BIPOLAR JETS LAUNCHED FROM ACCRETION DISKS. II. THE FORMATION OF ASYMMETRIC JETS AND COUNTER JETS

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Received 2013 April 30; accepted 2013 July 7; published 2013 August 8

ABSTRACT

We investigate the jet launching from accretion disks, in particular the formation of intrinsically asymmetric jet/counter jet systems. We perform axisymmetric MHD simulations of the disk–jet structure on a bipolar computational domain covering both hemispheres. We apply various models such as asymmetric disks with (initially) different scale heights in each hemisphere, symmetric disks into which a local disturbance is injected, and jets launched into an asymmetric disk corona. We consider both a standard global magnetic diffusivity distribution and a novel local diffusivity model. Typical disk evolution first shows substantial disk warping and then results in asymmetric outflows with a 10%–30% mass flux difference. We find that the magnetic diffusivity profile is essential for establishing a long-term outflow asymmetry. We conclude that bipolar asymmetry in protostellar and extragalactic jets can indeed be generated intrinsically and maintained over a long time by disk asymmetries and the standard jet launching mechanism.

Key words: accretion, accretion disks – galaxies: active – galaxies: jets – ISM: jets and outflows – magnetohydrodynamics (MHD) – stars: protostars

Online-only material: color figures

1. INTRODUCTION

Jets as highly collimated beams of high-velocity material are a ubiquitous phenomenon in astrophysics. They originate from young stars, microquasars, or active galactic nuclei (AGNs), and are thus found over a wide range of luminosity and spatial scale. The common understanding is that magnetohydrodynamic (MHD) forces are responsible for launching, accelerating, and collimating these jets (Blandford & Payne 1982; Pudritz & Norman 1983; Uchida & Shibata 1985; Pudritz et al. 2007). Furthermore, it is clear that accretion and ejection are related to each other—one efficient way to remove angular momentum from a disk is to eject it vertically into a jet or an outflow (Ferreira & Pelletier 1995; Li 1995; Casse & Keppens 2002; Fendt 2006).

Observational data have confirmed the coexistence of bipolar jets in most jet-forming regions. Jet and counter jet appear, however, typically asymmetric in shape with very few exceptions (see, e.g., Mundt et al. 1990; Ray et al. 1996, 2007). One of the best studied sources is RW Aur. Woitas et al. (2002) find jet/counter jet velocity differences of about 50% close to the source and point out that this could be hint at variations intrinsic to the central engine. On the other hand, the outflow velocities seem to vary within a few years, and also the velocity asymmetry. Similarly, the jet/counter jet mass and momentum fluxes differ substantially. This was confirmed by Hartigan & Hillenbrand (2009) who find again that the redshifted flow is slower than the blueshifted flow. Also Melnikov et al. (2009) confirm the velocity differences close to the source and also find differences in the mass fluxes: \( M_{\text{red}} = 2.6 \times 10^{-9} M_\odot \text{yr}^{-1} \) > \( M_{\text{blue}} = 2.0 \times 10^{-9} M_\odot \text{yr}^{-1} \). Another example is DG Tau, where recent studies find the blue being faster and less collimated (Podio et al. 2011). However, the mass fluxes were found to be similar for jet and counter jet.

Liu et al. (2012) report a velocity ratio of jet to counter jet of about 1.34 and suggest different mass loss rates for both jet components. A recent spectroscopic study of PV Cep detects recent jet outbursts that are asymmetric (Caratti o Garatti et al. 2013). While the mass fluxes are again similar, the outflow velocities as well as the electron densities differ. Ellerbrook et al. (2013) find similar mean velocities in both lobes of the jets observed, while they report asymmetric variations in mass outflow rates and local velocities, and suggest that the jet launching mechanism on either side of the disk is not synchronized.

One exception is the jet source HH 212 that ejects an almost perfectly symmetric bipolar structure (Zinnecker et al. 1998), suggesting that the causal origin of jet knots is located close to the central engine. On the other hand, if jets form intrinsically asymmetrically, and if thus symmetric jets would need special conditions to be formed, we may ask: what is this kind of “natural” ejection process and what are the additional conditions for symmetry?

In the case of extragalactic jets, asymmetries between jet and counter jet are observed regularly (see, e.g., Laing 1988 or Urry & Padovani 1995 for an early review); a typical example may be M87 (Perlman et al. 1999). For highly relativistic parsec-scale jets, beaming may play the dominant role, depending on the inclination angle. For kiloparsec jets, which are expected to move slower with about 0.1c, beaming must be weak; however, extinction by the galactic disk will affect the observed bipolar jet structure. Jet propagation on kiloparsec scales will also be affected by the potentially asymmetric medium in the host galaxy (Jeyakumar et al. 2005). On the large scale, the AGN jets look remarkably symmetric in spite of clear asymmetries on smaller scales; typical examples are Cyg A (Carilli & Barthel 1996) and 3C 175 (Bridle et al. 1994). Due to the detection method used (measuring radio synchrotron emission), a direct measurement of jet velocities or mass fluxes is not possible for extragalactic jets.

A small number of papers have addressed the topic of jet asymmetries theoretically. To our knowledge, the first paper
on this topic was by Wang et al. (1992) applying steady-state MHD models to the disk–jet structure. Chagelishvili et al. (1996) discuss the MHD origin of the one sidedness state MHD models to the disk–jet structure. Chagelishvili on this topic was by Wang et al. (1992) applying steady-

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2. MODEL SETUP

As discussed in detail in Paper I, we model the launching of MHD jets from slightly sub-Keplerian disks, which are initially in pressure equilibrium with a non-rotating corona. Here we extend our approach to simulations on a computational domain including both hemispheres. The simulations are performed in a 2.5-dimensional setup in cylindrical coordinates (thus applying three-dimensional axisymmetry). We apply the MHD code PLUTO (Mignone et al. 2007, 2012), solving the time-dependent resistive MHD equations, as described in Paper I. We use the CENO3 algorithm as a third-order interpolation scheme for spatial integration (Del Zanna & Bucciantini 2002) together with a third-order Runge–Kutta scheme for time evolution and an HLL Riemann solver. For the magnetic field evolution, we apply the constrained transport method (FCT) ensuring solenoidality \( \nabla \cdot \mathbf{B} = 0 \).

