Zoom-in Simulations of Protoplanetary Disks Starting from GMC Scales

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Abstract

We investigate the formation of protoplanetary disks around nine solar-mass stars formed in the context of a (40 pc)$^3$ Giant Molecular Cloud model, using RAMSES adaptive mesh refinement simulations extending over a scale range of about 4 million, from an outer scale of 40 pc down to cell sizes of 2 au. Our most important result is that the accretion process is heterogeneous in multiple ways: in time, in space, and among protostars of otherwise similar mass. Accretion is heterogeneous in time, in the sense that accretion rates vary during the evolution, with generally decreasing profiles, whose slopes vary over a wide range, and where accretion can increase again if a protostar enters a region with increased density and low speed. Accretion is heterogeneous in space, because of the mass distribution, with mass approaching the accreting star–disk system in filaments and sheets. Finally, accretion is heterogeneous among stars, since the detailed conditions and dynamics in the neighborhood of each star can vary widely. We also investigate the sensitivity of disk formation to physical conditions and test their robustness by varying numerical parameters. We find that disk formation is robust even when choosing the least favorable sink particle parameters, and that turbulence cascading from larger scales is a decisive factor in disk formation. We also investigate the transport of angular momentum, finding that the net inward mechanical transport is compensated for mainly by an outward-directed magnetic transport, with a contribution from gravitational torques usually subordinate to the magnetic transport.

Key words: accretion, accretion disks – magnetohydrodynamics (MHD) – protoplanetary disks – stars: formation – stars: low-mass – stars: protostars

1. Introduction

Since the first observations of exoplanets (Wolszczan & Frail 1992; Mayor & Queloz 1995), interest in modeling planet formation has grown enormously. All planets in the solar system and the majority of detected exoplanets move in nearly coplanar orbits, consistent with planets forming in circumstellar disks, and these disks are therefore commonly referred to as protoplanetary disks.

Observations of disks around young stellar objects (YSOs) strongly indicate that protoplanetary disks are connected to the early formation phase of the host star, where vigorous accretion from a thick envelope is controlling the dynamics of the disk-star system. Hence, protoplanetary disk should ideally be modeled in the context of star formation. Existing models typically start from more or less idealized initial conditions, often in the form of a spherical stellar core with a Bonnor–Ebert-like profile, and follow the collapse during star and protoplanetary disk formation (Machida et al. 2004, 2006, 2007; Tomida et al. 2010, 2013; Li et al. 2011; Machida & Matsumoto 2011; Seifried et al. 2011, 2012; Joos et al. 2012, 2013; Krumholz et al. 2013; Myers et al. 2013; Bate et al. 2014; Padoan et al. 2016). Although these cloud collapse models are useful to study specific cases of star and protoplanetary disk formation, they neglect the influence of the stellar environment. In fact, stars preferentially form as a result of the collapse of pre-stellar cores in filamentary structures of Giant Molecular Clouds (GMCs; Blitz 1993), and as illustrated by the recent simulations of Padoan et al. (2014), the accretion rates are in general heavily influenced by external factors, resulting in formation times that can vary with an order of magnitude or more for similar-mass stars. Therefore, modeling star formation under idealized assumptions, with isolated boundary and initial conditions, is an oversimplification.

In this paper we present a fundamentally different approach to modeling star and protoplanetary disk formation, accounting for the influence of the environment ab initio. In order to avoid the dependency of strongly idealized initial and boundary conditions, we start from GMC scales and follow the process of individual star formation and protoplanetary disk formation down to scales of 2 au, properly anchoring the formation of individual stars in the larger GMC context. We model the GMC in a cubic box of size (40 pc)$^3$, with a total mass of $\approx 10^5 M_\odot$. This is motivated by observations by, for example, Murray (2011), who found GMCs ranging between 5 and 200 pc in size, with masses from $10^3$ to $10^7 M_\odot$. Realistic turbulence in our GMC model is driven by feedback from supernovae, which are evolved self-consistently by allowing massive sink particles to explode and enrich the surrounding medium. The numerical method and the model are discussed in Section 2. In Section 3 we present the stellar accretion and disk formation processes, as well as discuss the importance of our results. Finally, we summarize in Section 4.

2. Methods

In this section we describe the setup of our zoom-in simulations, with focus on the refinement and accretion criteria, since these are the key technical parameters of our model. Further descriptions of the implementations of sink particles can be found in Padoan et al. (2014) and in Kuffmeier et al. (2016). Our approach is feasible owing to extensive application of adaptive mesh refinement (AMR), using a heavily modified version of the RAMSES code (Teyssier 2002; Fromang et al. 2006), which can handle up to 29 levels of refinement relative
to the box size (Nordlund et al. 2014). In Figure 1, we sketch the procedure for our simulations.

2.1. Evolving the Initial Conditions

The starting point of the entire simulation is a GMC model, which is represented in RAMSES by a cubic box of size (40 pc)$^3$, with periodic boundary conditions. The box is filled with a self-gravitating, magnetized gas. The average H$_2$ number density is initially 30 cm$^{-3}$, corresponding to a total mass in the box of approximately $10^5 M_\odot$. The initial magnetic field strength is 3.5 $\mu$G. Initially, the box was evolved with driven turbulence and star cluster particles, including supernovae, using the unigrid MHD STAGGER code (Kritsuk et al. 2011). After an initial period of burn-in, a snapshot was restarted with RAMSES, and refinement was added (see Vasileiadis et al. 2013). Potentially a real GMC fragment might be broken up by supernova explosions or other forms of feedback. In the experiment, this is prevented by the periodic boundary conditions, which can be interpreted as corresponding to embedment in a similar or larger amount of nearby gas. The assumed GMC lifetimes are in agreement with the “star formation in a crossing time” paradigm (Elmegreen 2000; Elmegreen & Shadmehri 2003), with recent numerical modeling (Padoan et al. 2016), and with observational estimates (Murray 2011). To initialize a turbulent state, solenoidal forcing in a shell of wavenumbers in the range of $1/L \leq k/2\pi \leq 2/L$ was initially applied (Padoan & Nordlund 2002), but this forcing was turned off when sufficient forcing from exploding supernovae had developed; no artificial forcing was active during the time interval reported on here. The induced velocity dispersion of the cold and dense medium is consistent with Larson’s velocity law (Larson 1969, 1981),

$$\sigma_u \propto 1.2 L^{0.4},$$

with $\sigma_u$ being the velocity dispersion in km s$^{-1}$ and $L$ the size in parsecs. In reference to the recipe of Franco & Cox (1986) for heating due to UV photons (Osterbrock & Ferland 2006), heating is quenched in dense gas, and by using an optically thin cooling function motivated by Gnedin & Hollon (2012), the heating and cooling processes are otherwise modeled schematically as optically thin.

Based on lifetimes interpolated from Schaller et al. (1992), the evolution of stars with more than 8 $M_\odot$ is followed until they explode as core-collapse supernovae. According to observations (Richardson et al. 2002), supernovae release approximately 1 foe, corresponding to $10^{51}$ erg of thermal energy. The integrated energy yields from the stellar winds are expected to be smaller, and we therefore omit their modeling (modeling winds from especially early-type stars in the current connection would also be extremely costly with current computer codes, since the time step must be kept correspondingly short in the entire model, forcing reduced spatial resolution and/or other compromises in other parts of the model). Initial temperatures of supernovae are of the order of 10$^6$ K, and they expand with initial velocities of the order of 1000–3000 km s$^{-1}$, thus requiring calculations with time steps of only a few weeks for a brief period after each supernova explosion.

2.2. Zoom-in Procedure

During the evolution of the GMC with RAMSES, structures of high density form as a consequence of convergent flows and compression due to shocks in cold dense molecular clouds, which are characterized by supersonic turbulence. If the compression is sufficient, self-gravity overcomes opposing pressure forces, resulting in collapse. When the gas density exceeds a critical value, well above the densities reached by turbulence alone, and conditions fulfill additional criteria such as convergence of flow, local potential minimum, etc. (see, e.g., Padoan et al. 2014, for details), the collapsing mass is transferred to a so-called sink particle. The sink particle represents a star, which only influences the environment through its gravity and by accreting surrounding gas.

The analysis of star formation and protoplanetary disk formation around single stars with RAMSES is performed in two steps, using AMR on top of a uniform grid with $128^3$ cells (corresponding to seven levels). First, star formation is followed during 4 Myr of evolution, with nine levels of refinement. That is, the highest level of refinement is $\ell = 16$, corresponding to a dynamic range of $2^{16}$ relative to the entire box size of (40 pc)$^3$. The minimum cell size is thus $\Delta x_{\min} = 2^{-\ell} \times 40 \text{ pc} \approx 126 \text{ au}$.

In the next step, the environment of a single sink with a final mass of 1–2 $M_\odot$ is followed in more detail, by zooming in on the particular object with an increased maximum refinement level of $\ell = 22$, corresponding to a minimum cell size of $\approx 2$ au, and by introducing an accordingly higher time cadence of output from the simulation (in the following called snapshots). The cadence is chosen to lie in the range $0.1–0.5$ kyr. This second stage is the key aspect of this study, given that it provides information about protostellar accretion, including the subsequent formation of protoplanetary disks. To minimize numerical diffusion, and to simplify subsequent analysis, we perform velocity transformations that keep the chosen sink particle stationary in the model’s frame of reference during the second simulation step. To avoid any tendency of coarse grid imprint on smaller scales that a sudden increase in resolution might have on an already-initiated collapse, we start the zoom-in simulations at a time shortly before the sink particle formed in the parental run. We also allow the AMR to adjust continuously and the AMR levels to appear sequentially according to the developing collapse.

