Magnetohydrodynamic Accretion–Ejection: Jets Launched by a Nonisotropic Accretion-disk Dynamo. I. Validation and Application of Selected Dynamo Tensorial Components

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Abstract

Astrophysical jets are launched from strongly magnetized systems that host an accretion disk surrounding a central object. The origin of the jet-launching magnetic field is one of the open questions for modeling the accretion–ejection process. Here we address the question of how the accretion-disk magnetization and field structure required for jet launching are generated. Applying the PLUTO code, we present the first resistive magnetohydrodynamic simulations of jet launching including a nonisotropic accretion-disk mean-field \( \alpha \Omega \) dynamo in the context of large-scale disk-jet simulations. Essentially, we find the \( \alpha_r \)-dynamo component determining the amplification of the poloidal magnetic field, which is strictly related to the disk magnetization (and, as a consequence, to the jet speed, mass, and collimation), while the \( \alpha_{\Omega} \) and \( \alpha_r \)-dynamo components trigger the formation of multiple, antialigned magnetic loops in the disk, with strong consequences for the stability and dynamics of the disk–jet system. In particular, such loops trigger the formation of dynamo-inefficient zones, which are characterized by a weak magnetic field and therefore a lower value of the magnetic diffusivity. The jet mass, speed, and collimation are strongly affected by the formation of dynamo-inefficient zones. Moreover, the \( \theta \) component of the \( \alpha \) dynamo plays a key role when the dynamo interacts with a nonradial component of the seed magnetic field. We also present correlations between the strength of the disk toy dynamo coefficients and the dynamical parameters of the jet that is launched.

Unified Astronomy Thesaurus concepts: Jets (870); High energy astrophysics (739); Active galactic nuclei (16); Magnetohydrodynamics (1964); Astrophysical fluid dynamics (101); Plasma astrophysics (1261); Magnetohydrodynamical simulations (1966); Young stellar objects (1834)

1. Introduction

Astrophysical jets, consisting of collimated high-speed outflows, are launched from a wide range of astrophysical objects such as young stellar objects (YSOs), microquasars, or active galactic nuclei (AGNs). Although these sources span orders of magnitude in terms of extension, timescales, and energy scales, it is commonly accepted that these jets are launched from systems that host an accretion disk surrounding a central object (Frank et al. 2014; Hawley et al. 2015; Pudritz & Ray 2019). There is further agreement on the key role of the large-scale magnetic field for the launching, acceleration, and collimation of jets. With the launch, we denote the transition between accretion and ejection, respectively, as the mass loading of outflows and jets. Quite a number of studies have investigated this launch process (see, e.g., Blandford & Payne 1982; Uchida & Shibata 1985; Casse & Keppens 2002; Fendt 2006; Zanni et al. 2007; Tzeferacos et al. 2009; Stepanovs & Fendt 2014; Stepanovs et al. 2014; Fendt & Gällmann 2018).

Still, the origin of the jet-launching disk magnetic field is not completely understood. Analytical models and numerical simulations have so far mostly assumed a predefined large-scale, open magnetic field structure, whose strength and configuration are essential parameters when understanding and determining the jet dynamics (Murphy et al. 2010; Stepanovs & Fendt 2016). For the origin of the accretion-disk magnetic field, one may consider an extended central stellar magnetic field, the advection of magnetic flux from the ambient interstellar medium, or a magnetic field that is generated by a dynamo process in the disk. The latter scenario is particularly interesting in order to generate jets from AGNs, which host as a central object a supermassive black hole, that, unlike stellar objects, cannot produce its own magnetic field.

Disk dynamos have already been suggested some decades ago (Pudritz 1981a, 1981b; Brandenburg et al. 1995), and there is extensive literature on dynamo theories and their applications to astrophysics (for reviews, see, e.g., Brandenburg & Subramanian 2005; Rincon 2019). Essentially, astrophysical dynamos are thought to have turbulent, thus small-scale, nature, while, on the other hand, one is interested in its large-scale effects on the dynamics of these systems. Overall, it is prohibitively expensive to model the turbulence involving the smallest scales, and, at the same time, to describe astrophysical systems on large scales. The disk turbulence that may lead to both a turbulent dynamo effect and also to a turbulent magnetic diffusivity is generally thought to be generated by the magnetorotational instability, MRI (Balbus & Hawley 1991).

For this reason, two paths to modeling dynamos have been pursued. These are (i) direct simulations, which study the natural amplification of the magnetic field (see, e.g., Gressel 2010; Bai & Stone 2013; Gressel & Pessah 2015; Hogg & Reynolds 2018; Riols & Latter 2018; Dhang et al. 2020) and (ii) the mean-field approach (see, e.g., Krause & Rädler 1980; Rüdiger et al. 1995; Campbell 1999; Rekowski et al. 2000; Bardou et al. 2001; Chabrier & Küker 2006), by which (ia) (semi)analytical solution can be derived or (iib) global numerical simulations covering the large scales of astrophysical systems can be run.

In this paper, we follow the second approach. Previous work in this field on large-scale disk-jet simulations and a possible
The origin of large-scale magnetic field in accretion disks has been performed by von Rekowski et al. (2003), Stepanovs et al. (2014), Fendt & Gaßmann (2018), and Dyda et al. (2018), essentially demonstrating that a mean-field dynamo-generated magnetic field can efficiently launch jet or outflows. Most recently, mean-field dynamos have also been considered in general relativistic magnetohydrodynamic (MHD) tori (Bucciantini & Del Zanna 2013; Bugli et al. 2014; Tomei et al. 2020) and disks (Vourellis & Fendt 2020).

In particular, we extend the work of Stepanovs et al. (2014) by studying the effects of a nonscalar mean-field dynamo. This has not yet been done for launching simulations. By prescribing a nonisotropic dynamo, in the induction equation, we are able to disentangle the dynamo effects shown in Stepanovs et al. (2014) and Fendt & Gaßmann (2018) in terms of the different components of the dynamo tensor.

Mean-field MHD theory arises from averaging the small-scale dynamics of a turbulent flow pattern, which is (in our case) affected by the central gravity and a subsequent rotation pattern, gas pressure gradients, and Lorentz forces. Both analytical theory (e.g., Rüdiger & Kichatinov 1993) and direct numerical simulations resolving the disk turbulence (e.g., Gressel 2010) have clearly detected an anisotropic nature of the dynamo tensor. As demonstrated by previous work, the different tensor components have different amplitudes and different impacts on the magnetic field components. We thus believe that the anisotropy of the dynamo will have an essential impact on the structure of the dynamo-generated magnetic field. This has not been shown before in numerical simulations of jet-launching disks.

The paper is organized as follows. In Section 2, we describe our model setup and the numerical approach. In Section 3, we investigate the different effects of the single components of a vectorial dynamo tensor on the jet-launching process. We summarize our paper in Section 4. In the appendices, we define our control volumes and provide test simulations comparing our new code to previous works.

2. Model Approach

2.1. MHD Equations

We solve the time-dependent, resistive MHD equations by applying PLUTO code (Mignone et al. 2007) version 4.3 on a spherical grid \( (R, \theta, \phi) \) assuming axisymmetry. We refer to \( (r, \phi) \) as cylindrical coordinates. The code integrates and solves numerically the set of MHD conservation laws, in particular for the conservation of mass,

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0,
\]

where \( \rho \) and \( \mathbf{v} \) are, respectively, the plasma density and the flow velocity; and the momentum conservation,

\[
\frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot \left[ \rho \mathbf{v} \mathbf{v} + \left( P + \frac{B \cdot B}{2} \right) I - BB \right] + \rho \nabla \Phi_\text{g} = 0,
\]

where \( P \) and \( B \) denote the gas pressure and the magnetic field, respectively. The central object of mass \( M \) provides the gravitational potential \( \Phi_\text{g} = -GM/R \). The energy is conserved through the equation

\[
\frac{\partial e}{\partial t} + \nabla \cdot \left[ \left( e + P + \frac{B \cdot B}{2} \right) \mathbf{v} - (\mathbf{v} \cdot B) B + \eta \mathbf{j} \times B \right] = \Lambda_{\text{cool}},
\]

where the total energy density is defined as

\[
e = \frac{P}{\gamma - 1} + \frac{\rho v^2}{2} + \frac{B^2}{2} + \rho \Phi_\text{g},
\]

with the polytropic index \( \gamma = 5/3 \).

