Magnetic Effects Promote Supermassive Star Formation in Metal-enriched Atomic-cooling Halos

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Abstract

Intermediate-mass black holes (with \(\geq 10^5 M_\odot\)) are promising candidates for the origin of supermassive black holes (with \(\sim 10^9 M_\odot\)) in the early universe (redshift \(z \sim 6\)). Chon & Omukai first pointed out direct collapse black hole (DCBH) formation in metal-enriched atomic-cooling halos (ACHs), which relaxes the DCBH formation criterion. On the other hand, Hirano et al. showed that magnetic effects promote DCBH formation in metal-free ACHs. We perform a set of magnetohydrodynamical simulations to investigate star formation in magnetized ACHs with metallicities \(Z/Z_\odot = 0, 10^{-5}, \) and \(10^{-4}\). Our simulations show that the mass accretion rate onto the protostars becomes lower in metal-enriched ACHs than in metal-free ACHs. However, many protostars form from gravitationally and thermally unstable metal-enriched gas clouds. Under such circumstances, the magnetic field rapidly increases as magnetic field lines wind up due to the spin of protostars. The region with the amplified magnetic field expands outwards due to the orbital motion of protostars and the rotation of the accreting gas. The amplified magnetic field extracts angular momentum from the accreting gas, promotes the coalescence of low-mass protostars, and increases the mass growth rate of the primary protostar. We conclude that the magnetic field amplification is always realized in metal-enriched ACHs regardless of the initial magnetic field strength, which affects the DCBH formation criterion. In addition, we find a qualitatively different trend from the previous unmagnetized simulations in that the mass growth rate is maximal for extremely metal-poor ACHs with \(Z/Z_\odot = 10^{-5}\).

Unified Astronomy Thesaurus concepts: Magnetohydrodynamical simulations (1966); Primordial magnetic fields (1294); Population II stars (1284); Star formation (1569); Protostars (1302); Supermassive black holes (1663)

1. Introduction

The formation of supermassive black holes (SMBHs) with mass \(\sim 10^9 M_\odot\) in the early universe at redshift \(z \geq 6\) is one of the outstanding issues in astrophysics (see Woods et al. 2019 for a review). As researchers find SMBHs in earlier epochs of the universe, the demands on theoretical models become greater. The number of observational samples will increase with upcoming observations, such as with the James Webb Space Telescope (JWST). The theoretical formation scenario for SMBHs should be updated accordingly.

One of the popular formation channels of SMBHs is the direct collapse black hole (DCBH) scenario (see Inayoshi et al. 2020 for a review). In metal-free atomic-cooling halos (ACHs), the gas component is cooled only by atomic hydrogen and can gravitationally contract while remaining at a high temperature of \(\sim 8000\) K. The high temperature leads to a large Jeans mass, suppression of cloud fragmentation, and also a high mass accretion rate that allows the protostar to increase its mass efficiently. The high mass accretion rate leads to an expanded envelope of the protostar with a relatively low surface temperature of a few thousand K (Hosokawa et al. 2012), hence limiting radiative feedback that can otherwise halt accretion. The result is the formation of a supermassive star (SMS). Once the SMS exceeds the mass threshold of the general relativistic instability, the star can collapse to form a massive black hole with a similar mass \((\sim 10^5 M_\odot)\); Umeda et al. 2016). This marks the formation of an intermediate-mass black hole (IMBH) with \(\sim 10^5 M_\odot\), a candidate seed of SMBHs. To explain the observed number density of high-redshift quasars (a few Gpc\(^{-3}\); e.g., Bañados et al. 2016), previous studies provided a number of conditions under which metal-free ACHs can be realized: \(H_2\)-dissociating ultraviolet radiation (e.g., Omukai 2001; Agarwal et al. 2012; Latif et al. 2013), high-velocity collisions (Inayoshi et al. 2015), baryon–dark matter streaming velocities (Tanaka & Li 2014; Hirano et al. 2017), dynamical heating due to a violent halo merger (Wise et al. 2019), and turbulent cold flows (Latif et al. 2022).

