GR-MHD Disk Winds and Jets from Black Holes and Resistive Accretion Disks

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Abstract

We perform GR-MHD simulations of outflow launching from thin accretion disks. As in the nonrelativistic case, resistivity is essential for the mass loading of the disk wind. We implemented resistivity in the ideal GR-MHD code HARM3D, extending previous works for larger physical grids, higher spatial resolution, and longer simulation time. We consider an initially thin, resistive disk orbiting the black hole, threaded by a large-scale magnetic flux. As the system evolves, outflows are launched from the black hole magnetosphere and the disk surface. We mainly focus on disk outflows, investigating their MHD structure and energy output in comparison with the Poynting-dominated black hole jet. The disk wind encloses two components—a fast component dominated by the toroidal magnetic field and a slower component dominated by the poloidal field. The disk wind transitions from sub- to super-Alfvénic speed, reaching velocities ≈0.1c. We provide parameter studies varying spin parameter and resistivity level and measure the respective mass and energy fluxes. A higher spin strengthens the \( B_c \)-dominated disk wind along the inner jet. We disentangle a critical resistivity level that leads to a maximum mass and energy output for both, resulting from the interplay between reconnection and diffusion, which in combination govern the magnetic flux and the mass loading. For counterrotating black holes the outflow structure shows a magnetic field reversal. We estimate the opacity of the innermost accretion stream and the outflow structure around it. This stream may be critically opaque for a lensed signal, while the axial jet funnel remains optically thin.

Unified Astronomy Thesaurus concepts: Accretion (14); Magnetohydrodynamics (1964); Black hole physics (159); Active galaxies (17)

1. Introduction

Astrophysical jets appear as linearly collimated structures of high speed that are typically found in young stellar objects, X-ray binaries, gamma-ray bursts, or active galactic nuclei (AGNs). The physical mechanisms that produce these jets (jet launching) have been studied extensively. A consensus has been achieved, that launching of relativistic jets requires the existence of an accretion disk around a black hole and a strong magnetic field.

Blandford & Payne (1982) have proposed that jets can be formed as a result of magnetocentrifugal acceleration of matter from the surface of an accretion disk (hereafter the BP mechanism). On the other hand, Blandford & Znajek (1977) suggested that relativistic jets can be launched from the magnetosphere of a black hole by extracting rotational energy (hereafter the BZ mechanism). An interesting question for AGN jets is which of these mechanisms is responsible for generating the jets we observe. One way to investigate and compare the efficiency of these mechanisms is to use magnetohydrodynamic (MHD) simulations. In the case of the BZ mechanism, the equations of MHD need to be solved in a general relativistic (GR, GR-MHD) context (Einstein 1915).

Despite the abundance of observational data, it is almost impossible to resolve the jet-launching area for more than a few sources. Applying VLBI, Doeleman et al. (2012) determined a jet base of M87 of approximately 5.5\( R_\odot \) (Schwarzschild radii), which may imply that the BZ mechanism is responsible for feeding energy into the jet. On the other hand, Boccardi et al. (2016) find for the launching region of the Cygnus A jet a scale of 227 \( R_\odot \), thus suggesting that at least part of the jet may result from a disk wind (BP mechanism). Only high-resolution radio observations of the jet base and the accretion disk may discriminate which of the two mechanisms is more involved in the launching of jets. Another unresolved problem connected to the launching question is the matter content of relativistic jets. BZ-driven jets are expected to be leptonic and mass loaded by pair production in the strong radiation field of the black hole–disk corona, while jets launched as disk winds would be fed with hadronic disk material.

Most recently, the long-lasting search for a direct proof for the existence of supermassive black holes succeeded when the Event Horizon Telescope Collaboration (EHTC) released the first striking pictures of the shadow of the central black hole in M87, observed with short-wavelength 1.3 mm (Akiyama et al. 2019a). Radiation from an asymmetric ring around the black hole was detected and identified as a signature of the photon sphere around a Kerr black hole (Akiyama et al. 2019b).

As for the launching area, the jet propagation has been extensively investigated as well. For the example of M87\textsuperscript{15 GHz Very Large Array observations find that the jet knots are moving (outward) with apparent velocities of about 0.5c (Biretta et al. 1995). More recently, for the same source radio observations by Asada & Nakamura (2012) find an indication of a change in the jet opening angle at \( 10^5 \) Schwarzschild radii distance from the central black hole such that the jet shape changes from parabolic to conical. A similar behavior was detected for the jet and counterjet of NGC 4261 (Nakahara et al. 2018).

A physical complete theory that will fully connect the AGN jet-launching mechanism with the observed behavior of the jet is still under development. The general approach is to perform GR-MHD simulations of the close environment of the central black hole and the accretion disk to investigate and compare
the launching mechanisms of relativistic extragalactic jets. In
the past 20 years a significant number of GR-MHD codes have
been developed and used to simulate rotating disks around
black holes and their resulting outflows (Koide et al. 1999; De
Villiers & Hawley 2003a; Gammie et al. 2003; Noble et al.
2006, 2009a; Del Zanna et al. 2007; Bucciantini & Del
Zanna 2013; McKinney et al. 2014; Zanotti & Dumbser 2015;
Porth et al. 2017).

Koide et al. (1999) studied the development of a relativistic
jet in a Schwarzschild spacetime and identified a magnetic-
driven and a gas-pressure-driven component. De Villiers &
Hawley (2003b) focused on the accretion process between the
disk and the black hole for different black hole spins.
McKinney & Gammie (2004) examined the energy flux in
the black hole horizon in an attempt to detect the BZ
mechanism. McKinney et al. (2012) tested magnetically
choked accretion flows and detected quasi-periodic oscillations
between the accreting inflow and the jet magnetosphere. In
Tchekhovskoy et al. (2010, 2011) the authors simulated
accretion flows into extreme Kerr black holes to measure the
energy extracted by the BZ mechanism. Radiative transfer in
combination with GR-MHD codes allows the study of spectra
of GR accretion disks (Noble et al. 2009b) or their evolution in
the super-Eddington limit (McKinney et al. 2014; Sadowski
et al. 2014). Recently, Nakamura et al. (2018) compared the jet
funnel seen in GR-MHD and force-free electrodynamic
simulations with VLBI data of M87, finding good agreement
concerning a parabolic jet shape.

With a number of codes available, it is possible to perform
comparison studies, as analytical test problems in GR-MHD do
not exist. A major breakthrough along these lines has been
achieved as an integral part of the EHTC studies, comparing a
set of GR-MHD codes (including HARM3D) in the ideal-MHD
limit, simulating a torus around a black hole (Porth et al.
2019). All codes produce very similar results, confirming the
robustness of the methods used.

In contrast to most of the GR-MHD simulations including
the above-mentioned code-comparison studies, one of the
specific features of our present study is that we follow the
evolution of a thin disk right from the start of the simulation.
Thin disks were first studied in a purely hydrodynamic
approach by Shakura & Sunyaev (1973) in the nonrelativistic
limit and by Novikov & Thorne (1973) for the general
relativistic case. As a seminal step forward, the α-viscosity was
invented as a mean driver of angular momentum exchange in
disks (Shakura & Sunyaev 1973). Launching simulations of
jets out of thin disks using nonrelativistic resistive MHD were
pioneered by Casse & Keppens (2002). Those simulations and
many follow-up studies essentially apply resistivity or mag-
etic diffusivity to allow matter to be accreted through the
magnetic field that threads the disk, and also disk material to be
loaded on the jet magnetic field, eventually leading the system
into an inflow–outflow structure in a quasi-stationary state. We
further refer to Zanni et al. (2007), who studied the efficiency
of the magnetocentrifugal acceleration mechanism for different
levels of resistivity (see also Sheikholeslami et al. 2012).

It thus seems essential to apply resistive MHD for disk–jet
launching also for the relativistic case. Magnetically diffusive
MHD codes for the relativistic case have been developed only
rather recently. Resistive MHD for special relativistic simulations
was pioneered by Komissarov (2007). Palenzuela et al. (2009)
applied an implicit–explicit solver for the resistive GR-MHD
equations in order to deal with the stiff part of the electric
field, allowing them to model magnetized rotating neutron stars
(Palenzuela 2013). A similar scheme was used by Dionysoiopoulou
et al. (2013) for the resistive version of the WHISKY code (Baiotti
et al. 2005), which was then to study collisions of binary neutron
stars (Dionysoiopoulou et al. 2015). The ideal-MHD code ECHO
(Del Zanna et al. 2007) was also extended to the resistive regime
considering as a fully covariant mean-field dynamo closure
(Bucciantini & Del Zanna 2013). Subsequently, Bugli et al.
(2014) investigated the evolution of a kinematic mean-field
dynamo in thick accretion disks. Porth et al. (2017) presented
a GR-MHD code particularly suited for black hole accretion and
Ripperda et al. (2019) evolved it further including resistivity and a
new inverse solver for the electric field.

In the present paper we have expanded the physics of the
parallel, 3D, conservative, GR-MHD code HARM3D (Gammie
et al. 2003; Noble et al. 2006, 2009a) by implementing resistivity
in the form of a magnetic diffusivity, following Bucciantini & Del
Zanna (2013) and Qian et al. (2017, 2018). This allows us to run axisymmetric (so-called 2.5D) simulations
of thin accretion disks around black holes in order to
investigate the detailed launching conditions that favor the
generation of relativistic jets. In particular, we are interested in
comparing the energy budget of the jet launched from the black
hole magnetosphere to the jets launched from the disk and to
compare the outflow mass fluxes to that of the disk accretion.
Compared to our previous works (Qian et al. 2017, 2018), we
can now take advantage of the parallelization of the code and
can aim for long-lasting simulation runs on larger domains and
with better grid resolution.

Our paper is structured as follows. Section 2 introduces the
basic theory of (resistive) GR-MHD. Section 3 includes a
summary of the initial setup, as well as the characteristic
properties of the simulations. Section 4 discusses our reference
simulation and the outflows it develops. Section 5 compares the
reference simulation with simulations of different black hole
rotation and levels of magnetic diffusivity. In Section 6 we briefly
discuss our results in light of the recently detected black hole
shadow in the center of M87. Finally, Section 7 summarizes our
work. In the appendices we present test simulations for the
implementation of resistivity and a test simulation of the thin-disk
setup using the GR-MHD code in the mildly relativistic limit.

2. Theoretical Background

Here we review the basic equations of resistive GR-MHD as
a basis for our implementation of resistivity (in the form of
magnetic diffusivity) in the formerly ideal GR-MHD code
HARM3D (Gammie et al. 2003; Noble et al. 2006, 2009a).

We adopt the signature of Misner et al. (1973) for the metric
(−, +, +, +) and use geometrized units where \( G = c = 1 \).
Greek letters run for 0, 1, 2, 3 (\( t, r, \theta, \phi \)) and Latin letters run
for 1, 2, 3 (\( r, \theta, \phi \)). Radii are expressed in units of the
gravitational radius, \( R_g = GM/c^2 \), while time is in units of
light-travel time \( t_g = GM/c^2 \). Vector quantities are denoted
with bold letters, while the vector and tensor components are
indicated with their respective indices.

Our code uses the “3+1” decomposition of the GR-MHD
equations, where the time component is separated from the
spatial components, which are expressed as 3D manifolds.
The spacetime is described by the metric \( g_{\mu \nu} \) in Kerr–Schild
coordinates with \( g = \det(g_{\mu \nu}) \). A zero angular momentum
observer (ZAMO) frame exists in the spacelike manifolds
moving only in time with velocity \( n_\mu = (-\alpha, 0, 0, 0) \), where \( \alpha = 1/\sqrt{1 - v^2} \) is the lapse function. The gravitational shift is \( \beta^i = \alpha^2 g^i \). The four-velocity of the fluid in the comoving frame is \( u^\mu = (u^t, 0, 0, u^\theta) \). The code solves the equations of resistive GR-MHD using a conservative scheme based on the previous works of Gammie et al. (2003) and Noble et al. (2006).

We extended the physics simulated by the code to the resistive GR-MHD regime by implementing magnetic diffusivity in the ideal GR-MHD version of the code, following the work of Bucciantini & Del Zanna (2013) and Qian et al. (2017). As a result, we were required to increase the number of variables from 8 to 11, adding the three components of the electric field. We denote our new resistive GR-MHD code with rHARM3D.