In total, we follow the evolution of the environment of nine sinks with higher resolution. In Table 1, we show the properties of the selected sinks. We refine cells if they fulfill certain criteria, e.g., on mass density and distance from the sink.
To maintain sufficient resolution in the inner part of the system of the order of $10^6$ cells on the highest level of refinement, we slowly change the density refinement ladder during the run. We illustrate the resulting distribution of cells refined to the highest level for the different sinks in Figure 2. For a more detailed description of the refinement strategy, please see Appendix A.

### 3. Results

In this section, we present and discuss the influence of the large-scale environment on the evolution of solar-mass sinks.

#### 3.1. Structure of Stellar Environments

To illustrate the structure around the different sinks in detail, we plot the density distribution in the three different planes of the coordinate system, together with the velocity field, in rectangular slices of size $l = 50$ kau in Figure 3. The slices illustrate the shape of the pre-stellar cores and their environments at the time of sink creation. At large distances from the sink (beyond $\approx 10$ kau), the slices show very low densities—less than $10^{-21}$ g cm$^{-3}$ for sinks 1 and 2. Especially compared to sinks 4–9, the density is distributed more homogeneously around sinks 1 and 2 than for the other sinks, where the pre-stellar cores are more perturbed. In contrast to what is assumed in many local molecular cloud collapse models, none of the pre-stellar cores appear as even roughly spherically symmetric.

The core around sink 1 is closest to resembling a classical pre-stellar core, flattened owing to the underlying magnetic field (Allen et al. 2003). However, even in the least disturbed cases the cores show strong velocity differences with respect to their low-density environment. Moreover, some of the sinks appear to be at the endpoints of filamentary arms about $\approx 100$ to $\approx 1000$ au in width, which reach out to distances beyond several tens of kau (see sinks 2, 3, 4, 5, 6, 7, and 8 in the $xy$-plane; 4, 5, 6, and 8 in the $xz$-plane; and 4, 5, and 9 in the $yz$-plane). This may be related to the gravitational collapse, which tends to concentrate mass at the end of filaments (Seifried & Walch 2015).

In Figure 4 we show the density distribution (average densities in spherical shells) in the radial direction for the snapshots closest to (generally within a few hundred years of) sink creation. We can see that depending on the sink, the densities can vary by several orders of magnitudes at distances beyond $10^3$–$10^4$ au. For comparison, we also plot the density profile of a singular isothermal sphere (SIS)

$$\rho_{\text{SIS}} = \frac{\sigma_v^2}{2\pi G r^2}.$$  

(2)

The averaged density profiles do not differ dramatically from a spherical isothermal Bonnor–Ebert density profile with $\rho \propto r^{-2}$ and a central flat core. Although the azimuthally averaged density profiles around different protostars may be similar, taking into account the perturbations of and around the pre-stellar cores (Figure 3), the actual stellar environment can nonetheless be significantly different owing to the filamentary structure. The densities at larger distances from sink 1 and 2 are lower than $10^{-21}$ g cm$^{-3}$, whereas the other sinks show higher densities beyond distances of about $10^3$ au. At radii smaller than $\approx 10^3$ au the sinks have a similar generic profile, with a flat density distribution close to the sink that decreases beyond a certain radius. We require a threshold value for the density to trigger sink creation, which corresponds to having the Jeans length resolved by a number of cells at the highest level of
refinement. Therefore, the inner flat core in the profile is approximately at the same level of several times $10^{-15}$ g cm$^{-3}$ just after sink formation for all sinks, except for sink 4, where we required slightly different conditions for sink creation, namely, a higher threshold density, corresponding to 25 instead of 50 cells per Jeans length. By using our method of sink creation, the lower the maximum level of refinement, the earlier the sink forms. Therefore, the time of sink creation forestalls the actual birth of the protostar (with an amount of time similar to the freefall time at the central densities, which is $<1000$ yr). Such a time delay is short compared to the evolution around young protostars for time intervals of the order of 100 kyr and therefore negligible.

### 3.2. Accretion Profiles

In this section, we investigate the evolution of the sinks for the first 1 Myr, based on the parental run, before we focus on the zoom-ins for a more detailed study of their evolution during the first 100 kyr.
around the sinks nor artificially remove a certain fraction of the accreting mass. Taking into account that protostellar mass loss is expected to be about 50% of the accreted mass (Zanni et al. 2007), the mass range is appropriate for our purposes. Given that we use an output cadence of 50 kyr in the low-resolution parental run, we can only roughly estimate the accretion onto the sink, but it is striking that the sinks accrete their masses on very different timescales. While sinks 1 and 2 accrete most of their mass within only about 100 kyr, some other sinks accrete only a fraction of the mass that they have after 1 Myr in that time, and in some cases undergo long periods with accretion rates exceeding $10^{-6} \ M_\odot \ yr^{-1}$ at times later than 800 kyr after sink creation. Considering the different accretion profiles, together with the fact that some of the selected sinks have only evolved for less than 1 Myr in the parental run, we are aware that some of the sinks have not yet finished their accretion process or may be fed at later times after quiet periods with infalling mass from large scales.

Comparing the different accretion profiles in the context of the stellar environment, we conclude that more isolated, less distorted pre-stellar cores (stars 1 and 2) correlate with a shorter accretion time of the protostar. Essentially, they consume the gas reservoir of their pre-stellar core, and the accretion process stops. This is in accordance with a classical picture of star formation from spherically symmetric isolated pre-stellar cores. The differences in duration of the accretion process are a consequence of the spatial extent of the pre-stellar core or its possible feeding filament(s). Taking into account that the freefall time scales as $t_{ff} \propto R^{3/2}$, infalling mass located at larger distances from the sink accretes later onto the sink. As seen in Figure 3, several pre-stellar cores are strongly perturbed, and many sinks lie along filaments of gas that extend to distances beyond $10^4$ au from the sink. These filaments can feed the sink during its accretion phase with material initially located far away from the sink, thus prolonging the accretion timescale compared to the classical scenario of a collapsing Bonnor–Ebert sphere of about $10^4$ au in size—even though the collapsing cores may in fact initially be smaller. As pointed out by Padoan et al. (2014), mass clumps at distances beyond $10^4$ au may eventually fall onto the system and thus cause accretion bursts, which potentially explain the luminosity problem observed for YSOs (Kenyon et al. 1990). We investigate the effect of the accretion profiles on the stellar evolution, including a detailed comparison with observations in a follow-up paper.

### 3.2.2. Accretion during the First 100 kyr of Evolution

Figure 6 shows the evolution of the accretion rates during the first 100 kyr based on the zoom-in simulations. We start by presenting the similarities for the different sinks before we focus on the differences in their accretion histories. All accretion profiles quickly peak to values in the range of $3 \times 10^{-5} \ M_\odot \ yr^{-1}$ for sink 4 and $6 \times 10^{-5} \ M_\odot \ yr^{-1}$ for sink 3, before the accretion rates decrease during the subsequent evolution. The quick rise within at most a few kiloyears seen for the sinks indicates the accuracy of the selected time of sink creation as discussed above. The accretion rate peaks with a small offset after $t = 0$, due to the flat density profiles around the different sinks at the time of their creation (Figure 4).

After the initial peak, the accretion rates of the sinks decrease in ways that differ between individual sinks. Sink 1 (blue dots) shows the most consistent continuously decreasing accretion rate, while other sinks settle to a nearly stable average accretion rate or have only slowly decaying profiles, after an initial drop in the rate. After 100 kyr, sink 1 has an accretion rate of only $\sim 10^{-7} \ M_\odot \ yr^{-1}$ and is therefore transitioning to a state that observationally would be classified as either class I or early class II. Sink 3 and sink 7 have very similar profiles, and, as for sink 1, they also show a continuously decreasing accretion rate. However, the change is more modest, starting at peak values of about a few times $10^{-5} \ M_\odot \ yr^{-1}$. They still show accretion rates of about $10^{-6} \ M_\odot \ yr^{-1}$ after $t = 100$ kyr. Sink 7 also shows a drop in the accretion, from a peak value of $4 \times 10^{-5}$ to about $6 \times 10^{-6} \ M_\odot \ yr^{-1}$ at approximately 20 kyr after sink creation. However, the accretion rate then only decreases slightly, to a minimum of $3 \times 10^{-6} \ M_\odot \ yr^{-1}$ after $t = 75$ kyr, after which the accretion profile starts to rise again, to about $8 \times 10^{-6} \ M_\odot \ yr^{-1}$ at the end of the simulation at $t = 100$ kyr. Although we do not account for the later accretion profile of sink 5 owing to its shorter simulated evolution, we can see that the profile is almost identical to the one for star 8 and also similar to the evolution of sinks 7 and 9. Such a period of re-increasing accretion rate is also seen for sink 6 in the time interval between $t = 27$ and 62 kyr. The accretion profiles for this sink, as well as for sink 9 and (especially) sink 4, are more intermittent and episodic.

We caution the reader that we are generally underestimating the amplitudes and extents of the bursts for two reasons. First, even our highest resolution of 2 au is coarse compared to the actual size of a protostar, and second, the snapshot cadences we used were at least 200 yr, with the effect that shorter bursts are
often missed in this plot. Therefore, we postpone a detailed analysis and comparison of the accretion histories with observations to future work. Nevertheless, the fact that some profiles show intermittent profiles may help to understand observations such as the ones by Safron et al. (2015) of a class 0 YSO (HOPS 383) with fluctuations of more than 30 in the mid-infrared. We stress that these early accretion bursts that occur within a few tens of thousands of years after stellar birth are different from the large-scale infalls, previously suggested by Padoan et al. (2014), explaining the luminosity variations of class I YSOs.