Recently, Chon & Omukai (2020) pointed out the possibility of IMBH formation in metal-enriched ACHs. Traditionally, researchers have considered that low-metallicity gas clouds were unsuitable for IMBH formation because the low gas temperature due to the fact that the metal and dust coolings produce numerous low-mass protostars, reducing the accretion rate onto each protostar by fragmentation. Chon & Omukai (2020) performed a set of hydrodynamical simulations to follow the star formation in ACHs with different metallicities \(Z/Z_\odot = 10^{-6} - 10^{-3}\). They found that for slightly metal-enriched ACHs with \(Z/Z_\odot \leq 10^{-4}\), thousands of stars are formed by frequent fragmentation of the gas cloud. However, most stars competitively merge into the most massive (primary) star (“supercompetitive accretion”). The mass growth rate of...
phase, as shown in Figure 1 (Hirano & Machida 2022). The resultant strong magnetic field promotes the coalescence of protostars and enhances mass accretion onto the primary protostar at the center of the metal-free ACHs (Hirano et al. 2021, hereafter Paper I). Evolution over cosmic time means that metal-enriched ACHs may initially have a stronger magnetic field than metal-free ones. Therefore, we speculate that magnetic effects might be more effective for metal-enriched ACHs than for metal-free ones. This paper investigates whether or not the magnetic effects can promote DCBH formation in metal-enriched ACHs.

2. Numerical Methodology

We perform a set of three-dimensional (3D) ideal magnetohydrodynamical (MHD) simulations of star formation in metal-enriched ACHs. The numerical simulation setup follows Paper I, but we switch the thermal evolution models of the ACHs according to the metallicities, as shown in Figure 2.

2.1. Numerical Methodology

We solve the ideal MHD⁵ equations (Equations (1)–(4) of Machida & Doi 2013) with the barotropic equation-of-state (EOS) tables. We adopt EOS tables of ACHs under strong Lyman–Werner radiation⁶, \( J_{31} = 10^3 \) for metallicities of \( Z/Z_c = 0, 10^{-5}, \) and \( 10^{-4} \) based on a simulation of chemical reactions (Omukai et al. 2008) as used in Chon & Omukai (2020). We assume a blackbody spectrum of Population II stellar sources with an effective temperature of \( 10^4 \) K to realize the strong Lyman–Werner radiation. We define the EOS model with \( Z/Z_c = 10^{-6} \) as the metal-free case (\( Z/Z_c = 0 \)) because the thermal evolution with \( Z/Z_c = 0 \) follows that with \( Z/Z_c = 10^{-6} \) (Omukai et al. 2008). Instead of a sink particle technique, we adopt a stiff-EOS technique to accelerate the time evolution while connecting magnetic field lines to high-density regions, which is necessary to reproduce the magnetic field amplification (Paper I, Hirano & Machida 2022). We set a threshold density \( n_{th} = 10^{16} \text{ cm}^{-3} \), which reproduces dense cores whose radius is roughly consistent with the approximate mass–radius relation of a rapidly accreting protostar with \( R_{\ast} \approx 12(M_{\ast}/100 M_\odot)^{1/2} \text{ au} \) (Equation (11) of Hosokawa et al. 2012). Note that the difference in radius between our model and that of Hosokawa et al. (2012) is about a factor of 10 at the maximum. The radii of the protostars formed in our simulations tend to be smaller than those in Hosokawa et al. (2012). Figure 2 shows the thermal evolution models defined by the resultant EOS tables.

We use our nested grid code (Machida & Nakamura 2015), in which rectangular grids of \( (n_s, n_s, n_s) = (256, 256, 32) \) are superimposed. The grid size \( L(l) \) and cell width \( h(l) \) of the \( l \)th grid are twice as large as those of the \((l+1)\)th grid (i.e., \( L(l) = 2L(l+1) \) and \( h(l) = 2h(l+1) \)). We use the index “l” to describe a grid level. The base grid \((l = 1)\) has grid size \( L(1) = 6.68 \times 10^7 \text{ au} \) and cell size \( h(1) = 2.61 \times 10^5 \text{ au} \). A finer grid is generated to resolve the Jeans wavelength of at least 32 cells. The finest grid \((l = 18)\) has grid size \( L(18) = 510 \text{ au} \) and cell size \( h(18) = 1.99 \text{ au} \).

