
KDEi...... & 258  & 36.1 & 2.70 & 0.60 & 0.017 & 255 & 32.8 & 2.69 & 1.52 & 0.18 & 1652 & 19.2 & 4.21 & 3.68 & 0.19 \\
\hline
SFP....... & 2308  & 427 & 72.3 & 30.2 & 0.071 & 2263 & 369 & 72.1 & 38.4 & 0.10 & 9729 & 759 & 289 & 116 & 0.15 \\
SF0....... & 1374  & 469 & 24.6 & 6.78 & 0.014 & 1346 & 441 & 24.5 & 6.88 & 0.016 & 6058 & 1455 & 113 & 24.7 & 0.017 \\
SFR....... & 741  & 285 & 8.76 & 1.36 & 0.005 & 726 & 272 & 8.73 & 1.67 & 0.006 & 3484 & 1102 & 46.0 & 0.019 & 0.002 \\
KDPlr...... & 291  & 10.2 & 0.88 & 0.87 & 0.085 & 286 & 8.33 & 0.88 & 1.52 & 0.18 & 1652 & 19.2 & 4.21 & 3.68 & 0.19 \\
KDPi...... & 151  & 5.73 & 0.48 & 0.32 & 0.056 & 148 & 4.79 & 0.47 & 0.54 & 0.11 & 835 & 10.8 & 2.34 & 1.52 & 0.14 \\
KDEi...... & 258  & 36.1 & 2.70 & 0.60 & 0.017 & 255 & 32.8 & 2.69 & 0.94 & 0.18 & 1652 & 19.2 & 4.21 & 3.68 & 0.19 \\
\hline
\end{tabular}
\caption{Jet Luminosity}
\end{table}

The jet energy flux has contributions from both the fluid
and the magnetic field terms in $T_{\tau\tau}$. Figure 4 shows that the fluid
term dominates, in part because of the importance of the rest-
mass contribution to the energy. The figure also shows that, for
the $a/M = 0.9$ model, matter and electromagnetic fields con-
tribute roughly comparable amounts to the angular momentum
flux in the jet. This is true for all the rotating hole models. An examination of the time average of the jet angular momentum
flux as a function of radius shows that in the three nonzero spin
models the electromagnetic angular momentum jet flux begins
(and peaks) near $r = 2M$. For the Schwarzschild hole model
(KDO), the angular momentum flux is reduced and does not
originate near the black hole; instead it seems to arise from
entrainment beyond $r = 10M$.

These data support the hypothesis that when the black hole
spins, the dominant source of angular momentum (and energy)
carried out in the outflow is the rotation of the black hole itself.
The near-constancy of the jet’s angular momentum flux from
very small radii outward in the Kerr hole models is telling;
there is little else inside $r = 2M$ besides the spin of the hole
itself that could supply such a large quantity of angular moment-

um. This conclusion is further supported by the fact, already
emphasized in Paper I, that in model KDE, the net angular momentum accreted by the black hole is so small that, if these
conditions persisted, $a/M$, the ratio of angular momentum to
mass in the black hole, would have to decrease.

To derive an effective jet efficiency that can be compared more
easily to the usual estimates of accretion efficiency in disks, we
compute a jet energy flux by summing the fluid and electromagnetic unbound energy fluxes and subtracting the contribution due
to the rest mass in the jet. We use a time average from $t = 3200M$
to the end of the simulation, thus omitting the outflow bursts that
result from the initial infall of material toward the hole, and cal-
culate the outflow values at $r = 100M$. An estimate of the lu-
minosity in the jet can then be obtained by taking the ratio of the
unbound energy flux $E_{\text{jet}}$ to the accretion rate $M$. These values are
shown in the column labeled $\eta_{\text{jet}}$ in Table 3. Next, these ratios can
be compared with the usual estimates for accretion efficiency
derived from the binding energy of the marginally stable orbit,
which are, for $a/M = 0, 0.5, 0.9$, and 0.998, $\eta_{\text{ms}} = 0.057, 0.082$, 0.16, and 0.32, respectively. This is listed in the column labeled
$\eta_{\text{jet}}/\eta_{\text{ms}}$. What these numbers reveal is that the energy expelled in
the jet rises from a modest supplement to the expected disk lu-
minosity in the zero-spin model to an amount comparable to the
disk luminosity in the highest spin example.