### 2.1. Basic GR-MHD Equations

A magnetized fluid in a general relativistic environment is described by the Maxwell equations (Maxwell 1865) in covariant form

\[
\nabla_\mu F^{\mu \nu} = 0, \quad \nabla_\nu F^{\mu \nu} = J^\mu,
\]

where

\[
F^{\mu \nu} = u^\mu e^\nu - e^\mu u^\nu + \epsilon^{\mu \nu \alpha \beta} u_\alpha b_\beta \tag{2}
\]

\[\ast F^{\mu \nu} = -u^\mu b^\nu + b^\mu u^\nu + \epsilon^{\mu \nu \alpha \beta} u_\alpha e_\beta \tag{2}\]

are the antisymmetric Faraday and Maxwell tensors, \( e^\mu \) and \( b^\mu \) are the electric and magnetic field in the fluid rest frame, and \( \epsilon^{\mu \nu \alpha \beta} \) is the Levi-Civita symbol

\[
\epsilon_{\alpha \beta \gamma \delta} = \sqrt{-g} [\alpha \beta \gamma \delta], \quad \epsilon^{\alpha \beta \gamma \delta} = -\frac{1}{\sqrt{-g}} [\alpha \beta \gamma \delta]. \tag{3}
\]

The magnetic and electric fields as measured by the normal observer are defined as \( B^i = n^a F^{\mu \nu} \delta_{ab} \) and \( E^i = n^a F^{\mu \nu} = -\alpha F^{\mu} \). The equations of motion for the magnetized fluid are

\[
\nabla_\mu T^{\mu \nu} = 0, \tag{4}
\]

where \( T^{\mu \nu} = T_{\text{fluid}}^{\mu \nu} + T_{\text{EM}}^{\mu \nu} \) is the stress-energy tensor, which can be split into a fluid part and an electromagnetic part. The fluid component can be written as

\[
T_{\text{fluid}}^{\mu \nu} = (\rho + u + p) u^\mu u^\nu + pg^{\mu \nu}, \tag{5}
\]

where \( \rho \) is the mass density, \( u \) is the internal energy, and \( p \) is the thermal pressure. Pressure and internal energy are connected through

\[
u = \frac{p}{\Gamma - 1}, \tag{6}
\]

where \( \Gamma \) is the polytropic exponent. The electromagnetic component can be written as

\[
T_{\text{EM}}^{\mu \nu} = (b^2 + e^2) \left( u^\mu u^\nu + \frac{g^{\mu \nu}}{2} \right) - b^\mu b^\nu - e^\mu e^\nu - u^\alpha e_\beta (u^\mu e^{\alpha \beta \gamma} + u^\nu e^{\alpha \beta \gamma}), \tag{7}
\]

These two components can be combined into the total stress-energy tensor

\[
T^{\mu \nu} = (\rho + u + p + b^2 + e^2) u^\mu u^\nu + \left( p + \frac{1}{2} (b^2 + e^2) \right) g^{\mu \nu} - b^\mu b^\nu - e^\mu e^\nu - u^\alpha e_\beta (u^\mu e^{\alpha \beta \gamma} + u^\nu e^{\alpha \beta \gamma}). \tag{8}
\]

### 2.2. From Ideal to Resistive GR-MHD

Resistivity enters the equations in the form of an (anomalous) magnetic diffusivity \( \eta = \eta(r, \theta) \) that is believed to be of turbulent nature. In ideal MHD, the electric field is given by Ohm’s law \( E + v \times B = 0 \). In the resistive regime Ohm’s law becomes

\[
E + v \times B = \eta J, \tag{9}
\]

or in covariant form in the fluid frame

\[
e^\mu = \eta j^\mu, \tag{10}
\]

where \( J \) is the electric current density. In the resistive environment, the electric field can no longer be calculated by the cross product of fluid velocity and magnetic field and new equations need to be formulated. By setting \( \eta = 0 \), we get back into the ideal case \( e^\mu = 0 \).

Furthermore, magnetic diffusivity puts a restriction in the time step of a numerical simulation, as the diffusive time step goes as \( dt_\eta = (\Delta x_i)^2 / \eta \), where \( \Delta x_i \) is the smallest cell size in any dimension of the grid. Thus, for high values of magnetic diffusivity we expect the diffusive time step to become lower than the MHD step and to effectively determine the evolution of the simulation.

### 3. Simulation Setup

This paper considers GR-MHD simulations of thin accretion disks that rotate differentially around a (rotating) black hole and are threaded by a poloidal magnetic field. Here we describe the initial conditions we use for our models, the boundary conditions, and other numerical details of the simulation.

#### 3.1. Numerical Grid

Depending on our problem setup, we apply a different numerical grid. The original grid of HARM applying modified Kerr–Schild coordinates is used for our test simulations of diffusivity and for the simulation in the mildly relativistic limit (see Appendices A and B).

For our science applications we decided to construct a stretched grid in order to shift the outer boundary condition as far out as possible. This grid is an extension of the original HARM grid and is based on the hyper-logarithmic grid of Tchekhovskoy et al. (2009). With that we may concentrate cells close to the black hole in the radial direction and concentrate cells close to the equatorial plane or the polar axis in the polar direction, allowing us to resolve the turbulent disk and the polar jet at the same time. Furthermore, with such a scheme the outer boundary is causally disconnected from the inner simulated area of interest close to the black hole or the disk.

In the hyper-logarithmic grid the radial coordinate is split into two parts. The first part follows a logarithmic scaling as in
the original HARM code (Gammie et al. 2003). Beyond a transition radius \( R_t \), the grid becomes substantially more scarce, up to the outer radius \( R_{\text{out}} \).

Physical and numerical radial coordinates translate as

\[
r(x) = \exp \left[ \frac{1}{2} x_1 + 4H(x - x_{11})^4 \right],
\]

where \( x_1 \) is the uniformly spaced numerical radial coordinate and \( x_{11} \) is the transition radius (corresponding to \( R_t \)). The function \( H = H(x - x_{11}) \) is a step function that is equal to unity for \( x > x_{11} \) and vanishes otherwise. In Figure 1 we show the relation between the numerical and the physical radial coordinates.

The physical and numerical polar coordinates are connected by

\[
\theta(x_2) = \theta_{\text{start}} + x_2 \theta_{\text{length}} - h_{\text{slope}} \sin(4\pi x_2),
\]

where \( \theta \) and \( x_2 \) are the physical and numerical polar coordinates, respectively, while \( \theta_{\text{start}} \) denotes the starting angle and \( \theta_{\text{length}} = \pi - 2\theta_{\text{start}} \) the angular length of the coordinate in radians. The factor \( h_{\text{slope}} \) governs how many grid cells are focused toward the equatorial plane and toward the symmetry axis. We note that these coordinates are slightly different from the original HARM code, where the choice of focusing coordinates and the increase of resolution for the polar coordinate are only possible toward the equatorial plane.

Our typical maximum resolution in the polar coordinate is \( \Delta \theta = 0.00625 \) along the polar axis and in the equatorial plane, while the minimum resolution is \( \Delta \theta = 0.025 \) at \( 45^\circ \). The radial coordinate is best resolved close to the horizon, where \( \Delta r_{r=2} = 0.02 \), and is radially decreasing with \( \Delta r_{r=10} = 0.1 \), \( \Delta r_{r=50} = 0.5 \), and \( \Delta r_{r=100} = 1 \).

### 3.2. Boundary Conditions

For our simulations we use outflow boundary conditions in the inner and outer radial boundary. The values of the primitive variables are copied from the boundary cells to the ghost cells. At the same time we make sure that there is no inflow from the boundaries by checking that the velocity is pointing outward at each boundary cell. As an extra measure in the inner boundary, we make sure we have 10 cells of our grid inside the black hole event horizon in order to prevent numerical effects from propagating outside of it.

Furthermore, one of the reasons we modified our numerical grid into the hyper-logarithmic version we described before is because we wanted to have the outer boundary as far away from the disk as possible. Before adopting the hyper-logarithmic grid, we had noticed a collimation effect in the magnetic field lines that we had deemed as artificial (see Appendix B in Qian et al. 2018). By selecting an outer radius of \( R_{\text{out}} = 10^4 \), we make sure that the outer boundary stays causally disconnected from the disk. In the axial boundary we impose axisymmetric boundary conditions where the vector values are being reflected along the small cutout in both axes.

### 3.3. Initial Conditions

The initial disk density distribution is described by a nonrelativistic vertical equilibrium profile, such as applied in Sheikhi Nezami et al. (2012),

\[
\rho(r, \theta) = \frac{\Gamma - 1}{\Gamma} \frac{r_{\text{in}}}{r} \frac{1}{\epsilon^2} \left( \sin \theta + \epsilon \frac{\Gamma}{\Gamma - 1} \right)^{\Gamma/(\Gamma - 1)},
\]

slightly modified to fit into our code. Here \( r_{\text{in}} \) is the initial inner disk and \( \epsilon = H/r \) is the initial disk aspect ratio as is defined by the vertical equilibrium of a disk with a local pressure scale height \( H(r) \). The pressure and internal energy are given by the polytropic equation of state \( p = K p^\Gamma \) and Equation (6), where \( K \) is the polytropic constant. For the polytropic exponent \( \Gamma \) we will use different values for different simulations as specified in the sections below.

Around the disk we prescribe an initial “corona.” For the choice of a polytropic index of \( \Gamma = 4/3 \), the disk has a finite outer radius much smaller than the outer radius of the stretched grid. Furthermore, the upper and lower disk surfaces do not follow lines of constant polar angle as implied by Equation (13). The initial coronal density and pressure are given by

\[
\rho_{\text{cor}} \propto r^{1/(1-\Gamma)}, \quad p_{\text{cor}} = K_{\text{cor}} r^{\Gamma}.\]

The coronal temperature is chosen to be much higher than the disk temperature, \( K_{\text{cor}} \gg K \). This implies a density jump between disk and corona but a pressure equilibrium along the disk surface. More specifically, for our simulation we chose \( K = 0.001 \) for the disk initial condition and \( K = 1 \) for the initial corona. The corona collapses instantly the moment the simulation starts, and part of it is also expelled by the initial ejections from the disk, meaning that the values are quickly replaced by the floor values of the simulation (see Section 3.5). However, the polytropic equation \( p = K p^\Gamma \) is not enforced in any step of the code except the initial condition. The code uses Equation (6) to connect pressure and internal energy, which means that entropy and temperature are free to change.

The disk is given an initial orbital velocity following Paczyński & Wiita (1980),

\[
\tilde{u}^\phi = r^{-3/2} \left( \frac{r}{r - R_{\text{PW}}} \right),
\]

where \( R_{\text{PW}} \) is a constant of choice, here equal to the gravitational radius \( R_g \). This approximation is applied in the \( \phi \)-component of the fluid velocity \( \tilde{u}^\phi \).

The disk is initially threaded by a large-scale poloidal magnetic field, implemented via the magnetic vector potential.
$A_\phi$ following $B = \nabla \times A_\phi$. In most cases we use the inclined field profile suggested by Zanni et al. (2007),

$$A_\phi(r, \theta) \propto (r \sin \theta)^{3/4} \frac{m^{5/4}}{(m^2 + \tan^{-2} \theta)^{5/8}}. \quad (16)$$

The parameter $m$ determines the initial inclination of the field, which plays an important role for the magnetocentrifugal launching of disk winds (Blandford & Payne 1982). The magnetic field strength is then normalized by choice of the plasma-$\beta = p_{\text{gas}}/p_{\text{magn}}$.

### 3.4. The Magnetic Diffusivity

The simulations presented in this paper apply a scalar function for the magnetic diffusivity that is constant in time. The diffusivity is assumed to be of turbulent nature, thus much larger than the microscopic resistivity, and thought to be generated by the magnetorotational instability (MRI; Balbus & Hawley 1991). In general, the magnetic diffusivity distribution is chosen such that it resembles a magnetized diffusive disk within an ideal-MHD wind and jet area.

We apply a magnetic diffusivity profile, as it is typically used in jet-launching simulations (see, e.g., Zanni et al. 2007; Sheikhnezami et al. 2012; Qian et al. 2018), i.e., a Gaussian profile along the polar angle with a maximum at the initial disk midplane

$$\eta(r, \theta) = \eta_0 \exp \left[ -2 \left( \frac{\alpha}{\alpha_\eta} \right)^2 \right], \quad (17)$$

where $\eta_0$ is the level of diffusivity along the equatorial plane, $\alpha = \pi/2 - \theta$ is the angle toward the disk midplane, and $\alpha_\eta = \arctan(\chi \times \epsilon)$ is the angle that measures the scale height of the diffusivity profile. The parameter $\chi$ compares the scale height of the diffusivity profile with the disk pressure scale height. This profile—as artificial as it might seem—focuses the high diffusivity values in the equatorial plane, allowing for a highly resistive material and for a lowly resistive to asymptotically ideal-MHD disk wind and jet. Since we take resistivity as a result of turbulence, we expect higher diffusivity in the highly turbulent areas of the interior of the disk.

Initially, we also considered an anisotropic resistivity profile with different values of $\alpha$ affecting the poloidal and toroidal components of the magnetic field. According to Ferreira (1997), such a profile would help stabilize the disk evolution reaching a stationary state. In contrast to Zanni et al. (2007), Murphy et al. (2010), and Sheikhnezami et al. (2012), who applied an anisotropic diffusivity in their simulations, in our case the disk loses its mass rather quickly, mainly due to the disk wind. This rapid mass loss is actually minimizing the stabilization effect by an anisotropic magnetic diffusivity. Furthermore, the initial ejections created by the absence of equilibrium between the disk and the black hole delay the reach of a stationary condition even more. Based on that, we decided that the introduction of anisotropic diffusivity would not contribute much in the stability of the disk.

When testing the performance of our code, we found that the simulations become more stable when we apply a low background diffusivity (1000 times lower than in the disk) along the rotational axis. We thus apply an exponentially decreasing profile along the axial boundary within six grid cells. As this axial diffusivity is confined within an opening angle of $\lesssim 3^\circ$, it does not affect the physics of the jet launching. We also apply an exponential decrease in the radial diffusivity profile from radius $r = 3$ toward the horizon, resulting in a smooth transition from the high disk diffusivity to the ideal-MHD black hole environment. Figure 2 shows the 2D distribution, as well as the radial and angular profiles of $\eta$.

For the magnitude of the magnetic diffusivity $\eta_0$ we apply a range of values, $\eta_0 = 10^{-10} ... 10^{-2}$ (in code units). These values correspond to some kind of standard parameters applied in the literature in diffusive MHD simulation in GR (Bucciantini & Del Zanna 2013; Bugli et al. 2014; Qian et al. 2017, 2018), in nonrelativistic simulations (Casse & Keppens 2002; Zanni et al. 2007; Sheikhnezami et al. 2012; Stepanovs & Fendt 2014), but have been modeled concerning strength and spatial distribution also by direct simulations, e.g., by Gressel (2010). Concerning the diffusive numerical time stepping and the strength of the numerical diffusivity, we refer to the discussion in our previous works (Qian et al. 2017, 2018).