We assume that the nonideal MHD effects are ineffective in metal-enriched ACHs with a high temperature of >100 K in this study, but discuss this in Section 4.3.

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\(^{6}\) \( J_{31} \) is the intensity at the Lyman–Werner bands normalized in units of \( 10^{-21} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ Hz}^{-1} \text{ sr}^{-1} \).
2.2. Initial Condition

We adopt the initial condition as a cloud with an enhanced Bonnor–Ebert (BE) density profile \( f \times n_{\text{BE}}(r) \) where \( f = 1.2 \) is an enhancement factor to promote cloud contraction. The initial cloud has a central density \( n_0(0) = f \times n_{\text{BE}}(0) = 1.2 \times 10^4 \text{ cm}^{-3} \), mass \( M_0 = 1.82 \times 10^3 M_\odot \), and radius \( R_0 = 2.09 \times 10^6 \text{ au} = 10.1 \text{ pc} \). We impose a uniform temperature \( T_0 = 7700 \text{ K} \) and a rigid rotation of \( \Omega_0 = 7.67 \times 10^{-15} \text{ s}^{-1} \) for the initial cloud. The resultant ratios of thermal and rotational energies to the gravitational energy of the initial cloud are \( \alpha_{\text{th}} = 0.7 \) and \( \beta_0 = 0.01 \), respectively. We add an \( m = 2 \) mode nonaxisymmetric density perturbation to the initial cloud with the amplitude of the perturbation as 10%. We do not include turbulence and do not consider a small-scale dynamo (e.g., Sur et al. 2010; McKee et al. 2020) because we only consider very weak fields that are significantly amplified by the rotation of protostars and accreting gas (Figure 1).

2.3. Model Parameters

We adopt two parameters for the gas clouds: (1) metallicities and (2) initial magnetic field strength. We adopt three different gas metallicities \( Z/Z_\odot = 0, 10^{-3}, \) and \( 10^{-4} \), which cover the metallicity range for “supercompetitive accretion” \( (Z/Z_\odot \leq 10^{-4}) \); Chon & Omukai 2020). We adopt five different initial magnetic field strengths \( B_0/G = 0, 10^{-12}, 10^{-10}, 10^{-8}, \) and \( 10^{-6} \). We adopt \( B_0/G = 10^{-12} \) at \( n_{\text{th}} = 10^3 \text{ cm}^{-3} \), extrapolated from the cosmological seed value \( (10^{-6} \text{ G at } n_{\text{th}} = 1 \text{ cm}^{-3}) \); Xu et al. 2008) through flux freezing during cloud compression (described by a power law \( B \propto n_{\text{th}}^{1/2} \)) as the minimum value. We impose a uniform magnetic field with the same direction as the initial cloud’s rotation axis in the whole computational domain. In Cartesian coordinates, the initial directions of the global magnetic field and the rotation axis are parallel to the \( z \)-axis.

Table 1 summarizes the simulation models adopted in this study. We define the model names by connecting the common logarithms of two model parameters. We adopt model Z4B12 with \( (Z/Z_\odot, B_0/G) = (10^{-4}, 10^{-12}) \) as the fiducial model. We expect that Z4B12 is the model for which it is most difficult to form an SMS because it has the lowest accretion rate (corresponding to the lowest gas temperature) and the weakest magnetic field among the models in this study. We can conclude that the magnetic effects support DCBH formation in any other model by confirming the formation of an SMS in model Z4B12.

2.4. Simulations

We assume that the region where the gas density exceeds \( n_{\text{th}} \) is a dense core that hosts a protostar. We analyze the number and masses of such dense cores to examine the fragmentation process. We define the epoch of formation of the first protostar \( (t_{\text{ps}} = 0 \text{ yr}) \) when the gas number density first reaches the threshold density \( n_{\text{max}} = n_{\text{th}} \). We terminate the calculation either when \( t_{\text{ps}} = 500 \text{ yr} \) or when the time step in the calculation becomes very short \( (dt \ll 0.01 \text{ yr}) \) because we cannot follow the further time evolution of the system. We could calculate the time evolution for about 200–500 yr after formation of the first protostar (Table 1).