In all models, if the rest-mass contribution is included, the
energy carried with the fluid is the largest contribution to the jet.
However, the relative importance of the electromagnetic part
rises sharply with increasing black hole spin. The ratio of the
jet Poynting flux to the unbound energy flux in the fluid jet (after
subtracting rest mass) is given in the last column of Table 3. As
these numbers show, the Poynting flux becomes fully half the
net fluid energy flux at $a/M = 0.9$ and 90% at $a/M = 0.998$.

Returning to the mass flux in the jet, we ask whether the mass
ejected in the unbound jet is significant compared to the rate at
which mass is fed into the near-hole region. The initial torus is,
conserve, the mass source for both the accretion flow and the
jets. The transport of angular momentum allows accretion to
occur, forming an accretion disk that extends toward the black
hole, where an inner torus forms near $1.5r_{\text{ms}}$. To quantify
the mass flow through the inner region, three quantities are of in-
terest: the bound flux through the radial shell at $r = 15M$ (the
inner edge of the initial torus), the amount of bound mass ac-
creted into the black hole, and the amount of unbound mass
mattered off by the jets. What remains after accounting for these
fluxes should be equal to the mass contained in the region inside

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|}
\hline
Model   & $a/M$ & $\eta_{\text{jet}}$ & $\eta_{\text{jet}}/\eta_{\text{ms}}$ & $E_{\text{(EM)}}/\left(E_{\text{(FL)}} - M\right)$ \\
\hline
KDO...... & 0.0  & 0.002 & 0.03 & 0.06 \\
KDI...... & 0.5  & 0.013 & 0.16 & 0.34 \\
KDP...... & 0.9  & 0.029 & 0.18 & 0.47 \\
KDE...... & 0.998 & 0.18  & 0.56 & 0.897 \\
\hline
\end{tabular}
\caption{Jet Luminosity}
\end{table}
The bound mass accreted into the black hole \( \Delta M_{\text{hole}} \) and the unbound mass carried off by the jets \( \Delta M_{\text{jet}} \) are taken from Table 2. The flux of bound matter in the main disk and corona through the shell at \( r = 15M \) \( \Delta M_{r=15} \) is computed by time-integrating the mass flux \( \rho U^r \) over the shell at \( r = 15M \) subject to the condition \(-hU^r < 1\). The final mass is computed at the end of the simulation by direct integration of the gas density inside \( r = 15M \). Note that the accretion flow at \( r = 15M \) is mainly established by the MHD turbulence in the main body of the disk, while the accretion rate into the hole and, to some degree, the outflow rate in the jets are determined by processes near the black hole, and those rates need not match the accretion rate supplied by the main disk. One consequence is the spin-dependent buildup of the inner torus, as noted in Paper I; this is again evident in the values presented in Table 4. Although \( M_{\text{final}} \) should equal the amount that passed through \( r = 15M \) minus the losses to the hole and the jet, the agreement in these numbers is not perfect. Two sources of error in this comparison are the entrainment of mass into the jet outside of \( r = 15M \) in all models and the need to evaluate the mass fluxes using the data dumps that are spaced at intervals of 80 Myr. It is clear, however, that unbound mass flux in the jets becomes a significant component of the mass flow through the inner region for high-spin holes.

4. Jets in Other Simulations

The KD simulations are, of course, computed from a specific initial condition; only the spin of the hole is varied in a significant way. Given that, it is important to consider how general the results are and how much they depend on the specific setup used in this study. In this section, we summarize the properties of jets that have been seen in other accretion simulations, including those we have run using the GRMHD code and those run using other codes. These simulations complement the analysis of the KD models by probing the role of initial conditions, numerical resolution, and relativistic effects.