Here we emphasize another important impact of physical resistivity: it suppresses the MRI (Fleming et al. 2000; Longaretti & Lesur 2010). Overall, we thus do not expect to detect any MRI being resolved in our disk structure.

As discussed in Qian et al. (2017), the diffusion rate will be of order $k^2 \eta$ (Fleming et al. 2000), with the wavenumber $k$. From Balbus & Hawley (1991) we know that the MRI grows only in a certain range of wavenumbers $k \in [0, k_{\text{max}}]$, in the linear MRI regime—depending on whether the numerical grid may resolve certain wavelengths and whether certain wavelengths will fit into the disk pressure scale height. Furthermore, there exists a wavenumber $k_{\text{MRI}}$ for which the MRI growth rate reaches a maximum (see Hawley & Balbus 1992 for the case of a Keplerian disk). A certain number of MRI modes can therefore be damped out when $k_{\text{MRI}} \eta$ is comparable to the maximum growth rate of MRI. Moreover, for a large enough $\eta$, it is even possible to damp out most of the MRI modes in the linear evolution of MRI.

In Qian et al. (2017) a thorough investigation of resistive effects on the accretion rate of an initial Fishbone & Moncrief (Fishbone & Moncrief 1976) torus was presented. They could show that for this setup for $\eta \lesssim 10^{-3}$ the MRI seemed to be completely damped, while for lower $\eta$ the onset of the MRI and...
thus of massive accretion was substantially delayed. This result was claimed to be consistent with Longaretti & Lesur (2010), demonstrating that the growth rate of the MRI substantially decreases with $1/\eta$ beyond a critical diffusivity.

In addition to the point that we do not expect the MRI to play a role in our simulations owing to the disk resistivity, we also note that we consider a thin disk that is thread by a strong magnetic field. Thus, angular momentum transport is dominated by the torque of the magnetic lever arm.

Other consequences of considering a magnetic diffusivity are physical reconnection of the magnetic field and also physical ohmic heating. Both effects are present in our simulations, and we will discuss their impact on the accretion–ejection system accordingly.

3.5. The Density Floor Model

As typical for any MHD code, rHARM3D cannot work in vacuum. This is a problem also for relativistic MHD codes, in particular for their inversion schemes, so usually a floor model is applied to circumvent numerical problems when the initial disk corona collapses.

Depending on the model setup, we apply a different floor model for the density and pressure. Note that in particular in our approach that applies a large-scale initial disk magnetic field, we potentially deal with a high magnetization $B^2/\rho$ and/or low plasma-beta $P/B^2$ at large radii. Thus, for simulations on a large grid, we choose a floor profile following a broken power law. For the density we apply

$$\rho_{fl}(r) \propto \left[ \left( \frac{r}{r_0} \right)^{\frac{1}{2} (1-\Gamma_1)} + \left( \frac{r}{r_0} \right)^{\frac{1}{2} (1-\Gamma_2)} \right], \quad (18)$$

while the internal energy follows

$$u_{fl}(r) \propto \left[ \left( \frac{r}{r_0} \right)^{\frac{1}{2} \Gamma_1 (1-\Gamma_1)} + \left( \frac{r}{r_0} \right)^{\frac{1}{2} \Gamma_2 (1-\Gamma_2)} \right], \quad (19)$$

with $\Gamma_1 = 4/3$ and $\Gamma_2 = 2$, and where $r_0$ marks the transition radius between the two power laws with typically $r_0 = 10 R_g$ (see Figure 3). With that we implement higher floor values for large radii in order to avoid a too high magnetization. Close to the black hole we apply the same floor profile as in the original HARM code.

3.6. Characteristic Quantities of the Simulations

Here we define a number of physical quantities, which will later be used to characterize the evolution in different simulations. The mass contained in a disk-shaped area between radii $r_1$ and $r_2$ and between surfaces of constant angle $\Theta_1$ and $\Theta_2$ is calculated as

$$M_{\text{disk}} = 2\pi \int_{r_1}^{r_2} \int_{\Theta_1}^{\Theta_2} \rho(r, \Theta) \sqrt{-g(R, \Theta)} \, d\Theta \, dr, \quad (20)$$

where $\sqrt{-g}$ is the square root of the determinant of the metric. The mass flux through a sphere of radius $R$ between angles $(\Theta_1, \Theta_2)$ is

$$M(R) = 2\pi \int_{\Theta_1}^{\Theta_2} \rho(R, \Theta) u^\theta(R, \Theta) \sqrt{-g(R, \Theta)} \, d\Theta. \quad (21)$$

Similarly, we calculate the mass flux in the $\theta$-direction considering the $u^\theta$ component and the area element $\sqrt{-g}$.

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**Figure 3.** Floor model used in the science simulations. The radial density (top) and internal energy (bottom) distributions are calculated as broken power laws.

$\sqrt{-g(R, \Theta)} \, d\Theta \, dr$. This is in particular used to calculate the disk wind mass flux from the disk surface, considering two surfaces with a constant opening angle $\Theta$ that is chosen to be similar to the initial disk density distribution. We thus obtain

$$M(\Theta) = 2\pi \int_{r_1}^{r_2} \rho(r, \Theta) u^\theta(r, \Theta) \sqrt{-g(r, \Theta)} \, dr. \quad (22)$$

The Poynting flux per solid angle is defined as

$$F_{\text{EM}}(r, \Theta) = -T_i^\theta$$

$$= \left[ (b^2 + e^2)(u^\theta u_t + \frac{1}{2} g^\theta_\theta) - b^\theta b_t - e^\theta e_t - u_j e^i b_i (u^\mu u^\nu g^\mu_\nu + u^\mu e^\mu_\nu g^\mu_\nu) \right]. \quad (23)$$

By integration along the polar angle we obtain the flux through a sphere of radius $R$,

$$E_{\text{EM}}(R) = 2\pi \int_0^\pi \sqrt{-g(R, \Theta)} \, F_{\text{EM}}(R, \Theta) \, d\Theta. \quad (24)$$

The corresponding electromagnetic flux is

$$F_{\text{EM}}(r, \Theta) = -T_i^\theta$$

$$= \left[ (b^2 + e^2)(u^\theta u_t + \frac{1}{2} g^\theta_\theta) - b^\theta b_t - e^\theta e_t - u_j e^i b_i (u^\mu u^\nu g^\mu_\nu + u^\mu e^\mu_\nu g^\mu_\nu) \right]. \quad (25)$$

By integration along the radius, we obtain the flux through a surface of constant angle $\Theta$,

$$E_{\text{EM}}(\Theta) = 2\pi \int_{r_1}^{r_2} \sqrt{-g(r, \Theta)} \, F_{\text{EM}}(r, \Theta) \, dr. \quad (26)$$

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The poloidal Alfvén Mach number is \( M_{A,p} = [h \mu_0^2 B_p^2]^{1/2} \) (see also Qian et al. 2018), where \( h \equiv \Gamma/(\Gamma - 1)(P/\rho) + 1 \) is the specific enthalpy of the fluid. Alfvénic Mach numbers \( M_{A,p} < 1 \) imply that the magnetic energy is dominating the kinetic energy of the fluid and that the dynamics of the outflow is most likely governed by the strong magnetic field in that area.

4. A Reference Simulation

In the following we will first describe the details of our reference simulation that will be used to compare our parameter runs for characteristic properties of the source. The reference simulation runs for 9000\( t_g \), corresponding to approximately 67 disk orbits at the initial inner disk radius. In Figure 4 we show the evolution of the density distribution, the poloidal magnetic field lines, and the poloidal velocity field up to time \( t = 8000 \).

4.1. Initial Conditions

The distributions for the initial density, pressure, angular velocity, and magnetic vector potential are given by Equations (6), (13), (15), and (16). For the disk rotation (15) we impose a factor of 0.95 in order to treat a sub-Keplerian disk. For the disk gas law we apply \( \Gamma = 4/3 \) and \( K = 0.001 \). For a Kerr parameter of \( a = 0 \), the horizon is located at \( r = 1.4358 \) and the innermost stable circular orbit (ISCO) at \( r = 2.32088 \).

For the numerical grid we choose a transition radius \( R_o = 200 \) and an outer grid radius \( R_{\text{out}} = 10^4 \). The initial inner disk radius \( r_0 = 7 \) is outside the ISCO in order to avoid possible initial ejections of gas, as the initial disk is not in force equilibrium within GR. At this radius the initial angular velocity of the disk is \( \Omega_{\text{ref}} \approx 0.047 \), thus slightly lower than the Keplerian value \( \Omega_K \approx 0.052 \), and corresponding to an orbital period of \( T_o \approx 135 \).

The initial corona is given by Equation (14) with \( K_{\text{cor}} = 1 \), resulting in a higher coronal entropy. The initial magnetic field structure follows Equation (16) with \( m = 0.6 \). The magnetic field strength is fixed by the choice of the plasma-\( \beta \) = 10 at the initial inner disk radius. The magnetic diffusivity profile is given by Equation (17) with \( \chi = 3 \) and \( \eta_0 = 0.001 \) (see Figure 2).

4.2. Evolution of Disk Mass and Disk Accretion

As the disk evolves, accretion sets in and the inner disk radius changes to lower values, extending down to \( r \approx 3 \) right outside of the ISCO after having completed more than 200 rotations at this radius. Since the shape of the disk changes constantly, it is difficult to measure the total disk mass. One option is to measure the mass within a disk area defined by the inner surface located at \( r = 3 \), an outer surface at \( r = 100 \), and the surfaces of constant opening angle of \( \theta \approx 80^\circ \) and \( \theta \approx 100^\circ \). The disk mass is then obtained by integrating the mass density in the disk area as specified above.

The evolution of the disk mass is shown in Figure 5. Since the disk is not in equilibrium, there is a rapid change in the innermost part of the disk that causes a small initial increase. We understand that the extra mass for the disk arises from the initial corona, which immediately starts to collapse, and by that squeezes and relaxes the inner disk until a quasi-equilibrium is reached at \( t \approx 300 \). After that, the disk mass decreases steadily until \( t \approx 5500 \), when the slope of the disk mass evolution changes. This is mainly due to changes in the disk outflow. Note that by the end of the simulation the disk has lost more than 80% of its initial mass.

Figure 6 shows the normalized accretion rates measured through three different radii, \( r = 2, 4, \) and \( 13 \), and integrated over the disk scale height. Close to the horizon, measured at radius \( r = 2 \), the accretion rate is first negligible, mainly because of the absence of disk material in that radius. After \( t = 3000 \), accretion rate increases. Note that by now the inner disk radius, located initially at \( r = 7 \), has moved closer to the black hole, populating that area with dense disk material. The enhanced accretion level is accompanied by substantial accretion spikes. However, the underlying base accretion rate seems to decrease as the disk loses mass. The accretion mass flux in the inner area is of the order of \( 10^{-4} \).

After \( t = 6000 \) and until the end of the simulation, the innermost area around \( r = 2 \) becomes almost empty again, with the exception of a thin stream of material that is connecting the disk with the black hole. It looks like that at this point in time all material close to the black hole has fallen into it but has not been replenished by disk material from larger radii. As a result, accretion at \( r = 2 \) is halted completely for a substantial period of time, until it is temporarily restarted by disk material that has newly arrived (accreted) from larger radii. This relaunch of accretion is indicated by the spikes in the accretion rate at late times.

Similar to \( r = 2 \), at radius \( r = 4 \) accretion is not significant until \( t = 1500 \), while it gradually increases afterward until \( t = 3000 \). In the following strong accretion phase \( (t \in [3000, 6000]) \), there is also a significant amount of material moving outward. In the inner part of the disk, just outside the ISCO, the gas is actually moving in both directions, radially inward and outward, thus indicating the turbulent character of the motion. The highly turbulent nature of the inner accretion flow is shown in Figure 7. The figure demonstrates the rapid change in density and velocity within short time. Note the strong gradient in velocity at the ergosphere (yellow line; dark blue indicates high infall speed).

Since the average accretion rate at radius \( r = 4 \) is similar to that measured at \( r = 2 \), we conclude that the accretion mass flux is conserved and, thus, no outflow is ejected from this area close to the black hole. Even farther out, at radius \( r = 13 \), the accretion process looks quite different. The accretion rate is again of an order of magnitude similar to smaller radii. The accretion spikes that are seen at lower radii now are replaced with much broader time periods of high mass accretion, indicating a slower change to the accretion rate.

However, we still detect a few accretion spikes during the third phase of evolution. In fact, the accretion spikes that are observed at \( r = 13 \) are subsequently followed by spikes at \( r = 4 \) and \( r = 2 \). We measure a time delay between the spikes at \( r = 13 \) and \( r = 4 \) varying between \( \Delta t = 75 \) and \( \Delta t = 40 \). The time delay between the spikes at \( r = 4 \) and \( r = 2 \) is \( \Delta t \approx 10 \). An approximate average accretion velocity can be defined by dividing the distance traveled by the fluid by the time delay of the spikes. For the three major spikes appearing at radius \( r = 2 \) at \( t = [5630, 6070, 7920] \) we measure a similar velocity

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5 The disk mass and mass flux are normalized by the mass of the initial disk material included in the disk area as specified above.