Compared to the mass evolution of the sinks in the parental run, the mass evolutions of the sinks in the zoom-in runs show significant differences. First of all, the sinks accrete much less mass than in the parental run, which is mainly a consequence of the accretion efficiency parameter acc_eff being 0.5 instead of 1 as in the parental run. In retrospect, a lower acc_eff could have been used also in the parental run, to compensate for the missing outflow due to the unresolved winds and jets. Simply multiplying the sink mass in the zoom-in runs according to our choice of the accretion efficiency by a factor of 2 is not sufficient, because a higher acc_eff would cause a larger sink mass and thus a deeper gravitational potential. This would cause additional accretion, as the sink would be able to gravitationally attract mass from larger distances that is gravitationally unbound for a lower sink mass. From smaller test runs evolved with both recipes we have found that acc_eff = 0.5 corresponds to reducing the sink mass by one-third in an identical run with acc_eff = 1. Sink 1 is the most massive of the plotted sinks after 100 kyr in the parental run and is almost three times as massive as sink 3 and about twice as massive as sink 7 and 9, although it has accreted less mass than these sinks in the zoom-in runs.

By tracing the history of the gas with tracer particles that are passively advected with the gas motion, we showed in Kuffmeier et al. (2016) that a small fraction of the gas located within 100 au from the protostar at time \( t \) after sink creation was located beyond \( 10^4 \) au at the time of sink creation. Given that local core-collapse models usually use a homogeneous mass distribution similar to the theoretical profile of a Bonnor–Ebert sphere, we expect that the majority of low-mass sinks accrete on timescales that are longer than predicted by these models. Furthermore, the diversity shows the limitations of an isolated core-collapse model as an initial condition, when aiming to comprehend the entire process involved in sink creation, especially the occurrence of late mass infall through filaments.

### 3.2.3. Diversity in the Angular Distribution of Accretion

In contrast to the case of a nonmagnetized, rotationless collapsing sphere, we find that the gas does not accrete uniformly. Instead, gas accretes along accretion channels and accretion sheets, and some parts of the gas even flow in the outward direction.

In Figure 7, we illustrate this spatial heterogeneity of the accretion process by showing the angular distribution of the accretion rates at a radius of 50 au, 10 kyr after sink creation, for sinks 1, 2, 4, 5, 6, 7, 8, and 9. In order to highlight the infall and outward motion of material, as well as the differences in magnitude, we use a symmetric log color scale. That means that values between \( \pm 10^{-11} M_\odot \text{ au}^{-2} \text{ yr}^{-1} \) are plotted according to a linear color scale, while we use a logarithmic scale for values below or above the threshold values. Furthermore, we choose cutoff values of \( \pm 10^{-10} M_\odot \text{ au}^{-2} \text{ yr}^{-1} \) to emphasize that red corresponds to outward motion of gas and blue to infall. When comparing the projections of the sinks with each other, we find differences in the extent of the accretion channels and sheets, as well as the amount of gas moving away from the sink. Sinks 5, 7, and 8 only show weak outward motions of gas in contrast to the other four sinks, where a significant surface area is dominated by the outward motion of gas. In contrast to what is found in models of collapsing envelopes with uniform density, we do not see a pure bipolar outflow at 50 au at this stage, and distribution of outflowing and infalling gas does not show a clear morphology. Nevertheless, we occasionally see outward motions approximately perpendicular to the disk plane with speeds similar to the Kepler speed at distances of a few au. These speeds are consistent with magnetocentrifugally driven winds launched on scales of our highest resolution of 2 au.

The strength of the magnetic fields around the protostar is probably overestimated in our models, due to the absence of nonideal MHD effects (Krasnopolsky et al. 2010, 2011; Li et al. 2011; Tsukamoto et al. 2015; Hennebelle et al. 2016; Wurster et al. 2016), and a significant part of the gas initially moving outward may be caused by magnetic interchange instabilities that occur because of high magnetic pressure close to the protostar, resulting in the emergence of strong and expanding magnetic loops (Li et al. 2014). Considering the high densities at the center, we expect ohmic dissipation to have the strongest effect in reducing the field strength. Nevertheless, we speculate that a buildup of high magnetic pressure, such as seen in our simulations, might occur also to some extent during the real formation phase of protostars. Other groups carrying out ideal MHD simulations of a local cloud collapse see this effect as well for their turbulence-free simulations (Seifried et al. 2011; Joos et al. 2012). Seifried et al. (2013) do not detect the occurrence of the magnetic interchange instability, when including turbulence in their simulations. We think that the instability does not occur in their case because the resolution is lower than in our study (8 au compared to 2 au in our simulations). The lower resolution has the effect that the magnetic pressure does not pile up as much as seen in our case. Recent simulations by Masson et al. (2016), accounting for dissipation effects from ambipolar diffusion and ohmic resistivity, show that nonideal effects can circumvent the occurrence of the magnetic interchange instability. However, the strengths of the nonideal MHD coefficients depend on the degree of ionization of the gas, which may be highly variable between protostars for several reasons. First, the efficiency of shielding by dust depends crucially on how the dust is distributed in space. In cases where the dust distribution is very filamentary, with the region in between with much lower dust densities, the effective absorption will be much less than if the distribution were uniform. Second, we point out that the efficiency of shielding even for a uniform dust-to-gas ratio strongly depends on the grain size distribution. Moreover, the cosmic-ray intensity may depend on the location (Padovani et al. 2009, 2013, 2014), e.g., because of “magnetic bottle” effects, and because of the distance in spacetime from acceleration regions such as SN shock fronts. Finally, if the cosmic-ray intensity is efficiently reduced in dense regions, the main ionization source becomes short-lived radionuclide \( ^{26}\text{Al} \) (Umeyabashi & Nakano 2009; Cleeves et al. 2014; Tomida et al. 2015), and given that the abundance of \( ^{26}\text{Al} \) can differ
among stars by orders of magnitude (Vasileiadis et al. 2013), this can also contribute significantly to variations of the ionization degree in space and time.

In the case of sink 7 we find that, even at \( t = 75 \) kyr, the gas falls toward the sink through sheets from all spatial directions. However, the infall rates relax to more modest values and accretion appears to happen mainly through two sheets with a small distance from each other at \( t = 100 \) kyr. In contrast, the Hammer projections for sink 6 (not shown) show slightly broader sheets, with higher accretion rates, as well as a significant contribution of gas moving in the outward direction. We point out that the enhanced infall rates for sink 6 are not surprising, taking into account that the density over several hundred au in the \( xy \)-plane from the sink is enhanced by up to a factor of 10 (Figure 4) at the time of stellar birth, while the density is about the same or only slightly higher in the other coordinate-axis planes. Consequently, there is more mass available that can accrete onto the star–disk system. Nevertheless, the pure difference in the absolute amount of mass cannot account for the significant outward motions observed for sink 4 compared to the nearly pure infall morphology around sink 7.

3.3. Disk Formation and Evolution

An increasing number of observations of class 0 objects reveal the existence of circumstellar disks in the midst of their formation process. They are not perfectly symmetric, and their profiles differ from static thin standard accretion disks (van der Marel et al. 2013). Due to the violent processes during protostar formation, the identification of a disk depends on the definition of the term “disk.” The fundamental property that causes disk formation is the conservation of angular momentum during the infall, which determines the size of the disk. The left panel of Figure 8 shows the average specific angular momenta in the stellar surroundings of the nine sinks close to their individual \( t = 0 \). For a \( 1 M_\odot \) star, the specific angular momentum at 1 au is about \( 4.5 \times 10^{19} \) cm\(^2\)s\(^{-1}\), and the plot illustrates that at this time the average specific angular momentum of essentially all mass inside about 1000 au has low enough angular momentum to reach orbits inside 1 au. Remarkably, the distribution over radius is similar for all sinks, scaling roughly linearly with radius. Since we are resolving each level of the refinement ladder with a large number of cells and we have of the order of 50 cells per Jeans length, it is unlikely that this scaling is significantly influenced by numerical resolution effects.

To compare the evolution of the level of rotation around the different sinks in our simulation, we plot the sum of the specific angular momentum of the gas within a distance of 100, 1000, and 10,000 au from the eight sinks of the advanced zoom-ins. The symbols belong to the same sinks as in Figure 5.

Figure 7. Hammer projection of the accretion at a radius of 50 au from the sink at time \( t = 10 \) kyr for sinks 1, 2, 4, 5, 6, 7, 8, and 9 that were evolved with a minimum cell size of 2 au. The colors represent inflow and outward motion according to the color bar. In order to illustrate both positive and negative values, we decided to use a linear range in between \( \pm 10^{-13} M_\odot \) au\(^{-2}\) yr\(^{-1}\) and a logarithmic scale beyond.

Figure 8. From left to right: specific angular momentum as a function of radius from the nine host sinks at \( t = 0 \), and evolution of total specific angular momentum of the gas located within a distance of 100, 1000, and 10,000 au from the eight sinks of the advanced zoom-ins. The symbols belong to the same sinks as in Figure 5.
around sink 1 peaks at a value of about $5 \times 10^{18} \text{ cm}^2 \text{s}^{-1}$ at $t \approx 10 \text{ kyr}$ and then drops below $10^{17} \text{ cm}^2 \text{s}^{-1}$ after $t \approx 100 \text{ kyr}$. In contrast, sink 8 has a much flatter accretion profile, and the specific angular momentum within 100 au is roughly constant of the order of $10^{19} \text{ cm}^2 \text{s}^{-1}$ since $t \approx 40 \text{ kyr}$. Also, the initial strength of specific angular momentum at distances beyond 1000 au can differ by more than an order of magnitude.