4.1. Jets in the SF Simulations

The SF models (DH03) consist of initial constant-\( \ell \) tori containing weak poloidal field loops. Three black hole spin values were considered: zero spin (SF0), prograde \( a/M = 0.9 \) (SFP), and retrograde \( a/M = -0.998 \) (SFR). The main differences between the SF and KD models are that the initial torus is hotter and the initial pressure maximum is closer to the black hole in the SF models. In addition, the SFR model represents an example of a counterrotating black hole. We find that jets are produced in all three SF models, including retrograde model SFR. This is noteworthy, since the marginally stable orbit for this model, at \( r = 9M \), is quite distant from the black hole, and one might expect that the thin, radially extended plunging inflow would not be conducive to jet formation. The base of the jet in the SFR model occurs just outside the static limit, i.e., well inside the marginally stable orbit, and the configuration of the base and injection region are similar to that of the SFP model. However, the SFR funnel-wall jet is considerably weaker than its prograde counterparts, and the funnel wall is more flared out; consequently, the SFR jet has a relatively large opening angle near its base.

It was noted in DH03 that these disk models generate rapid accretion, a fact that is reflected in the elevated rates shown in Table 2; e.g., 34% of the initial torus mass is accreted for model SF0, whereas only 16% is accreted for model KD0. The greater accretion rate in the SF models is most likely connected to their greater heat content and therefore greater thickness. To the extent that magnetic stresses grow as longer wavelength magnetorotational modes become unstable, greater disk thickness promotes faster accretion. The fractional unbound mass ejection rates in the SF models are also elevated. This could arise from at least two possible causes: the greater rate at which mass is fed into the inner regions of the accretion flow, making more mass available to be expelled in the jet, and the greater amount of energy released by accretion, some of which can be utilized for driving outflows. One can regard energy drawn from the rotation of the black hole as catalyzed by accretion, so that the rate of rotational energy extraction is also proportional to the accretion rate.

4.2. \( \Gamma = 4/3 \) Equation of State

Since gas pressure acceleration plays a role in the acceleration of the jet, it is of interest to see how the jet might vary with the adiabatic index, \( \Gamma \). Model KDPi is similar to KDP except that it employs the softer \( \Gamma = 4/3 \) equation of state. From Table 2 we see that this model yields a slightly weaker jet than the corresponding \( \Gamma = 5/3 \) simulation; only 0.21% of the initial torus mass is ejected in the jet, compared to 0.30% for KDP. A similar reduction occurs for energy and angular momentum. However, the accretion rates are similarly affected, resulting in comparable efficiencies for models KDP and KDPi. Any dependence of the jet on the adiabatic index in this simulation seems to be modest. It is possible that more general (and realistic) equations of state that include, for example, resistive dissipation of magnetic field energy and heat losses by generation of photons could lead to greater changes.

4.3. Initial Toroidal Field

In this simulation, the Schwarzschild (KD0) torus is initialized with a moderate \( (\beta = 10) \) toroidal field. The grid resolution is \( 128 \times 128 \times 64 \). This model eventually generates an accretion flow produced by MRI-driven turbulence. What is notable for the present purposes is that this model generates essentially no jet. Several factors may contribute to suppressing the jet, but there are two in particular that stand out. First, the accretion rate is considerably smaller in this simulation, and the flow does not form a substantial inner torus, so there is no large pressure gradient to push material against the funnel wall. Second, and perhaps more important, there is no significant poloidal magnetic field in the funnel. In the KD models, an extended poloidal field is generated by the plunging accretion flow, and this field is ejected into the funnel when that flow reaches the black hole. This becomes the radial funnel field, whose polarity is established by the initial conditions. In the toroidal field model, there is no initial north-south field mirror symmetry and no large-scale, systematic poloidal field lines generated by the inflow. Thus, the toroidal field model provides an example of jet formation switched off by the absence of a high-pressure inner torus and a significant poloidal field in the funnel.

| Model | \( \Delta M_{\text{hole}} \) | \( \Delta M_{\text{jet}} \) | \( \Delta M_{r=15} \) | \( M_{\text{final}} \) |
|-------|-----------------|-----------------|-----------------|-----------------|
| KD0   | 0.161           | 0.00147         | 0.184           | 0.0194          |
| KDI   | 0.141           | 0.00275         | 0.171           | 0.0259          |
| KDP   | 0.062           | 0.00361         | 0.101           | 0.0281          |
| KDE   | 0.038           | 0.01400         | 0.115           | 0.0545          |
4.4. Numerical Resolution

The values shown in Table 2 for low-resolution model KDPr are comparable to those of model KDP. The amount of material ejected through the jet is very similar for mass, energy, and angular momentum. However, the corresponding accreted quantities are systematically lower, yielding slightly higher efficiencies for model KDPr. As noted in Paper I, the increased numerical resolution in the high-resolution KD models better captures the turbulent and highly dynamical accretion flow. However, the constancy of the ejected rates between the two models suggests that the present numerical resolution is adequate to capture the qualitative behavior of the jets.