3. Results

This section shows the dependence of the simulation results on two parameters. First, Section 3.1 presents the simulation results of the fiducial model (Z4B12) and compares it with the corresponding nonmagnetized model (Z4B00). Next, Section 3.2 confirms the dependence of the magnetic effects on the metallicities \( (Z) \). Finally, Section 3.3 shows the dependence on the initial magnetic field strength \( (B_0) \). In all metal-enriched models, we confirm the magnetic field amplification as in the metal-free ACHs and the increase in the mass accretion rate onto the primary protostar due to the transfer of angular momentum by magnetic effects.

3.1. Fiducial Model

Figure 3 compares the simulation results of the fiducial model (Z4B12) and the unmagnetized model (Z4B00) during the first 500 yr of the protostar accretion phase. We place three different panels at each epoch: mass spectrum of the dense cores \( (N_{\text{core}}) \); maps of the distributions of the gas number density on the \( z = 0 \) and \( y = 0 \) planes \( (n) \); and maps of the absolute magnetic field strength on the \( z = 0 \) plane \( (|B|) \). When the first protostar forms at the center of the collapsing gas cloud \( (t_{\text{ps}} = 0 \text{ yr}) \), the maps of gas density of models Z4B12 and Z4B00 are identical because the magnetic field strength in the vicinity of the protostar is too weak \( (|B| \sim 10^{-7} \text{ G}) \) at most, as in Figure 4(a)) to affect the dynamics of the collapsing gas cloud.

During the first 100 yr after formation of the first protostar, the region around the cloud’s center becomes gravitationally unstable due to the high mass accretion rate. Fragmentation occurs and forms many protostars (see the mass spectrum in Figure 3). The rapid spin of protostars winds up magnetic fields and amplifies the local magnetic field strength according to \( B = B_0 \exp(5t) \) (Figure 4). During the first 100 yr, the magnetic field strength in the region of about 100 au around the primary protostar is amplified to \( |B| \sim 10^{-4} \text{ G} \) (Figure 5(a)). The peak at \( R \sim 60 \text{ au} \) corresponds to the magnetic field amplified by another core. This amplified magnetic field is strong enough to affect the gas dynamics and decrease the number of dense cores. Figure 4(b) shows two density ranges where the magnetic field is noticeably amplified: a high-density region in the circumstellar disk \( (n \gtrsim 10^{13} \text{ cm}^{-3}) \) and a low-density region just above the disk \( (n \sim 10^{10} \text{ cm}^{-3}) \). The rotation of the dense cores amplifies the magnetic field in the disk, which in turn amplifies the low-density region above (see also the bottom panels of Figure 3). The presence of two peaks is transient and becomes less noticeable over time as the magnetic field amplifies throughout the region at the disk’s outer edge.

We continue the simulations until \( t_{\text{ps}} = 500 \text{ yr} \). The amplified “seed” magnetic field increases the amplification rate according to the induction equation, \( \partial B/\partial t = \nabla \times (\mathbf{v} \times B) \), in which the amplified seed field \( B \) contributes to further increase of the magnetic field. As a result, the amplified region of the magnetic

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\(^{7}\) As the calculation progresses, several small (low-mass) cores are ejected from the center and move to the region with coarse numerical resolution. Such cores are dissolved numerically. We note that the number of low-mass cores is underestimated in the mass spectrum.