The electric current density is determined by Ampère’s law \( \mathbf{J} = \nabla \times \mathbf{B} \). As shown e.g., by Casse & Ferreira (2000a, 2000b), Zanni et al. (2007), and Tzeferacos et al. (2013), cooling may play a role in the jet-launching process as both density and velocity are subjected to cooling effects. For the sake of simplicity, as in Sheikhlazemi et al. (2012) and Stepanovs & Fendt (2014), the cooling term is set to be equal to the ohmic heating, which is, therefore, instantly radiated away.

The magnetic field evolution is determined by the induction equation. Here, we have implemented into the code a mean-field dynamo term (Krause & Rüdiger 1980),

\[
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B} + \tau_\text{dyn} B - \mathbf{J}),
\]

following mainly the approach of Stepanovs et al. (2014). The tensors \( \tau_\text{dyn} \) and \( \eta \) describe the \( \alpha \) or effect of the mean-field dynamo and the magnetic diffusivity.

2.2. Numerical Setup

As we solve the nondimensional MHD equations, no intrinsic physical scales are involved. All of the primitive MHD variables, i.e., \( \rho, \mathbf{v}, P, \mathbf{B} \), as well as the length scale and timescale, are normalized to their value at the initial inner disk radius \( R_{\text{in}} \). Thus, velocities are normalized to \( v_{\text{K, in}} \) corresponding to the Keplerian speed at \( R_{\text{in}} \). As a consequence, the time unit is given in units of \( t_{\text{in}} = R_{\text{in}} / v_{\text{K, in}} \), and therefore the quantity \( 2 \pi t_{\text{in}} \) corresponds to one revolution at the inner disk radius. In the following, all times are measured in units of \( t_{\text{in}} \), implying that \( t = 2000 \) (in short) corresponds to \( t = 2000 t_{\text{in}} \).

The computational domain covers a radial range of \( R = [1, 100] R_{\text{in}} \) and an angular range of \( \theta = [10^{-8}, \pi/2-10^{-8}] \approx [0, \pi/2] \). A stretched grid is applied in the radial direction considering \( \Delta R = R \Delta \theta \). The domain is discretized with a number of \( [N_R \times N_\theta] = [512 \times 128] \) grid cells, which allows the initial disk height \( H = 0.2 r \) with 16 cell to be resolved.

For the resolution study (see the Appendix to Mattia & Fendt 2020, hereafter Paper II), we have discretized the domain with \( [N_R \times N_\theta] = [1024 \times 256] \) and \( [N_R \times N_\theta] = [256 \times 64] \) grid cells, namely 32 and 8 cells per disk height, respectively.

Our scale-free simulations may be applied and thus scaled to a variety of jet sources. We apply the same physical scaling as described in previous works (Zanni et al. 2007; Tzeferacos et al. 2009; Sheikhlazemi et al. 2012; Stepanovs & Fendt 2014). For an astrophysical scaling of our normalized quantities for typical jet systems we refer to Table 1.

For spatial integration, we use the piecewise parabolic interpolation method (see Mignone 2014). Time integration is achieved through a third-order Runge–Kutta scheme, while for the flux computation a Harten–Lax–van Leer Riemann solver is employed (Toro 2009). To preserve the solenoidal condition of the magnetic field, the method of upwind constrained transport...
Table 1

|            | YSO | BD  | AGN | (Units) |
|------------|-----|-----|-----|---------|
| $R_{in}$   | 0.1 | 0.01| 0.05| au      |
| $M_{in}$   | 1   | 0.05| 10^3| $M_s$  |
| $\rho_{in}$| 10^{-10} | 10^{-13} | 10^{-12} | g cm^{-3} |
| $v_0$      | 94  | 66  | 6.7 × 10^7 | km s^{-1} |
| $B_0$      | 15  | 0.5 | 1000 | G      |
| $t_0$      | 1.7 | 0.25| 0.5  | days   |
| $M_0$      | 3 × 10^{-5}| 2 × 10^{-10} | 10  | $M_s$ yr^{-1} |

(Londrillo & del Zanna 2004) is applied. In order to achieve stability, we choose a Courant–Friedrichs–Lewy time stepping with CFL = 0.4 < 1/√$N_{dim}$. This may be a challenge for our diffusive MHD simulations, in particular for high resolution as the diffusive time step goes as $\tau_\alpha = (\Delta R)^2/\eta$.

2.3. Initial Conditions

The simulations start with a very weak initial seed field, thus with a very low disk magnetization, defined as the ratio between the magnetic pressure and the thermal pressure

\[ \mu_{in} = B_{in}^2/P_{in} = 10^{-5} \]

measured at the disk midplane. Therefore, the initial structure of the accretion disk can be obtained as a solution of the hydrostatic equilibrium between thermal pressure gradients, gravity, and centrifugal force (Zanni et al. 2007; Stepanovs & Fendt 2014), ignoring the Lorentz force (Stepanovs et al. 2014; Fendt & Gaßmann 2018).

\[ \nabla P + \rho \nabla \phi_k - \frac{1}{R} \rho v_0^2 (e_R \sin \theta + e_\theta \cos \theta) = 0. \quad (6) \]

This equation can be solved by assuming that all of the quantities X scale as power laws, X = $X_0 R^{k X} f_k(\theta)$, where $X_0$ is the corresponding quantity evaluated at the innermost radius of the disk (midplane). For the sake of clarity, we summarize here the pertinent formulas. Self-similarity requires that every characteristic speed should scale as the Keplerian velocity, $\propto R^{-1/2}$. Combining this assumption with a polytropic gas, $P \propto \rho^{\gamma}$, the power-law coefficients are $\beta_\rho = -1/2$, $\beta_p = -5/2$, and $\beta_\rho = -3/2$ (as in the self-similar solution of, e.g., Blandford & Payne 1982). A key parameter to describe the initial disk structure is the ratio between the isothermal sound speed and the Keplerian velocity at the disk midplane of the inner radius $c_s = v_0/\sqrt{\gamma R \rho} \propto R^{-1/2}$. For an initially thin disk, $c_s = z/R \propto 0.1$, and solving for the $z$ component of Equation (6) with $\rho_{in} = 1$ at the inner disk radius, we obtain

\[ F_p = \left[ \frac{2}{5 \epsilon} + \left( \frac{1}{\sin \theta} \right) \right]^{5/2} \quad (7) \]

and where we have chosen $P_{in} = 0.01$. Following the polytropic relation assumed before, the disk pressure is defined by $P_p = F_p^{3/5}$.

Following Stepanovs & Fendt (2014) and Stepanovs et al. (2014), by solving the radial component of Equation (6), we obtain the disk rotation profile. Outside the disk, we define a hydrostatic corona,

\[ \rho_c = \rho_{c,in} R^{1/(1-\gamma)}, \quad P_c = \frac{\gamma-1}{\gamma} \rho_{c,in} R^{7/(1-\gamma)}, \quad (8) \]

with $\rho_{c,in} = 10^{-3} \rho_{in}$. All simulations are initialized with a purely radial magnetic field vanishing outside the disk, defined by the vector potential

\[ B = \nabla \times A_{e_\phi} = \nabla \times \left[ \frac{B_{p,in}}{r} \exp(-8(z/H)^2) \right] e_{\phi}. \quad (9) \]

Here, $B_{p,in} = \epsilon \sqrt{2/\mu_{in}}$ defines the strength of the initial poloidal magnetic field and $\mu_{in} = 10^{-5}$ is the initial magnetization along the disk midplane. The quantity $H = 2\epsilon r$ represents the geometric disk height, which is twice the initial pressure scale height.