4.5. Axisymmetric Relativistic Simulations

To the extent that our results reflect physical processes likely to occur in accretion systems, they should be reproducible and seen in other global models. Although global simulations of this type are still uncommon, jets similar to what we report here have been observed by McKinney & Gammie (2004). They have carried out an axisymmetric GRMHD simulation of a torus around an $a/M = 0.938$ black hole, along with a series of variations on this fiducial model, including both $\Gamma = 5/3$ and 4/3, different spin values $a/M$, different initial magnetic field strengths and topologies, and different grid resolutions. They find the same general late-time flow structures as those characterized in Paper I and described here, including a conical mass outflow along the funnel edge, along with a significant electromagnetic flux in the funnel. The unbound material in the outflow begins just inside the equipartition contour $\beta = 1$. The funnel is evacuated and contains a large-scale radial field. In an $a/M = 0.9$ simulation they find an outward Poynting flux in the funnel that has a time-averaged value of $E^{(E)} / M = 0.028$; for an $a/M = 0.5$ model, this ratio is $10^{-3}$. In our models the corresponding ratios of Poynting flux in the jet to rest-mass accretion rate are 0.013 and 0.0044 for those $a/M$ values, respectively (Table 3). McKinney & Gammie (2004) also observe a strong increase in funnel Poynting flux with increasing black hole spin. Because their simulations make use of a quite distinct numerical algorithm and greater (albeit two-dimensional) grid resolution, the overall agreement of the results is gratifying.

Another highly suggestive result of McKinney & Gammie (2004) is their observation that the strongest electromagnetic jets are found in models that have an initial net vertical field. This reinforces the conclusion that much of the Poynting flux portion of the jet depends on the ability of the accretion flow to establish a net poloidal field in the funnel. It may also be related to the generic conclusion of shearing-box studies that a net vertical field enhances the saturation magnetic field strength resulting from the MRI (e.g., Sano et al. 2004). This line of thought is further supported by the work of Koide et al. (2000), who performed two axisymmetric GRMHD simulations that began with a preexisting vertical magnetic field for both a prograde and a retrograde orbiting equatorial matter disk around an $a/M = 0.95$ Kerr hole. Although the Koide et al. simulations could be run for only very short amounts of time, they found that jets with $r_{\text{jet}} \sim 0.4c$ are produced in both cases. In the counterrotating disk the dominant component of the jet is the electromagnetic flux, powered primarily by the spin of the hole rather than by the disk. They also reported that both magnetic pressure and Maxwell stress forces are of comparable importance in accelerating the jet.

4.6. Jets in Pseudo-Newtonian Simulations

There have been several investigations of black hole accretion in the pseudo-Newtonian approximation (Hawley & Krolik 2001, 2002; HB02; Kato et al. 2004). The initially poloidal simulation of Hawley & Krolik (2002) is very close in character to the KD0 model, because it began from a hydrostatic initial condition with pressure maximum at $r = 20M$ and a purely poloidal magnetic field, and the pseudo-Newtonian potential was designed to mimic the principal properties of the Schwarzschild metric. HB02 was similar in design but differed in that the initial condition was a constant angular momentum torus centered at $r = 200M$. The work of Kato et al. (2004) was similar to that of HB02, but for two different initial field strengths and nonzero initial external gas pressure surrounding the accretion disk. Like HB02, the accreting matter was centered at a relatively large distance, in this case 80M.

In Hawley & Krolik (2002), although a high-temperature and (sometimes) magnetically dominated region is formed within a conical region near the axis, no true outflow is created; all the matter remains bound. It is possible that the failure to form an outflow in this simulation was due to its cylindrical grid. This choice of symmetry necessitated a cutout around the axis. The magnetic field could be absorbed by this cutout, preventing the buildup of a poloidal field that takes place in our GR simulations. In addition, we find here that jet initiation takes place at high latitudes very close to the event horizon; yet in the Hawley & Krolik simulations, the axial cutout eliminated a sizable range of polar angles in the region immediately outside the event horizon.