The differences in extent of protoplanetary disks and the initial variation in the rotational velocity profiles around the sinks at larger radii raise the question to what extent these two properties are connected. Considering that gas falls toward the sink, potentially from larger distances, the specific angular momentum at larger distances determines the rotational structure of the accreting gas. The specific angular momentum within a spherical shell of 100 au around the sinks quickly increases during the first $\sim 5 \text{ kyr}$ for sinks 4, 6, and 9, as may be seen in the initial rise at the very left in the second panel of Figure 8. Below, we describe in more detail that circumstellar disks quickly form around these sinks.

Similar to Frimann et al. (2016), we quantify the rotational extent of the gas motion by investigating the angle $\alpha$ (Brinch et al. 2007, 2008),

$$\alpha = \arctan \left( \frac{\langle v_r \rangle}{\langle v_\phi \rangle} \right).$$

Theoretically, the angle can vary between $\pi/2$ for pure infall and $-\pi/2$ for pure outward motion, with a value close to 0 representing pure rotation. As shown by Frimann et al. (2016), it serves as an adequate measure of the rotational support of the gas around the sink, and thus as a valid first proxy of a possible disk. In Figure 9, we show the evolution of $\alpha$ within spheres of radii of 50, 100, 300, and 1000 au around the sinks. The decrease of $\alpha$ with time for all of the sinks shows that the relative amount of rotation increases for all sinks during their evolution (although sink 7 has not evolved long enough to show a clear decrease). However, just as with the accretion profiles of the different sinks (see Figure 6), the curves evolve remarkably differently.

We interpret low $\alpha$ values as an indicator for the early formation of rotationally supported disks around these sinks. The gas around the remaining sinks does not show such low values. In particular, the velocity profile around sink 7 is significantly dominated by the infall of material, even after 100 kyr, with an $\alpha$ of more than 0.5. This indicates that no rotationally supported disks have formed around them.

To illustrate the structure around the sinks in our study, we show slices in the plane perpendicular to the mean angular momentum vector at $t = 50 \text{ kyr}$ for sinks 1, 4, 5, 6, 7, and 9 in Figure 10. The images reveal the variety in disk formation for the different stellar environments, as already seen in the evolution of $\alpha$. Similarly to the images of the pre-stellar cores (Figure 3), the images show filamentary arms feeding the forming protoplanetary disk. These filaments are not necessarily aligned with the disk plane, but more often approach the disk from various angles. Moreover, in agreement with recent observations from ALMA and the Subaru Next Generation Adaptive Optics (HiCIAO), the disks show evidence of large-scale features, such as spiral arms or inflowing gas streams (Yen et al. 2014; Liu et al. 2016).

### 3.4. Angular Momentum Transport

An important aspect in protoplanetary disk studies is how the angular momentum is transported from the environment to the star–disk system. To study the flow of angular momentum, we consider a cylindrical test volume with the height equal to the diameter ($h = 2R$) and calculate the angular momentum flux through the cylinder wall (the “radial direction”) and through the top and bottom of the cylinder (the “vertical
direction”). We refer the reader to Appendix C for the detailed calculation. Three terms contribute to the transport of angular momentum in the vertical and radial directions. The magnetically induced transport is associated with the Maxwell stress $-\mathbf{B} \otimes \mathbf{B}$. In the vertical direction it is

$$F^B_v(R) = \pm \int_0^R \int_0^{2\pi} r \, d\phi \, \frac{B_\phi(r, \phi, \pm h/2) B_\phi(r, \phi, \pm h/2)}{4\pi},$$

(4)

while in the radial direction it is

$$F^B_r(R) = -\int_{-h/2}^{h/2} R \, d\phi \, R \, \frac{B_\phi(R, \phi, \pm h/2) B_\phi(R, \phi, \pm h/2)}{4\pi}.$$  

(5)

The mechanical flux of angular momentum is associated with the Reynolds stress $\rho \mathbf{v} \otimes \mathbf{v}$. In the vertical direction it is

$$F^v_v(R) = \pm \int_0^R \int_0^{2\pi} r \, d\phi \, \rho \, \mathbf{v}_\phi \times (\mathbf{r}, \phi, \pm h/2) \mathbf{v}_\phi(r, \phi, \pm h/2) \mathbf{v}_\phi,$$

(6)

(while $\pm$ indicates the signs at top and bottom, respectively), while in the radial direction it is

$$F^v_r(R) = \int_{-h/2}^{h/2} R \, d\phi \, R \, \rho \, \mathbf{v}_\phi \times (\mathbf{r}, \phi, \pm h/2) \mathbf{v}_\phi(r, \phi, \pm h/2).$$

(7)

Finally, we have the contribution associated with the gravitational potential $\nabla \Phi \nabla \Phi$, accounting for angular momentum transport through spiral arms and similar nonaxisymmetric structures. In the vertical direction it is

$$F^g_v(R) = \pm \int_0^R \int_0^{2\pi} r \, d\phi \, \left(\frac{\nabla \Phi \Phi}{4\pi G}\right)_{(\phi, \pm h/2)}(\nabla \Phi \Phi)_{(\phi, \pm h/2)},$$

(8)

while in the radial direction it is

$$F^g_r(R) = \int_{-h/2}^{h/2} R \, d\phi \, R \, \left(\frac{\nabla \Phi \Phi}{4\pi G}\right)_{(\phi, \pm h/2)}(\nabla \Phi \Phi)_{(\phi, \pm h/2)}.$$ 

(9)

In contrast to Joos et al. (2012), we consider the total value of the signed fluxes, rather than splitting each contribution into negative and positive parts. In this manner, we average out the natural spatial fluctuations induced by turbulence and obtain net fluxes (by definition positive in the outward direction).

We compute the fluxes of angular momentum within cylinders of different radii ($r = 28\, \text{au}$, $r = 80\, \text{au}$, $r = 128\, \text{au}$, and $r = 180\, \text{au}$), with tops and bottoms at $+h$ and $-h$, with in one case $h = r$, and in another case $h = 2r$. Using these cylindrical control volumes, we compared the different contributions relative to one another and also compared their radial and vertical components. In Figure 11, we illustrate the time evolution of the different components around sink 4.

As may be seen from Figures 11 and 12, angular momentum is—consistent with expectations—predominantly transported inward by the mechanical Reynolds flux, while the gravitational acceleration and the magnetic Maxwell stress generally account for transport in the outward direction. Transport induced by the magnetic stress is typically stronger than the transport caused by gravity by a factor of a few to several, though strong spiral arms can occasionally cause enhancements of the gravitational component. Comparing the contributions in the radial and vertical directions in Figure 12, we find that the strongest transport of angular momentum transport occurs in the vicinity of the disk midplane in the radial direction, consistent with the fact that the fluctuating radial component of the magnetic field, which tends to outline trailing spirals, is strongest near the disk midplane. The vertical Maxwell component is strongest at larger radii. This indicates that angular momentum is not transported in a narrow jet. However, this is a consequence of the resolution given that the minimum cell size of 2 au is too coarse to resolve the launching of narrow jets. The flux in the radial direction near the top and bottom is similar to the vertical flux through the top and bottom, respectively. This indicates that the angular momentum flux is oriented roughly at 45$^\circ$ there. Another interesting aspect concerns the asymmetry of the fluxes in the vertical direction. In contrast to idealized core-collapse setups without turbulence, but in line with our results presented above, the angular momentum is transported heterogeneously. Even when averaging over periods of a few times 10$^4$ yr—corresponding to about 100 orbital times at 50 au—we clearly see that more angular momentum is transported on one side than the other. While the resolution is insufficient to study outflows via jets and winds quantitatively, these asymmetries are nevertheless consistent with recent observations of asymmetric outflows (e.g., IRAS 03292+3039; Tobin et al. 2015). We interpret them as a natural consequence of the underlying filamentary and chaotic nature of the large-scale accretion flows.

### 3.5. Disk Size Evolution

An important quantity in understanding the disk evolution around the different protostars is the time evolution of the disk sizes. $\alpha$ is a good indicator of the relation of rotational to radial velocity, but early disks might already form, although the gas has a strong radial velocity contribution. We thus estimate the disk size in the following way. We first calculate the total angular momentum vector $L_{100}$ inside a sphere of 100 au around the protostar. We then calculate the azimuthal velocity
for all cells that are located inside a cylinder of 1000 au in radius, ±8 au in height, and with the radial direction being perpendicular to $L_{100}$. Afterward, we estimate an average azimuthal velocity $v_{\phi}$ for all cells that are located between radii $r$ and $r + dr$. Altogether we consider 100 radial bins with $dr$ increasing exponentially with increasing radius. Finally, we determine the disk size as the radius where $v_{\phi}/v_K$ (with $v_K$ being the Kepler speed) drops below a threshold value $a$, though we do not take into account velocities inside a distance of seven cells (14 au) from the sink in order to avoid potentially low rotational velocities that are induced by the sink parameters. Theoretically, a thin rotating gaseous disk with the radial structure described by power laws has azimuthal velocities

$$v_{\phi} = v_K \left(1 - \mathcal{O}\left(\frac{h}{r}\right)^2\right)^{1/2}$$

(Armitage 2007). The small deviation of the rotational speed of the gas from the Kepler speed is induced by the decrease of gas pressure with radius in the disk. For thin disks, we thus expect $v_{\phi}$ to be nearly equal to $v_K$. However, considering that early disks can be rather thick, and taking into account the violent accretion process from the outside, we relax the lower velocity limit somewhat and choose a threshold value of $a = 0.8$. Based on this method, we plot in the top panel of Figure 13 the evolution of the disk sizes $r_{\text{disk}}$ during protostellar evolution for the different protostars. In general, we can see that the disk sizes increase during the evolution. However, similar to the time evolution of $\alpha$, we see strong differences among the different protostars. While some disks extend to several hundred au, others only extend to at most a few tens of au. This is in agreement with differing disk sizes depending on the initial density profile in local core-collapse models (Machida et al. 2014). We find a significant time variation in the disk size during the evolution, indicating that disk formation is an intermittent process, such as expected when accounting for turbulence. Part of this intermittency is most probably the result of some of the disks being marginally Toomre unstable, as discussed below, but this is of course ultimately related to the time variation and strength of mass infall from the environment onto the disk. We expect this intermittency to decline with time, when the infall rate decreases and the replenishment time increases in the disks.