6 With the exception of the floor density accretion.

7 This is also the time sequence for our data dumps, so we cannot provide a higher time resolution for the pattern speed of the spikes.
from radius $r = 4$ to radius $r = 2$ of 0.2 for all three spikes. For the pattern speed of the spikes from $r = 13$ to $r = 4$ we measure velocities of 0.12, 0.225, and 0.16, respectively. These values derived for the pattern speed agree well with the radial velocity that we observe in this area of the disk.

At late stages of the simulation (between $t = 6000$ and 9000) we notice a decline in the accretion rate at all three radii. This is accompanied by the opening of a larger gap between the horizon and the inner disk, meaning that the inner disk radius moves out. At this time, the disk has already lost 70% of its mass. During this period, the disk accretion becomes disconnected from the black hole. We interpret this as follows. Due to the decrease of density and pressure (following accretion and ejection of disk material), this area becomes magnetically dominated. The strong magnetization leads to the structure of a magnetically arrested disk (MAD; see Narayan et al. 2003). When the magnetic flux is advected to the black hole, the magnetization in this area decreases again, and accretion restarts (see Figure 8).

Figure 4. Snapshots of our reference simulation. Shown is the density distribution (in log scale) at times $t = 0$, 1000, 2000, 4000, 6000, and 8000. The poloidal magnetic field lines are shown (white lines), while the poloidal velocity field is represented by the black arrows.
Finally, in Figure 9 we display the mass fluxes vertical to the surfaces of constant opening angles \((\theta = 80^\circ, 100^\circ)\) that approximate the surfaces of the initial disk density distribution. We see that the mass fluxes of accretion and radial outflow along the disk are comparable. However, both are dominated by the vertical mass loss from the disk surface. The low accretion rate is comparable to an MAD structure (Tchekhovskoy et al. 2011), and due to the strong disk magnetic field, a strong outflow is launched, but at the same time the accretion rate decreases. Obviously, also the strength of magnetic diffusivity plays a role (see our comparison study below). We may conclude that most of the mass that the disk is losing is due to the strong disk wind that is launched.

We note that since a substantial disk wind is present during the whole simulation, the wind mass-loss rate is changing. The wind mass flux increases until about \(t = 3000\) and then decreases again until \(t = 6000\). In the late stages of the simulation the wind mass flux is highly variable. These two different phases of wind ejection seem to correspond to similar phases in the disk accretion, visible in Figure 6 (middle panel), which shows large variations in the mass accretion rate, or in Figure 5, which indicates a change in the disk mass evolution at \(t = 3000\).

In order to double-check our mass flux integration, we have measured the total mass loss of the disk with two different methods. First, we integrate all mass flux, leaving the surfaces of the disk area as specified previously. Second, we calculate the mass loss from the mass evolution of the disk (see Figure 5). Figure 10 compares the time evolution of the two measurements. Essentially, both show excellent agreement, confirming our methods to determine the evolution of the disk.

The mass loss remains negative for the majority of the simulation, with small exceptions of momentary mass increase especially in the later stages. On average, we have a mass-loss rate of the order of \(10^{-2}\)\(M_d\). The rate of mass loss, however, changes a lot, following a repeating pattern similar to the one appearing in the vertical mass flux from the disk surface, demonstrating that the large mass loss is due to the disk wind. Based on Figure 9, if we integrate over time, we find that out of a total of 85% of the disk mass lost during the simulation, approximately 73% is from the disk wind, 10% is from accretion to the black hole, and 2% is across the outer disk radius.

4.3. Outflow from the Black Hole Magnetosphere

The most prominent feature of our reference simulation (as visible in Figure 4) is the outflow that develops from the area around the black hole. It starts around \(t = 1000\) with the advection of magnetic flux toward the black hole. The field lines that enter the ergosphere are being twisted and turned along the toroidal direction, creating eventually a jet toward the polar direction, according to the BZ mechanism (Blandford & Znajek 1977). Up to \(t = 3000\), this jet has been fully developed, and it enters a quasi-steady state until the end of the simulation \((t = 9000)\), even though its strength still depends on the advection of the magnetic flux and, through that, on the accretion rate of the disk. The jet is identified by a parabolic-shaped funnel of high-velocity fluid that originates from the area around the black hole and moves almost parallel (in the later stages) to the symmetry axis toward the outer parts of our domain.

The jet funnel can be seen clearly in Figure 11, where we plot the \(z\)-component of the fluid frame velocity and the Lorentz factor at time \(t = 4000\). The jet seems to consist of fast-moving inner parts with \(\gamma \approx 1.8\), a moderately fast moving envelope with \(\gamma \approx 1.5\), and the outer part where the Lorentz factor values stay below \(\gamma = 1.3\). The fast-moving inner parts
seem discontinuous, and we can clearly distinguish 2–3 knots of high velocity in larger radii \( r > 50 \), while closer to the black hole the high values of Lorentz factor seem to have a more continuous distribution (see Figure 12).

We select four radii, \( r \approx 4, 12, 52, \) and 75, where high-velocity knots appear. In Figure 13 we see how the radial velocity, the mass flux, and the electromagnetic energy flux (Poynting flux) per solid angle are distributed along the polar angle in these radii. In general, the Poynting flux distribution follows the high-velocity areas, proving that the jet funnel has a strong electromagnetic component. The mass flux in the funnel area does not show a significant increase in comparison with the disk wind area and the disk, where the mass density is
Figure 12. Magnetohydrodynamic accretion–ejection structure close to the black hole, \( r < 15 \) at \( t = 4000 \). Shown is the Lorentz factor (left), the poloidal Alfvén Mach number (log scale; second from left), the plasma-\( \beta \) (log scale; second from right), and the magnetization \( \rho B^2 \) (log scale; right). The high Alfvén Mach number indicates flows that are dominated by kinetic energy, regardless of the highly magnetized area. The red semicircle marks the horizon, the yellow line marks the ergosphere, and the green line indicates the radius of the ISCO.

beginning, the flux remains almost constant, but after \( t = 3000 \) it drastically increases, indicating the development of a strong jet. This is mainly due to the advection of magnetic field and energy toward the black hole, along with the mass accretion from the disk. Beyond \( t = 6000 \)—due to the high disk mass loss—the variability in the accretion rate triggers the Poynting flux, leading to strong variations in the funnel and in most of the disk wind.

For the disk wind, at small radii the associated Poynting flux shows a steady increase with time. However, this is again an artifact due to integration area that cannot follow the bent geometry of the funnel flow. Also, the base of the funnel flow is partly extending beyond the chosen integration domain (limited to \( 25^\circ \)). It is thus not accounting for the initial funnel Poynting flux, but contributing to the Poynting flux we measure for the “wind.” This is indicated clearly in Figure 13. At larger distances, the Poynting flux remains at low levels, now following the true geometry of the disk wind. For the \( B_\phi \)-dominated disk wind the Poynting flux has very low but still positive values in the outer radii.

Table 1 shows the time-averaged Poynting flux measured at radius \( r = 100 \) for the three previously mentioned angular regions. As expected, the higher values of Poynting flux are detected in the jet funnel, about two times larger than the corresponding flux in the disk wind. The \( B_\phi \)-dominated disk wind also drives a Poynting flux about six times larger than the flux in the \( B_\phi \)-dominated disk wind. In total, the electromagnetic energy output of the disk is led mainly by the Poynting-dominated jet from the black hole, where we also detect the highest velocities.

This seems to contradict earlier results (Qian et al. 2018) indicating a disk wind substantially contributing to the total electromagnetic flux. We think that the reason for this difference is mainly the shorter live time of the simulation in Qian et al. (2018), in particular for the simulation with high spin. This is visible in Figure 14, where we see that for early times \( t \approx 500 \) the Poynting flux of the \( B_\phi \)-dominated wind (green curve) dominates the inner jet.

considerably higher, since the accelerated material consists primarily of floor density values.

Figure 12 shows the Lorentz factor, the poloidal Alfvén Mach number, the plasma-\( \beta \), and the magnetization over an area of \( 15R_\text{g} \) at time \( t = 4000 \). The highly magnetized funnel coincides with the high-velocity area of the jet. It starts as a sub-Alfvénic flow right outside of the ISCO; however, even though the area of the funnel is highly magnetized (plasma-\( \beta \approx 1, B^2/\rho \approx 1 \)), the flow is accelerated quickly to super-Alfvénic speed, indicating that it is dominated by kinetic energy.

4.4. Evolution of the Poynting Flux

We now examine the electromagnetic energy fluxes (Poynting flux) of our reference simulation. In Figure 14 we show the evolution of the integrated Poynting flux through a surface at radius \( r = 100 \). We further split our integration domain into the following three areas. The first region is between \( 0^\circ < \theta < 25^\circ \) and mainly covers the funnel region hosting the relativistic jet from the black hole magnetosphere. The disk wind area (covering larger polar angles) is split into two more regions (see also our Section 4.5): a region between \( 25^\circ < \theta < 65^\circ \), where the \( B_\phi \)-dominated disk wind evolves, and a region between \( 65^\circ < \theta < 80^\circ \), where the poloidal magnetic field dominates.\(^8\) The chosen separation does not exactly follow the direction of the funnel, as the geometry of the funnel flow changes with time. However, it is a good approximation for the average location of the funnel, especially in higher radii. Note that even though the majority of the (bent) funnel jet is inside the opening angle we have just defined, at \( r \approx 2–4 \) it is rooted closer to the equatorial plane, resulting in very low values of Poynting flux measured for the launching region and higher values for the disk wind regions (see also Figure 13).

The three phases of the disk evolution can also be seen in the evolution of the Poynting flux of the jet funnel. In the

\(^8\) Section 4.5 discuss the different types of disk wind extensively.
between the rotational axis and the jet funnel, the peaks in Lorentz factor and electromagnetic energy flux indicate the launching mechanisms, especially in the context of both AGNs and YSOs. Numerous works have investigated the launching mechanisms, especially in the nonrelativistic regime. In this section, we continue the analysis of the disk outflows, extending their study to (physically) larger grids of higher resolution.

4.5. The Accretion Disk Wind

The origin of accretion disk winds has been studied in the context of both AGNs and YSOs. Numerous works have investigated the launching mechanisms, especially in the nonrelativistic regime (Casse & Keppens 2002; Zanni et al. 2007; Sheikhnezami et al. 2012; Stepanovs & Fendt 2014). It has become clear since the seminal work of Ferreira (1997) that the magnetic resistivity is a key parameter for the investigation of the disk wind since it allows the gas to penetrate the magnetic field lines and thus allows for both (i) advection toward the black hole and (ii) mass loading the disk wind.

In a strong disk magnetic field, magnetocentrifugally accelerated outflows can be driven once the material is lifted from the disk plane into the launching surface, usually located around the magnetosonic surface. Qian et al. (2017, 2018) have extended the study of disk winds to the general relativistic regime. However, they have found that—in contrary to nonrelativistic disks—it is mainly the pressure gradient of the toroidal magnetic field that launches the disk winds, while the energy output by the disk wind can indeed be comparable to the BZ outflow launched by the BH. In addition (or rather as a consequence), disk winds from relativistic disks are quite turbulent and do not evolve in the smooth outflow structures that are known from nonrelativistic cases. In this section we continue the analysis of the disk outflows, extending their study to (physically) larger grids of higher resolution.

4.5.1. General Overview

In Figures 15 and 16 we present the velocity structure, the Alfvén Mach number, and the plasma-β for different areas of the disk wind. In order to emphasize the dynamic range of the disk wind, we restrict the velocity plots to $v_p < 0.1c$.

The plots of radial velocity (Figure 15, left, and Figure 16, top) nicely demonstrate the wind launching surface where the radial velocity changes sign, thus indicating the transition from accretion to ejection. The total poloidal velocity vectors start from inside the disk, where accretion dominates, and then continue across zero-velocity surface into the disk wind. The radial disk wind velocity increases as the wind leaves the disk surface, reaching up to $\mu' = 0.1c$ and following the magnetic field lines. Our vectors clearly demonstrate the connection between disk accretion and wind ejection.

In Figure 16 we show the poloidal Alfvén Mach number $M_{A,p}$. The Alfvén surface is located slightly above the disk surface (which we defined by $u' = 0$), implying that the fluid leaves the disk surface with sub-Alfvénic speed, $M_{A,p} < 1$. 

Figure 13. Comparison of the angular distribution of mass flux (red), Poynting flux per solid angle (green), and Lorentz factor (blue) for the reference simulation at $t = 4000$ at four radii, $r = 4, 12, 52,$ and 75. Negative mass flux indicates accretion toward the black hole. The BZ-driven jet funnel is clearly distinguished by the peaks in Lorentz factor and electromagnetic energy flux. For increasing radii, the mass flux increases, demonstrating the matter-dominated disk wind. In low radii, between the rotational axis and the jet funnel, the floor density material falls toward the black hole.

Figure 14. Evolution of the total Poynting flux for our reference simulation at radius $r = 100$. We split our domain into three regions. The first is between $0^\circ < \theta < 25^\circ$, expressing the Poynting flux from the relativistic jet funnel (red). The second is between $25^\circ < \theta < 65^\circ$, which includes the $B_T$-dominated disk wind (green). The third is between $65^\circ < \theta < 80^\circ$, which includes the $B_P$-dominated disk wind (blue).
In particular, the wind from the inner disk carries a so-called magnetic tower already in the altitude of \(0^\circ < \theta < 25^\circ\). This is a major difference from the nonrelativistic launching regime, in which the magnetic field dominates the poloidal magnetic field. We will first describe the inner wind component.