2.4. Boundary Conditions

The boundary conditions are identical to those of Stepanovs & Fendt (2014). We report them in this section for convenience. Along the rotational axis and the equatorial plane, the standard symmetry conditions are applied. The inner radial boundary is divided into two different areas. One is the area that is suited for disk accretion located at $\theta > \pi/2 - 2\epsilon$; the other is the coronal area at $\theta < \pi/2 - 2\epsilon$, where we choose $2\epsilon \approx \arctan(2\epsilon)$. The boundary condition along the inner radial boundary is essential for stabilizing the corona against collapse to the central object. While $v_\theta = 0$ along the inner disk boundary, the radial velocity follows a power law $v_R = v_{R,in} R^{-1/2} \leq 0$, where the inequality is imposed in order to enforce the boundary behaving as a “sink.” Along the coronal area, we prescribe a constant inflow velocity into the domain $v_p = 0.2$ (in units of the Keplerian speed at $R_{in}$) in the radial direction, which could be interpreted astrophysically as a stellar wind.

From previous jet formation simulations (see, e.g., Ouyed & Pudritz 1997) we expect the terminal jet speed to reach the Keplerian velocity at the inner disk. For $v_\phi$, we prescribe a power law across the inner boundary (for both the disk and the coronal boundary),

\[ v_\phi = v_\phi_{in} R^{-1/2}. \quad (10) \]

The boundary conditions for the poloidal magnetic field along the inner radial boundary obey the divergence-free condition. The method of constrained transport requires only the $\theta$-component of the magnetic field to be defined along the boundary, while the radial component is recovered from the Maxwell equations. At the outer boundaries, both $B_\phi$ and $B_\theta$ follow a power law:

\[ B_{\phi,\theta} = B_{\phi,\theta,\text{out}} R^{-1}, \quad (11) \]

while $B_R$ is again recovered using the solenoidality condition. This is compatible with a constant gradient condition. For $B_{\phi,\theta}$ this implies in particular the conservation of the electric current across the boundary.

Along the inner boundary, we prescribe $B_\phi = 0$ toward the coronal region, while we again adopt a power law $\propto R^{-1}$ for the boundary area toward the inner disk. Along the inner radial disk boundary, we prescribe the poloidal magnetic field
inclination, choosing an angle
\[ \varphi = 70^\circ \left[ 1 + \exp \left( -\frac{\theta - 45^\circ}{15^\circ} \right) \right]^{-1}, \] (12)
where \( \varphi \) is the angle of the magnetic field with respect to the disk surface. Note that here again we solve for the divergence-free condition of the magnetic field, recovering the solution with the inclination prescribed.

These boundary conditions are slightly different from those of Stepanovs et al. (2014) and Fendt & Gaßmann (2018), as we do not suppress the advection of the magnetic flux from the inner disk toward the axis. This inner boundary condition is known to be quite critical for the numerical stability of the simulation, as it is time dependent, thus implementing a feedback loop from the cells of the active domain. While the boundary condition as defined in the papers mentioned above was chosen because it was found to be less prone to numerical instabilities, for the present paper, we decided to release that boundary condition and allow the magnetic flux to be advected across the boundary toward the axial region. The advection of flux toward the axis has some impact on the structure of this innermost area, but it does not change the structure and the evolution of the surrounding disk jet, which is our major focus. We also think that the advection of magnetic flux toward the axis is a more physical boundary condition.

Across the inner and outer boundaries, both the density and pressure are extrapolated by a power law,
\[ \rho = \rho|_{R_{in},R_{out}} R^{-3/2}, \quad P = P|_{R_{in},R_{out}} R^{-5/2}, \] (13)
where \( R_{in} \) and \( R_{out} \) are the inner and outer radii of the domain. Along the outer boundaries, we apply the standard PLUTO outflow (zero-gradient) conditions. In addition, we still prescribe \( v_R \) to be nonpositive in the disk region and nonnegative in the coronal region.

### 2.5. The Dynamo Model

For a thin disk, the nondiagonal components of the mean-field dynamo tensor are negligible. In our approach, we consider the explicit form of the dynamo terms following Rüdiger et al. (1995) and Rekowski et al. (2000),
\[ \sigma_{\text{dyn}} = (\alpha_R, \alpha_\theta, \alpha_\phi) = -\sigma_0 c_s F_\alpha(z), \] (14)
where \( c_s \) is the adiabatic sound speed at the disk midplane and \( F_\alpha(z) \) is a profile function,
\[ F_\alpha(z) = \begin{cases} \sin \left( \frac{z}{H} \right) & z \leq H \\ 0 & z > H \end{cases} \] (15)
(Bardou et al. 2001), which confines the dynamo action within the accretion disk. Note that Rüdiger et al. (1995) applied a slightly different profile, namely a linear function \( F_\alpha(z) = z/H \) in the disk. We prefer the approach of Bardou et al. (2001) which effectively avoids the discontinuity at the disk surface and is thus better suited for a simulation that also includes the disk corona.

As in Stepanovs et al. (2014), we choose a radial dependence of the dynamo \( \alpha \propto R^{-1/2} \), as this profile also follows the sound speed. Note, however, that compared to our former simulations, in the present setup, the radial profile of the dynamo is not necessarily constant in time. As the sound speed is included in the dynamo tensor, along with the disk sound speed, the dynamo tensor is also updated every time step. We will demonstrate that this variation only has a minor impact on the overall evolution of the system. However, it represents a more consistent approach and is furthermore in agreement with the analytical models of mean-field dynamo theory (Rüdiger & Kichatinov 1993; Rüdiger et al. 1995).

### 2.6. The Diffusivity Model

For the magnetic diffusivity tensor, we assume a diagonal structure (as for the dynamo). We adopt an \( \alpha \)-prescription as typically applied in our previous work (Stepanovs et al. 2014),
\[ \eta = (\eta_R, \eta_\theta, \eta_\phi) = \eta_0 \alpha_s c_s H F_\eta(z), \] (16)
where \( c_s \) is the adiabatic sound speed at the disk midplane, \( H \) is the initial disk pressure scale height, while \( \alpha_s \) is the dimensionless parameter of the turbulence (Shakura & Sunyaev 1973). Thus, the diffusivity is assumed to essentially have a turbulent nature, most probably caused by the MRI (Balbus & Hawley 1991). Again, we define a profile function,
\[ F_\eta(z) = \begin{cases} 1 & z \leq H \\ \exp \left[ -2 \left( \frac{z}{H} \right)^2 \right] & z > H \end{cases} \] (17)
which confines the diffusivity within the disk region. Note that the magnetic diffusivity, or resistivity, respectively, is motivated here by the disk turbulence, and thus is much stronger than the microscopic value.

In the literature of jet-launching simulations (see, e.g., Jactemin-Idle et al. 2019) without a mean-field dynamo, the magnetic diffusivity is usually computed as
\[ \eta = \eta_0 v_A H F_\eta(z), \] (18)
where the two model approaches described above coincide if \( \alpha_s = \sqrt{2\mu_D/\gamma} \), where \( \gamma \) is the polytropic index and \( \mu_D \) is the magnetization computed at the disk midplane. This model approach, in the following denoted as the standard diffusivity model, is, however, not used in this paper. One reason is that we want to avoid the accretion instability from occurring (Campbell 2009). This can be avoided when the feedback between the magnetization and magnetic diffusivity is chosen to be stronger than \( \alpha_s \propto \sqrt{\mu} \) (see our previous work Stepanovs & Fendt 2014). Note that we already have a feedback loop on the magnetic diffusivity, as the growth of the magnetic field is naturally related to the mean-field dynamo.

We therefore apply the so-called strong diffusivity model that we have previously invented (Stepanovs & Fendt 2014; Stepanovs et al. 2014),
\[ \alpha_s = \frac{2}{\gamma} \left( \frac{\mu_D}{\mu_0} \right)^2, \] (19)
where \( \mu_0 = 0.01 \). Because the initial magnetic field does not intersect the disk midplane, for the quantity \( \mu_D \), we calculate the ratio between the average total magnetic field (vertically averaged at a certain radius) in the disk and the gas pressure at the disk midplane (Stepanovs et al. 2014). As demonstrated previously, this approach allows a more stable evolution of the
The accretion instability is the disk mass loss which increases magnetization, which increases the mass loss, and so on.

However, in order to be able to evolve the long-term properties of the various dynamo models, we choose the diffusive quenching over the standard quenching. The standard quenching was found to be prone to the accretion instability for our setup. Furthermore, the diffusive quenching works much more smoothly compared to the standard quenching and leads to the same result. We note that we do not put any lower bounds on the turbulence level $\alpha_{\text{ss}}$. This may affect, via the dynamo alpha $\alpha_0$, the critical dynamo number, above which we expect an effective magnetic field amplification. The study of physically more self-consistent feedback models for dynamo quenching will be the subject of our future work.

