MAGNETOHYDRODYNAMIC SIMULATIONS OF A ROTATING MASSIVE STAR COLLAPSING TO A BLACK HOLE

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ABSTRACT

We perform two-dimensional, axisymmetric, magnetohydrodynamical simulations of the collapse of a rotating star of 40 $M_\odot$, in light of the collapsar model of gamma-ray bursts. Considering two distributions of angular momentum, up to $\sim 10^{15}$ cm$^2$ s$^{-1}$, and the uniform vertical magnetic field, we investigate the formation of an accretion disk around a black hole and the jet production near the black hole. After material reaches the black hole with high angular momentum, the disk forms inside a surface of weak shock. The disk reaches a quasi-steady state for stars whose magnetic field is less than $10^{10}$ G before the collapse. We find that the jet can be driven by the magnetic fields even if the central core does not rotate as rapidly as previously assumed as long as the outer layers of the star have sufficiently high angular momentum. The magnetic fields are chiefly amplified inside the disk due to the compression and the wrapping of the field. The fields inside the disk propagate to the polar region along the inner boundary near the black hole through the Alfvén wave and eventually drive the jet. The quasi-steady disk is not an advection-dominated disk but a neutrino cooling–dominated one. Mass accretion rates in the disks are greater than 0.01 $M_\odot$ s$^{-1}$ with large fluctuations. The disk is transparent for neutrinos. The dense part of the disk, which is located near the black hole, emits neutrinos efficiently at a constant rate of $<8 \times 10^{51}$ ergs s$^{-1}$. The neutrino luminosity is much smaller than those from supernovae after the neutrino burst.

Subject headings: accretion, accretion disks — gamma rays: bursts — methods: numerical — MHD — supernovae: general

1. INTRODUCTION

During the collapse of a massive star of greater than 35–40 $M_\odot$, the stellar core is considered to promptly collapse to a black hole (Woosley & Weaver 1995; Heger et al. 2003). Stellar material of greater than several solar masses may fall into the black hole at extremely high accretion rates ($>1 M_\odot$ s$^{-1}$; Woosley & Weaver 1995; Heger et al. 2003). If the star has sufficiently high angular momentum before the collapse, an accretion disk is likely to form around the black hole (MacFadyen & Woosley 1999). Jets are suggested to be launched from the inner region of the accretion disk near the black hole through magnetic and/or neutrino processes. Gamma-ray bursts (GRBs) are expected to be driven by the jets. This scenario of GRBs is called the collapsar model (Woosley 1993). Supported by the accumulating observations implying an association between GRBs and the deaths of massive stars (e.g., Galama et al. 1998; Hjorth et al. 2003; Zeh et al. 2004), the collapsar model seems most promising.

In light of the collapsar scenario, MacFadyen & Woosley (1999) and MacFadyen et al. (2001) have performed two-dimensional hydrodynamic simulations of rotating massive stars during collapse after the formation of a black hole. Taking account of the viscous heating with the $\alpha$-prescription (Shakura & Sunyaev 1973) and neutrino cooling, they showed that an accretion disk forms around the black hole if the angular momentum of the progenitor has the right distribution. The disk cools chiefly via advection, and the structure of the disk is shown to be well described by a steady, one-dimensional model of an advection-dominated accretion flow (ADAF; Popham et al. 1999). However, jets cannot be produced in the hydrodynamic simulations of MacFadyen & Woosley (1999). They suggested that collimated outflows can be produced through magnetohydrodynamic (MHD) and/or neutrino processes (MacFadyen & Woosley 1999; MacFadyen et al. 2001). Jet propagation through the stellar envelope has been investigated via hydrodynamic simulations (Aloy et al. 2000, 2003; Zhang et al. 2003). The jets are shown to be ultrarelativistic and collimated in the envelope, although the jets are assumed to be initiated from the central region of the collapsing star, where the thermal energy is deposited at a rate of about $10^{51}$ ergs s$^{-1}$.

In order to investigate jet production in a collapsar, Proga et al. (2003) have performed MHD simulations of the stellar collapse of a rapidly rotating 25 $M_\odot$ progenitor, whose iron core is assumed to be a 1.7 $M_\odot$ black hole and whose magnetic field is assumed to be radial (monopole-like) and uniform. They considered neutrino cooling processes in an optically thin regime and resistive heating, whose properties are highly uncertain. In their relatively short ($\leq 0.28$ s) simulation of a part of the collapsar (about 10–5000 km), relativistic jets of up to 0.2c are revealed to be launched and collimated magnetically. Mizuno et al. (2004a) have also shown that relativistic jets ($\leq 0.3c$) are formed in general relativistic MHD calculations of a collapsing massive star of 15 $M_\odot$. Faster jets can be produced near a black hole with larger rotational parameters (Mizuno et al. 2004b). It should be emphasized that their simulations were performed with a very small computational domain of 360 km $\times$ 360 km and for a very short duration of $\sim$5 s for a 3 $M_\odot$ black hole. We note that MHD calculations have been performed for a collapsing star with relatively small mass in the previous studies (Proga et al. 2003; Mizuno
et al. 2004a, 2004b), although the formation of a black hole is assumed.