### 4.5.2. \(B_p\)-dominated Disk Wind

Considering the strength of the magnetic field components, we see that the toroidal field dominates the poloidal magnetic field. This is shown in Figure 16, where we plot the ratio \(|B_p|/|B_t|\). In particular, the wind from the inner disk carries a toroidal field 10 times larger than the poloidal component. We believe that this results from the fact that at this time the innermost part of the disk has completed a larger number of orbits: at time \(t = 4000\) and at \(r = 5\) we have almost 50 orbits, compared to about 18 at \(r = 10\) and only 10 at \(r = 15\). Hence, simply the twist of the originally poloidal magnetic field may induce such a strong toroidal field component. If the simulation would evolve further, we expect this area of a toroidally dominated magnetic field to grow along the disk.

We also need to compare the magnetic pressure to the gas pressure. This is done in Figure 16, where we present the distribution of plasma-\(\beta\) in the area of the disk. Inside the disk, we find plasma-\(\beta > 100\) (as prescribed by our initial condition), but as we move away from the disk surface, the plasma-\(\beta\) quickly starts decreasing to values below 0.01. This finding supports the idea of a magnetic-pressure-driven disk wind.

Interestingly, we find that the disk wind separates into two components considering the plasma-\(\beta\). There is an inner component of the disk wind that develops from the innermost part of the accretion disk (\(r \lesssim 10\)). This wind component has a rather high gas density and pressure, resulting in high poloidal plasma-\(\beta\) and low magnetization, \(B^2/\rho \sim 0.0001\). The second wind component originates from larger radii, and it is dominated by the poloidal magnetic field. We will first describe the inner wind component.

![Reference simulation sim0] Figure 15. Reference simulation sim0. Shown are the radial velocity (left), with superimposed contours of the vector potential (black lines), and the Alfvén Mach number (right), both at time \(t = 4000\).

However, it quickly accelerates to super-Alfvénic velocity. This is a major difference from the nonrelativistic launching simulations we have cited above, where the extension of the sub-Alfvénic regime is more comparable to the self-similar solution described by Blandford & Payne (1982), in which the flow in the area close to the disk is magnetically dominant, with matter accelerated along the field lines by the magnetic stress (or so-called magnetocentrifugally). The flow then consecutively passes the Alfvén and the fast-magneto sonic surface, before it becomes collimated by magnetic tension.

The fact that this mechanism may work as well for relativistic jets has been suggested by numerical simulations by Porth & Fendt (2010), however without considering the launching process out of the accretion disk. In our reference simulation, the picture is quite different, with an Alfvén surface much closer to the disk surface. The flow reaches super-Alfvénic speed of \(M_{A,p} > 5\) already in the altitude of \(z < 10\) from the disk midplane. Thus, we conclude that we do not find evidence for a large BP-driven region of the disk wind, and the outflow is most probably driven by the magnetic pressure gradient of the toroidal field, thus as a so-called magnetic tower (Lynden-Bell 1996).

### Table 1

| Run | \(\alpha\) | \(\eta_0\) | \(\langle M_{\text{vis}}\rangle\) | \(\langle M_{\text{B}}\rangle\) | \(\langle M_{\text{p}}\rangle\) | \(\langle M_{\text{EM,fun}}\rangle\) | \(\langle M_{\text{EM,P}}\rangle\) | \(\langle M_{\text{EM,fun}}\rangle_{\text{r}}\) | \(\langle M_{\text{EM,fun}}\rangle_{\text{t}}\) |
|-----|-----------|------------|----------------|-----------------|----------------|----------------|----------------|----------------|----------------|
| sim0 | 0.9 | 0.001 | \(-0.75\) | 9.72 | 4.15 | 1.02 (25) | 2.40 (58) | 0.73 (16) | 4.89 | 2.38 | 0.38 |
| sim1 | 0 | 0.001 | \(-1.59\) | 6.20 | 1.83 | 0.26 (14) | 1.02 (56) | 0.55 (30) | 0.44 | 0.55 | 0.23 |
| sim2 | 0.5 | 0.001 | \(-1.57\) | 7.51 | 2.88 | 0.67 (23) | 1.57 (54) | 0.64 (22) | 2.87 | 1.29 | 0.32 |
| sim3 | 0.9 | 0.001 | \(-1.27\) | 5.77 | 4.88 | 0.96 (20) | 3.48 (71) | 0.44 (9) | 3.00 | 4.52 | 0.21 |
| sim4 | 0.9 | 0.01 | \(-0.53\) | 5.61 | 3.17 | 0.79 (25) | 1.94 (61) | 0.44 (14) | 1.82 | 1.93 | 0.19 |
| sim5 | 0.9 | 0.0001 | \(-1.24\) | 12.8 | 3.70 | 1.43 (39) | 1.81 (49) | 0.46 (12) | 4.11 | 1.93 | 0.24 |
| sim6 | 0.9 | \(10^{-10}\) | \(-1.24\) | 11.6 | 3.04 | 1.37 (45) | 1.33 (44) | 0.33 (10) | 4.17 | 1.72 | 0.25 |

Note. The average mass fluxes in units of \(10^{-5}\) are measured over the whole simulation period and are normalized by the initial disk mass. The average Poynting fluxes are in code units of \(10^{-5}\). The average vertical wind mass flux is integrated along the radius vector along \(80^\circ\) or \(100^\circ\) up to \(r = 100\). The average radial wind flux is integrated along a spherical surface at \(r = 100\). Note that simulations sim3, sim5, and sim6 end before \(t = 6000\). The columns show from left to right the simulation run ID; the spin parameter \(\alpha\); the maximum diffusivity \(\eta_0\); the average accretion rate at \(r = 2\) \(\langle M_{\text{acc}}\rangle\); the average vertical mass flux \(\langle M_{\text{B}}\rangle\); the total radial mass flux \(\langle M_{\text{r}}\rangle\) \((0^\circ < \theta < 80^\circ)\); the average mass flux in the jet funnel \(\langle M_{\text{fun}}\rangle\) \((0^\circ < \theta < 25^\circ)\); the average mass flux in the \(B_p\)-dominated disk wind \(\langle M_{\text{fun}}\rangle\) \((25^\circ < \theta < 65^\circ)\); the average mass flux in the \(B_p\)-dominated disk wind \(\langle M_{\text{t}}\rangle\) \((65^\circ < \theta < 80^\circ)\); the electromagnetic energy flux in the jet funnel \(\langle \dot{E}_{\text{EM,fun}}\rangle\) \((0^\circ < \theta < 25^\circ)\); the electromagnetic energy flux in the \(B_p\)-dominated disk wind \(\langle \dot{E}_{\text{EM,fun}}\rangle\) \((25^\circ < \theta < 65^\circ)\); the electromagnetic energy flux in the \(B_p\)-dominated disk wind \(\langle \dot{E}_{\text{EM,P}}\rangle\) \((65^\circ < \theta < 80^\circ)\); Values in parentheses show the percentage of each individual radial mass flux over the total radial mass flux.
similar structure for the overall disk wind in high plasma-β simulations. These outflows are dominated by the toroidal magnetic field component, also known as tower jets (see below), and are accelerated by the vertical toroidal magnetic field pressure gradient. However, in our simulations we notice that this turbulent outflow layer has a certain, rather narrow opening angle. If we assume that the extent of this layer defines a characteristic length, we may also assume that the extension of this structure in the $f$-direction may be similar, possibly hinting to a series of outflow tubes around the disk.

Interestingly, Britzen et al. (2017) have recently suggested that such turbulent loading of jet channels may happen in M87, leading to large-scale episodic wiggling of the overall jet structure.

In Figure 16 (top right) we show the $z$-component of the velocity, where we can distinguish a number of “branches” with values higher than in the adjacent area. These branches are actually part of the $B_{\phi}$-dominated disk wind. They seem to stay connected to the surface of the disk from where they are originally launched and then continue through the $B_{\phi}$-dominated wind following the poloidal magnetic field lines. The footpoint of the branches coincides with highly magnetized disk areas. This might explain the acceleration within the branches—on the other hand, when this material enters the $B_{\phi}$-dominated wind, the plasma-β increases without weakening the acceleration. We note that the strong $V_z$ component pushes the disk wind material toward the boundaries of the funnel outflow. As for an alternative scenario, we may think of a magnetic-pressure-driven radial outflow that drags the poloidal field with it, thus stretching it into a radially aligned poloidal field distribution.

4.5.3. $B_{\phi}$-dominated Disk Wind

We now discuss the second wind component that originates in the outer, main body of the disk. Here, for radii $r \geq 10$, the $|B_{\phi}/B_p|$ ratio decreases with radius and the poloidal field starts to dominate. This outer wind becomes launched almost parallel to the magnetic field lines (see velocity streamlines and poloidal field lines in Figure 16), and it retains that direction as well for larger distances. The vertical velocity component is substantially lower compared to the inner disk wind, implying a weaker acceleration despite the higher magnetization. When comparing the local escape speed with the local poloidal velocity of the disk wind, we find that the disk wind is launched
with sub-escape velocity. However, the wind becomes further accelerated to \( u_{\infty} > u_{\text{esc}} \) and becomes eventually fast enough to escape the gravity of the black hole.

In Figure 16 (first panel), we notice that in the area where the disk wind develops the wind tends to follow the radial direction in general. However, in Section 4.2 above we quantified the launching of the disk wind as mass flux escaping the disk surface in the polar direction (\( \theta \)-component of the velocity). Thus, after being launched vertically from the disk surface, the wind further develops into a kind of radial outflow. This overall picture connecting between the launching in polar direction and the radial outflow can be verified by calculating the mass fluxes through the respective boundaries.

### 4.5.4. Connecting the Vertical and Radial Disk Wind

Following the considerations of the disk wind toward the end of Section 4.2, we show in Figure 17 the disk wind mass fluxes, only in this case we are interested in the situation at smaller radii, where the disk outflow is stronger. We calculate the mass fluxes vertical to surfaces of constant opening angle of \( \theta = 80^\circ, 100^\circ \) that approximate the opening angle of the initial disk density distribution. We integrate the mass fluxes in the range \( r = [4, 100] \), where we also separate between infall (motion toward the disk surface) and outflow (motion away from the disk surface), thus providing the net vertical fluxes. We find that the disk wind seems to increase for \( 0 < t < 3000 \) while it decreases for \( 3000 < t < 6000 \); overall we measure an average mass flux of \( \dot{M} = 9.72 \times 10^{-5} M_{\odot} \). The variations in the mass flux during the second phase are much stronger; this is consistent with a similar behavior in the accretion rate (see Section 4.2 and Figure 6).

A similar behavior is observed in the radial mass fluxes. These increase or decrease with radius, depending on the phase during the simulation. Simultaneously, this is visible as an increase or decrease, respectively, in the mass load of the disk wind. Note that the latter we can also observe by measuring the difference in the two mass fluxes. The overall time-averaged mass flux is \( \dot{M} = 4.56 \times 10^{-5} M_{\odot} \) across a spherical surface at \( r = 100 \). The difference in the two mass fluxes is deposed as mass in the area of the disk wind increasing its density. Taking into account this mass sink, as well as all mass fluxes through the surfaces of the integration area, we find a good agreement between the radial and the disk wind fluxes for small time intervals. The remaining difference is due to the jet funnel that is constantly loaded by the floor model for the density and that naturally contributes to the radial mass fluxes and also increases the mass load in the radial wind.

Our detection of a \( B_{z} \)-dominated disk wind confirms the results of Qian et al. (2018), who interpreted their results in terms of a tower jet (Ustyugova et al. 1995; Lynden-Bell 1996). However, the whole disk wind in Qian et al. (2018) is entirely dominated by the \( B_{z} \), while in our simulation it is restricted to the disk wind from the inner disk only. As our new simulations have a higher resolution, Qian et al. (2018) may have not been able to resolve the inner part of the disk wind properly.

#### 4.5.5. Magnetic Reconnection and Ohmic Heating

Since the disk evolves in a resistive environment, we expect the generation of ohmic heating, which will affect the internal and magnetic energy in the disk. As we do not use radiative transfer, we cannot directly compare the energetics of ohmic heating with the emitted radiation.

However, we can attempt an estimation of the generated heating. For the reference simulation, we calculated an approximation of ohmic heating as \( \eta j^2 \) and compared it with the internal and magnetic energy of the fluid. We separated the area into two parts—the first one is from \( r = 5 \) to \( r = 20 \), and the second one is from \( r = 20 \) to \( r = 50 \). Since the resistivity is concentrated to the accretion disk (and thus the ohmic heating), we also constrain the area between \( 5^\circ \) above and below the equatorial plane. The ohmic heating is mostly generated from the inner part of the disk, as the magnetic field gradients \( j \sim B \) are largest over there.

We find that up to time \( t = 5000 \) ohmic heating generates a total energy of \( 1.5 \times 10^{-4} \) (in code units). This is somewhat higher than the total magnetic energy in this disk area but substantially lower than the internal energy of the disk. At larger radii, from \( 20 < r < 50 \) the ohmic heating rate is even lower, making it overall negligible in comparison with the magnetic and internal energy.

Another physical mechanism that contributes to the heating of our fluid is magnetic reconnection. It has been shown (de Gouveia Dal Pino & Lazarian 2005; de Gouveia Dal Pino et al. 2010) that in AGNs the magnetic reconnection episodes that occur mostly in the inner disk and the black hole magnetosphere can heat up the disk material and at the same time accelerate the ejected disk wind.