3. A Toy Model for an Anisotropic Mean-field Dynamo

This section aims to disentangle the effects that are physically caused by the different components of the dynamo tensor (the three components of a vector in our case) in order to gain a detailed understanding of the physical process of the field amplification in action.

3.1. Anisotropic Dynamo and Diffusivity Coefficients

Our aim is to generalize the dynamo models applied previously (von Rekowski et al. 2003; Stepanovs et al. 2014; Fendt & Gaßmann 2018). These works applied a scalar (thus isotropic) $\alpha$ coefficient. Here, we apply the anisotropy of the dynamo, assuming that the coefficients $\tilde{\alpha}_{\Gamma}$, as described in Section 2.5, are not necessarily the same (Rüdiger et al. 1995).

The role of diffusivity has been widely discussed in the literature (Zanni et al. 2007; Sheikhnezami et al. 2012). Here, we assume

$$\eta_0 = \left(\frac{1}{2}, \frac{1}{2}, 1\right) \eta_0,$$

where $\eta_0 = 0.165$ recovers the reference values of Stepanovs et al. (2014). In order to have a direct comparison with the simulations of Fendt & Gaßmann (2018), we set the dynamo tensor components as

$$\tilde{\alpha}_{\Gamma} = (\phi, \psi, \chi) \alpha_0,$$

with $\alpha_0 = 0.775$. Setting $\psi = \phi = \chi = 1$, we recover the reference simulation of Fendt & Gaßmann (2018).\footnote{Note that $\alpha_{\text{ss}}$ as well as the dynamo tensor (now also considering sound speed) are now defined differently. Thus, the coefficients $\alpha_0$ and $\eta_0$ are not defined in the same way.}

The strength of the dynamo coefficients $(\phi, \psi, \chi)$ is summarized in Table 2. From this set of simulation runs, we will consider a sample of eight exemplary runs in order to disentangle the influence of the different components of the alpha tensor on the magnetic field structure and the disk and jet evolution.

3.2. Evolution of the Magnetic Field

Figure 1 shows for the different parameter runs the density distribution of the disk-jet structure, together with the magnetic field geometry (as contour lines of the vector potential). We point out that in all cases but simulation $\eta_0 B$ (which is described more in detail in Section 3.4), the initial magnetic field has the radial structure described in Section 2.3.

Overall, we see that in all simulation runs, the magnetic field in the inner region close to the rotation axis that has been
generated by dynamo action shows a large-scale open geometry. Together with a substantial strength, this magnetic field structure is able to eject disk material into an outflow with a high degree of collimation. On the other hand, we also see that the very field structure depends on the choice of the dynamo tensor, and thus the strength of the tensor components. The choice of different coefficients \((\phi, \psi, \chi)\) in our toy dynamo model leads to a different magnetic field configuration.

### 3.2.1. A Supercritical Poloidal Dynamo

The induction equation tells us that the dynamo action governed by \(\alpha_\phi\) is the only way to increase the poloidal magnetic field up to the strength that is required for jet launching. Even for \(\alpha_\phi = 0\) the toroidal magnetic field is still dynamo-amplified through the \(\Omega\) effect and also the \(\alpha_R\) dynamo component. However, the dynamo does not lead to a substantial amplification of the poloidal magnetic field. Therefore, the latter cannot increase and stays confined within the disk. Neither the strength nor the launching angle that is required to produce a Blandford–Payne outflow can be reached.

On the other hand, as a consequence of the quenching model applied, the magnetic diffusivity still increases as the toroidal magnetic field grows. As a consequence, the poloidal magnetic field still evolves, even if the field is not enhanced.

![Figure 1. Magnetohydrodynamic evolution of the toy model dynamo simulations (see Table 2) at \(t = 4000\). Shown are the density distribution (color, in logarithmic scale) and magnetic field lines (white lines). The poloidal magnetic field is represented by the contour lines of the vector potential \(A_\phi\). The dashed lines indicate a negative polarity of the poloidal magnetic field.](image-url)

| Run ID | \(\phi\) | \(\psi\) | \(\chi\) | \(\eta_\psi\) | Comment |
|--------|---------|---------|---------|-------------|---------|
| Scalar | 1.0     | 1.0     | 1.0     | 30          | As Stepanovs et al. (2014) |
| \(\phi_A\) | 1.0     | 1.0     | 2.0     | 10          | Strong amplification |
| \(\phi_B\) | 1.0     | 1.0     | 0.5     | 10          | Weak amplification |
| \(\phi_C\) | 1.0     | 1.0     | 0.1     | 10          | Very weak amplification |
| \(R_A\) | 2.0     | 1.0     | 1.0     | 10          | Magnetic loops at \(R \approx 40\) |
| \(R_B\) | 0.75    | 1.0     | 1.0     | 10          | Magnetic loops at \(R \approx 20\) |
| \(th_A\) | 1.0     | 5.0     | 1.0     | 10          | Multiple loops in \(R \in [15, 80]\) |
| \(th_B\) | 1.0     | 0.1     | 1.0     | 4           | Magnetic loops at \(R \approx 15\) |

*Note.* The magnetic diffusivity distribution is the same with \(\eta_\psi = 0.165\). The run time of the simulations is \(\eta_\psi\) in units of 1000.
Different evolution of the magnetic field topology. That is, for \( \phi < 0.8 \) or \( \phi > 1.5 \), a second magnetic loop is formed that is antialigned to the loop structure induced farther in. These loops, characterized by a reversal in the toroidal field, are substantially different from the ones described previously and play a significant role in the evolution of the magnetic field and of the disk–jet system (see our discussion below). We point out that the antialigned magnetic loops can also be formed when considering a scalar dynamo tensor, when the scalar \( \alpha_0 < 0.6 \) (Fendt & Gaßmann 2018).

As shown in Section 3.2.1, the coupling between the toroidal magnetic field and the dynamo tensor component \( \alpha_\phi \) is the main mechanism responsible for the generation of the poloidal field. For \( 0.8 < \phi < 1.5 \), the toroidal magnetic field, being amplified by the \( \Omega \)-effect from the radial weak seed field, shows a monotonic behavior (after being amplified). As the system evolves, the poloidal field is amplified over the whole accretion disk.

By looking at the spatial and temporal numerical derivatives of the toroidal field, we find that because of the highly anisotropic character of \( \alpha_\theta \), some “dynamo-inefficient zones” are formed. These are areas of vanishing poloidal field strength, but, in addition, in such zones also the toroidal magnetic field cannot be amplified. The number and the location of these zones, where the dynamo is not efficient, depend on the strength of the three dynamo components and not exclusively on \( \alpha_R \).

Furthermore, for \( \alpha_\theta > 3 \), we find that the toroidal field shows multiple dynamo-inefficient zones. On the other hand, the dynamo-inefficient zones of case \( th_A \) remain confined in the accretion disk. This is illustrated in the top panels of Figure 3, where we show the disk magnetization at the same evolutionary time, \( t = 4000 \). The difference among the three simulation runs \( Sc, R_B, \) and \( th_A \) is clearly visible.

For the case of the scalar dynamo, the local disk magnetization is only weakly dependent on the radius. It is relatively low along the midplane and increases toward the disk surface. This is understandable as the disk gas pressure decreases with altitude while the poloidal field remains rather constant vertically.

For simulation run \( R_B \), for which \( \alpha_R = 0.75 \), we see that a dynamo-inefficient zone has developed around radius \( R \approx 23 \). Typically, these zones seem to be anchored at the disk midplane. As they are balanced by a low magnetic pressure, they vertically extend while preserving the total pressure equilibrium.

For simulation run \( th_A \), for which \( \phi = 1 \) and \( \psi = 5 \), we find multiple dynamo-inefficient zones along the accretion disk. Note that due to their proximity, these zones are able to connect and reconnect.

Because the coupling between the toroidal field and the \( \alpha_\phi \) component of the dynamo tensor is the only way to dynamo-amplify the radial field component, the radial field that is amplified from the toroidal also has different polarities.