\(^{8}\) The induction equation \( \partial B/\partial t = \nabla \times (\mathbf{v} \times B) \) can be simplified to \( \partial B/\partial t = (\mathbf{v} \times B)/L = \Omega B \) where \( \mathbf{v} \) and \( L \) are the typical timescale and typical length scale, and \( \Omega \) is the typical angular velocity, which should correspond to the angular velocity of the spin or orbital motion of protostars. Assuming \( \Omega \) to be constant and integrating the simplified equation \( \partial B/\partial t = \Omega B \), we have \( B = B_0 \exp(\Omega t) \), where \( B_0 \) is a constant of integration. This simplification or estimation holds as long as magnetic feedback can be ignored. The angular velocity \( \Omega \) decreases due to the magnetic tension force when the magnetic field (or Lorentz force) is strong. Since we initially assumed very weak fields, exponential growth can occur.
field strength extends outwards due to the orbital motion of protostars and the rotation of accreting gas (Figures 4(c) and 5).

To clarify the importance of the magnetic field, we plot the plasma beta (ratio of thermal to magnetic pressure) in Figures 4(d)–(f) and 6. Figure 6 shows that the plasma beta on the equatorial plane is $\sim 10^2$–$10^3$ around the center, indicating that the magnetic pressure is $\sim 0.1$–$1\%$ of the thermal pressure (see also Figures 4(d)–(f)). Although the magnetic pressure is lower than the thermal pressure, it can affect the gas dynamics.
The magnetic effects enhance “supercompetitive accretion” more in the magnetized model than in the unmagnetized model. The transfer of angular momentum from the accreting gas increases the mass accretion rate and promotes fragmentation. Then, the transfer of the orbital angular momentum of the number of protostars facilitates the coalescence of low-mass protostars into the primary one. The more protostars form, the more the rotational motion amplifies the local magnetic field. Thus, the magnetic effects positively affect the growth of an SMS in this case.

Figure 7(b) compares the time evolution of the primary protostellar mass between the magnetized (Z4B12; red solid line) and unmagnetized (Z4B00; red dashed line) models. The early histories of mass growth are almost identical until $t_{ps} \sim 240$ yr. After that epoch, the mass growth is faster for the magnetized model than for the unmagnetized one. In the magnetized model, the mass of the primary star suddenly rises due to the runaway coalescence of almost all low-mass protostars (Figure 7(c)) at $t_{ps} \sim 400$ yr. After that, the number of protostars increases to $N_{core} \sim 50$ due to the high mass accretion rate. The intermittent fragmentation and coalescence allow the primary star to gain mass more efficiently.

### 3.2. Dependence on the Metallicity

Next, we examine the dependence of the magnetic effects on metallicity ($Z$). The gas temperature decreases as the metallicity increases. Thus the mass accretion rate should be lower in the cloud with higher metallicity (e.g., Hosokawa & Omukai 2009). The unmagnetized models (dashed lines in Figure 7(b)) represent this trend as the mass growth rate of the primary protostar decreases with metallicity. At the end of the simulations at $t_{ps} = 500$ yr, the masses of primary protostars reach about a few $10^5 M_\odot$, which is consistent with the

| Model | $Z$ | $B_0$ | $t_{col}$ | $t_{ps}$ |
|-------|-----|-------|-----------|----------|
| Z4B00 | $10^{-4}$ | 0 | 1,177,683 | 500 |
| Z4B12 | $10^{-12}$ | 1,177,683 | 500 |
| Z4B10 | $10^{-10}$ | 39 | 250 |
| Z4B08 | $10^{-8}$ | 15 | 500 |
| Z4B06 | $10^{-6}$ | 38 | 290 |
| Z5B00 | $10^{-4}$ | 0 | 1,206,806 | 500 |
| Z5B12 | $10^{-12}$ | 34 | 500 |
| Z5B10 | $10^{-10}$ | 39 | 250 |
| Z5B08 | $10^{-8}$ | 15 | 500 |
| Z5B06 | $10^{-6}$ | 38 | 290 |
| Z0B00 | $0 \sim 10^{-6}$ | 0 | 1,209,600 | 500 |
| Z0B12 | $10^{-12}$ | 98 | 500 |
| Z0B10 | $10^{-10}$ | 72 | 500 |
| Z0B08 | $10^{-8}$ | 108 | 360 |
| Z0B06 | $10^{-6}$ | 98 | 190 |

Note. Column (1): model name. Columns (2) and (3): parameters $Z$ (the metallicity) and $B_0$ (the initial magnetic field strength) of the initial cloud. Column (4): elapsed time when the first protostar forms (corresponding to $t_{ps} = 0$ yr). Column (5): elapsed time at the end of the simulation after formation of the first protostar.
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previous unmagnetized simulation (see Figure 4(a) of Chon & Omukai 2020).