2.1. Numerical Setup—Initial and Boundary Conditions

We apply the same initial disk structure and boundary conditions as before; however, we extend the disk–outflow system across the equatorial plane into both hemispheres. Above and below the disk, respectively, a hydrostatic corona is prescribed in pressure balance with the disk gas pressure (thus, implying a density jump and an entropy jump from the disk to the corona; see Paper I).

For the numerical grid, we have applied a few options (see Table 1). The subgrid across the equatorial plane is equidistant in order to allow for a proper treatment for magnetic diffusivity in PLUTO (see below). Above and below the equidistant subgrid, we usually attach a stretched subgrid in order to reach large distances from the disk. However, we have one simulation on a large equidistant grid in order to test jet launching in the case of a vertically extended profile of diffusivity.

We run simulations applying a different grid resolution. With the highest-resolution grid, we resolve the disk scale height at the innermost radius with up to \( \pm 10 \) grid cells.

2.1.1. Units and Normalization

We apply the same code units and normalization as in Paper I. Throughout this paper, distances are expressed in units of the inner disk radius \( r_i \), while \( p_{d,i} \) and \( \rho_{d,i} \) are the disk pressure and density at this radius, respectively.\(^2\) For the sake of comparison to previous papers, we may assume \( r_i = 0.1 \) AU for a protostellar jet and \( r_i = 10 \) Schwarzschild radii for an AGN jet. Velocities are measured in units of the Keplerian velocity \( v_{K,0} \) at the inner disk radius. Thus, by assuming smaller inner disk radii, the outflow speed will become larger. Time is measured in units of \( t_i = r_i/v_{K,i} \), which can be related to the Keplerian orbital period \( t_{K,i} = 2\pi r_i/v_{K,i} \). The (initial) disk aspect ratio \( \epsilon \) is the ratio of the isothermal sound speed to the Keplerian speed, both evaluated at disk midplane, \( \epsilon = c_s/v_{K} \). Pressure is given in units of \( p_{d,i} = \rho_{d,i} v_{K,i}^2 \). The magnetic field is measured in units of \( B_i = B_{K,i} \). We adopt \( v_{K,i} = 1 \), \( \rho_{d,i} = 1 \) in code units.

\(^2\) The index “i” refers to the value at the inner disk radius at the equatorial plane at time \( t = 0 \).
The main goal of this paper is to investigate the symmetry of bipolar jets launched from a diffusive accretion disk. It is therefore essential to carefully check the numerical setup, in particular the internal boundary conditions describing the sink, in order to prevent numerical artifacts generating asymmetries.

Our boundary conditions are, in general, equivalent to the conditions we have applied previously (see Figures 1 and 2 in Paper I). This holds in particular for the outflow boundary conditions (modified from the original code; see Porth & Fendt 2010; Porth et al. 2011), the radial and vertical outer boundaries, and the axisymmetry boundary condition along the rotation axis. For the outer disk radius, we apply an outflow boundary condition. We feed the inner jet launching area by accretion from the outer disk areas. Since compared to Paper I the physical extent of the computational domain is somewhat reduced, also the mass reservoir for disk accretion is smaller, which limits the disk mass evolution to comparatively shorter timescales. The sink boundary conditions allow the mass and angular momentum of the accreting material to be absorbed.

Obviously, for the present investigation the equatorial-plane boundary condition is omitted. The disk itself may now evolve into a warped structure, breaking the intrinsic hemispheric symmetry of the disk and the outflow. A disk midplane, if it exists, will not be necessarily along the equatorial plane. Further consequences are, e.g., that the electric currents can now flow across the disk midplane.

2.1.3. Initial Conditions

We apply the same basic initial conditions as in Paper I. In addition, a few extensions to this setup were made in order to break the symmetry and to govern an asymmetric evolution.

As in Paper I, the initial magnetic field is prescribed by the magnetic flux function \( \psi \) following Zanni et al. (2007),

\[
\psi(r, z) = \frac{3}{4} B_{z,0} r^2 \left( \frac{r}{r_1} \right)^{3/4} \frac{m^{5/4}}{(m^2 + (z/r)^2)^{5/8}},
\]

where \( B_{z,0} \) measures the vertical field strength at \((r = r_1, z = 0)\). The (initial) field tension is determined by the parameter \( m \). We apply \( m = 0.4 \) as in Paper I.

For the disk density and pressure, we apply the same distribution as Equations (6) and (7) of Paper I. However, we also run models with an initially asymmetric disk structure, considering a pressure scale height in the upper disk hemisphere \( \epsilon = \epsilon_{\text{up}} = 0.15 \) different from the pressure scale height in the lower disk hemisphere \( \epsilon = \epsilon_{\text{down}} = 0.10 \).

The initial disk corona follows the same distribution as in Paper I. In order to compare intrinsic effects of asymmetric bipolar launching with external effects, we have, however, also run comparison simulations with an initially symmetric disk, but with a disk corona of different density/pressure in the upper and lower hemispheres, respectively.

2.2. Magnetic Diffusivity

In order to extend the setup of Paper I into a truly bipolar configuration, we need to reconsider the model for the magnetic diffusivity \( \eta(r, z) \), which was by definition symmetric. An obvious constraint is that the prescription of diffusivity should not a priori influence the symmetry of the system. Since the disk structure may now evolve asymmetrically between the two hemispheres, the magnetic diffusivity must be able to follow such a disk evolution. Naturally, in asymmetric disks the disk midplane does not follow the equatorial plane.

In general, we assume the diffusivity to be anomalous and of turbulent origin. Since we do not resolve the disk turbulence from first principles (e.g., by resolving the magnetorotational instability, MRI), we apply a parameterized diffusivity distribution effectively following in principle an \( \alpha \)-prescription. We neglect geometrically more complex distributions such as, e.g., MRI-inactive dead zones of turbulent diffusivity or a more detailed local treatment of MRI turbulence (Gammie 1996; Pessah et al. 2007; Gressel 2010; Fromang et al. 2013; Lesur et al. 2013; Bai & Stone 2013).

As detailed in Paper I, we assume a diagonal diffusivity tensor with the non-zero components \( \eta_{\phi\phi} \equiv \eta_\phi \) and \( \eta_{rr} = \eta_{zz} \equiv \eta_\theta \), where we denote \( \eta_\phi \) as the poloidal magnetic diffusivity and \( \eta_\theta \) as the toroidal magnetic diffusivity.
as the toroidal magnetic diffusivity. The anisotropy parameter \( \chi = \eta_\|/\eta_\perp \) quantifies the different strength of diffusivity in poloidal and toroidal directions. Here, we apply \( \chi = 3.0 \) for all simulations.