### 5. Comparison Study

We will now compare our reference run \( \text{sim0} \) with a number of simulations that apply different physical parameters such as black hole spin, magnetic field strength, or magnetic diffusivity (see Table 1).

#### 5.1. Accretion–Ejection and Black Hole Rotation

We now discuss how the dynamical evolution of accretion–ejection interrelates with the black hole rotation, i.e., the Kerr parameter \( a \).

We first concentrate on the disk accretion. Figure 18 shows the disk accretion rates at \( r \approx 2 \) for the simulation runs \( \text{sim1}, \text{sim2}, \text{sim0} \), each normalized with the mass of the respective initial disks.

While for \( \text{sim0} \) the accretion rate in the first stages of the evolution \( (t \in [0, 3000]) \) is constant and very low, for...
slower-rotating black holes the accretion rate shows a noticeable increase. Also, this first stage, which looks different from the later evolution, lasts longer in the case of \( a = 0.9 \). We think that this is due to the fact that the horizon \((r = 2)\) and the ISCO \((r = 6)\) are located closer to the initial disk radius. Therefore, it takes less time to bring disk material to the ISCO, from which it falls to the horizon. At later stages, all simulations show a similar behavior, with only the accretion spikes in \( \text{sim}0 \) being slightly stronger.

On average, for the duration of the simulation, the normalized accretion rate at \( r = 2 \) for the Schwarzschild black hole is slightly higher, \( \langle M \rangle = -1.59 \times 10^{-5} M_{\text{disk},0} \), while for the case of \( a = 0.9 \) we find \( \langle M \rangle = 7.49 \times 10^{-6} M_{\text{disk},0} \). Specifically, the three systems accrete 18.4\%, 17.3\%, and 11\% of their initial disk mass into the black hole for the duration of the simulations.

In Figure 19 we compare the disk wind that is launched from the disk surface. The disk mass flux is in general positive with a few exceptions\(^9\), meaning that there is a substantial mass injection from the disk into the outflow.

Following the same method as described toward the end of Section 4.5.4, we measure a normalized mass flux for the disk wind of \( \langle \dot{M} \rangle = 6.2 \times 10^{-5} M_{\text{disk},0} \) for the case of \( a = 0 \), a flux of \( \langle \dot{M} \rangle = 7.41 \times 10^{-5} M_{\text{disk},0} \) for the case of \( a = 0.5 \), and a flux of \( \langle \dot{M} \rangle = 9.72 \times 10^{-5} M_{\text{disk},0} \) for the case of \( a = 0.9 \). This implies that the three accretion–ejection systems accumulate a mass loss of 49\%, 62.3\%, and 80.7\% of their initial disk mass by the disk wind. The cases \( a = 0 \) and \( a = 0.9 \) differ by almost 30\% in the disk wind mass flux. For the radial fluxes there is a similar increase by 178\% between the simulations applying \( a = 0 \) and \( a = 0.9 \) (see Table 1). Thus, as an overall trend we find that the disk wind mass flux increases for higher black hole spin.

We understand that this is due to the ejection of mass that is launched from the innermost radii of disk accretion for high \( a \) (see Figure 6, middle panel). These ejections, and thus positive radial mass fluxes inside the disk, do not appear for the cases of low spin \( a = 0 \) and \( a = 0.5 \), for which accretion dominates, and which result in an overall lower disk wind ejection rate (see in Figure 19). There is also the interplay between the evolution of the disk structure and the distribution of magnetic diffusivity. As the ISCO radius is affected by the Kerr parameter, the disk is located completely inside the high diffusivity area for \( a = 0 \), while part of the inner radii has lower diffusivity for the case of \( a = 0.9 \).

Note that the radius \( r = 3 \) is just outside the ISCO for simulation \( \text{sim}0 \) but inside the ISCO for \( \text{sim}1 \) and \( \text{sim}2 \), which we think explains why no ejection is visible in the case of the latter two simulations. In order to check this hypothesis, we also measured the mass flux at \( 1R_g \) and \( 2R_g \) outside of the ISCO for each of our simulations. Only in simulation \( \text{sim}0 \) does a positive mass flux from this radius appear, subsequently contributing to the increased mass flux we measure in the disk corona.

We further investigate the radial mass fluxes through a surface of radius \( r = 100 \). We find that the increase in the mass flux is much higher than in the vertical fluxes. We have also analyzed the radial mass flux of the disk wind by comparing the fluxes in three domains of the outflow (see Figure 19 for numerical values). The innermost flow area is from \( 0^\circ \) to \( 25^\circ \), and it indicates the mass flux in the Poynting-dominated jet. The adjoining area from \( 25^\circ \) to \( 65^\circ \) covers the \( B_\phi \)-dominated wind launched in the innermost disk. The third domain from \( 65^\circ \) to \( 80^\circ \) contains the mass flux from the \( B_\theta \)-dominated disk wind. Obviously, we also include the fluxes from the lower hemisphere.

We recognize that our choice for the limits in the polar angle will not always coincide perfectly with the physical part of the flow we want to study. This holds especially in the earlier and later times of the simulations when both the jet and the disk wind are strongly evolving, either further being developed (early) or dying off because of the disk mass loss (late). For the Poynting-dominated jet, the floor density model that dominates this area obviously determines most of the mass flux.

Comparing the simulations, we find that the relative contribution of the \( B_\theta \)-dominated disk wind to the overall mass flux is similar for simulations \( \text{sim}1 \), \( \text{sim}2 \), and \( \text{sim}0 \)—even though in absolute values the wind mass flux increases with black hole spin. The relative contribution of the \( B_\phi \)- and \( B_\theta \)-dominated disk winds, however, depends on the black hole spin. In the case of a Schwarzschild black hole the \( B_\phi \)-dominated disk wind contributes 65\% to the total disk wind mass flux, while for the case of \( a = 0.9 \) the contribution is at 77\%. For the counterrotating black hole the contribution increases to 89\%, while it shows the strongest wind also in absolute values. We conclude that the black hole rotation not only increases the disk outflow mass flux in general but also contributes substantially in the \( B_\phi \)-dominated disk wind as it is generated from the inner part of the disk.

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\(^9\) Most of the negative flux occurrences appear in the late stages of \( \text{sim}1 \).
Figure 20. Poynting flux and black hole spin. Comparison of the radial Poynting fluxes at \( r = 100 \) for simulation runs applying a Kerr parameter \( a = 0 \) (sim1, red), \( a = 0.5 \) (sim2, green), \( a = 0.9 \) (sim0, blue), and \( a = -0.9 \) (sim3, magenta). Note that simulation sim1 is scaled 10 times lower than the others.

Finally, we compare the Poynting flux in our simulations. Figure 20 shows the time evolution of the Poynting flux through a surface at \( r = 100 \) in the area of the funnel flow for the four different cases of black hole spin. There is a clear trend that the Poynting flux from the jet funnel increases with spin parameter. The highest Poynting flux appears in the reference simulation with \( a = 0.9 \). For simulation sim1 the flux is substantially (factor 10) lower than for the simulation with a rotating black hole. Also, in sim1 the absence of black hole rotation results in a relatively higher flux from the disk wind. A question arises on what drives the Poynting flux from a nonspinning black hole. We believe that this Poynting flux is driven by the rapidly rotating (infalling) material that is just outside the horizon in a fashion similar to the BZ mechanism. The magnetic field lines are twisted by the rotating disk, creating a jet with smaller electromagnetic energy flux.

5.1.1. A Counterrotating Black Hole

We now investigate how a counterrotating black hole affects the overall jet launching. It has been suggested that the efficiency of the BZ process in prograde systems is slightly higher compared to retrograde black hole—torus systems (Tchekhovskoy & McKinney 2012). Here we extend this analysis for resistive GR-MHD and for thin accretion disks. We have set up simulation run sim3 with a negative Kerr parameter \( a = -0.9 \), but otherwise identical to our reference simulation.

A first comparison shows the accretion rate at radius \( r = 2 \) (see Figure 18) and the disk wind mass flux (see Figure 19) for both simulations. For \( a = -0.9 \) the ISCO is located at \( r \approx 8.7 \). As a result, since the inner radius of the initial disk is located farther in at \( r = 7 \), accretion toward the black hole starts immediately with a sudden infall of the disk area inside ISCO. Furthermore, the disk immediately looses a substantial fraction of mass, about 30\% until \( t = 300 t_g \). Afterward, the disk structure adjusts such that its inner radius remains outside the ISCO and the normal—slow—accretion begins as soon as angular momentum is removed from the disk material.

All simulations start with an initial setup with \( B_\phi = 0 \). However, by the rotation of the footpoints of the field lines (accretion disk or spacetime) a toroidal field is induced. In the prograde simulations, the \( B_\phi \) in the disk wind and the black hole magnetosphere have the same sign since both the disk and the black hole rotate in the same direction. At the equatorial plane \( B_\phi \) changes sign (see Figure 21, left panel), since the magnetic field lines are anchored at infinity.

In contrast, for the case of retrograde black hole rotation, simulation sim3, the \( B_\phi \) in the black hole magnetosphere and in the outflow launched from there is induced with the opposite sign compared to the disk wind (see Figure 21, right panel), resulting in another boundary layer with \( B_\phi = 0 \) appearing between the jet funnel and the disk wind.

While we expect (and find) the black-hole-driven outflow to have a different sign for negative Kerr parameter, we would expect the disk wind to have \( B_\phi \) with the same sign for positive and negative Kerr parameter, again with \( B_\phi = 0 \) and a change of sign at the equatorial plane. However, to our surprise, we find that in the disk area close to the inner disk radius the \( B_\phi \) changes sign three times (instead of only once; see Figure 21). In fact, the \( B_\phi \) in the wind above the disk area is directed opposite to the \( B_\phi \) below the disk surface.10 Along the disk surface \( B_\phi = 0 \).

This also affects the poloidal component of the magnetic field (mainly the radial component), as it is visible purely from the shape of the field lines. The change of sign in \( B_\phi \) close to the equatorial plane is intrinsically connected to the type of

10 Of course similar for the upper and lower hemisphere, respectively.
accretion: Figure 22 shows the radial velocity for simulation sim3 and clearly indicates that inside the disk some material is moving outward, while accretion happens along the surface layers of the disk. For the case of prograde rotation, accretion is mainly along the equatorial plane. This unexpected behavior, however, does not affect the overall accretion rate.

For simulation sim3 with $a = -0.9$ we find—similar to the prograde case—an outgoing Poynting flux, which is indicative of Blandford–Znajek launching. The Poynting flux in the funnel area increases with time, with a time average value of $\langle E_{\text{EM}} \rangle = 3 \times 10^{-3}$ at radius $r = 100$. For comparison, the Poynting flux at $r = 100$ for the prograde simulation sim0 is $\langle E_{\text{EM}} \rangle = 4.89 \times 10^{-3}$. Furthermore, the Poynting flux from the disk wind appears to be stronger than the one from the funnel having a time average of $\langle E_{\text{EM}} \rangle = 4.52 \times 10^{-3}$ at $r = 100$. We do not find significant differences in the electromagnetic energy emitted within the funnel flow between the prograde and retrograde simulations.

It would have been interesting to follow the retrograde setup for a longer time, but the simulation stopped at $t \approx 5500$, most probably due to the high mass loss and also the complex magnetic field and velocity structure.

Although we find for the retrograde black hole rotation a few remarkable and also unexpected features that can be astrophysically interesting, we do not want to overinterpret, as we think that the retrograde case is not likely realized in nature. Retrograde black hole rotation may be realized by galaxy mergers with accompanied binary black hole mergers, but not from pure disk accretion. Similarly, counterrotating black hole–disk systems may be expected from specific initial conditions for neutron star mergers and thus may affect the subsequent gamma-ray burst activity.

### 5.2. Impact of Magnetic Diffusivity

MRI is thought to be the main driver of turbulence in accretion disks (Balbus & Hawley 1991, 1998). The feasibility of MRI has been demonstrated also in GR-MHD simulations (Penna et al. 2010; McKinney et al. 2012). Overall, turbulence results in a dissipative effect for the magnetic field, which we express through a mean magnetic diffusivity, in analogy to the $\alpha$-effect for turbulent viscosity (Shakura & Sunyaev 1973).

In contrast with ideal MHD, the disk material is now able to move across the magnetic field (lines) while accreting toward the black hole. The advection of magnetic flux is reduced owing to the weaker coupling between magnetic field and mass. It is thus worth investigating the effect of diffusivity on the accretion–jet/ejection mechanism and the launching of outflows and jets. As described above, we have implemented a background diffusivity fixed in time and space that mainly follows the disk structure (see Section 3.4).

In the following we focus on varying the strength of the disk magnetic diffusivity. Further studies considering the scale height or the radial profile need to be done, as it has been worked out for nonrelativistic studies of jet-launching simulations (see, e.g., Sheikhnezami et al. 2012; Stepanovs & Fendt 2014).

We have run three further simulations that are identical to our reference simulation but consider $\eta_0 = 10^{-2}$ (sim4), $10^{-4}$ (sim5), and $10^{-10}$ (sim6), respectively (see Table 1). We observed that a higher magnetic diffusivity stabilizes the simulation run, and simulation sim4 runs until $t = 15,000$. Simulations with lower diffusivity levels were terminating earlier, however still providing enough information for a comparison.

In Figure 23 we compare the accretion rate at radius $r = 2$ for different levels of magnetic diffusivity. For simulation sim4 with the highest level of diffusivity we notice an almost constant (in comparison with the other simulations) accretion rate without any spikes. Still some spikes start appearing after $t = 9000$ when we plot the long-term accretion evolution of sim4, even though the background accretion does not change much. Overall, for this simulation we cannot identify the three phases of accretion rate we found in the reference simulation, even with the longer simulation time.