By contrast, in HB02 a structure similar to that of KD0 was seen: a narrow outflow of unbound gas moving outward along the boundary between the corona and the funnel. At a fixed height above the equatorial plane, the density is roughly constant in the corona but then drops rapidly toward the axis. Gas pressure also drops through the jet body but levels off in the funnel. Magnetic pressure dominates in the corona, but, unlike in the KD models, the magnetic field in the funnel is not strong and $\beta$ remains greater than 1. The field strength in the funnel has likely been reduced by the axial boundary condition, which again consists of a cut-out cylinder surrounding the axis with outflow boundary conditions. The axial boundary condition seems to have had little effect on the overall flow, which resembles what is seen in KD0, but did prevent the buildup of the field within the funnel. The primary jet acceleration mechanism in the HB02 model is the high gas pressure found near the centrifugal barrier. Jet collimation is due to the magnetized corona surrounding the jet. Kato et al. (2004) found a magnetic field configuration similar to the one reported in Paper II for the Schwarzschild hole simulation: a radial field in the funnel surrounded by a predominantly toroidal field. The jet that develops is driven by strong toroidal field pressures and propagates in a $\beta \sim 1$ environment. The jet is collimated in large part by pressure from the surrounding corona.

In sum, we find that the crucial distinction is not between Newtonian and relativistic treatments but between the Schwarzschild and Kerr metrics. Pseudo-Newtonian dynamics can, in many respects, be a good approximation to dynamics in a Schwarzschild metric, and in both sorts of simulations we find similar-appearing weak outflows. On the other hand, the effects of black hole rotation cannot be reproduced in this way. The dramatic strengthening of the outflow with increasing $a/M$ is an intrinsically relativistic effect.

5. CONCLUSION

In Paper I of this series we described the overall structure seen in a series of simulations of black hole accretion flows around Kerr holes of different spin. Unbound outflows, or jets, are a natural outcome of accretion in our self-consistent MHD disk simulations. In this paper we have made a more detailed examination of these jets.
The jets that emerge in our simulations originate in the region near the black hole, where the inner accretion torus, the corona, and the centrifugal funnel wall meet. There are two components to the unbound outflows: a hot, fast, tenuous outflow inside the axial funnel and a colder, slower, massive jet confined to the funnel wall by a centrifugal barrier. Matter is largely excluded from the funnel by a centrifugal barrier; gas pressure gradients and Lorentz forces combine to push matter from the coronal and inner torus regions of the accretion flow against the “wall” formed by this centrifugal barrier. This process can occur even in purely hydrodynamic accretion flows, if those flows are sufficiently hot (Hawley 1986; HB02). But when there is a magnetic field present, and especially when the black hole rotates rapidly, additional forces come into play. Field lines tied to matter in the ergosphere (see, e.g., Fig. 7 in Paper II) exert strong electromagnetic forces that accelerate matter outward along the edge of the cone. This picture resembles that envisioned by Blandford & Payne (1982) for a disk wind, but it is restricted to a relatively narrow hollow cone because it is the only location in which there are rapidly rotating poloidal field lines.

Electromagnetic fields, which fill the funnel interior, contribute significantly to both the energy and the angular momentum outflow, and their relative importance grows with greater black hole spin. The key ingredient is the presence of a radial field within the funnel. Although our simulations begin without such a field, a magnetic field is injected into the funnel as the accretion flow first reaches the black hole early in the simulation. The resulting field builds until it is in rough pressure balance with the surrounding corona, and it stays approximately the same from then on. The poloidal field lines extend into the ergosphere, and this permits frame dragging to tap into the rotational energy of the black hole, increasing the flux of energy in the funnel outflow and the funnel-wall jet in the models with a spinning black hole. Not surprisingly, given the way magnetic fields strengthen with increasing black hole spin (Paper II), the “efficiency” of the jet relative to the mass accretion rate is a rapidly rising function of $a/M$.