Figure 12. Time-averaged fluxes in the vertical (left and right panels) and radial (middle panel) directions during the period from 15 to 45 kyr for sink 4. The left panel shows the vertical flux as a function of radius at a constant height of 180 au, the middle panel shows the vertical flux as a function of height at a constant radius of 180 au, and the right panel shows the vertical flux as a function of radius at a constant height of $-180$ au. Blue solid corresponds to the magnetic component, green dashed to the mechanical component, cyan dot-dashed to the gravitational component, and the black dotted line represents the total flux.

Figure 13. Disk size (top panel), disk mass (middle panel), and disk-to-stellar mass ratio (bottom panel) as a function of time around the different protostars after sink creation. The symbols belong to the same sinks as in Figure 5.
The differences are also reflected in the diversity of disk masses for the different disks during their evolution. In order to compare disk masses, we add up the mass of the cells that are located within \( r_{\text{disk}} \), within a maximum \( \frac{h}{r} \) ratio of 0.2, and within a maximum vertical distance from the midplane of \( h = 8 \text{ au} \) (middle panel in Figure 13). We find that the disk masses are highest (\( \sim 1\% M_\odot \)) around the sinks where \( \alpha \) approaches 0 most quickly and remains close to \( \alpha = 0 \) (sinks 4 and 6). Moreover, we compare the disk masses with the mass of the host sink in the bottom panel of Figure 13. We notice two trends. First, early disks tend to have large disk-to-stellar mass ratios of up to \( \sim 1 \), because of the short duration of mass accretion by that time. Second, after about 50 kyr the ratios vary between 0.1\% and 10\% for the stable disks in our simulation. Such differences of three orders of magnitude in disk-to-star ratios are remarkable and stress the influence of the stellar environments on star–disk properties—even more so when accounting for the cases where disks do not even form.

In order to constrain the surface density profiles around the sinks, we compute the mass located in cylindrical shells of size \( \Delta r = 10 \text{ au} \) and \( h = 500 \text{ au} \) in height, with the vertical \( z \)-axis being aligned with the total angular momentum vector in a sphere of \( r = 100 \text{ au} \). Afterward, we divide the masses by the corresponding disk surfaces to estimate the column densities. When looking at the surface density profiles around the sinks after 50 kyr (Figure 14), we find that the profiles of the most evident disks (sinks 4, 6, and 9) show a dependence in between \( \Sigma \propto r^{-1} \) and \( \Sigma \propto r^{-1.5} \), with a profile closer to \( \Sigma \propto r^{-1} \) in the inner part. The overall surface densities of the two more massive disks (sinks 4 and 6) are larger by about a factor of two compared to the minimum-mass solar nebula (MMSN; Weidenschilling 1977; Hayashi 1981). However, in the case of the weaker disk around sink 9, the surface density profile is 10–20 times lower than the MMSN. According to our rotational velocity criterion, sink 8 also hosts a disk-like structure. Its surface density profile is much flatter and more perturbed than for the other sinks. However, we caution that the run around this sink was carried out with lower resolution—a minimum cell size of 8 au instead of 2 au. Also, even though the surface density profile for the higher resolved disks is rather smooth, we stress that abundant fluctuations (see Figure 10) are by construction averaged out in a radial profile.

Considering that the surface densities for two of our disks are above MMSN values, we investigate whether the disks are gravitationally unstable. To constrain that, we estimate the azimuthally averaged Toomre parameter (Toomre 1964) based on cylindrical shells of height 100 au,

\[
Q = \frac{c_s \Omega}{\pi G \Sigma},
\]

where \( c_s \) is the sound speed, \( \Omega \) is the orbital frequency, and \( G \) is the gravitational constant. For reasons of simplicity, we assume perfect rotation inside the shell such that

\[
\Omega = \sqrt{\frac{GM}{r^3}},
\]

where \( M \) is the mass of the sink and \( r \) is the radius of the shell.

We illustrate Toomre’s \( Q \) inside the disk around sink 6, which is one of the two massive disks, during its evolution in a contour plot (left panel of Figure 15). We find that \( Q \) drops below 1 at early times, but we stress that at such early times no disk is present, and hence the Toomre criterion should not be applied owing to its dependence on the Keplerian orbital angular velocity \( \Omega \). At later times, \( Q \) is larger than 1, but occasionally very close to 1. If we applied more refinement and were able to resolve the disk inside 2 au, we might expect \( Q \) to be even lower at some locations in the disk. However, considering only the pure hydrodynamical \( Q \) parameter is an oversimplification. Instead, we also have to take into account potential magnetic support of the disk. In Figure 15 is also shown the magnetic Toomre parameter (Kim & Ostriker 2001)

\[
Q_{\text{mag}} = \frac{\sqrt{(c_s^2 + v_A^2)} \Omega}{\pi G \Sigma},
\]

where \( v_A = B / \sqrt{4\pi \rho} \) is the Alfvén velocity. As for the sound speed, we take the mass-weighted average value inside the column for \( v_A \). The toroidal magnetic fields inside the disk contribute to the pressure support of the disk, partly stabilizing it against gravitational collapse. As seen in the right panel of Figure 15, compared to \( Q \), \( Q_{\text{mag}} \) is generally slightly enhanced in the disk, but still only marginally larger than 1 at some locations at times when a disk is present.

Moreover, these values are spatially averaged, and there may be denser clumps within the shell that lead to \( Q \) values lower than 1, and thus to gravitational collapse inside the disk. In Figure 16 we show the distribution of both \( Q \) parameters inside the disk at \( t = 50 \text{ kyr} \) around sink 6. Although we find spiral structures with low \( Q \) values (left panel in Figure 16), we clearly see the additional support provided by the magnetic field (right panel in Figure 16) preventing collapse at least at this point in time.

Nevertheless, the results show that \( Q_{\text{mag}} \) can well be in the marginal range. Taking into account nonideal MHD effects inside the disk may further reduce the magnetic field strength and thus the magnetic disk support. Several works investigated the effects of nonideal MHD in idealized symmetrical disks in detail (e.g., Lesur et al. 2014; Bai 2015; Gressel et al. 2015). The importance of nonideal effects remains to be investigated for our disks. For computational reasons, we leave both tasks, the study of disks with higher resolution and the investigation of nonideal MHD effects, to future work, but already at this point in time.

![Figure 14](image_url) Column density around the different sinks at \( t = 50 \text{ kyr} \). The lines belong to the same sinks as in Figure 4. The black solid line illustrates the MMSN relation of \( \Sigma = 1700 \frac{M_\odot}{\text{au}} \cdot (r/\text{au})^{-3/2} \), and the black dotted line shows a surface density profile of \( \Sigma = 1700 \frac{M_\odot}{\text{au}} \cdot (r/\text{au})^{-1} \).
point we can conclude that some stars may host disks that become massive enough to be gravitational unstable, in particular in the early embedded phase of the evolution.

### 3.6. Mass-to-flux Ratio

An often-considered quantity in analytical and numerical core-collapse models is the ratio between mass and magnetic flux threading a sphere of uniform density, in short, the mass-to-flux ratio (e.g., Allen et al. 2003; Hennebelle & Fromang 2008). It is commonly given as

\[
\mu = \frac{M_{\text{core}}}{\Phi_{\text{core}}} = \frac{M_{\text{core}}}{\int_A B \, dA} / \left( \frac{0.13}{\sqrt{G}} \right),
\]

with \( M_{\text{core}} \) being the mass enclosed in the core, \( \Phi_{\text{core}} \) the magnetic flux through the sphere with cross section \( A \), \( \left( \frac{M}{\Phi} \right)_{\text{crit}} \) the critical mass-to-flux ratio (Mouschovias & Spitzer 1976), and \( G \) the gravitational constant. However, as shown above, stars do not form from pre-stellar cores of uniform density. Instead, the cores are significantly distorted, with feeding filaments such as seen in Figure 3. Therefore, the assumption of an isolated spherical core as the initial condition for modeling star formation is an oversimplification. Accounting for these asymmetries also shows the difficulty of the concept of an initial mass-to-flux ratio. For cores that are highly distorted and continuously fed by accreting filaments the mass budget is not fixed, and hence a changing mass-to-flux ratio does not necessarily imply diffusion of mass relative to the magnetic field.