For the diffusivity profile, we have investigated several options which we will discuss in the following. Table 2 compares the parameter setups for the different magnetic diffusivity distributions applied in our simulations.

2.2.1. A Global Magnetic Diffusivity Prescription

A first option for the diffusivity distribution is to mirror the profile applied in Paper I along the equatorial plane,

\[
\eta_\| (r, z) = \alpha_{m,1} f_1 (r, z) \equiv \alpha_{m,1} \nu_{A,0} H_0 \exp \left(-\frac{2 z^2}{H_0^2}\right),
\]

(2)

with the Alfvén speed \( \nu_{A,0} \equiv \nu_A(r, z = 0) \) and the disk thermal scale height \( H_0 \equiv H(r) = \epsilon r = c_s(r, z = 0)/v_k(r, z = 0) \), measured at the midplane and at time zero. Essentially, the diffusivity profile (Equation (2)) is geometrically tied to the equatorial plane, potentially inducing hemispheric symmetry of the disk and the outflow.

As in Paper I, we allow for a diffusivity scale height \( H_0 \) larger than the thermal scale height \( H_0 \) (see the discussion in Paper I). In fact, Gressel (2010), who investigated the MRI-induced turbulence of accretion disks by high-resolution box simulations, finds an increasing level of turbulence with disk height. His simulations attest a maximum level of turbulence at about two to three disk pressure scale heights, which are in nice agreement with our model approach. More recent simulations by Beckwith et al. (2011) and Simon et al. (2013) indicate similar scale heights for the turbulent stresses.

In Equation (2), both \( \nu_A \) and \( c_s \) can be chosen time-independent (as in Paper I) or evolving in time (as, e.g., in Murphy et al. 2010). For comparison, in the present bipolar reference simulation (denoted by rfr-no-f1) we apply a magnetic diffusivity profile \( f_1 (r, z) \) fixed in time.

2.2.2. A Local Prescription of Magnetic Diffusivity

In the present paper, we investigate an asymmetric evolution of the disk–jet structure in both hemispheres. Therefore, we need to apply a new, truly bipolar setup in which the leading parameters for the diffusivity profile are no longer defined with respect to the equatorial plane. It is clear that for an asymmetric disk evolution, a hypothetical disk midplane will not coincide with the equatorial plane. Therefore, in order to consider the spatial evolution of the asymmetric disk structure, it is essential to apply a local prescription of diffusivity.

Among other options, one possibility is to relate the magnetic diffusivity to the local pressure or density. One may assume that a magnetic diffusivity profile will follow a power law \( \eta \sim \rho^{1/3} \). In this case, the diffusivity is proportional to the local sound speed \( \eta \propto c_s \). We may generalize the power law:

\[
\eta_\| = \alpha_{m,2} f_2 (r, z) \equiv \alpha_{m,2} \rho^{\Gamma}.
\]

(3)

The profile with \( \Gamma \approx 1/3 \) results in a relatively broad vertical diffusivity profile, wider than the previously used exponential profile (Equation (2)). It has been shown that this will impact the jet acceleration and collimation (Fendt & Čemeljić 2002), maybe more than it affects the launching mechanism itself. For comparison, we have also applied different power laws, such as \( \Gamma = 2/3 \), or even steeper profiles. The drawback of the simple profile (Equation (3)) is that the diffusivity decreases with radius. The outer disk material is therefore strongly coupled to the magnetic field, leading to a superefficient angular momentum removal, rapid accretion, and, thus, a short lifetime of the outer disk. On the other hand, in the case of MRI-driven turbulence, the MRI activity is expected to cease for large radii, where “dead zones” for active accretion may exist (Gammie 1996; Beckwith et al. 2011; Simon et al. 2013).

Therefore, we have favored another option for the magnetic diffusivity which allows for a disk diffusivity increasing with radius,

\[
\eta_\| (r, z) = \alpha_{m,3} f_3 (r, z) \equiv \alpha_{m,3} \frac{H_\| (r, z)}{H_0} \frac{1}{1 + \frac{H_\| (r, z)}{H_0}}.
\]

(4)

Thus, we apply a density-weighted “local disk scale height” \( H_\| (r, z) \) mainly following the local sound speed in the gas flow,

\[
H_\| (r, z) = \rho (r, z)^{\nu_k} \left[ \frac{\rho (r, z)}{\nu_k (r)} \right]^{\nu_k (r)},
\]

(5)
Figure 1. Time evolution of the bipolar jet–disk structure for reference simulation \textit{rfr-no-f1} applying a fixed-in-time and fixed-in-space magnetic diffusivity distribution (Equation (2)). Shown are the evolution of mass density (color) and the poloidal magnetic field (contours of the poloidal magnetic flux $\Psi(r, z)$) for the dynamical time steps $t = 0, 100, 1000, 2000, 3000$.

A color version of this figure is available in the online journal.

and a quenching term in order to avoid numerically problematic diffusivities. We typically choose $\tilde{H}_0 = 0.5$, $\sigma_\rho = 3/2$, $\sigma_r = 3/2$, or $\tilde{H}_0 = 0.1$, $\sigma_\rho = 3/2$, $\sigma_r = 5/2$ (simulation run \textit{gda-sh-f3}).

Essentially, both prescriptions for the magnetic diffusivity (Equations (3) and (4)) allow for a smooth transition from accretion to ejection, and, also to follow the changes in the local disk structure. As the outflow density decreases along the streamlines, also the outflow diffusivity decreases. For low densities toward the asymptotic outflow, the ideal MHD will be approached. A physical motivation can be the following. Turbulently diffusive disk material is lifted from the disk and is further accelerated along the outflow while, however, the turbulent motions decay. In Paper I, we have estimated the scale height, where the turbulence will be damped to be several disk thermal scale heights (see the discussion above citing Gressel 2010).