Because we included physical resistivity, the magnetic field is able to reconnect and to change its topology within the accretion disk. In particular, instead of one magnetic loop that is visible (see Figure 3, bottom-right panel), now more magnetic loops arise (see Figure 3, bottom-left panel). On the other hand, the reversal of the toroidal field is associated with a maximum in the tensor component \( \alpha_\phi \), which undergoes a reversal at smaller radii (bottom-left panel of Figure 3).
Compared to the results of Fendt & Gallaßmann (2018), here we find that the reconfiguration of the magnetic field structure does impact the jet evolution on a weaker level. We believe that this is mainly due to the midplane boundary condition, which is absent in the previous paper. In particular, here we enforce symmetry between the upper and lower hemispheres, which can be violated in a bipolar setup. However, the reversal of the toroidal and radial field components that directly define the disk magnetization still play a key role in the disk-jet evolution. We note that disk magnetization is the main ingredient of the diffusivity model for the resistive disk evolution.

Because the dynamo-inefficient zones correspond to zones of low diffusivity, as a result the accretion process can be affected. In fact, accretion can be suppressed across such zones, leading to underdense and overdense regions (compared to the simulations without multiple loops). We find that these underdense/overdense regions are strongly related to the existence of a vertical field. We experienced numerical issues when underdense zones are located too close to the inner boundary, for example, unphysical values of the fluid density or the fluid pressure.

Here we need to comment briefly on the “dead zones” that have been proposed for protoplanetary disks. Although the dynamo-inefficient zones we detect in our simulations may look similar to these dead zones, the physical processes involved are not the same. Dead zones in protoplanetary disks have been proposed by Gammie (1996) on the basis of a lack of coupling between matter and magnetic field due to an insufficient degree of ionization. This lack of coupling would not allow the MRI to operate, and, as a consequence, also accretion is unlikely to happen, due to the lack of angular momentum exchange. As a result, a layered accretion is expected on a theoretical basis, which could indeed be realized in numerical simulations (Fleming et al. 2000; Fleming & Stone 2003). Also, resistivity was found to play an essential role in suppressing the MRI (see, e.g., Sano et al. 2000; Fromang et al. 2002; Flock et al. 2012). Dead zones in protoplanetary disks are also thought to be responsible for creating transition disks (Pinilla et al. 2016).

It is interesting to note that for both the protoplanetary dead zones and for our dynamo-inefficient zones, resistivity plays a leading role. However, for the first approach, it is the resistive decoupling that suppresses the MRI (and would subsequently suppress the dynamo action of the MRI), while for our models the dynamo-inefficient zone is formed as result of a minimum of the magnetic diffusivity.

Finally, we note that as the dynamo-efficient zones basically result from the feedback of the magnetic field on the magnetic diffusivity, a change in the quenching model—from the diffusive quenching to the standard quenching—may affect the exact location and width of the dynamo-inefficient zones.

### 3.2.3. Amplification of the Magnetic Field

The majority of our parameter runs apply a supercritical dynamo $\alpha_\phi > \alpha_{\text{crit}}$ (Table 2). The resulting magnetic field strength and geometry support a collimated outflow. In Figure 2 we show the time evolution for the disk poloidal magnetic energy, integrated from $R = 10$. For comparison, the case of an isotropic dynamo is shown.

The three different dynamo tensor components play a different role in the amplification of the poloidal magnetic field. The $\phi$ component of the dynamo is the main ingredient that amplifies the poloidal magnetic field in the disk, while the $R$ and $\theta$ components determine the formation of the dynamo-inefficient zones that, subsequently, also determines the poloidal magnetic field structure.

The $\phi$ component of the dynamo tensor essentially already influences the very early stages of the disk-jet evolution—a higher strength of $\alpha_{\phi}$ leads to a faster and stronger amplification, as we can see by comparing the “$\phi$” simulations to the isotropic model in Figure 2.

The other dynamo components ($\alpha_R$ and $\alpha_\theta$) become important only once the poloidal field has been amplified to substantial strength and through the presence (or absence) of the dynamo-inefficient zones. In particular, where a dynamo-inefficient zone is built up in the inner disk, it already triggers,
in units of $t_{\text{in}}$, the temporal evolution of the system on short timescales ($\approx 100$ after its formation). A dynamo-inefficient zone located farther out plays a minor role during the early phase of the disk evolution.

We emphasize that the evolution of the disk magnetic field is strictly correlated with the existence of dynamo-inefficient zones, because these features lead to the formation of multiple antialigned magnetic loops in the disk (see Figure 1 and Section 3.2.2). A higher strength of the dynamo tensor component $\alpha_R$ leads to—on average—a higher amplification of the toroidal field. However, once the dynamo is quenched by magnetic diffusivity, the magnetic field strength decreases to the magnitude that we recovered in the isotropic dynamo simulation. Therefore, we interpret the effect of a higher $\alpha_R$ to be a more rapid amplification of the poloidal field. On the other hand, a lower $\alpha_R$ leads to a slower toroidal (and therefore poloidal) field amplification.

We find a different behavior when a dynamo-inefficient zone (only one) is forming, which extends beyond the accretion-disk surface. As discussed in Section 3.2.2, the reversal of the toroidal field corresponds to a spatially stationary point in the $\theta$ component of the magnetic field. As a result, the poloidal magnetic energy is higher than for the isotropic dynamo model, simply because in the dynamo-inefficient regions of the disk, the vertical field component becomes stronger.

On the other hand, this increase in the vertical component of the magnetic field is partially suppressed in the presence of multiple magnetic zones, compared to the case of an isotropic dynamo tensor. Our understanding of this effect is that the existence of quite a number of field reversals (which effectively decrease the local magnetic energy) more than compensates for the induction of a vertical field component (which would lead to a decrease of the local magnetic energy).

### 3.2.4. The Dynamo Number

The dynamo number is usually quoted as a measure of dynamo activity. Only dynamos with a supercritical dynamo number evolve rapidly and work efficiently against magnetic diffusivity and finally lead to a strong, saturated poloidal magnetic field. The dynamo number can therefore tell us when and where the growth of the magnetic field reaches saturation. In Figure 4, we compare the dynamo number as a function of time and radius for different cases.

We first show the dynamo number for different strengths of the tensor component $\alpha_\phi$ (top panel). We see that as $\chi$
decreases, the amplification of the poloidal field is weaker and also slower, as also indicated by Figure 2. These differences in the magnetic field evolution are reflected in the dynamo number. In the time evolution of the dynamo number for all simulations, we can clearly distinguish three evolutionary stages.4

We may first define an (i) initial phase (indicated in blue) during which the dynamo number is almost infinite, simply because the diffusivity is still low (as implied by the quenching triggered by the magnetic diffusivity). Then comes a (ii) dynamo phase (indicated in white) that is characterized by a strong competition between dynamo action and diffusive quenching. During this phase, we recognize magnetic loops being present, surviving from the early stages ($t \lesssim 500$ in the inner disk region) of dynamo evolution. In a subsequent (iii) final phase (indicated in red), these magnetic loops have been washed out or have been broken up, and a quasi-steady state of the dynamo evolution is reached. The timescale when the final phase is reached depends on the radius (and thus on the dynamical timescale that is defined by the disk rotation at this radius). In the inner radii, the final stage is reached around $t \lesssim 500$, while in the outer disk regions, it is reached only at $t \gtrsim 5000$. In this final phase, dynamo action and diffusive quenching are fully balanced.

Note that in the inner disk region the second dynamo phase is missing because of the rapid evolution of the dynamo. Here, the magnetic energy reaches the saturation level already very early, with a timescale of the first two phases being much smaller.

Considering now the effect of different levels of dynamo anisotropy we find the following results. For larger $\chi$, we do not find a second phase at larger radii because the magnetic field is amplified on a shorter timescale. In addition to that, for larger $\chi$, the first phase has a shorter lifetime at every radius.

Looking at the innermost parts of the accretion disk, a larger $\chi$ leads to an overall smaller dynamo number at the stage of the quasi-steady state. This is a consequence of the quadratic dependence on the disk diffusivity (see Equation (19)) that balances and quenches the mean-field dynamo effect. For the latest evolutionary stages, we notice that although this happens at different times, for each choice of $\chi$, the simulation also reaches its steady state at a larger radius. This is an indicator of a faster evolution of the magnetic field for larger $\chi$.