Nevertheless, we estimate a proxy of the mass-to-flux ratio to compare the relative magnetization around the different sinks at the time of their creation and to investigate how they might influence the formation of disks and their properties (see Figure 17). We compute the mass-to-flux ratio according to Equation (14) at different distances and \( t = 0 \) by accounting for all the mass within the given distance from the forming sink and by integrating over the absolute magnetic flux through the spherical test surface. The plot shows a general trend of decreasing mass-to-flux ratio with increasing radius. This correlation is caused by the fact that gravitational collapse already started at \( t = 0 \), with the effect that mass already has piled up close to the center of the collapsing core. We are aware that the implementation of accretion onto sink particles causes diffusion of mass across field lines, deviating from the assumption of ideal MHD. Therefore, we plot the mass-to-flux ratios at \( t \approx 0 \), when the sink mass is almost zero and the diffusion effect is negligible. We point out that one cannot avoid some magnetic diffusion in numerical MHD simulations, but that the magnetic diffusivity is certainly underestimated in our case, compared to models that explicitly account for nonideal MHD effects. Considering the mass-to-flux ratios only on scales of the pre-stellar core in the range of \( 10^3 \)–\( 10^4 \) au, we find supercritical mass-to-flux ratios that lie in the range between \( \approx 2 \) and \( \approx 10 \), in agreement with observations (Falgarone et al. 2008; Girart et al. 2009; Beuther et al. 2010). With respect to disk formation, we find that mass-to-flux ratios at a few to several au distances are largest around sinks 6 and 8, which also are the sinks that show the most significant disk formation. This sounds intuitively reasonable, given that a higher mass-to-flux ratio implies less magnetic field strength and therefore potentially weaker magnetic braking. However, the mass-to-flux ratios around sinks 4 and 9 are in contrast the lowest, but as seen in Figure 13, disks form around both sinks. We interpret this result to mean that a spherical mass-to-flux ratios is not a very precise proxy for the formation of disks from pre-stellar cores, due to the effects of turbulent motions. As discussed above, the pre-stellar cores are not spherical in shape, and apparently the differences in the distribution of specific angular momentum inside the cores have a stronger effect on disk formation than the mass-to-flux ratios.

### 3.7. Angular Momentum and Magnetic Field Misalignment

The fact that the Sun contains more than 99.9% of the mass in the solar system, but less than 1% of the total angular momentum of the entire system, has been one of the major puzzles in models of solar system formation and thus of star formation in general. In classical core-collapse models of star formation, it was thus suggested that angular momentum can be efficiently transported by magnetic fields through magnetic braking. However, studies and simulations of magnetized spherical core collapse revealed that this mechanism is in fact effective enough to suppress the formation of circumstellar...
disks altogether. One mechanism that could help prevent the so-called “magnetic braking catastrophe” is ambipolar diffusion, as suggested by Mouschovias (1977). However, more detailed studies by Mellon & Li (2009) and Li et al. (2011) showed that ambipolar diffusion will rather cause a strong magnetic field at small circumstellar radii, leading to efficient accretion shocks.

Another suggestion of how to reduce the effect of magnetic braking is changing the angle between the total angular momentum vector and the total magnetic field vector of the collapsing core. Theoretical analysis indicated that magnetic braking is most efficient when the angular momentum and the magnetic field vectors are perpendicular, and comparatively weaker for a parallel configuration (Mouschovias & Paleologou 1979). More recent results (Joos et al. 2012; Krumholz et al. 2013) confirmed the fundamental idea that disk formation can be suppressed for certain angles, but, in contrast to the theoretical prediction by Mouschovias & Paleologou (1979), found that magnetic braking acts strongest in the case of a parallel alignment of the two vectors and weakest in a perpendicular configuration. As pointed out by Joos et al. (2012), the reason is that magnetic field lines are dragged toward the center of the collapse, such that an initial parallel field develops sections that are nearly perpendicular to the angular momentum vector, which act as efficient “lever arms.” Simulations of the collapse of spherical magnetized cores including the effect of turbulence by Seifried et al. (2013) also showed that circumstellar disks can form around young protostars without relying on increased magnetic diffusion due to the nonideal MHD effects. Although they generally agree with the result that a misalignment of the magnetic field with respect to the angular momentum vector facilitates disk formation, their conclusion is that this misalignment is rather the consequence of a more fundamental quantity present in GMCs, namely, turbulence. In our zoom-in simulations, we have the unique possibility to investigate the angular momentum transport in a consistent setup, rather than imposing ad hoc conditions on idealized core-collapse models.

In Figure 18 we show the evolution of the angle between the total angular momentum vector and the total magnetic field orientation within a radial distance of 100 au from the sinks. The plot illustrates the significant fluctuations of the angle around all sinks during their evolution, and thus we conclude that in the early phases of protostellar evolution, where the majority of the protostar is assembled, the relative orientation...
of the two quantities is highly dominated by the underlying turbulence and reflects the infall of gas clumps from different locations and with different angular momentum. We also see that in the case of sinks that quickly evolve a strong rotationally supported circumstellar velocity profile, the angle typically varies around 90° with ±50° during the formation period.

In some cases, the angle between the angular momentum vector and magnetic field vector can still vary significantly even during late phases, such as seen for sink 8 around 60 kyr. This particular change from an angle of about 20° to nearly 180° is most probably related to the peak in the accretion profile shortly after 60 kyr. Considering that the rise in the accretion profile is caused by an infalling clump of mass, it is not surprising that this clump can account for a significant shift in the alignment between magnetic field and angular momentum. Comparing the angle between the total angular momentum vector and the total magnetic field vector of the gas around sinks with α does not show a clear correlation between the level of rotation and the angle between the mean magnetic field vector and the total angular momentum vector, indicating that in realistic situations it is an oversimplification to focus on alignment/misalignment.

3.8. Replenishment Time

To improve our understanding of the dynamics around the sink, we plot where the mass that is located within a radial distance of 100 au from the sink and that has not accreted onto the sink was located at the time of stellar birth in Figure 19. One can see that for increasing times the gas stems from regions that were initially farther away from the sink, consistent with the classical model of inside-out collapse. However, comparing the different sinks at similar late times, a fraction of the gas may stem from very different distances (e.g., less than 5000 au for sink 1, but beyond 10^4 au for sink 10). We interpret this as a sign of spatial variations of the pre-stellar cores and particularly as evidence of infall through large-scale filaments feeding the cores. Looking more closely at the individual profiles, the origin of the gas around sink 6 is striking. Initially, it shows the same trend of an increasing original distance with increasing time, but when comparing t = 50 kyr with t = 100 kyr, almost no difference can be seen except for a lower mass contribution from the gas that was initially closely located to the sink. This feature may be caused by the fact that the gas orbits inside the disk for a long time without any addition of new gas. However, the significant accretion rates of more than 10^{-6} M_☉ yr^{-1} in the early phases of star formation (Figure 6) and the substantial disk masses around some of the protostars (and of sink 6 in particular) indicate that the gas inside the protoplanetary disk gets replenished rather quickly.

To illustrate the short replenishment time, we plot in Figure 20 the fraction of tracer particles that are located within 100 au from the sink at t = 50 kyr during the evolution of the stellar surrounding. For the sinks that show no or only weak signs of disks, the tracer particles move rapidly in radial direction as indicated by the sharp peak in the plot. The sinks with stronger signs of disks show broader peaks, but even in the case of the strongest disk (sink 6, magenta line), most of the gas only remains in that region for at most a few times 10^3 yr. This is in approximate agreement with the expected replenishment times of the disks defined as

\[ t_{\text{repl}} = \frac{m_{\text{disk}}}{m_{\text{accc}}}. \]  

3.9. Impact on Planet Formation

We have shown that circumstellar disks form at different times after stellar birth. In some cases disks of significant size already form less than a few times 10^4 yr after stellar birth, while in other cases they have not formed within the first 100 kyr. Assuming that these disks are indeed protoplanetary disks, the difference in disk formation times suggests that planet formation may occur at significantly different times after protostellar birth.

Current models of planet formation assume power-law profiles for disk properties that depend on radius and height. Motivated by the MMSN assumption, these power laws are time independent. However, as shown in this study and in previous works, as well as in agreement with observations (Tobin et al. 2015), protoplanetary disks may extend to radii of a few to several tens of au already in the early phase of protostellar evolution. Together with observations of gaps in the dust distribution of disks that are younger than 1 Myr (HL Tau; Maury et al. 2014; ALMA Partnership et al. 2015) and commonly assumed average disk ages of about 2 Myr (for a critical analysis of potential underestimating disk ages due to selection biases, please refer to Pfalzner et al. 2014), this raises the question whether planet formation already happens in the first few megayears, perhaps starting already in the first few hundred kiloyears after protostellar formation. At this stage, material moves rapidly in the radial direction through the disk, and the protoplanetary disks might still be fed with new material as discussed above. Therefore, the assumption of the disk mass as the mass reservoir for planet formation may be an oversimplification. Instead, the mass reservoir for planet formation might in practice be significantly larger, even for low disk masses, when accounting for the effects of rapid radial motions inside the disks and external infall onto the disk. Hence, we emphasize that planets may have a much larger mass reservoir, when properly accounting for all the mass that travels through the protoplanetary disk during the period of planet formation.

4. Conclusions

In this study, we investigated the first ~100 kyr of protostellar formation and evolution with a 2 au minimum cell size in nine cases, where 1–2 M_☉ stars (sink particles) formed when the minimum cell size was 126 au. Accounting for the physical environments in which the stars were embedded, we found that the accretion process onto protostars is heterogeneous in multiple ways, namely,

1. in time,
2. in space,
3. among protostars.