By considering the magnetic effects (solid lines of Figure 7(b)), the masses of the primary protostars at $t_{ps} = 500\,\text{yr}$ become about 10 times greater than those in the unmagnetized models for all metallicity models examined. Note that the mass infall rate far from the central region is almost the same in the magnetized and unmagnetized models. On the other hand, the mass infall rate near the center is higher in the magnetized models than in the unmagnetized ones. Thus, the magnetic field affects the mass accretion rate near the center.9 In the early accretion phase ($t_{ps} < 20\,\text{yr}$), the amplification rate of the magnetic field strength is higher with decreasing metallicity (Figure 7(a)). The origin of this early difference appears more clearly in the “seed” magnetic field amplified by the spin of protostars (Figure 7(a)), as the lower metallicity causes a higher accretion rate and the central high-density gas region fragments into many protostars (Figure 7(c)). After that, all magnetized models will eventually undergo significant magnetic field amplification. The transport of angular momentum by the amplified magnetic field promotes gas accretion and protostellar mergers. After only 500 yr, the primary protostars obtain a mass $\sim 10^3\,M_\odot$.

The mass growth rate is highest in the model with extremely low metallicity (Z5B12) rather than the metal-free model (ZOB12), as shown in Figure 7(b). The increase in metallicity has both positive and negative effects on mass growth: increasing the number of protostars, which contributes to exponential magnetic field amplification, and decreasing the mass accretion rate due to the lower temperature of the gas cloud. The mass growth history of the primary protostar is similar in the unmagnetized models ZSB00 and ZOB00 (Figure 7(b)). Therefore, we find that the effect of positive feedback wins out at the very low metallicity value $Z/Z_\odot = 10^{-5}$, but as the metallicity increases, the negative feedback becomes as important. We see no significant further net feedback effect.

Figure 8 shows 2D maps of density and magnetic field strength around the primary protostar at the end of the simulation ($t_{ps} = 500\,\text{yr}$). In the magnetized models, there are many protostars formed by the burst in fragmentation at $t_{ps} = 400\,\text{yr}$ (Figure 7(c)). Some protostars have escaped farther away from the cloud center (the primary protostar) in the magnetized models (Figures 8(b) and (e)) than in the unmagnetized models (Figures 8(a) and (d)). This is due to the increased efficiency of N-body interactions since the transport of angular momentum due to magnetic effects causes more protostars to fall into the cloud’s central region. We conclude that the rapid accreting gas flow moves protostars away from the cloud’s center in the magnetized models rather than in the unmagnetized models.

The structures of the infalling-accreting envelopes in edge-on views ($y=0$ planes) depend on two parameters. The envelope height decreases with metallicity because of decreasing temperature or thermal pressure. The inner diameter of the envelope (or the diameter of the thin disk) decreases with increasing magnetic field strength due to the transfer of angular momentum.

### 3.3. Dependence on the Initial Magnetic Field Strength

Finally, we examine the dependence on the initial magnetic field strength ($B_0$). Figure 9 summarizes $B$–$n$ diagrams at the same epoch ($t_{ps} = 100\,\text{yr}$) for all models. We confirm the amplification of the magnetic field in all models. There are two distinct density ranges of the amplification: a high-density region on the circumstellar disk ($n \gtrsim 10^{13}\,\text{cm}^{-3}$) and a low-density region above the disk ($n \sim 10^{11}\,\text{cm}^{-3}$) (as explained in Section 3.1).

Figure 10 summarizes the time evolution of properties of the protostar. The mass growth is faster in the magnetized models than in the unmagnetized models. However, the mass growth rate does not always increase with $B_0$. The gas infall (or negative radial) velocity towards the center is large in the amplified region due to the efficient removal of angular momentum by the magnetic field. The infall velocity, and hence the mass growth rate of the protostar, is naively considered to increase with initial magnetic field strength. However, the opposite results are often observed, in which a stronger initial magnetic field slows protostellar mass growth (Figures 10(d)–(f)).