For lower levels of diffusivity the evolution of the accretion rate has more similarities to simulation sim0. We identify similar phase changes to those we detected in our reference simulation; however, unfortunately the simulations stop before they reach a timescale that is comparable to that of the reference simulation. Even in this case, though, for sim5 the second phase starts at $t \approx 1600$, while for sim6 it starts at $t \approx 1100$; however, it is not as clear as in the reference simulation.

For the vertical flux of the disk wind we observe a similar behavior—a larger disk wind mass flux resulting for lower levels of diffusivity (see Figure 24). It therefore seems that high diffusivity reduces the efficiency for the magnetic field to a launch disk wind. This is straightforward to understand and has been observed in nonrelativistic simulations (Sheikhnezami et al. 2012): for a magnetic driving of outflows (Blandford–Payne or magnetic pressure driven) a strong coupling between magnetic field and matter is essential.
For the radial mass flux we detect a different behavior. A high radial mass flux appears for the reference simulation with $\eta_0 = 0.001$, while for both higher and lower diffusivity levels the mass flux decreases to approximately similar levels. The area where we find the $B_{\nu}$-dominated wind has a lower diffusivity level than the equatorial plane, but for simulation $\text{sim4}$ it is still significant enough to weaken the wind. The area of the $B_{\nu}$-dominated wind increases with the increase of diffusivity.

Finally, we investigate the Poynting fluxes for the different levels of diffusivity. Figure 25 shows the Poynting flux through the jet funnel at radius $r = 100$ for various $\eta$. The flux increases in time for all cases; however, comparing simulation $\text{sim4}$ (largest $\eta$) with the reference simulation, the increase is much slower. Simulations $\text{sim5}$ and $\text{sim6}$ show again very similar behavior, following the trend we observed in the accreting and vertical mass fluxes. Also, in the case of $\text{sim4}$ the flux from the disk wind is slightly stronger than the flux from the jet funnel.

The previous findings hint at preferred levels of diffusivity (or a preferred level of turbulence) that supports the launching of a disk wind. For higher diffusivity, the coupling between matter and field may not be efficient enough for launching, while for lower levels of diffusivity the mass loading becomes inefficient.

What is the mechanism behind these findings of a threshold value for the magnetic diffusivity of $\eta = 10^{-3} \ldots 10^{-2}$ where the flow becomes smooth and never MAD-like? We believe that it is the interplay between magnetic reconnection, magnetic diffusion, and ohmic heating that governs the disk structure at these scales. Magnetic reconnection destroys magnetic flux that is needed to launch strong outflows. It also generates turbulence to the flow. We would thus expect a high resistivity to weaken the outflow launching. On the other hand, a higher resistivity enables a more efficient mass loading of the outflow. Thus, a smaller resistivity would decrease the mass load of the outflow but potentially may produce outflows with higher speed (for the same magnetic flux available). Ohmic heating of the launching area would, in contrast, increase the mass loading (in classic MHD steady-state theory the mass load is determined by the sound speed at the launching radius).

Overall, our simulations seem to follow these trends. For low resistivity, resistive mass loading becomes less efficient, assisted by low ohmic heating. For high resistivity, reconnections weaken the outflow. For a critical resistivity in between, outflow launching becomes most efficient.

\section{6. A Black Hole Shadow?}

Motivated by the recent detection of a black hole shadow in the jet-launching core of M87, we here discuss a few features of our simulation that can possibly be interrelated with these new findings. As we do not consider radiation in our simulations, we cannot provide emission maps of direct or lensed radiation. However, we can estimate the opacities in our disk outflow system and thus the visibility of the innermost central region around the black hole. Obviously, if the black-hole-surrounding medium is opaque, a black hole shadow cannot be seen from a distant observer. This is interesting because the accretion structure and the metric depend on quantities that are not really known, that is, the black hole spin, the black hole mass, and the accretion rate. Our question here is whether we could derive some general features that allow us (or not) to observe a signal lensed into the photon sphere as claimed for M87.

Considerably one of the most critical points in detecting a black hole shadow is the structure of the accretion flow close to the black hole. Here we expect to find differences when considering a (turbulent) magnetic diffusivity for accretion or ideal MHD. One reason is that the coupling between matter and magnetic field is different. The other reason is that physical resistivity allows for reconnection of the magnetic field. This is a particularly strong effect along the equatorial plane close to the horizon—the reason being the advection of large-scale magnetic flux from the disk that connects to the horizon, leading to a field reversal across the midplane.

We therefore calculate the surface density (in code units or gravitational units, denoted by the overline),

\begin{equation}
\Sigma = \int_{\Sigma(r)} p(\rho) d\Sigma,
\end{equation}

along the accretion stream connecting the inner disk and the horizon for our different models. Figure 26 shows the (axisymmetric) surface density distribution close to the horizon. There is a clear trend toward higher surface densities (yellowish, in code units) for increasing black hole spin parameter when compared at a certain radius.

For example, at $r = 4$ the surface density increases by a factor of 100, when comparing spin parameters of $a = 0, 0.5, 0.9$, respectively (see Figure 26). The optical depth would differ by the same factor, if (!) assuming the same disk density scaling $\rho_0$ for all simulations. Note, however, that for different spin also the metric changes and comparing physical variables at a fixed radius is not necessarily meaningful. For $a = 0.9$ the ISCO is inside $r = 4$ and the yellow structure in Figure 26 still resembles a rotating accretion disk, while for a Schwarzschild black hole this radius is inside the ISCO.

We may thus better compare the surface densities at the radius of the respective photon spheres (dashed line) where most of the lensed radiation originates. Essentially, we find that $\Sigma$ is similarly small for all three spin parameters (blueish colors). Also, the $\Sigma$ inside the ISCO is small for all three cases shown (see blueish colors inside the dotted circles). That again implies that the photon orbit—and thus the photons lensed into this orbit—could be visible in all these cases, supposing that the optical depths are not extremely high (see below). Note that here we do not consider radiation from the disk or the outflow, but only the visibility of a hypothetical emission of photons that were lensed into the photon ring. That photon ring would...
always be located within the radius of disk accretion and thus potentially in a low-density area. In order to calculate the opacities of the disk–jet material, we need to know the physical densities. However, since the disk mass in our simulations acts as a test mass on the Kerr metric, we cannot specify the density in physical units without further assumptions. Rescaling our normalized accretion rate (in code units)\(^{11}\) of \(\dot{M} \approx 0.1\) to astrophysical units assuming typical AGN accretion rates of

\[
\dot{M} \equiv \rho_0 R_g^2 \frac{c M}{\sin i} \approx 0.01 \frac{M_\odot}{\text{yr}}^{-1},
\]

we can constrain the disk densities to

\[
\rho_0 = 4 \times 10^{-14} \frac{\text{g}}{\text{cm}^3} \left[ \frac{M}{0.01 \frac{M_\odot}{\text{yr}}^{-1}} \right] \left[ \frac{M_{BH}}{5 \times 10^9 \frac{M_\odot}{\text{yr}}} \right]^{-1} \left( \frac{\dot{M}}{0.1} \right). \tag{29}
\]

The rescaled astrophysical density \(\rho = \rho_0 \bar{\rho}\), where \(\bar{\rho}\) in code units follows from our simulations. For comparison, for M87 the EHT Collaboration derived an accretion rate of \(\dot{M} \approx 2.7 \times 10^{-3} \frac{M_\odot}{\text{yr}}^{-1}\) (Akiyama et al. 2019b) assuming a black hole mass of \(M = 6.2 \times 10^9 \frac{M_\odot}{\text{yr}}\). Earlier estimates suggested \(\dot{M} \approx (0.2-1) \times 10^{-3} \frac{M_\odot}{\text{yr}}^{-1}\) (Feng et al. 2003).

We may now estimate the optical depth of the disk material applying opacities for Thomson scattering \(\kappa(r) = \sigma_T n_e(r) = \sigma_T p(r)/m_p\). We integrate vertically over a sufficiently large geometrical scale height \(h(r)\) across the midplane and find the optical depth of the inner accretion stream:

\[
\tau(r) = \int_{-h(r)}^{h(r)} \kappa(r) d\Sigma = \frac{\sigma_T}{m_p \rho_0} \int_{-h(r)}^{h(r)} \bar{\rho}(r) d\Sigma
\]

\[
= 0.1 \Sigma(r) \left( \frac{\rho_0}{4 \times 10^{-14} \frac{\text{g}}{\text{cm}^3}} \right) \left( \frac{M_{BH}}{5 \times 10^9 \frac{M_\odot}{\text{yr}}} \right)
\]

\[
= 0.1 \Sigma(r) \left( \frac{M}{0.01 \frac{M_\odot}{\text{yr}}} \right) \left( \frac{M_{BH}}{5 \times 10^9 \frac{M_\odot}{\text{yr}}} \right) \left( \frac{\dot{M}}{0.1} \right)^{-1}, \tag{30}
\]

with \(\Sigma(r)\) again being the surface density in code units (see Figure 26).

\(^{11}\) Note that the accretion rates given in the previous figures \(\propto 10^{-4}\) are normalized by the initial disk mass, in order to be able to compare different simulations.

Radiation that is lensed into the photon sphere could then be observed if the accretion stream at this radius is optically thin, \(\tau \ll 1\). Taking the \(\Sigma \approx 10^{-2}\), we find that for the area close to the photon orbit the latter is actually the case for all three spin parameters and for the normalization used in Equation (30). Considerably higher accretion rates, say, \(\dot{M} \approx 10 \frac{M_\odot}{\text{yr}}^{-1}\), may lead to an opaque situation, which is also the case for smaller black hole masses, say, \(M \approx 10^8 \frac{M_\odot}{\text{yr}}\).

What the image of the very central region actually looks alike is beyond the scope of our paper, as we do not consider the radiation from the gas in our simulations. Thus, what we see as a large “hole” in the surface density toward the center of the Schwarzschild simulation (see Figure 26, left panel) may actually be bright owing to radiation from hot gas falling toward the horizon. How that emission would compare to the lensed signal at the photon sphere we cannot tell. Simulations by the EHT Collaboration suggest that the signal from the photon sphere is dominating.

A remaining question is the visibility of the innermost region when observed from high inclination, meaning along viewing angles close to the rotational axis. In the case of M87 the line of sight is about 20°. In our models, as well as in almost all simulations in the literature, this is the area of the funnel flow of low (floor) density. For our reference simulation, when integrating for the surface density along radial directions, we find the distribution shown in Figure 27.

We find that when implying the same normalization as above, the optical depth along the line of sight close to rotational axis is low. As expected, the disk material blocks the radiation along light paths close to the equatorial plane. Interestingly, we also see that the disk wind clearly contributes to the opacity and may thus, depending on the physical density (meaning accretion or ejection rate), also block the view toward the central black hole.

Overall, we find that for reasonable disk densities, \(\rho_0 = 10^{-14} \frac{\text{g}}{\text{cm}^3}\), the line of sight toward the central black hole and its shadow is not blocked by the outflow material for viewing angles below 25° to the rotational axis. For a line-of-sight inclination larger than 25° a massive disk wind blocks the view toward the center, where \(\tau(\theta) > 1\).

7. Summary

In this paper, we have extended the newly developed resistive GR-MHD code eHARM (Qian et al. 2017) to the parallel version HARM3D in order to apply our models of jet launching from thin accretion disks surrounding a black hole to
longer timescales and larger spatial scales, also considering a higher numerical resolution.

In our model the disk is threaded by inclined open poloidal field lines. In general, our simulation results demonstrate how the magnetic field strength, the disk magnetic diffusivity, and the black hole spin influence the MHD launching of disk winds and the Blandford–Znajek jet from the black hole. Essentially we are able to compare the strength and power of both jet components for different Kerr parameters. In the following we summarize our results.

1. Our implementation of resistivity is based on Qian et al. (2017) following Bucciantini & Del Zanna (2013). We tested the code by simulating the decay of the magnetic field and comparing the evolution with the analytic solution. We find a perfect match for diffusivities $\leq 10^{-2}$.

2. As a test for the model setup we run a set of GR-MHD simulations in the mildly relativistic limit. Here only the disk evolution was simulated with the inner grid boundary far from the black hole. Strong disk outflows were found, similar to the magnetocentrifugally driven outflows observed in the nonrelativistic simulations in the literature.

3. As a reference simulation we applied the code for a setup considering a black hole with Kerr parameter $a = 0.9$, together with a disk magnetic diffusivity profile that follows a Gaussian distribution. We have investigated the physics of the accretion–ejection mechanism between the disk and the launched wind while focusing somewhat on the nature of the outflows and the development of the disk wind.

4. We provide a detailed study of the MHD characteristics of the disk wind structure. A thin disk exists until accretion and disk wind have depleted the initial mass reservoir of the disk. We resolve the disk surface where accretion of material is turned into ejection. The Alfvén surface of the disk wind is close to the disk surface—the disk wind is thus launched with sub-Alfvénic speed but quickly accelerated to super-Alfvénic velocities. The counterrotating disk seems to develop a different accretion mode with layered accretion in the upper disk levels.