2.3. Parameter Runs

A summary of all parameter runs is given in Table 2 listing the following simulations. The reference run with no symmetry breaking and diffusivity profile $f_1$, denoted by \textit{rfr-no-f1}, a run with global disk asymmetry set by the initial disk scale height and diffusivity profile $f_1$, denoted by \textit{gda-sh-f1}, a run with local disk asymmetry set by a local overpressure inserted height and diffusivity profile $f_1$, denoted by \textit{lda-op-f1}, a run with no disk asymmetry injected in an asymmetric ambient gas and diffusivity profile $f_2$, denoted by \textit{na-ism-f2}, a run with global disk asymmetry set by the initial disk scale height and diffusivity profile $f_3$, denoted by \textit{gda-sh-f3}, and similarly, but with high resolution, the run \textit{gda-sh-f3-hr}. In the lower part of the table, we have listed a number of additional simulations with similar parameters which are briefly discussed in this paper, but for which we do not show plots. These are simulations that have been run to further test the diffusivity profile $f_2$ and are denoted by \textit{tst-...-f2} (see below), or further approve the simulations with local injection, denoted by \textit{tst-...-f1}.

3. TEST CASE—LAUNCHING SYMMETRIC BIPOLAR JETS

We first present jet launching simulations resulting in symmetric bipolar jets. The first example is the evolution of jets following a magnetic diffusivity description fixed in time and space (simulation \textit{rfr-no-f1}). This case serves as a reference simulation for this paper and also allows for a comparison with the one-hemispheric simulations of Paper I. This inflow–outflow evolution is shown in Figure 1.

As we see, the outflow evolves perfectly symmetrically in both hemispheres for almost 3000 dynamical time steps, until the outer disk starts to deviate from symmetry due to numerical effects. However, even for these late time steps, the inner disk, which is the main jet launching area, is still highly symmetric.

In particular, this can be seen in the time evolution of the mass fluxes. Figure 2 shows the accretion rate and the ejection rates of the two hemispheres. The ejection rates are integrated from $r = 1.5$ to $r = 10$ at $z = \pm 3H(r)$, while the accretion rate is integrated from $z = -3H$ till $z = 3H$ at $r = 10$. It is not surprising that due to the symmetric diffusivity profile which remains fixed in time together with symmetric initial and boundary conditions, we find a symmetric evolution of both
the disk and the outflows. The symmetric bipolar jet structure we obtain does mainly serve as a test case for the numerical setup, in particular for the sink boundary conditions. While the symmetric evolution is expected on physical grounds, even a little numerical failure in the setup would lead to asymmetry quickly. The reference simulation demonstrates that our setup is properly defined without such failures. Naturally, in the very quick time evolution of the mass fluxes for reference simulation rfr-no-f1. Shown are the evolution of the accretion rate and the mass ejection rates from the upper (solid line) and lower (dashed line) disk surfaces (all in code units). Ejection rates are measured in the control volumes with \( r_1 = 1.5 \) and \( r_2 = 10.0 \), while the accretion rate is vertically integrated at \( r = 10.0 \).

Comparing to the one-hemispheric simulation in Paper I (reference run case1), the mass fluxes we measure now are quite similar. As an accretion rate for the bipolar simulation we measure \( M_{\text{acc}}(t = 2000) \approx 0.03 \), while for the previous reference case1 it was \( M_{\text{acc}}(t = 2000) \approx 0.015 \). The outflow rates for the bipolar simulation are \( M_{\text{eject}}(t = 2000) \approx 0.01 \) in each direction, which agree with the \( M_{\text{eject}}(t = 2000) = 0.007 \) in Paper I nicely. However, due to the smaller grid extension, the disk mass reservoir is smaller, leading to a faster decay of the disk mass (see the outer disk structure at \( t = 3000 \)). Therefore, disk accretion and mass ejection decay faster as well.

4. BIPOLAR JETS FROM ASYMMETRIC DISKS

Here we present simulations in which we have disturbed the internal hemispheric disk symmetry. We have applied two options for disturbing the disk symmetry—either by prescribing a global asymmetric initial state or by injecting a localized overpressure at certain time. In this section, we apply a global, symmetric magnetic diffusivity prescription (Equation (2)), constant in time. The next section will deal with simulations using a local diffusivity description.

4.1. Asymmetric Disks

Obviously, in our simulations the disk asymmetry is put in by hand. It is therefore interesting to discuss possible physical processes leading to a disk asymmetry at first hand. A self-consistent treatment of the asymmetric disk evolution is beyond the scope of this paper and will probably require a higher resolution of a three-dimensional treatment in order to resolve possible disk instabilities.

Warping can be present whenever a misalignment is present in the disk or a flat disk becomes unstable by external forces (Ogilvie & Latter 2013). Processes leading to an asymmetric disk evolution may be distinguished in effects external to the disk or disk internal effects, such as a central stellar magnetic field (Peiffer & Lai 2004; Romanova et al. 2004, 2013; Lai & Zhang 2008), tidal effects by a companion star (Larwood et al. 1996; Facchini et al. 2013), or a black hole (Bardeen & Petterson 1975; Pringle 1992), or radiation pressure from the central source (see, e.g., Maloney et al. 1996, 1998; Wang & Li 2012; Wu et al. 2013).

4.2. An Asymmetric Disk Scale Height

Option I is to prescribe an initial disk structure with a global pressure asymmetry between the two disk hemispheres. We achieve this by applying a different thermal disk scale height for the initial disk in each hemisphere. In simulation gda-sh-f1, we have applied \( \epsilon = H/r = 0.1 \) for the upper hemisphere and \( \epsilon = H/r = 0.15 \) for the lower hemisphere. Consequently, we have a density and a pressure jump across the equatorial plane, \( \Delta P/P = 0.2 \).

The disk turns into an asymmetric evolution right from the beginning, evolving into a warped structure along the midplane (Figure 3). A series of warps are visible along the disk with warp amplitudes of a few local disk scale heights. After about 1000 dynamical time steps, the warp amplitudes start to decrease—first along the inner disk, while the outer disk is still in a warped state.

The disk asymmetry is reflected in the jet evolution. Along with the initial asymmetric disk evolution, the jets launched from the inner disk are asymmetric as well. This is clearly visible in the poloidal magnetic field structure (Figure 3), but also in the mass fluxes we measure. We also note a different timescale in the jet propagation. The upper jet reaches the grid boundary earlier than the lower jet which is delayed by about \( \Delta t = 10\% \) (note that this obviously depends on the grid size, and happens rather early at \( t < 50 \)).

Figure 4 shows the time evolution of the mass fluxes. Comparing the two jet fluxes we find that at early stages, \( t < 700 \), the lower outflow carries about 80% of the mass load of the upper outflow. From \( t \approx 700–1500 \), an asymmetric inflow–outflow system is established. The differences in mass flux are now about 10%. After \( t \approx 1500 \), the variation in the outflow rates decreases, and the inner disk has established a symmetric structure.