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9 It should be noted that we confirmed that some models show a high mass accretion rate of $\sim 100\,M_\odot\,\text{yr}^{-1}$. However, we calculated the cloud evolution only for $\leq 500\,\text{yr}$ after the formation of first protostar. We need further calculations to evaluate the long-term mass accretion rate correctly.
uniformly throughout the cloud, but proceeds first around the center of the gas cloud where the rotation is fastest. There are strong magnetic fields spread to the outer area where the rotation is slower. Efficient gas accretion occurs when the gas and dense cores are sufficiently present in the cloud.
amplification region. If most of the gas (and other cores) in the amplified region accretes onto the primary core, the rotation-driven amplification of the magnetic field would not work well because of the deficit of the accreting gas. In this case, the magnetic field amplification stagnates until the gas is supplied from the area outside the amplification region. Figures 10(g)–(i) show strong oscillations in the number of cores, especially in magnetized models, indicating violent fragmentation following the rapid coalescence of cores. Figures 10(f) and (i) show the synchronous change in mass growth rate and the number of cores.

Figures 10(d)–(f) confirm that the mass growth is maximized with $Z_{\odot}/Z_{\odot} = 10^{-5}$ for the different values of $B_0$.

Figures 10(j)–(l) also show that the rotational velocities stay low enough that the primary star attains a rotational velocity that is about 0.1 times the Keplerian velocity, regardless of the model parameters. This is consistent with the requirement for SMSs to have low surface velocities of only 10%–20% of the Keplerian value to continue accreting and reach their final masses (Haemmerlé et al. 2018). In addition, the figures indicate that the magnetic field has a subdominant role in determining the rotation of the protostar.

4. Discussion

4.1. Evaluation of the Final Stellar Mass

We show that magnetic effects enhance the mass growth rate in metal-enriched ACHs. One of the limitations of this study is that the mass of a dense core with $n \geq n_{th}$ is the upper limit for
the protostellar mass formed inside the dense core. We have to confirm whether a large amount of gas can accrete efficiently onto the protostar in future work. We note that the rotation-driven amplification of the magnetic field is more efficient inside the unresolved region of this study, $n > n_0 = 10^{16}$ cm$^{-3}$ because the angular velocity $(\Omega = v/r)$ increases with decreasing distance from the central protostar. Then the magnetic effects could also support efficient mass accretion onto the protostar in future work.

In addition, the bursts of fragmentation and merger of protostars also promote the formation of massive stars by preventing protostellar radiative feedback. Figure 4(d) of Chon & Omukai (2020) showed that about 3000 stars formed in $10^5$ yr after the first protostar. The magnetic effects enhance the “supercompetitive accretion” and cause some fraction of (or all) stars to fall onto the primary star. Such a highly variable mass accretion rate is advantageous for DCBH formation because the frequent periods of high accretion rate can inflate the protostar and keep the effective temperature low enough to suppress the radiation feedback that can stop mass accretion onto the protostar (e.g., Sakurai et al. 2016).

4.2. Mass Growth from IMBH to SMBH

After forming IMBHs in metal-enriched ACHs, the issue is whether a high accretion rate can be realized to grow from IMBHs to SMBHs. Yajima et al. (2017) and Park et al. (2022) investigated super-Eddington accretion onto IMBHs in low-metallicity environments. Regan et al. (2020) showed that $2/3$ of low-metallicity ACHs ($Z/Z_\odot \geq 10^{-3}$) have a high accretion rate ($\geq 0.1 M_\odot$ yr$^{-1}$) by analyzing the Renaissance simulation data. Recently, Chon et al. (2021) reported the metallicity criterion for DCBH–SMBH growth as $Z/Z_\odot \sim 10^{-5}–10^{-3}$. In this way, there is a growing chance for DCBHs to become SMBHs in low-metallicity environments.