Although this mechanism is reminiscent of the scenario proposed by Blandford & Znajek (1977) in that energy is extracted by the magnetic field from a rotating hole, the role of matter is much more important here than in the original Blandford-Znajek picture. Much of the energy flux in the funnel outflow and funnel-wall jet comes from the fluid enthalpy, not just the electromagnetic component. In addition, in the classical Blandford-Znajek picture, the density was so low everywhere that the field was taken to be “force-free”; here, by contrast, most of the mass and energy flow occurs in the funnel-wall jet, where the ratio of magnetic field energy density to matter energy density is generally less than unity. In this regard the picture is similar to the Punsly & Coroniti (1990) ergosphere-driven wind model.

The specific conclusions drawn from our simulations can be more broadly interpreted by comparing and contrasting with the jets observed in other simulations (for both rotating and non-rotating black holes), with general relativity, and with only a pseudo-Newtonian approximation.

Jets are generated by both prograde and retrograde accretion flows. Two retrograde simulations, the SFR model (DH03) and one by Koide et al. (2000), provide an interesting contrast. The SFR jet was considerably weaker than that seen in the KD models, whereas the Koide et al. jet was comparable in strength to their prograde case. Koide et al. (2000) argue that their retrograde model provides a clear demonstration of jet power derived from the black hole spin energy. The contrast with SFR likely results from their use of a preexisting magnetic field in the simulation, whereas in SFR the funnel field must grow from a self-consistent MRI-driven dynamo within the disk.

Our simulations have shown that preexisting large-scale poloidal fields are not essential for jet production. Nevertheless, poloidal magnetic fields are key for both jet acceleration and jet collimation. We find that there is no jet production when no significant poloidal field is injected into the funnel, as was the case when an initially toroidal magnetic field configuration was used. Although poloidal magnetic fields appear to be essential, gas pressure in the disk also plays a role in jet formation. For practical reasons, all global simulations that include the accretion disk within the computation domain must use relatively thick, hot disks. In such disks gas pressure alone can accelerate gas away from the hot inner torus, but any angular momentum remaining within that gas will keep it out of the funnel. The SF models provide examples of jets formed from hotter, thicker inner tori. The elevated accretion rates seen in the SF models are also accompanied by higher ejection rates. We see only a weak dependence on the equation of state, but the role of the vertical thickness and temperature of the inner torus in jet formation is a topic requiring further study.

In the KD models, jet collimation occurs near the hole because of the confining pressure of the surrounding corona, and magnetic fields play an important role in establishing and maintaining that corona. Similar collimation was seen in the pseudo-Newtonian HB02 simulation; even though the jet acceleration is pressure-driven and funnel fields are not important, the confining corona remains magnetically dominated.

Finally, we must ask about the quantitative reliability of these simulations when applied to jets in nature. Despite the limited radial domain studied within our simulation volume, we are beginning to connect to astrophysical scales that are now solvable with VLBI. For example, the width of the jet in M87 (Junor et al. 1999; Ly et al. 2004) at a distance of $\sim 100$ gravitational radii from the central engine is comparable to that produced by the mild collimation seen in the jets in our simulations.

That outflows readily form seems to be a robust conclusion. It follows merely from the fact that the region around the axis must always have very low density if most matter has significant orbital angular momentum, while the magnetic field readily penetrates the evacuated cone if it has any poloidal component. However, we have not fully demonstrated the expected scaling of jet strength with accretion rate, nor have we investigated the effects of more realistic (i.e., strongly nonadiabatic) thermodynamics.

On the other hand, we have shown that the energy content of the jets is an extremely strong function of black hole spin. Crudely speaking, the jet efficiency (i.e., energy ejected after subtracting rest-mass energy, relative to the rate of rest-mass accretion) increases by an order of magnitude from a nonspinning black hole to a slowly spinning one and by another order of magnitude from slow spin to rapid spin. When the black hole rotation is truly rapid ($a/M = 0.998$), the efficiency of energy production in the outflow (0.18) is comparable to the radiation efficiency predicted by classical disk theory (0.32). Angular momentum loss in that case is also substantial; the angular momentum expelled in the unbound outflow exceeds the angular momentum accreted onto the black hole.