The same behavior is also visible in the radial velocity distribution. Figure 5 shows the complex velocity field of the

Note, however, the axisymmetric setup.
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Figure 3. Time evolution of the bipolar jet–disk structure for simulation gda-sh-f1 applying a fixed-in-time diffusivity profile (Equation (2)) and evolving from an asymmetric initial state with different thermal scale heights for the upper and lower disk hemispheres, $\epsilon_{\text{up}} = 0.15$ and $\epsilon_{\text{down}} = 0.1$. We show the evolution for time $t = 0, 100, 1000, 2000$ of the mass density (colors) and the poloidal magnetic field (lines), i.e., contours of the poloidal magnetic flux $\Psi$.

(A color version of this figure is available in the online journal.)

Figure 4. Time evolution of the mass fluxes for simulation gda-sh-f1. Shown are the evolution of the accretion rate and the mass ejection rates from the upper (solid line) and lower (dashed line) disk surfaces (all in code units). Ejection rates are measured in the control volumes with $r_1 = 1.5$ and $r_2 = 10.0$ (middle) while the accretion rate (left) is integrated at $r = 10.0$ (left). For comparison, we show the asymptotic (vertical) mass fluxes through the jets, integrated from $r_1 = 2.0$ to $r_2 = 50.0$ at $z = 75, 75$ (right).

inner disk. Accretion starts first in the innermost disk regions. The disk asymmetry is well reflected in the $v_r$ distribution. However, after $t = 1000$ the system turns to a symmetric geometry. We also see the highest accretion velocities in the upper disk layers. One is tempted to indicate this as layered accretion; however, due to the higher density in the lower disk layers, the radial mass flux is in fact equally distributed over the disk height.

We believe that the disk evolution into a symmetric state is due to two reasons: firstly by the restoring force of the symmetric gravitational potential and secondly by the symmetry and the time-independent prescription of the magnetic diffusivity. The latter is interesting since it is an indirect effect only, as the magnetic diffusivity does not provide a force term in the MHD equations which could directly reconfigure the disk structure. The diffusion timescale $\tau_\eta = \Delta Z^2 / \eta \approx H^2 / \eta \approx 10^2$ is about 100 dynamical timescales for $\eta \approx 0.1$, $H \approx 1$ at radius $r = 10$.

So far we have concentrated on the inflow–outflow dynamics close to the jet launching region. We now consider the evolution further downstream of the outflows. This is interesting as it is this part of the outflow which is in principle accessible by the observations. Figure 4 (right) shows the mass fluxes of jet and counter jet far from the source, integrated from $r = 2$ to $r = 50$ at $z = 75$ and $z = -75$, respectively. We find that the mass flux asymmetry of about $5\%$ at the launching region propagates to the asymptotic region, where we find a similar mass flux.
Figure 5. Disk evolution for the jet-disk system gda-sh-f1. Shown is the radial velocity $v_r(r,z)$ for the dynamical time steps $t = 40, 100, 1000, 2000$ (from top) indicating the asymmetric evolution of the disk and the disk wind. (A color version of this figure is available in the online journal.)

difference. The time lag between the launching region and the asymptotic domain of about $\Delta t = 300 = 1200 - 900$ can be explained by the propagation period of the accelerating outflow of velocities $v_z \simeq 0.2-0.8$ propagating a distance $\Delta z = 75$. The maximum jet velocity is achieved along a narrow cone between $r = 10$ and $r = 20$ (at $z = 100$). The outflow velocity of the lower jet is about $1.2$, while for the upper jet it is $1.1$ times the Keplerian speed at the inner disk radius. The bulk of the mass flux in the outflow is, however, confined within a layer between $r = 15$ and $r = 30$. This bulk flow is moving with $v_z \simeq 0.4$, but has a mass flux a factor four higher compared to the high-speed jet within the narrow cone discussed before.

Figure 6. Time evolution of the inner disk accretion for simulation lda-op-f1. Shown is the evolution of the radial velocity (color) for the dynamical times $t = 400, 404, 420, 500, 2800$. Simulation lda-op-f1 applies a fixed-in-time diffusivity distribution (Equation (2)). An overpressure is added at $t = 400$ in the upper disk hemisphere at radius $r = 12$, lasting for a few rotations. (A color version of this figure is available in the online journal.)
4.3. A Localized Disk Asymmetry in One Disk Hemisphere

We start simulation 1da-op-f1 with a symmetric initial disk structure. As for reference simulation rfr-no-fa1, the outflows evolve into a symmetric jet–counter jet structure (Figure 7). However, at $t = 400$, when a quasi steady state of the inner inflow–outflow structure is reached, we disturb the symmetry of the disk structure by inserting a localized overpressure in the upper disk hemisphere (see Figure 6). This injection is localized within a box of size $(\Delta r \times \Delta z) = (1.5 \times 0.4)$ located at $(r, z) = (11.25, 1.2)$, and is switched on for $\Delta t = 20$, corresponding to 0.1 of an orbital period at this radius. The injected material has on average a 20 times higher density and a 2000 times higher pressure compared to the ambient disk material.

The injected material disturbs the disk symmetry as it expands across the disk. The disturbance is slowly advected into the jet launching region. However, we observe that the disk asymmetry decays faster than it is advected along the disk into the jet launching area at small disk radii. This is easy to understand. The expansion happens roughly with sound speed, while the advection happens with a subsonic speed $v_r \simeq \epsilon v_{\text{Kep}}$. Figure 6 shows the disk accretion velocity evolution. The expansion following the initial injection first penetrates the whole disk in the vertical direction, before the overdensity is advected slowly inward. Thus, the symmetry of the inner disk is not really affected by the injection in the upper hemisphere. The dominant outflow from the inner disk is only weakly affected. When the size of the launching area reaches the injection site, the asymmetry has almost disappeared.

A slight structural asymmetry is visible (see the sixth flux surface contour in Figure 7) which has propagated from the disk surface to a distance $z \simeq 70$ at this time.