Note that the dynamo number can also be used as a tracer to identify the dynamo-inefficient zones. As the latter correspond to a minimum in the magnetic diffusivity, here the dynamo number will have a sudden growth. On the other hand, the dynamo-inefficient zones are not only zones where the toroidal magnetic field has a minimum, but they are also zones where the toroidal field cannot be amplified. For such reason, the general application of the dynamo number as a measure of dynamo activity can be misleading, because its sudden growth (corresponding with the field reversal) does not necessarily lead to further magnetic field amplification.

This is shown in Figure 4, where we display in the bottom panels the dynamo number for the simulation runs that result in the generation of dynamo-inefficient zones. In contrast, the upper panels show simulations that do not lead to dynamo-inefficient zones. The figure nicely demonstrates a similar evolution of these simulations up to radii where the dynamo-inefficient zones have been established when a quasi-steady state is reached.

Interestingly, the dynamo-inefficient zones—representing a minimum in the toroidal and radial magnetic field components—do not directly affect the dynamo activity farther out. Outside the field-reversal zone, a saturation of the magnetic field can be reached. This is in particular visible when comparing the two right panels (runs $\phi_C$ and $th_B$).

Looking at the dynamo number in more detail, we understand why case $R_A$ and case Scalar (isotropic dynamo) are almost indistinguishable (Figure 2, left). The dynamo-inefficient zone that is present in case $R_A$ appears only at later stages of the evolution, as it is located at about $R \approx 40$ while the magnetic field in the ambient parts of the disk is amplified only on a longer timescale. In contrary, the dynamo-inefficient zone of $R_B$ is already formed earlier at $t \approx 1000$, and therefore, a different evolution of the poloidal disk magnetic field takes place and also on a shorter timescale. The time evolution of cases $th_A$ and $th_B$ will be discussed below (see Section 3.4).

### 3.3. Dynamics of Accretion–Ejection

So far we have mainly investigated the evolution of the magnetic field structure that is generated by the accretion-disk dynamo, applying different model assumptions for the dynamo tensor. Obviously, the difference in the field structure—difference in strength and geometry—will have a strong impact on the dynamics of the accretion disk and the disk wind or jet. In this section, we want to discuss the dynamical evolution of the accretion–ejection structure and compare the results for different dynamo models.

#### 3.3.1. Accretion and Ejection Rate

As pointed out in the previous sections, the dynamo tensor components that amplify the toroidal field ($\alpha_R$ and $\alpha_\phi$) work on longer timescales than the $\phi$ component of the dynamo tensor (which amplifies the poloidal magnetic field). Also, a larger dynamo component $\alpha_\phi$ leads to a higher magnetic diffusivity. In turn, this leads to a higher accretion rate, as shown in the top panel of Figure 5, as the disk diffusivity enables the replenishment of the disk matter that is lost from the inner disk (by accretion or ejection) to the outer disk regions.

On the other hand, the ejection rate only weakly depends on the strength of the $\phi$ dynamo, especially in the early stages of the evolution ($t \approx 100$), as shown in the bottom panel of Figure 5. While the inner regions reach a quasi-steady state for $t \gtrsim 500$, the ejection rate decreases until it reaches a quasi-constant level. This magnitude is higher for larger $\chi$, mostly because of the enhanced accretion rate.

We find that the ratio between the accretion and accretion rate is higher for lower $\chi$. This can be understood as follows. A higher strength of $\alpha_\phi$ leads effectively to a stronger and faster amplification of the magnetic field. A larger $\chi$, which is itself a consequence of applying an anisotropic dynamo tensor, leads to a stronger disk magnetization. Because of the diffusive quenching we apply (see Equation (19)), a higher disk

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4 Here, we point out that, as opposed to the simulations described in the lower panel of Figure 4, the top panel is marked by the absence of the multiple magnetic loops described in Section 3.2.2. For this reason, the three evolutionary stages we prefer to define consider both time and space.

5 Note that the disk gas pressure, in the absence of dynamo-inefficient zones, is subjected to only very small changes during the temporal evolution of the accretion disk.
magnetization implies a higher disk magnetic diffusivity, which in turn supports higher accretion rates.

For example, Figure 5 shows that for $\chi = 2.0$, about $\leq 50\%$ of the accretion mass flux becomes ejected. For $\chi < 0.5$, all of the matter accreted becomes ejected into the jet structure. This result is in nice agreement with resistive nondynamo-launching simulations (Zanni et al. 2007; Sheikhnezami et al. 2012), which showed a correlation between the disk magnetic diffusivity and the accretion–ejection rate ratio.

Once the poloidal field has become dynamo-amplified, the $R$ and $\theta$ components of the dynamo tensor can play a major role in the magnetic field evolution and, thus, also in the dynamics of accretion–ejection as they potentially induce dynamo-efficient zones in the disk. For simulations in which no dynamo-efficient zones emerge, differences in the toroidal magnetic field do not really impact on the poloidal field components, even on longer timescales.

On the other hand we find that a toroidal field reversal and the subsequent formation of multiple antialigned loops (and the corresponding dynamo-efficient zones) in the disk leads to a decrease in the accretion rate. The reason is the diffusive quenching we apply. At the locations where the toroidal field vanishes in the disk, the magnetic diffusivity also has a minimum (because of the low disk magnetization; see Equation (19)). A low diffusivity lowers the accretion efficiency.

We point out that the increase in the poloidal magnetic energy shown in Figure 2 is a value integrated over a control volume. Therefore, even if the overall magnetic energy is high, the formation of zones of low magnetic diffusivity leads to a decrease in the overall accretion rate. As a consequence, the disk mass that is lost by accretion and ejection cannot be efficiently replenished, therefore the accretion rate decreases with time. The ejection rate is also affected, but at later times. The most immediate consequence of the lower accretion rate is the formation of underdense and overdense zones in the accretion disk.

The radial distance of a dynamo-inefficient zone from the inner disk radius is strictly correlated with the timescale at which we observe a decrease in the accretion rate. This is the case, for example, for simulations $R_B$ and $R_A$ (see Figure 5, left). While in the former case the dynamo-inefficient zone already leads to a decrease in the accretion rate at time $t \approx 2000$, the latter case shows no difference from the simulation applying an isotropic dynamo tensor until time $t \approx 4000$. Note that the dynamo-inefficient zone is formed only at $t \approx 4000$, and therefore can impact the accretion and ejection rates only on a longer timescale (see Figure 4).

3.3.2. Jet Speed and Collimation

An immediate consequence of a variation in the dynamo tensor components is the jet kinematics. As pointed out by Stepanovs & Fendt (2016), a higher poloidal disk magnetization will lead to a stronger jet, for example, in terms of mass flux and velocity. We know from simulations applying a scalar dynamo model (Fendt & Gaßmann 2018) that the terminal jet speed is correlated with the strength of $\alpha_{0i}$; in particular, a stronger dynamo leads to a faster jet. Note that these properties—jet speed, mass flux, or collimation—are global properties and thus accessible in principle by observations, different from the intrinsic local conditions in the disk such as turbulence and dynamo action.

As for the evolution of the magnetic field, the three components of the dynamo tensor also have a different impact on the jet kinematics. When considering different magnitudes of the dynamo $\chi$, from our simulations we find a correlation similar to the one discovered in Stepanovs & Fendt (2016). It is a fact that a stronger $\phi$ component of the dynamo results in a stronger amplification of the poloidal magnetic field. As a direct consequence, as the midplane pressure shows only a very weak dependence on the dynamo model, a larger $\chi$ leads to a higher poloidal disk magnetization (see Figure 6). Consequently, with a higher disk magnetization, more magnetic energy is available to accelerate the outflow.

We show the terminal jet speed, here computed as the maximum speed at $R = 100$, as a function of the magnetization in Figure 6. This figure indicates a very clear trend, as proposed.
by Stepanovs & Fendt (2016). In addition, it demonstrates again the gain in magnetization for different parameters for the dynamo parameter. We find that the maximum jet speed reaches the Keplerian velocity at the inner disk radius. However, the maximum speed decreases for smaller $\chi$. This is shown also in Figure 7, where we compare the distribution of the jet poloidal velocity for different simulation runs.