Although these simulations show how strong outflows can be created in some circumstances, they do not solve all the problems of astrophysical jet creation. The most serious limitation is that the outflow speeds we find are at most mildly relativistic, whereas we observe much more strongly relativistic jets in numerous places, including Galactic black hole binaries, active galaxies, and gamma-ray bursts. It is possible that truly relativistic
jets might result from stronger magnetic fields in the near-hole region (Meier 2003). How such fields might be created is an important question for further investigation.

The work of J. F. H. and J.-P. D. was supported by NSF grant AST-0070979, NSF ITR grant PHY02-05155, and NASA grant NNG04-GK77G. J. H. K. and S. H. were partially supported by NSF grants AST-0205806 and AST-0313031. The simulations were carried out on the Bluehorizon system of NPACI. We thank Charles Gammie and Jonathan McKinney for the discussion of their results and Henrique Schmitt and Tracy Clarke for additional guidance. J. H. K. is also grateful to the Institute of Astronomy, Cambridge, for its hospitality while this work was completed and to the Raymond and Beverly Sackler Fund for support during his visit there.

REFERENCES

Balbus, S. A., & Hawley, J. F. 1998, Rev. Mod. Phys., 70, 1

Begelman, M. C. 2001, in ASP Conf. Ser. 250, Particles and Fields in Radio Galaxies, ed. R. A. Liang & K. M. Blundell (San Francisco: ASP), 1

Blandford, R. D. 2002, in Lighthouses of the Universe, ed. M. Gilfanov, R. A. Sunyaev, & E. Churazov (Berlin: Springer), 381

Blandford, R. D., & Payne, D. 1982, MNRAS, 199, 883

Blandford, R. D., & Znajek, R. 1977, MNRAS, 179, 433

De Villiers, J. P., & Hawley, J. F. 2003a, ApJ, 589, 458

———. 2003b, ApJ, 592, 1060 (DH03)

De Villiers, J. P., Hawley, J. F., & Krolik, J. H. 2003, ApJ, 599, 1238 (Paper I)

Hawley, J. F. 1986, in Radiation Hydrodynamics in Stars and Compact Objects, ed. D. Mihalas & K.-H. Winkler (New York: Springer), 369

Hawley, J. F., & Balbus, S. A. 2002, ApJ, 573, 738 (HB02)

Hawley, J. F., & Krolik, J. H. 2001, ApJ, 548, 348

———. 2002, ApJ, 566, 164

Hirose, S., Krolik, J. H., De Villiers, J. P., & Hawley, J. F. 2004, ApJ, 606, 1083 (Paper II)

Junor, W., Biretta, J. A., & Livio, M. 1999, Nature, 401, 891

Kato, Y., Mineshige, S., & Shibata, K. 2004, ApJ, 605, 307

Koide, S., Meier, D. L., Shibata, K., & Kudoh, T. 2000, ApJ, 536, 668

Krasnopolsky, R., Li, Z.-Y., & Blandford, R. D. 1999, ApJ, 526, 631

Livio, M. 2002, Nature, 417, 125

Ly, C., Walker, R. C., & Wrobel, J. M. 2004, AJ, 127, 119

McKinney, J. C., & Gammie, C. F. 2004, ApJ, 611, 977

Meier, D. L. 2003, NewA Rev., 47, 667

Meier, D. L., Eddington, S., Godon, P., Payne, D. G., & Lind, K. R. 1997, Nature, 388, 350

Misner, C. W., Thorne, K. S., & Wheeler, J. A. 1973, Gravitation (San Francisco: Freeman)

Ouyed, R., & Pudritz, R. E. 1997, ApJ, 484, 794

Punsly, B., & Coroniti, F. V. 1990, ApJ, 354, 583

Romanova, M. M., Ustyugova, G. V., Koldoba, A. V., Cheeketkin, V. M., & Lovelace, R. V. 1997, ApJ, 482, 708

Sano, T., Inutsuka, S., Turner, N. J., & Stone, J. M. 2004, ApJ, 605, 321

Shibata, K., & Uchida, Y. 1985, PASJ, 37, 31

Stone, J. M., & Norman, M. L. 1994, ApJ, 433, 746