That part of the outflow which originates from the initially asymmetric part of the disk (where the blob was injected) starts indeed asymmetric. As the mass flux launched from larger disk radii is small compared to the jet launched from the inner part, the total jet mass fluxes into the two hemispheres differ only marginally. Figure 8 shows the time evolution of the mass fluxes. The ejection rates are integrated along the disk surface between $r = 1.5$ and $r = 10$, thus inside the radius where the overpressure is injected. Still the expanding material disturbs the launching site such that we see a 5%-10% effect in the outflow mass fluxes. This effect is somewhat delayed from the time of injection since the disturbance needs time to be advected inward. The asymptotic outflow rates, integrated from $r = 2$ to $r = 50$ at $z = \pm 75$, show only marginal differences, however. Note that the flux surface passing through $(r, z) = (50, 75)$ anchors at $r = 5$ in the disk, and thus in a weakly disturbed region.

In summary, although the initially symmetric disk is clearly disturbed by the asymmetric injection, the asymmetry decays faster than it propagates to the inner jet launching site. Thus, the asymptotic, collimated jet, which originates in the inner disk area, is only marginally asymmetric. The main mass flux is launched along the innermost field lines.

5. A LOCAL MAGNETIC DIFFUSIVITY MODEL AND BIPOLAR JET LAUNCHING

In order to allow the disk and the jet to follow a truly asymmetric evolution, without the restoring effects of a symmetric, global magnetic diffusivity distribution, we have applied a local description of diffusivity which follows the local evolution of the disk and the outflow (see Table 2). The simulation discussed here also applies an initial disk asymmetry.

The results of this section are somewhat preliminary, as a physically self-consistent parameterization for a local magnetic (turbulent) diffusivity is not available. In order to investigate the main features of the local approach, we have applied two diffusivity distributions. One follows Equation (3), the other one refers to Equation (4). The main difficulty one faces with a local description is a feedback mechanism such that low densities...
Disk accretion

Jet mass fluxes close to disk

Jet mass fluxes at top boundary

Figure 8. Time evolution of the mass fluxes for simulation lda-op-f1. Shown are the evolution of the accretion rate and the mass ejection rates from the upper (solid) and lower (dashed) disk surfaces (all in code units). Ejection rates (middle) are measured in the control volumes with \( r_1 = 1.5 \) and \( r_2 = 10.0 \), while the accretion rate (left) is integrated at \( r = 5.0 \). For comparison we show the asymptotic (vertical) mass fluxes through the jet, integrated from \( r_1 = 2.0 \) to \( r_2 = 50.0 \) at \( z = -75, 75 \).

Figure 9. Time evolution of the inner disk structure for simulation run gda-sh-f3. Here the diffusivity profile is prescribed as a local Equation (4). Shown are the evolution of magnetic diffusivity (color) and the poloidal magnetic field (contours of the poloidal magnetic flux \( \Psi(r,z) \)) for the dynamical time steps \( t = 50, 500, 1000 \) (clockwise from upper left). For comparison, the diffusivity profile of the reference simulation is shown (lower left). (A color version of this figure is available in the online journal.)

lead to a lower diffusivity, thus stronger matter–field coupling, thus more efficient angular momentum removal and a faster accretion, leading to even lower densities.

Moreover, the local prescription (Equation (3)) gives a diffusivity profile decreasing with radius, contrary to the usual choice in the literature, Equation (2), which is also realized in Equation (4).

We find that the feedback mechanism is most efficient in simulations with a strong coupling between diffusivity and density, such as a diffusivity profile (Equation (3)). These are simulations tst-...-f2 in the second part of Table 2 for which we do not show visualizations. In the extreme cases, we have observed a feedback mechanism strong enough to lead to a temporary gap opening over the innermost disk radii. As a result, the accretion in this area is strongly episodic, and subsequently also the mass ejection. We will present a detailed investigation of this feedback mechanism in a forthcoming paper.

Applying a more sophisticated prescription of the magnetic diffusivity following Equation (4), the overall accretion–ejection evolution is more similar to the picture established by the simulations using a global diffusivity profile. The fundamental difference is, however, the longer
lasting and more persistent asymmetry in the disk and the outflows.

Figure 9 shows the time evolution of magnetic diffusivity and the magnetic field for a simulation with a more sophisticated local setup of magnetic diffusivity following Equation (4). This is simulation *gda-sh-f3*. We clearly see how the magnetic diffusivity follows the structure of the disk and the outflow. The disk “warping” is seen in the diffusivity distribution as well. Mass loading and matter–field coupling depend on the local disk properties (defined by Equation (4)). Since the diffusivity profile is broader in the vertical direction, mass loading and angular momentum removal are more efficient, establishing higher accretion and outflow rates—in agreement with our results in Paper I.

The outflow mass fluxes develop a clearly asymmetric structure. Figure 10 shows the mass fluxes integrated from $r = 1.5$ to $r = 10$, along a surface $z = 0.3r$ parallel to the initial disk surface. Jet and counter jet injection mass flow rates differ substantially—by about 30% over at least 1000 time steps. The asymptotic jet mass fluxes integrated from $r = 2$ to $r = 40$ along a surface $z = \pm 50$ are somewhat lower (due to the fact that some of the material injected into the outflow leaves the grid in a radial direction) and still differ at a 15% level.

So far we have discussed the outflow mass fluxes. Observationally interesting is also the velocity structure and evolution of the outflow. Figure 11 shows the velocity evolution of the outflow launched in simulation *gda-sh-f3*. We show the time evolution of the mass-weighted jet velocity, integrated along $z = 50.0$ for the inner, fast jet for $r < 20.0$, and the surrounding slower outflow for $r > 20.0$. Such a distinction can be motivated by Figure 11 (bottom) which shows the profile of the vertical velocity across the jet (at $z = \pm 50$). Note that the mass-weighted velocity is somewhat smaller than the actual velocities. The fast jets, launched from the inner part of the disk, show time-dependent and bipolar asymmetric behavior. The velocities reach maximum values just below the Keplerian rotation speed at the launching site. The velocity asymmetry between jet and counter jet is even more intriguing than the asymmetry seen in the mass fluxes. Interestingly, the velocity asymmetry changes sign at about $t = 850$, a feature not visible in the mass fluxes in Figure 10. It would be very interesting to perform longer-lasting simulations to see whether the velocity asymmetry reverses at later times.

We have performed another simulation with the same setup but higher resolution (simulation *gda-sh-f3-hr*; see below) clearly confirming the results discussed in this section.
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Figure 12. Time evolution of the bipolar jet–disk structure for simulation na-ism-f2. Here the outflow is launched in hemispheres with different density contrast. Shown are the evolution of mass density (color) and the poloidal magnetic field (contours of the poloidal magnetic flux \(\Psi(r, z)\)) for the dynamical time steps \(t = 40, 100, 500, 1000, 3000\).