Accretion is heterogeneous in time, in the sense that accretion rate time profiles vary significantly. Initially, accretion rates are of the order of 10^{-5} to 10^{-4} M_☉ yr^{-1} and generally decrease during the subsequent evolution. However, the protostars in our simulations sometimes underwent periods of increased accretion, in which the accretion rates were enhanced by factors of a few to several. We saw evidence for
event amplitudes increasing with increasing spatial resolution, and we note that event maxima may be missed unless the snapshot cadence is correspondingly increased.

Accretion is heterogeneous in space, in the sense that mass accretion onto the star–disk system is filamentary, acting through accretion channels and accretion sheets, rather than in the form of a smooth, space-filling infall of mass.

Last, the accretion processes differ among protostars, both quantitatively and qualitatively, dependent on physical properties such as density, magnetization, and the strength of turbulence. These properties are typically determined by the dynamics of the environment on length scales similar to or larger than $\sim 10^4$ au, i.e., on scales that characterize the dynamical fluctuations in the general turbulence of the GMC.

We thus conclude that the diversity in the large-scale stellar environment profoundly influences the formation and evolution of protoplanetary disks. If the magnetization of the surrounding gas is not too large, protoplanetary disks can form as early as a few thousand years after star formation. In cases where the magnetization of the collapsing gas is sufficiently large (low mass-to-flux ratios), no disk of more than a few au in size may form around the star. We suggest that the main reason why the magnetic braking catastrophe is avoided in many cases is the reduction of magnetic braking caused by turbulence (see also Figure 19. Distribution of where the gas that has accreted onto the different sinks was located after times indicated in the legends at $t = 0$ kyr. The panels show from left to right and top to bottom the distribution for sinks 1, 3, 4, 5, 6, 7, 8, and 9.)
Seifried et al. (2013). The fact that protoplanetary disks form in our simulations, even though we do not account for nonideal MHD effects, suggests either that nonideal MHD effects are not important in situations with realistic turbulence, or that disks form even more frequently than seen in our study. Limited numerical resolution (see Appendix A) is not likely to play a significant role, since accretion rates and frequencies of disk formation were similar in early versions of these zoom-in simulations, where the number of cells per Jeans length was significantly smaller (Nordlund et al. 2014).

We also studied the setting of parameters in our sink particle recipe. On the one hand, the average mass accretion profiles of the sinks are rather robust to a broad range of settings, but on the other hand, changes in the settings can have significant effects on the exact conditions under which protoplanetary disks form. Choosing settings that favor disk formation has the side effect that massive clumps of gas may accrete onto the sink once in a while, thus causing periods with significant accretion bursts. We thus conclude that, at least when the accretion process is modeled with transition from flows to sink particles occurring at scales of a few au, there is a significant uncertainty from parameter dependencies. However, these uncertainties do not change the overall physical conclusion, that the outcome of star and disk formation strongly depends on the pre-stellar environment.

We conclude that protoplanetary disk formation is a ubiquitous process in GMCs, with rotationally supported disks forming as a result of the specific angular momentum of the collapsing gas, even though the process is counteracted by magnetic braking. An important factor ensuring that the formation of disks is not entirely suppressed is the presence of turbulence (in both flows and magnetic fields) inherited from the large-scale dynamics.

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### Appendix A

#### Refinement and Sink Accretion Parameters

To limit the computational costs only to the regions of interest, we apply geometrical refinement in our zoom-in runs, where a cell is only allowed to be resolved to a location-dependent maximum level of \( \ell = 1 + \text{max}(r) \). We then mark cells for refinement if they exceed level-dependent threshold values for one of the refinement criteria, such as number of cells per Jeans length, steep gradients in velocity, magnetic field magnitude, gas density, or gas pressure. The maximum allowed refinement level decreases with increasing distance from the protostar according to a refinement ladder (as shown in the left panel of Figure 21). Such a limitation on the refinement (referred to as “geometrical refinement” in the RAMSES user guide) is thus only a necessary, but not a sufficient, condition for cells to be refined. Our most important sufficient condition for refinement is the density (Jeans length) criterion. A cell is marked for refinement to a certain level if the density exceeds a level-dependent threshold value. We illustrate the chosen refinement ladders for the different runs in the right panel of Figure 21.

With respect to density gradients, cells are generally refined if the difference between neighboring cells exceeds a factor of 20. However, this only applies up to level of 20, and steeper density gradients are allowed for cells above level 20. Similarly, refinement is triggered by gradients in the speed and the magnetic field strength. We have found that refining on gradients helps in maintaining waves and features in the flow and in resolving shocks, in particular in the vicinity of the protostars. In Table 2 the parameters for refinement due to gradients in velocity and magnetic field strength in the different runs are listed. The first column shows the number of the sink. The second column indicates the maximum level at which cells are refined based on gradients in the velocity. The value in parentheses indicates the time in kyr after sink creation at which the maximum level was increased from 20 to 22. Refinement is triggered when the relative velocity gradient differs by more than a factor of \( f = 5 \) compared to the sound speed using the following inequality to trigger refinement:

\[
\sum_{i=x,y,z} \max \left( \frac{2|v_i^+ - v_i^-|}{c_i^2 + c_i^2}, \frac{2|v_i^+ - v_i^-|}{c_i^2 + c_i^2} \right) > f^2,
\]

where \( c_i \) is the sound speed and \( -\), \( 0 \), \( + \) indicate cells at \(-1\), \( 0 \), \( +1 \) cell distance from the center cell. The third column indicates whether cells are selected for refinement if the magnetic field strength between neighboring cells differs by more than a factor of 3. We point out that although these settings differ slightly between runs, their effect is only minor. The most important criterion for refinement is the density.

Besides the different refinement criteria, the remaining numerical parameter that is different in the runs is listed in the fourth column: it is the threshold value of the fast magnetosonic velocity at which the code switches from the HLLD solver to a more diffusive LLF (Lax–Friedrichs) solver.

![Figure 20. Number of tracer particles located within 100 au distance from the sink at \( t = 50 \) kyr, compared to the number of particles within 100 au. Blue corresponds to sink 1, green to sink 3, red to sink 4, cyan to sink 5, magenta to sink 6, yellow to sink 7, black to sink 8, and blue dashed to sink 9.](image-url)
This is to avoid developing exceedingly large magnetosonic speeds in otherwise uninteresting places, which has aggravating effects on the time step.

Appendix B
Sink Parameter Comparison

To assess the robustness of our results, we investigate the impact on the accretion process of setting different parameters in our sink particle model. To do this, we carried out a parameter study with the sink settings given in Table 3, for zoom-in simulations around sink 1. The table provides an overview of the parameters that are important for accretion onto the sink. The third column displays the density limit that has to be exceeded by a cell in order to accrete mass onto the sink. The fourth column gives the radius in cells inside which mass is allowed to be accreted onto the sink \( r_{\text{acc}} \). In the fifth column, we give the accretion rate \( \text{acc}_{\text{rate}} \) from a cell onto the sink. \( \text{acc}_{\text{rate}} \) is a prefactor, which is related to the assumed angle between the angular motion and the radial motion of the gas very close to the sink, so that the amount of mass removed becomes proportional to \( \text{acc}_{\text{rate}} \rho v_K \), where \( v_K \) is the Keplerian speed. The sixth column gives the fraction of the accreting mass that is added to the sink. The remaining mass is removed from the box, to mimic the mass lost in outflows. According to observations—in particular, the correspondence of core mass function to initial mass function in particular, the correspondence of core mass function to initial mass function between one-half and two-thirds of the envelope mass is lost in winds and due to envelope mass that is entrained in the wind, but it is nontrivial to measure how much. We therefore experiment with
setting the accreted fraction to both 50% and 100%, to gauge the importance.

We started each run from shortly before the time of sink formation on level 22 and evolved the runs for up to 100 kyr. In the left panel of Figure 22 we illustrate the accretion profiles of the different runs. In run 1a (blue dots), the accretion rate quickly increases to its maximum accretion rate of several times $10^{-3} \, M_\odot \, yr^{-1}$ at the very beginning of the accretion process. Afterward, the accretion rate drops more or less continuously—with a small shoulder around 20 kyr—to an accretion rate of $10^{-2} \, M_\odot \, yr^{-1}$. In run 1b (red triangles), we apply the same setting except for increasing acc_eff from 50% to 100%. The overall characteristic of a steep initial increase followed by a continuous relaxation is the same. As expected, the higher accretion efficiency parameter causes a higher accretion rate during the evolution. In run 1c (cyan squares), we decreased acc_rate to $10^{-3}$ and used an otherwise identical setting as for 1b. The differences are very subtle, showing that lowering the instantaneous accretion rate just lets the sink particle accrete more slowly or across more cells, even though sometimes the lower rate results in a significant pileup. Run 1c shows stronger fluctuations than 1a, including episodes of accretion bursts, while generally following accretion rates of similar strength as 1b. In run 1d (magenta asterisks), we decreased r_acc to 7.5 cells, with other parameters the same as in 1c. Similar to 1c, the accretion rate shows fluctuations of up to a factor of 2. The fluctuations—especially at later times—do not overlap with the ones in 1c and seem to be somewhat stronger. In run 1e (yellow plus signs), we increased rho_fr from $10^{-4}$ to $10^{-1}$ with respect to 1d and find a profile that shows a smoother evolution. In run 1f (green circles), we decrease r_acc to four cells, with other parameters as in 1e. We again find fluctuations in the accretion of mass onto the sink. Moreover, we see a striking difference, in that the accretion profile drops more steeply after about 30 kyr than seen for the other sinks with acc_eff = 1. After about 50 kyr and apart from the significant fluctuations, the strength of the accretion rate is close to the accretion rate for the run with acc_eff = 0.5 (1a). Finally, we carry out a comparison run 1g, with the same settings as in 1f, but with acc_eff = 0.5 (black crosses). We find that the profile is very similar to run 1a (both in absolute accretion rate and in the shape of the profile) for about the first 30 kyr and then shows stronger fluctuations. Run 1g does not show the significant drop in the accretion rate that was seen in 1f.