The two other dynamo tensor components affect the evolution of the disk magnetization in terms of the generation (or nongeneration) of magnetic loops and/or dynamo-inefficient zones. Because minima in the magnetic field strength do only have a very minor impact on the overall disk poloidal magnetic energy (and therefore on the disk poloidal magnetization; see Figure 5), a difference in $\phi$ does not necessarily lead to a different jet. The main reason why the jet dynamics are not substantially changed, at least in the early stages of the jet formation and propagation, is that the magnetic field structure remains very similar in the innermost disk regions (see Figure 3). This is actually the field structure that is responsible for launching the strongest—and also collimated—jet component.

On the other hand, the dynamo-inefficient zones lead to a different disk mass distribution (see Figure 5), which naturally affects the evolution of the whole disk–jet system. In particular, we observe a more turbulent configuration of the poloidal magnetic field (see Figure 7), which leads to the ejection of a slower and less massive jet (i.e., with a smaller ejection rate, as shown in the bottom panel of Figure 5). The latter has been proposed already by Fendt (2006).

Another observable is the jet collimation as an imprint of the overall jet dynamics. There are several options on how to best define jet collimation. For example, in Fendt (2006), Pudritz et al. (2006), and Sheikhnezami et al. (2012), the degree of collimation has been computed as the ratio of the (normalized) mass fluxes in the axial and lateral directions, respectively. Another option is the pure opening angle. Here we choose a different way to measure the jet collimation quantitatively, taking advantage of the spherical coordinates we applied. More specifically, we compute the opening angle of the jet flow for which the jet has its maximum velocity (or mass flux). Comparing the angle obtained for different (spherical) radii, we obtain a gradual change that in particular demonstrates the process of collimation.

What we find from our dynamo simulations is essentially that the jet degree of collimation shows only a weak dependence on
the strength of the dynamo component \(\alpha_\phi\). This is maybe expected as we know that collimation depends on the profile of the disk magnetic field rather than its strength (Fendt 2006; Pudritz et al. 2006). Therefore, no significant differences are found in the jet collimation for a substantially isotropic dynamo, while an anisotropic dynamo in general leads to a lower degree of jet collimation (see Figure 7, left).

Another feature that impacts the degree of jet collimation is the presence of magnetic islands, and thus magnetized vortices. These loops severely disturb the accretion–ejection structure, enhance the turbulence in the outflow flow, and also affect the efficiency of mass ejection.

The toroidal magnetic field, which plays a leading role in jet collimation, is affected by \(\alpha_\theta\) and \(\alpha_\phi\). In particular, the existence of zones where the mean-field dynamo does not work efficiently leads to a more turbulent configuration of both the poloidal and toroidal magnetic fields (see Figure 7, right). Note, however, that jets also self-generate a substantial toroidal field that usually supports collimation (Blandford & Payne 1982). Here, the turbulent injection and the turbulent field structure hinder a regular jet toroidal field. Thus, a weak or nonisotropic dynamo will produce a less collimated jet (see again Figure 7, right). To summarize, the dynamo-inefficient zones lead to a more turbulent evolution of both the magnetic field and the hydrodynamical quantities, resulting in a more turbulent and less collimated jet structure.

3.4. Early Evolution

Because the target of this toy model is to investigate the effects of the different dynamo components on the launching process, we now discuss the impact of the tensor component \(\alpha_\theta\) in more detail. This mostly relates to the very initial evolution of the simulation.

A first result is that for \(0 < \psi < 3\), the evolution of the disk–jet system shows no difference when compared with a scalar dynamo. A likely explanation we find comes directly from the choice of the initial configuration of the magnetic field in combination with the induction equation. As the seed field is purely radial, there is no \(B_r\) component that can be coupled by a dynamo process. Therefore, in the initial evolutionary states, no contribution can be provided from \(\alpha_\theta\). As the system evolves, diffusive quenching takes place quite rapidly, leading to a quasi-steady state. Eventually, the dynamo effects are counterbalanced by magnetic diffusivity, and the component \(\alpha_\theta\) plays a minor role, just because they are weak and had no time to evolve.

However, when increasing \(\alpha_\theta\), as for simulation run \(th_{-A}\), its dynamo effect on the temporal evolution becomes stronger. The most important difference to the scalar dynamo simulations is the formation of multiple dynamo-inefficient zones within the accretion disk. As the magnetic field can be amplified only between the dynamo-inefficient zones, this further leads to multiple regions in the disk where the magnetic diffusivity does not grow (see Figure 4).

The reason why the early temporal evolution is mostly dominated by the other two dynamo components essentially depends on the initial magnetic field configuration. On one hand, this might look unphysical, as the long-term dynamo amplification of the magnetic field should not depend on its initial structure. On the other hand, a weak field seed must be present in order to initialize a mean-field dynamo effect.

Essentially, a toroidal initial magnetic field leads to the same results (see also Stepanovs et al. 2014). Similar to the case of a radial initial field, component \(\alpha_\phi\) is not involved in the initial temporal evolution of \(B_\phi\) and therefore is able to play a role only when the magnetic field has already saturated. Thus, the field evolution generated from a purely toroidal initial field leads to results similar to those obtained from a radial seed field.

This is in contrast to simulations starting from a vertical seed field. We find a strong impact on the evolution of the system because of the strong shear between the rotating disk and the nonrotating (at \(t = 0\)) corona (Fendt & Gaßmann 2018). In addition, this is amplified by the \(\alpha_\theta\)-dynamo effect of the magnetic field.

This can be nicely seen by our simulation \(th_1B\) applying a vertical seed magnetic field that is derived from a constant vector potential \(A_\phi = 10^{-5}\) and is applying an anisotropic dynamo with \(\psi = 0.1\). Here, the vertical initial field is able to affect, through the mean-field dynamo, the magnetic field evolution and amplification. A dynamo-inefficient zone is formed around \(R \approx 15\). A collimated outflow is launched, although the overall jet structure shows less collimation compared to the simulation with isotropic outflow (with a radial initial field).

4. Conclusions

We presented MHD dynamo simulations in the context of large-scale jet launching. Essentially, a magnetic field that is amplified by a mean-field disk dynamo is able to drive a high-speed jet. All simulations were performed in axisymmetry, treating all three vector components for the magnetic field and velocity. We applied the resistive code PLUTO 4.3 (Mignone et al. 2007); however, we extended it by implementing an additional term in the induction equation that considers the mean-field dynamo action.

Extending our previous works on mean-field dynamo-driven jets (Stepanovs et al. 2014; Fendt & Gaßmann 2018), here we essentially investigated the effects of a nonscalar dynamo tensor. We applied:

(i) various (ad hoc) choices for the dynamo tensor components (this paper, Paper I), but also

(ii) an analytical model of the turbulent dynamo theory (Rüdiger et al. 1995) that incorporates both the magnetic diffusivity and the turbulent dynamo term, connecting their module and anisotropy by only one parameter, the Coriolis number \(\Omega^*\) (see Paper II).

In particular, we obtained the following results:

(1) We disentangled different effects of the dynamo tensor components concerning the magnetic field amplification and geometry. We find that the strength of the amplification is predominantly related to the dynamo component \(\alpha_\phi\). The stability of the disk and the launching process can be affected by reconnection events. The field geometry that favors reconnection is mainly governed by the dynamo components \(\alpha_\phi\) and \(\alpha_\theta\).

(2) We find that the component \(\alpha_\phi\) is strongly correlated to the amplification of the poloidal magnetic field, such that a stronger \(\alpha_\phi\) results in a more magnetized disk, which then launches a faster, more massive, and more collimated jet. In contrast, the amplification of the poloidal field depends substantially on the existence of dynamo-inefficient zones,
which, subsequently, affect the overall disk-jet evolution, and thus accretion and ejection.

(3) We find that both a stronger dynamo component $\alpha_\theta$ and also a radial component $\alpha_R$ defined by $\phi < 0.8 \lor \phi > 1.5$, respectively, lead to the formation of dynamo-inefficient zones. The formation of dynamo-inefficient zones can also be triggered by a vertical component of the initial magnetic field, even for a weak dynamo component $\alpha_\theta$. A strong $\alpha_\theta$ component triggers the formation of the dynamo-inefficient zones predominantly in the inner disk region. Those loops in general lead to a different evolution of the disk dynamics, as these zones are dynamo-inefficient and prevent the accretion of material from the outer regions of the accretion disk to the inner disk that loses mass by accretion and ejection.