(A color version of this figure is available in the online journal.)

Figure 13. Time evolution of the mass fluxes for simulation na-ism-f2. Shown are the evolution of the accretion rate and the mass ejection rates from the upper (solid line) and lower (dashed line) disk surfaces (all in code units). Ejection rates (middle) are measured in the control volumes with \(r_1 = 1.5\) and \(r_2 = 10.0\), while the accretion rate (left) is integrated at \(r = 5\). For comparison, we show the asymptotic (vertical) mass fluxes through the jet, integrated from \(r_1 = 3.0\) to \(r_2 = 40.0\) at \(z = \pm 80\).

6. IMPACT OF THE ENVIRONMENT

This section is devoted to the question whether jet asymmetries are due to external or internal interaction. It seems clear that the conditions in the ambient gas the jet is penetrating will certainly affect the expansion and the collimation on the large scales.

Here we consider a model setup, where a bipolar outflow is launched intrinsically symmetrically from a (symmetric) disk into an asymmetric star–disk corona (simulation run na-ism-f2). We do this by prescribing a different initial coronal density (and thus pressure), simply choosing \(\delta_{up} = 10^{-4}\) and \(\delta_{up} = 10^{-3}\).

Figure 12 shows the time evolution of density and magnetic field for such a setup. As expected, the initial outflow is highly asymmetric, more than the cases where the outflows are disturbed by intrinsic disk asymmetries. However, the asymmetry is clearly transient. As soon as the initial coronal material is swept out of the computational domain, the accretion–ejection system returns to hemispheric symmetry. This is easily visible in the mass flux evolution (Figure 13), showing a drastic asymmetry during the first hundreds of time steps. Again it is interesting to see various time lags in the flow evolution: firstly, the time lag between the launching time of the initial asymmetry and the arrival time of asymmetric features in the asymptotic region. Secondly, the time lag between the time step when the disk–outflow system has returned to symmetry, and the time when the outflow symmetry has reached the asymptotic jet.

Note that in these simulations, the disk structure itself stays rather symmetric, unlike the simulations discussed above, where we see the disk evolving into a warped structure. However, we can see back reaction of the asymmetric outflow onto the disk structure in this approach. The two asymmetric jets drive a different electric current system in both hemispheres. Both the current systems are connected within the disk—subsequent MHD forces therefore slightly distort the symmetric...
hydrodynamic disk structure. The deviation from symmetry is, however, not as strong as for the previous cases where we start with an initially asymmetric disk.

Clearly, on large timescales, the structure of the external medium may affect the jet symmetry on the propagation scales, even if the jet symmetry close to the disk is not perturbed. This is expected for protostellar jets originating deep in molecular clouds or AGN jets penetrating the central regions of a galaxy.

7. RESOLUTION STUDY

Finally, we present example results of our resolution study in brief. Figure 14 shows the inner part of the disk at a dynamical time step $t = 130$ when the inner disk evolution is rather violent. We compare simulation gda-sh-f3 applying our standard resolution of $\Delta r = 0.0417$, $\Delta z = 0.0333$ in the disk, with simulation gda-sh-f3-hr, applying a three times higher resolution, $\Delta r = 0.0165$, $\Delta z = 0.00824$. Although the higher resolution run shows somewhat more substructure, such as internal shocks or a more peaked diffusivity distribution for $r < 3$, the main features of the disk dynamics are just the same. In particular, the disk height is similar, as well as the structure and opening angle of the magnetic flux surfaces. Also the radially structured features of the outflow close to the disk are very similar in both hemispheres. The overall mass fluxes measured in the simulations are similar as well; however, the jet–counter jet asymmetry is somewhat more pronounced in the high-resolution run.

We note that high-resolution simulations are particularly difficult to manage in diffusive MHD since the CF time stepping condition $\Delta t \leq (\Delta x)^2/\eta$.

8. OBSERVATIONAL RELEVANCE

Here we summarize our numerical results concerning their observational relevance. We have found that asymmetries in an accretion disk do propagate into the outflow being launched and result in an asymmetric outflow. We have presented possible physical processes discussed in the literature leading to a disk asymmetry, i.e., a disk warping. To us, asymmetric disks seem to be a natural consequence of the dynamic disk evolution. However, in our model setup the asymmetry is put in by hand as an initial condition.

Depending on the kind of asymmetry, a time lag between disk asymmetry and the onset of outflow symmetry is found. The time lag is triggered by the timescales of advection and diffusivity in the disk, and also (in addition) by the timescales for jet acceleration and propagation.

We observe differences in mass flux between jet and counter jet of up to 30%. Similarly, we found a velocity difference of up to 40%. Our numerical values may be compared to observations of RW Auriga detecting outflow mass fluxes of $M_{\text{red}} = 2.6 \times 10^{-9} M_\odot \text{yr}^{-1}$ for the red and $M_{\text{blue}} = 2.0 \times 10^{-9} M_\odot \text{yr}^{-1}$ for the blue lobe (Mel’nikov et al. 2009). This asymmetry in the mass fluxes is confirmed by Ellerbroek et al. (2013).

Considering their claim that the launching mechanism on either side of the accretion disk is not synchronized, we would interpret the launching process in our model indeed as asymmetric, but synchronized by the diffusive dynamical evolution of the disk warps. Also the velocity differences observed, e.g., the velocity ratio of 1.34 in DG Tau (Podio et al. 2011; Liu et al. 2012), are close to our numerical values.

It is clear that our numerical numbers are model dependent; however, they seem to be generic to all of our models applying different parameters. What is also important is that these number values change in time. Velocity differences of 20%–30% may lead to propagating shocks along the jet. For the timescale of these velocity fluctuations, we find a typical value of about $500–1000$ dynamical timescales, which is equivalent to say 100–200 inner disk Keplerian periods or about five years (for $r_i = 0.1 \text{ AU}$ and $M_* = 1 M_\odot$). This timescale is in the range of the observed 10–50 yr time intervals indicated by the jet knot separation; however, without further detailed investigation, these values should not yet be applied to individual sources.