4.3. Nonideal MHD Effects

We calculate the ideal MHD equations in this study. If nonideal MHD effects become efficient and significantly

\[ B_0 [\text{G}] = 10^{-12} \]

\[ Z (Z_\odot) = 0 \]

\[ 10^{-8} \]

\[ 10^{-10} \]

\[ 10^{-12} \]

\[ \log_{10} (n [\text{cm}^3]) \]

\[ \log_{10} (m [M_\odot]) \]

\[ \log_{10} (t_{100} \text{yr}) \]

\[ \text{Figure 9. Phase diagrams of the absolute magnetic field strength for all models at } t_{100} = 100 \text{ yr after formation of the first protostar.} \]
dissipate the magnetic field within the disk, then rotation-driven amplification would not be efficient. Higuchi et al. (2018) showed that nonideal MHD effects are ineffective in metal-free star-forming clouds. We can assume that nonideal MHD effects are also ineffective in metal-free ACHs because the gas temperature is always higher in metal-free ACHs than in typical metal-free clouds (with $\sim 10^3 M_\odot$) at any stage of star formation with different gas densities. On the other hand, no one has yet examined the effects of nonideal MHD on metal-enriched ACHs. This is for future work.

In metal-enriched ACHs with $Z/Z_\odot = 10^{-5}$ and $10^{-4}$, we assume that nonideal MHD effects are negligible because the gas temperature, which has a minimum at about 200 and 100 K (Figure 2), respectively, is hot enough. The more metal-enriched ACHs with $Z/Z_\odot = 10^{-3}$ cool to below 100 K (Figure 1 of Chon & Omukai 2020). In such a low-temperature cloud, nonideal MHD effects could work. In addition, the mass growth rate in an unmagnetized ACH with $Z/Z_\odot = 10^{-3}$ is about an order of magnitude smaller than that in one with $Z/Z_\odot < 10^{-4}$ (Chon & Omukai 2020). This parameter ($Z/Z_\odot = 10^{-3}$) does not seem to lead to the formation of IMBHs.

4.4. Numerical Resolution

As described in Section 2.1, we resolve the Jeans wavelength with 32 cells to generate a new finer grid. However, we relaxed the condition to resolve the Jeans wavelength after generating the highest level grid. In the finest grid, the Jeans length was not sufficiently resolved in a high-density region. We need about ten times higher spatial resolution to resolve the Jeans length. Thus, we may overestimate/underestimate the number of fragments. Note that when a massive core or protostar exists, we may need to use other criteria to resolve fragmentation (see Equations (13)-(18) of Matsumoto & Hanawa 2003).

Most of the large or massive fragments in the simulation are reasonably resolved, with about 10 cells inside the large fragments. However, smaller fragments are not well resolved and may be transient in the simulation. High-resolution simulations resolving small fragments or clumps are extremely demanding computationally.
We expect that the number of protostars and amplification timescale should be changed with high-resolution simulations since small fragments may be resolved and therefore survive in the simulation. However, we also expect that the amplification of the magnetic field shown in this study is not qualitatively changed as long as the orbital and spin motions of protostars are resolved. The spatial resolution should be improved and we can address this fully in future studies.

5. Conclusion

We confirm the rotation-driven amplification mechanism of the magnetic field during the accretion phase of DCBH formation in metal-enriched ACHs. Although increasing the metallicity reduces the mass accretion, the collapsing central region is gravitationally unstable and fragments to form many protostars. The protostars drive magnetic field amplification to promote subsequent gas accretion and coalescence of stars. The effect of magnetic amplification is efficient independent of the initial magnetic field strength in the dynamically weak regime of initial magnetic field we study. The amplification mechanism of the magnetic field works even in metal-enriched ACHs with a primordial field strength $B_0/G = 10^{-12}$ at $n_H = 10^4$ cm$^{-3}$. The mass growth rate is maximal for extremely metal-poor ACHs with $Z/Z_\odot = 10^{-5}$, which is a qualitatively different trend to that found in the unmagnetized simulations. We conclude that the exponential magnetic field amplification is realized in metal-enriched ACHs and relaxes the DCBH formation criterion.

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