The most striking difference is the effect of the efficiency parameter, which separates the accretion profiles into two groups. The sinks that accumulate 100% of the mass selected in their surroundings (1b–1f) show higher mass accretion rates than in the case where 50% of that mass is removed from the box (run 1a and run 1g). Generally, the profiles then evolve in a similar manner, with some smaller fluctuations and more or less strongly evolved brief burst periods (e.g., run 1e at about $t = 50$ kyr) until the end of the simulation. However, there is one particular exception (run 1f), which evolves in a similar fashion to the other runs with acc_eff = 100% until about $t = 35$ kyr before the accretion rate drops more steeply until at about $t = 50$ kyr, and finally follows more the evolution of the sinks with low acc_eff, though with significantly stronger bursts than run 1a. Both the stronger amount of fluctuations in the accretion profiles and the fall-off of run 1f can be understood by studying the velocity profiles around the sinks.

In the right panel of Figure 22, we illustrate the evolution of $\alpha$ for the different sinks and find that for some settings a significant rotational gas motion evolves, whereas the velocity profile for other settings is mostly infall dominated during the entire evolution. The accretion profile of run 1a evolves rather smoothly and calmly, because there is no sign of significant disk formation for this setting of the sink parameters. The sinks with a surrounding disk show more intermittent accretion profiles, because clumps of mass rotating in the disk may eventually fall into the sink. Furthermore, the drop-off in the accretion profile in run 1f coincides with the drop in $\alpha$, and thus with the formation of a disk. The fact that a disk starts to form means that infalling mass does not accrete directly onto the star, but starts to build up a disk first. Therefore, the accretion profile drops more significantly compared to the runs where no disk or only a weak disk forms. We interpret the fact that disk formation occurs more strongly for high acc_eff as a direct consequence of the deeper gravitational potential induced by the more massive sinks.
From the results we conclude that the general profile of the sink accretion is robust to variations of the sink parameter settings over longer timescales. However, changes (particularly of acc\_eff) affect the disk formation process and therefore also the accretion process on small timescales. If the sink forms a circumstellar disk, the sink accretes mass in occasional bursts, causing a more intermittent accretion profile but affecting the long-term accretion profile only modestly. The parameter study illustrates the difficulties in choosing physically correct parameter settings. We selected the settings from run 1a, which shows the weakest sign of disk formation for the zoom-in simulations, for the sinks that formed on the highest level of refinement. This was a conservative choice in terms of determining the efficiency of disk formation, because choosing a higher accretion efficiency increases the probability of disk formation. Therefore, the frequency of disk formation we find is likely to be a lower limit and would have been higher if we had chosen settings more favorable for disk formation.

In comparison to effects due to the choice of sink particle parameters, limited numerical resolution is not likely to play a significant role, since accretion rates and frequencies of disk formation were largely similar in early versions of these zoom-in simulations, where the number of cells per Jeans length was significantly smaller (Nordlund et al. 2014).

**Appendix C**

**Angular Momentum Conservation**

Here we summarize how the change in angular momentum in a volume is related to flux densities through the surface of the volume. Our derivation is similar to that given in Joos et al. (2012), but more compact, since we use a conservative formulation of the Euler equation. It is also coordinate independent and contains pressure terms, which canceled out for the specific geometry of the volume and angular momentum component considered by Joos et al. (2012).

The Euler equation in conserved form may be written as

\[
\frac{\partial \rho v}{\partial t} = - \nabla \left( \rho v \otimes v - \frac{1}{4 \pi} B \otimes B \right) + \frac{1}{4 \pi G} \nabla \Phi \otimes \nabla \Phi - \nabla P_{\text{tot}},
\]

where the total pressure is

\[
P_{\text{tot}} = P + \frac{B^2}{2 \pi} + \frac{(\nabla \Phi)^2}{8 \pi G}
\]

and we have used a conserved formulation of the gravitational force (Jiang et al. 2013). The total angular momentum inside a volume with surface \(S\) is

\[
L = \int r \times \rho v \, dV.
\]

The time evolution of the total angular momentum can be calculated by integrating the time evolution of the angular momentum density \(L = r \times \rho v\), which is directly related to the Euler equation. To transform from \(\rho v\) to \(L\), we use the tensor identity

\[
r \times \nabla \cdot (X \otimes Y) = \epsilon_{ijk} r_j \partial_i [X_k Y_i] e_i
\]

\[
= \partial_i [X_k \epsilon_{ijk} r_k Y_i] e_i - \epsilon_{ijk} X_k Y_i e_i
\]

\[
= \nabla \cdot (X \otimes (r \times Y)) - X \times Y,
\]

where we have written it out in component form too (using the Einstein summation convention), to explicitly specify which tensor index the divergence applies to. The flux terms in the Euler equation are all symmetric tensors, and as a consequence the last term in Equation (20) disappears, leaving a pure divergence. The pressure term may be written as

\[
r \times \nabla P_{\text{tot}} = - \nabla \times (P_{\text{tot}} r).
\]

Transforming the volume integral into surface integrals, we find the time evolution of the total angular momentum expressed as a function of the flux through the surface

\[
\frac{dL_z}{dt} = - \int_S (r \times \rho v) (v \cdot n) - \frac{1}{4 \pi} (r \times B) (B \cdot n) + \frac{1}{4 \pi G} (r \times \nabla \Phi) (\nabla \Phi \cdot n) \, dA - \int_S P_{\text{tot}} r \times n \, dA,
\]

where \(n\) is the normal vector to the surface element \(dA\).

If our control volume is a sphere, the pressure term disappears, since \(n = e_z\). By rotational symmetry this is also the case for the \(x\)- and \(y\)-component. In the case where the control volume is a cylinder the pressure term also disappears, but only for the \(z\)-component of the angular momentum. If instead a cubical test volume is used, the pressure term has to be included when computing the time derivative of any component of the angular momentum.

In cylindrical coordinates, the change in the \(z\)-component of the angular momentum may be split up into a contribution from the cylinder wall, assumed to be at \(r = R\), and contributions from the top and bottom layer at \(z = \pm h/2\). The contribution from the cylinder wall is

\[
\frac{dL_z}{dt} \bigg|_{\text{cyl}} = - \int_{-h/2}^{h/2} dz \int_0^{2\pi} \int_0^R R \, d\phi \, \rho \nu_v v_r
\]

\[
- \frac{1}{4 \pi} B_r B_z + \frac{1}{4 \pi G} (\nabla \Phi)_{\phi} (\nabla \Phi)_{\phi},
\]

where all variables are evaluated at the fixed radius \(R\). The contributions from the top or bottom layer at fixed height \(\pm h/2\) are

\[
\frac{dL_z}{dt} \bigg|_{\text{top/bot}} = \mp \int_0^R dr \int_0^{2\pi} \int_0^{\pi/2} r \, d\phi \, \rho \nu_v v_z
\]

\[
- \frac{1}{4 \pi} B_r B_z + \frac{1}{4 \pi G} (\nabla \Phi)_{\phi} (\nabla \Phi)_{\phi},
\]

where the sign differs according to whether it is top or bottom layer, respectively, that is calculated.

In spherical coordinates the \(z\)-component of the angular momentum may be written as

\[
\frac{dL_z}{dt} \bigg|_{\text{ sph}} = - \int_0^{\pi} R \, d\theta \int_0^{2\pi} R \sin \theta \, d\phi \, R \sin \theta \, \rho \nu_v v_r
\]

\[
- \frac{1}{4 \pi} B_r B_z + \frac{1}{4 \pi G} (\nabla \Phi)_{\phi} (\nabla \Phi)_{\phi},
\]

where all variables are evaluated at the fixed radius \(R\) and the coordinates are \((r, \phi, \theta)\), with \(\theta\) being the polar angle.
Appendix D

Grid Alignment

Given the turbulent motions in the GMC, the mean angular momentum vector of the collapsing core is expected to be randomly oriented with respect to the grid in our simulation. Figure 23 shows the orientation of the mean angular momentum vector calculated within a sphere of 100 au radius around the sinks during the evolution of the zoom-in runs. At $t = 0$, we find a variation of different orientations of the mean angular momentum vector as expected from the underlying turbulence. However, with evolving time we see a tendency of alignment of the mean angular momentum vector with one of the coordinate axes for the runs that formed protoplanetary disks. This alignment is caused by numerical effects in our grid code RAMSES as pointed out by Hopkins (2015). Certainly the presence of grid alignment makes a detailed study of the disk dynamics challenging. The artificially aligned disk may induce additional torques on infalling gas and in this way enhance the presence of grid alignment and may influence the outcome of our runs to some degree, but the detected differences in the accretion and protoplanetary disk formation process are too large to be solely explained by numerical effects, whereas they are in line with the “inherited” variations from the stellar environments.

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Figure 23. Evolution of the angle between the total angular momentum vector of the gas within 100 au from the sink and the three coordinate axes (left panel: x-axis; middle panel: y-axis; right panel: z-axis). The symbols belong to the same sinks as in Figure 5.