For AGN jets, the timescale for jet fluctuations corresponds to 2 yr (for $r_i = 10 R_s$ and $M_{\text{BH}} = 10^8 M_\odot$). This is remarkably similar to the ejection times observed, e.g., in 3C 120 (León-Tavares et al. 2010) or 3C 390.3 (Arshakian et al. 2010).
The main message is that we obtain long-term fluctuations in jet velocities and mass loss rates, on timescales much larger than the typical kinematic timescale of the jet launching area, i.e., the Keplerian time in the range of about 10 days for protostars or 3 days for AGN jets. These timescales are set by the dynamical evolution of the underlying diffusive disk.

While the jet/counter jet asymmetry dies out in the long run when applying a simple model of a global diffusivity, the local diffusivity model—which is more physical in our opinion—leads to long-term asymmetries in the outflow properties which can be of relevance for seemingly one-sided outflows such as, e.g., HH34 or M87. Clearly, other processes such as foreground extinction, relativistic beaming, or the structure of the ambient medium will definitely play a role for the observational appearance of jets.

9. CONCLUSIONS

We have presented results of MHD simulations investigating the launching of jets and outflows from a magnetically diffuse disk in Keplerian rotation. The time evolution of the accretion disk structure is self-consistently taken into account. The simulations are performed in axisymmetry applying the MHD code PLUTO.

Based on Paper I that studied how magnetic diffusivity and magnetization affect the disk and outflow properties, such as the mass and angular momentum fluxes, jet collimation, or jet radius, the present paper investigates the hemispheric symmetry of outflows. We have set up a numerical scheme that can treat the outflows from an equatorial disk in the bipolar direction. The setup has been carefully checked against numerical artifacts triggering asymmetry in the disk–outflow evolution.

We disturb the disk–outflow bipolar asymmetry by applying several approaches. Bipolar symmetry is reached in the long term by gravitational forces, in particular the component orthogonal to the disk, and (in some cases) a symmetric magnetic diffusivity prescription.

In particular, we have obtained the following results.

1. A test case with a symmetric setup gave a perfectly symmetric bipolar evolution of the disk and outflow for several thousand rotations, clearly approving our numerical setup for this kind of study. The measured mass fluxes compare well with the one-hemispheric simulations in Paper I. For the very long-term evolution, numerical effects probably enhanced by the outer disk boundary condition start to disturb the symmetry of the outermost disk. The results presented in this paper are not influenced by numerically triggered asymmetries.

2. We then applied various options to disturb the outflow bipolar symmetry. First we prescribed an initially asymmetric disk applying a different pressure scale height in both disk hemispheres, $\epsilon_{\text{up}} = 0.1$, $\epsilon_{\text{down}} = 0.15$. In general, the disk structure evolves into a warped configuration, with warp amplitudes of a few initial disk scale heights. Electric currents are driven across the equatorial plane. The mass fluxes of the resulting outflows differ by about 10% over less than 1000 time steps, until a symmetric launching is achieved again. The outflow rates at larger distance from the source differ similarly; however, asymmetry lasts longer, just because of the propagation time of the material launched from the disk.

3. We then started the simulation with symmetric initial setup, but disturbing the symmetric disk structure by a localized (in time and space) injection of gas in overpressure at the time when a symmetric outflow is already well established. As for (2), a asymmetric outflow arises after the injection, however, after a time delay due to the propagation time of the asymmetric injection along the disk toward the main jet launching area. In the long term, a symmetric outflow is reestablished.

4. We then investigated the disk–jet evolution applying a local prescription of magnetic diffusivity $\eta = \eta(r, z, t)$ following the local disk evolution. We find that—as discussed in Paper I—the distribution of magnetic diffusivity affects the disk evolution (and subsequently the outflow evolution) essentially, as it governs the coupling between matter and field, and thus the angular momentum evolution, the disk accretion, the magneto-centrifugal acceleration, and also the mass loading of the outflow. We first applied a simple power-law diffusivity distribution as motivated by earlier papers.

5. For a more sophisticated magnetic diffusivity distribution following the density-weighted local sound speed, and an increasing diffusivity with radius, the accretion evolution is more persistent. We were able to follow the accretion–ejection process over many thousands of dynamical time steps. The most interesting result is that the bipolar asymmetry of jet and counter jet is long lasting. We find that in this case the warped structure of the disk survives many dynamical timescales also in the inner disk. We interpret this as due to the lack of a symmetric diffusivity distribution. There are the same restoring forces of gravity as in the case of a symmetric diffusivity profile. However, the matter is more directly coupled to the distorted magnetic field which in principle opposes the return to symmetry. In the end, we observe a persistent difference in the jet–counter jet mass fluxes—up to 30% and lasting longer than the simulation run time.

6. In order to compare between internal and external effects, we also investigated the launch of a jet outflow from a symmetric disk into an asymmetric ambient gas distribution. As probably expected, the initial outflow is strongly asymmetric with 20% different mass fluxes for jet and counter jet; however, as soon as the outflow has penetrated the asymmetric corona (i.e., when the outflow has left the grid and the initial condition has swept out of the domain), the outflow returns to symmetry (after about 2000 dynamical time steps). In comparison to the simulations discussed, the initial disk structure is symmetric and does not evolve into a warped structure.

Future studies should run for a longer time and on a larger computational domain in order to be more comparable to the observationally resolved features. Although we find it not essential for the present study, it would be helpful to apply an outer disk boundary condition providing a steady (and symmetric) mass inflow into the disk. An essential point is resolution. Our resolution study gave similar results for the cases investigated; however, the stratification of the inner disk is still not very well resolved, in particular the surface layer of the jet launching disk. This holds in particular for the disk surface where the jet launching actually happens and where the disk material is loaded onto the outflow.

Concerning the local prescription of magnetic diffusivity following profile $f_2$, we have observed a seemingly strong
feedback mechanism leading to superefficient accretion with weak outflows, and a strongly decaying disk mass may hint at interesting applications for unsteady jets. These features deserve further attention and will be the focus of a forthcoming paper.

We thank Andrea Mignone and the PLUTO team for the possibility to use their code. We acknowledge a detailed and extensive report by an expert referee, which helped to improve the presentation of the paper. S.S. acknowledges the hospitality by the Max Planck Institute for Astronomy. The simulations were performed on the THEO cluster of the Max Planck Institute for Astronomy. This work was financed by the SFB 881, subproject B4, of the German science foundation DFG, and partly by a scholarship of the Ministry of Science, Research, and Technology of Iran.

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