5. Two different types of disk winds were identified. The first one arises from the inner part of the disk $r \leq 10$ and is dominated by the toroidal magnetic field component, while carrying a large part of the mass flux. This type of disk wind has many similarities with the wind investigated by Qian et al. (2018), where it was identified as a tower jet (Lynden-Bell 1996). In contrast to Qian et al. (2018), we observe a second type of disk wind. This feature is launched from the larger radii and is dominated by the poloidal magnetic field. So far we believe that this is mainly due to the fact that the outer disk is less evolved in comparison with the inner part. The $B_p$-dominated disk wind shows higher radial mass flux even though it is not as highly magnetized.

6. We compare the accretion rates for different black hole spin parameters. For the same level of magnetic diffusivity ($\eta_0 = 0.001$), we find that for increasing spin the accretion rate decreases close to the horizon. At the same time, the accretion rate increases, and with it the mass flux of the launched disk wind in both the polar direction (launching) and the radial direction (acceleration) increases as well. This result is in contrast with previous works (Qian et al. 2018), where the connection between accretion and disk wind was much stronger.

7. We compare the accretion rates for different levels of magnetic diffusivity. For the same black hole spin we find that increasing diffusivity lowers the accretion rate and results in a decrease in the mass flux of the disk wind launched from the disk surface. The radial mass fluxes show only small differences that do not allow us to say beyond any doubt if they are affected by the changes in diffusivity. Definitely, a weaker coupling between matter and magnetic field, induced by the increase in magnetic diffusivity, affects both accretion rate and mass loading of the wind in a similar way.

8. The electromagnetic energy flux that is carried by different parts of the outflow is dominated by the flux of the jet funnel. This flux in the jet funnel is highly affected by the black hole rotation, as this part of the outflow is driven by the BZ mechanism. We find that the disk and the Poynting-dominated outflows are strongly connected, as the level of magnetic diffusivity does affect the electromagnetic flux in the jet as well—in spite of the fact that the diffusivity close to the horizon is negligible. Similar to the peak in the mass fluxes for the disk wind, the Poynting flux reaches a peak value for $\eta_0 = 10^{-3}$. We believe that this critical level for the resistivity is a result of the interplay between reconnection decreasing the magnetic flux launching the outflow and magnetic diffusion and ohmic heating, both increasing the mass flux.

9. The simulation of a counterrotating black hole revealed an interesting feature. The retrograde rotation induces additional field reversals in the toroidal component of the magnetic field in the inner disk area. In this case, the accretion is supported mainly from the surface material of the inner disk area (where $B_p = 0$), though without significantly affecting the accretion rate itself.
(10) Motivated by the recent discovery of the M87 black hole shadow, we calculate the optical depth of the innermost accretion flow and the outflow structure around it. We find that for high accretion rates, \( M \approx 10 M_\odot \) yr\(^{-1}\) for a black hole mass of \( M \approx 10^9 M_\odot \), the innermost accretion stream may be opaque for the lensed signal, while the jet outflow launched from the disk and black hole will most probably remain optically thin.

In summary, we have compared the efficiency of GR-MHD jet launching for a sample of combinations of the accretion disk magnetic diffusivity and the black hole spin. We find a substantial mass loading of the disk wind that accelerates up to 0.1c. The low-density, high-velocity jet funnel generated by the BZ mechanism can be affected by the resistive, turbulent environment of the accretion disk. The two components of the disk wind follow the same trend even though its strength is not suppressed by a high disk diffusivity but continuously supported by mass loading. Besides magnetic diffusivity, also reconnection and ohmic heating govern the strength of the disk wind.

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Appendix A

Test Simulations Considering Magnetic Diffusivity

In order to verify the implementation of magnetic diffusivity into \textsc{Harm}3D, we have performed two test simulations. Our tests are similar to those applied by Qian et al. (2017). In the first test we follow the diffusion of a parallel magnetic field in a rectangular box for different values of the magnetic diffusivity and compare it with the time-dependent analytic solution to the diffusion equation. The second test problem is a classic shock tube that allows us also to check how magnetic diffusivity affects the shock-capturing abilities of the code.

A.1. Diffusive Decay of a Vertical Field

The setup for the simulations treating the diffusive decay of a vertical field considers a hydrostatic gas distribution located in an almost rectangular box that is threaded by a weak magnetic field. A uniform magnetic diffusivity is applied for the whole box and is the only parameter affecting the magnetic field evolution. Applying different levels of magnetic diffusivity, we compare the simulated evolution of the magnetic field with the analytic solution. As in Qian et al. (2017), we find an almost perfect match.

A.1.1. Numerical Setup

The box simulations are performed in a \( 256^2 \) grid in a small sector of our domain space, along the equatorial plane, extending by \( \Delta r \) in radius and \( \Delta \theta \) in latitude. By choosing a large-enough radius \( r_0 \) to place the box \( (\Delta r \ll r_0) \), we establish that its shape is as close as possible to a perfect square, with a side length \( r \in [r_0 - \Delta r/2, r_0 + \Delta r/2] \) and a latitude \( \Delta \theta \) that corresponds to a \( z \)-direction side \( \Delta z \), where \( z = r \sin(\pi/2 - \theta) \) and \( \theta \in [\pi/2 - \Delta \theta/2, \pi/2 + \Delta \theta/2] \).

A relativistic gas is applied in the area of the box with a polytropic index of \( \gamma = 4/3 \). The gas is in hydrostatic equilibrium with radial profiles of density \( \rho(r) = C r^\alpha \) and pressure \( p(r) = \beta \rho^\gamma \), where \( \alpha = 1/(1 - \gamma) \), \( \beta = 1/(1 - \alpha) \), and \( C \) denotes a proper normalization constant. This profile balances the gravitational force at large distances from the black hole (where GR effects are negligible). The magnetic field is uniform in the \( \theta \)-direction. In the \( r \)-direction it follows a (time-dependent) Gaussian profile

\[ B_\theta(r, \tilde{t}) = \frac{1}{\sqrt{\tilde{t}}} \exp\left(-\frac{(r - r_0)^2}{4\eta \tilde{t}}\right). \]  

We apply a very high plasma-\( \beta \), \( \beta_{\text{HM}} = 10^3 \), in order to establish a weak magnetic field that does not initiate any advection of magnetic flux. The time variable \( \tilde{t} = t_0 + t \) is connected with the code running time \( t \) with the parameter \( t_0 \), which basically normalizes the Gaussian profile. Finally, we are using outflow boundary conditions in all four boundaries of the box.

A.1.2. Simulation Runs

We have placed the simulation box far from the black hole at a radius \( r_0 = 300.5 \) with a side length of \( \Delta r = 1 \approx \Delta \xi \). At this distance the shape of our box is quite close to square as \( \Delta r \ll r_0 \). We follow the magnetic field evolution as given initially by Equation (31) along the equatorial plane. We run a series of simulations for different strengths of magnetic diffusivity. In Figure 28 we compare the simulation results (solid lines) with the analytic solution (dashed lines).

For simulation box12 with \( \eta = 10^{-12} \) there is barely any change in the magnetic field distribution and the simulation perfectly matches the ideal-MHD limit. Note that for the very high resolution applied in these simulations, also the numerical diffusivity is low.\(^{12}\) As we increase the magnetic diffusivity (box4, box3, box2) to the values of \( \eta = 10^{-2} \), the magnetic field decays—faster for higher diffusivity. Overall, the initial field distribution decays following exactly the analytical solution.

However, for high levels of the magnetic diffusivity, \( \eta > 0.1 \), the code fails. In this case the magnetic field has completely lost its initial Gaussian distribution, which poses a limit in the values of diffusivity we are allowed to use in our simulation.

A.2. Diffusive Shock Tube Test

Following Qian et al. (2017), we perform a series of tests with our resistive code based on the classical 1D shock tube test that demonstrates the shock-capturing capability of the code. We employ a computational domain that extends for \( x \in [298.75, 302.25] \) in the limit of Minkowski spacetime using Cartesian coordinates with 4000 cells to reduce the effect of numerical diffusion. The initial condition of the test follows the setup of Dumbser & Zanotti (2009) and Bucciantini & Del Zanna (2013). We implement a discontinuity in the density of

\(^{12}\) See Qian et al. (2017) for an assessment of the numerical diffusivity of \textsc{Harm}2D.
the gas, in the gas velocities, and in the magnetic field; thus,

\[
\begin{align*}
(r, p, v^s, v^c, B^s, B^c) &= (1.08, 0.95, 0.4, 0.3, 0.2, 0.1, 0.3) \\
&= (1.08, 0.95, 0.4, 0.3, 0.2, 2.0, 0.3, 0.3)
\end{align*}
\]  

(32)

for \( x < 300.5 \) and

\[
\begin{align*}
(r, p, v^s, v^c, B^s, B^c) &= (1.0, 1.0, -0.45, -0.2, 0.2, 2.0, -0.7, 0.5) \\
&= (1.0, 1.0, -0.45, -0.2, 0.2, 2.0, -0.7, 0.5)
\end{align*}
\]  

(33)

for \( x > 300.5 \). The initial electric field is set to the ideal-MHD value. The boundary condition at the ends of the tube is fixed to the initial values (Dirichlet boundary conditions). For the equation of state we choose a polytropic index \( \gamma = 5/3 \).

In Figure 29 we show the evolution of the discontinuity in gas density and horizontal velocity for different values of magnetic diffusivity. We note that in all cases we see the distinct features that result from the breaking of the initial discontinuity and the velocity values describe accurately the behavior of the gas density. The left-going rarefaction wave has a negative velocity and moves faster than the compound wave that follows it. The contact discontinuity propagates with the same speed as the compound wave, which also appears in the density distribution, while the discontinuity moves slowly away from its initial position at \( x = 0 \). The discontinuity is followed by a slowly moving shock front and a fast-moving rarefaction wave, both with positive velocities.

Figure 28. Diffusive decay of a vertical magnetic field. Evolution of the \( \theta \)-component of the magnetic field for simulation runs applying four values for the magnetic diffusivity, \( \eta = 10^{-12} \) (top left), \( \eta = 10^{-13} \) (top right), \( \eta = 10^{-3} \) (bottom left), and box2 with \( \eta = 10^{-2} \) (bottom right). Each color represents a different time step \( \tau \) in the simulation. Solid lines show simulation results, while dashed lines show the analytical solution.
Figure 29. Time evolution of the classic 1D shock tube test. Shown are gas density (left) and horizontal velocity (right) for different levels of magnetic diffusivity, $\eta = 10^{-12}, 10^{-3}, 10^{-2}, 10^{-1}$ (from top to bottom). The initial conditions for $\rho_0$ and $V_{x0}$ are denoted by black lines, while colored lines denote the evolution for five consecutive time steps.
that for vertical magnetic diffusivity. We see that for transition have detected strong outflows from the innermost disk area (Casse & Keppens 2002; Zanni et al. 2007; Murphy et al. 2010; Sheiknazami et al. 2012; Stepanovs et al. 2014; Stepanovs & Fendt 2016), numerically confirming analytical derivations in steady state by Blandford & Payne (1982) and Ferreira (1997).

In order to test our general relativistic simulation setup, it is therefore interesting to do a comparison simulation toward the nonrelativistic limit, thereby placing the inner grid boundary and the inner disk boundary at a radius farther out. This excludes almost all general relativistic effects from the simulation, in particular any influence from a central black hole. For this simulation the inner disk radius has been placed at the location of the inner boundary at \( r = 40 \), well outside the marginally stable orbit. The other simulation parameters were \( \eta = 0 - 10^{-5} \), \( a = 0 \), \( \beta = 50 \), \( K = 10^{-3} \), and \( \Gamma = 5/3 \).

Figure 31 shows the evolutionary state of such a mildly relativistic simulation at time \( t = 6000 \). We see a clear disk outflow along the poloidal field lines anchored in the disk. The maximum speed reached by this configuration is somewhat larger than \( 0.1c \) at a distance of \( r = 150 \) from the launching point. This corresponds to the orbital speed at the launching radius of \( r = 40 \) and thus what we expect from the classical Blandford–Payne magnetocentrifugal acceleration in the nonrelativistic limit (see publications cited just above; see in particular Figure 1 in Zanni et al. 2007).

Note that in comparison to the nonrelativistic simulations cited above, the evolutionary time step is rather low. Nonrelativistic simulations have been performed until several hundred thousand rotational periods of the inner disk (Stepanovs et al. 2014). For our mildly relativistic simulation the time unit is the light-crossing time over the gravitational radius \( t_g \), and thus \( t = 6000 \) corresponds to about 50 inner disk orbits only. However, already a substantial disk wind is launched, as the disk evolution time is only several orbits and also the outflow kinematic timescale is much shorter than the disk evolutionary timescale.

Note also the axial flow of low density along the rotational axis. The density distribution follows mainly from our floor model, applying also a relatively high internal energy, which leads to a pressure-driven axial flow. While this seems kind of artificial, it can be motivated by the existence of a central wind (stellar wind or Blandford–Znajek-driven jet) and helps to stabilize the central area against collapse.

With these simulations we have therefore proven the applicability of our setup for magnetocentrifugally driven disk winds in a mildly relativistic setup. While we expected to see similar trends also for the general relativistic simulations, those simulations show in fact much more violent and variable characteristics for the disk wind (see the discussion in the main text above).

**Appendix B**

**Mildly Relativistic Limit**

Nonrelativistic jet-launching simulations of the disk outflow transition have detected strong outflows from the innermost...
Figure 31. Test case of a mildly relativistic simulation with an inner disk radius located at the inner boundary at $r = 40r_g$. Shown is the mass density distribution overlaid with poloidal magnetic field lines, the mass density distribution overlaid with poloidal velocity vectors, and the vertical velocity distribution overlaid with poloidal magnetic field lines (from left to right), all at time $t = 6000$.

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