(4) We investigated how the actions of the three different dynamo components affect the jet structure, respectively. We find that the strength of the magnetic field has a minor influence on the jet speed and mass; however, the field geometry, in particular the disk magnetic field profile, matters a lot. For lower $\alpha_\theta$, or in the presence of dynamo-inefficient zones within the accretion disk, the magnetic field follows a different configuration (with a larger-scale magnetic structure compared to a more turbulent structure), which immediately affects the jet structure and collimation.

(5) We disentangled a clear correlation between the anisotropy of the dynamo tensor and the large-scale motion of the jet. In particular, dynamos working with a larger $\alpha_\theta$ produce a magnetic field that is able to drive faster jets. The reason is that these dynamos lead to stronger disk magnetization and thus provide more magnetic energy for launching. This result nicely couples to correlations between the disk magnetization and various parameters of the jet dynamics as found by Stepanovs et al. (2016).

(6) We investigated the formation of the so-called dynamo-inefficient zones within the accretion disk and their effect on the disk–jet connection. In particular, such zones are related to a toroidal field reversal with zero derivative, which leads to the formation of multiple loops in the disk. As a consequence, the poloidal magnetic field (in both the disk and the jet) follows a more turbulent evolution, forming, e.g., reconnecting magnetic loops, which affect the overall jet launching, the jet mass loading, and subsequently the jet propagation. These zones result from certain conditions for the dynamo action, i.e., certain combinations of the dynamo tensor components.

So far, we have looked for nonisotropic dynamos applying (ad hoc) choices for the dynamo tensor components. In our follow-up paper (Paper II), we will apply an analytical model of turbulent dynamo theory (Rüdiger et al. 1995) that incorporates both the magnetic diffusivity and the turbulent dynamo term, connecting their module and anisotropy by only one parameter, the Coriolis number $\Omega^*$.

We thank Andrea Mignone and the PLUTO team for the permission to use their code. All simulations were performed on the ISAAC cluster of the Max Planck Institute for Astronomy. We acknowledge many helpful comments by an anonymous referee that led to a clearer presentation of our results.

Appendix A

Test Simulations and Comparison to the Literature

In order to validate our implementation of the mean-field dynamo tensor in the latest version of PLUTO, we performed comparison simulations to the reference simulation of Fendt & Gaßmann (2018), now restricted to one hemisphere.

Note that while in Stepanovs et al. (2014) and Fendt & Gaßmann (2018) the dynamo term was simply coupled with the magnetic diffusivity, here, because of its hyperbolic nature, the $\alpha$ tensor is coupled with the standard hyperbolic MHD flux terms, with a correction due to the solenoidal condition of the magnetic field. Some minor differences in the magnetic field evolution seem to arise from the different implementation schemes; however, the overall evolution of the system shows very small differences in the strength of the physical processes at work.

Our simulation runs until $t = 30,000$, corresponding to $\approx 5000$ inner disk rotations. Figure A1 shows the evolution of the density and of the magnetic field lines. We can distinguish three different zones of evolution—the innermost disk, the outer disk, and the corona. The temporal evolution is in very good agreement with Fendt & Gaßmann (2018), evolving the same features.

Throughout the inner disk region, the magnetic field lines have the typical open field lines inclined with respect to the disk surface. This configuration is particularly favorable for a Blandford–Payne-driven outflow. The outer disk region is filled with magnetic loops, which are pushed outwards by the magnetic pressure gradient and thereby diffuses through the disk until it is filled with magnetic energy and a local steady state is reached.

In contrast to Fendt & Gaßmann (2018), we find that the poloidal magnetic energy saturates toward a somewhat lower level, but this is simply because our computational domain is smaller. Integrated over the whole disk, Fendt & Gaßmann (2018) find a saturation magnitude of $\approx 2 \times 10^{-3}$ (in code units), while here we reach a saturation value of $\approx 1.2 \times 10^{-3}$ (assuming that the lower hemisphere follows the same evolution as the upper hemisphere).

On the other hand, the accretion and ejection rates saturate at similar magnitude, and the accretion–ejection ratio also agrees with our previous studies (Fendt & Gaßmann 2018). This again strongly supports our conclusion that the different implementation schemes are identical.
This allows our test simulations to be compared to the reference simulation of Stepanovs et al. (2014).

We again fitted the simulation data points with a power law in order to extrapolate the power-law index $\beta_q$ and compare it with the radial distribution at $t = 0$. We find that the disk rotation remains Keplerian with $\beta_q = -1/2$. However, the radial profile of the density distribution changes substantially from $\beta_q = -3/2$ to $\beta_q = -4/3$ up to $R \approx 30$, while for larger radii, the power index is $\beta_q = -5/4$. As the total mass flux is conserved, the ejection of matter immediately changes the accretion rate over the disk and is thus related to the changes in the profiles of the mass fluxes. The radial (accretion) velocity follows a power-law index $\beta_{v_R} = -2/5$. Because we are reaching a longer run time than Stepanovs et al. (2014), we are now able to get rid of the oscillations and also the reversal found by Stepanovs et al. (2014) in the outer disk regions (as their magnetic field was not yet diffused across the whole accretion disk).

The power-law coefficient of the sound speed changes during $t = 0$ and $t = 10,000$ from $\beta_{c_s} = -1/2$ to $\beta_{c_s} = -3/7$, which tells us that the mean-field dynamo only slightly changes its strength due to the disk sound speed through the disk–jet evolution (see Equation (14)). This change does not lead to any strong net effect on the temporal evolution of the disk–jet system, therefore we again find a difference from our previous results (Stepanovs et al. 2014; Fendt & Gaßmann 2018).

Also, the angular magnetic field component $B_\theta$ follows the same power law, namely $\beta_{B_\theta} = -5/4$. Note, however, that we do not find the decrease in the outer disk regions ($R \geq 40$) as found in Stepanovs et al. (2014) simply because of our longer simulation time.

Overall, by quantifying the essential dynamical properties of our simulation results, we find perfect agreement with the previous results that are based on a numerically different implementation of magnetic diffusivity and mean-field dynamo.

### Appendix B

#### Control Volumes and Fluxes

Here we define how we integrate global quantities that are used throughout the paper. The accretion rate is calculated by integrating the net radial mass flux through the disk, defined by an opening angle $\theta_S \equiv \arctan(2H/r)$,

$$M_{\text{acc}}(R) = 2\pi R \int_{\pi/2}^{\pi/2-\theta_S} \rho v_R R d\theta,$$  

(B1)

while the ejection rate is calculated by integrating the outflow in the vertical direction (through the disk surface),

$$M_{\text{eje}}(R; \theta_S) = \int_{R_{\text{in}}}^{R} \rho v_\theta(R) 2\pi R dR,$$  

(B2)

respectively. The magnetic disk energy (poloidal or toroidal) is integrated from a radius of choice $R$ to the outer radius $R_{\text{out}}$, and from the disk midplane to the disk surface, defined by $\theta_S$. We thus consider the disk magnetic energy outside $R$ for our considerations,

$$E_{\text{mag}} = \int_{R}^{R_{\text{out}}} \int_{\pi/2-\theta_S}^{\pi/2} \frac{1}{2} B^2 \sin(|\theta|) 2\pi R^2 d\theta dR.$$  

(B3)

The so-called disk magnetic field (and also the disk magnetization) is simply calculated as the average value of the magnetic field, at
each radius, within the initial disk defined by $\theta_i \equiv \arctan(H/r)$,

$$B_{\text{disk}}(R) = \frac{1}{\theta_i} \int_{\theta_i}^{\pi/2} B(R, \theta) d\theta$$  \hspace{1cm} (B4)

while the so-called disk diffusivity is the average value of the diffusivity at a certain radius within the initial accretion disk,

$$\eta_{\text{disk}}(R, t) = \frac{1}{\theta_i} \int_{\theta_i}^{\pi/2} \eta(R, \theta) d\theta.$$  \hspace{1cm} (B5)

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