On the other hand, jet production via neutrino annihilation is examined with a one-dimensional disk model (Popham et al. 1999). MacFadyen & Woosley (1999) have estimated the energy deposition rate due to neutrino annihilation to be \( \sim 5 \times 10^{50} \) erg s\(^{-1}\) using the disk model (Popham et al. 1999) and the mass accretion rates obtained from their hydrodynamic simulations. However, the deposition rate highly depends on the structure of the disks. Although the outer parts of a disk related to GRBs are ADAFs, as shown in MacFadyen & Woosley (1999), Popham et al. (1999), and Proga et al. (2003), the disk becomes a neutrino cooling-dominated flow (NDAF) in inner regions where neutrino cooling is more efficient than advective cooling (MacFadyen & Woosley 1999; Popham et al. 1999). The properties of NDAFs are extensively investigated using one-dimensional, height-integrated disk models with detailed microphysics (Kohri & Mineshige 2002; Yokosawa et al. 2004) and two-dimensional simulations (Lee & Ramirez-Ruiz 2006). The disks could be convection-dominated flows (Narayan et al. 2001). Which type of an accretion disk is realized in GRBs could depend on the size of the disk and mass accretion rates through the outer boundary of the disk (Narayan et al. 2001).

In this paper, we perform Newtonian MHD simulations of a collapsing massive star of 40 \( M_\odot \) to investigate the formation of an accretion disk and the production of jets from the disk in a collapsar. We consider magnetized stars with both rapidly and slowly rotating cores, which are assumed to collapse to a black hole promptly. Our simulations cover from the iron core to an oxygen-rich layer of the collapsing star, or 50–10,000 km, and are performed for a long term up to \( \sim 4 \) s, which is much longer than the previous studies (Proga et al. 2003; Mizuno et al. 2004a, 2004b), to examine the long-term evolution of the collapsing star, in particular, properties of the disk and jets during the late phase of the collapse. GRBs related to the final stage of a massive star have a long duration \( >2 \) s. Therefore, such long calculations are important for understanding the relation between long GRBs and the disk/jet system formed in the collapsing star.

In \S~2 we present basic MHD equations of the collapsing star, a numerical code, and the initial conditions of the star. Model parameters and numerical results are shown in \S~3. We discuss our numerical grid to check the resolution of the magnetorotational instability (MRI) and compare our results with those obtained in previous works (Proga et al. 2003; Mizuno et al. 2004a) in \S~4. Finally, we summarize our results in \S~5.

2. NUMERICAL METHOD AND INITIAL CONDITIONS

We carry out Newtonian MHD calculations of the collapse of a rotating massive star of 40 \( M_\odot \). In this study, the core of the star is assumed to collapse to a black hole promptly, although the prompt formation of the black hole depends on not only the progenitor mass (or the core mass) but also the equation of state (EOS) and the angular velocity distribution inside the core (e.g., Sekiguchi & Shibata 2004, 2005). Calculations are performed over the region from 50 to 10,000 km of the star. Fluid is freely absorbed through the inner boundary of 50 km, which mimics the surface of the black hole. The black hole mass, \( M_\odot \), is initially set to be that of the central region of the progenitor \(< 50 \) km \((0.001 \) \( M_\odot \)) and is continuously increased by the mass of the infalling gas through the inner boundary at \( r_m \),

\[
\Delta M = \Delta t 4 \pi r_m^2 \int_0^{\pi/2} \rho v_r \sin \theta \, d\theta,
\]
during time step \( \Delta t \), where \( \rho \) and \( v_r \) are the density and radial velocity estimated at the inner boundary, respectively.

General relativistic hydrodynamic simulations show that black holes form at 0.1–0.2 s after the onset of collapse (Sekiguchi & Shibata 2005). In our models, the core is absorbed through the inner boundary and the black hole mass reaches about 2 \( M_\odot \) at the free-fall time, \((3\pi/32G\rho_c)^{1/2}\sim 0.1 \) s, after the onset of the collapse, as we see below. Here \( \rho_c \) is the core density of \( \sim 10^{16} \) g cm\(^{-3}\). Hence, our procedure can approximate the collapsing phase to a black hole at an appropriate level. We note that Proga et al. (2003) have performed an MHD simulation of the period after the formation of a black hole of 1.7 \( M_\odot \).

2.1. Input Physics and Numerical Code

To calculate the structure and evolution of the collapsing star, we solve the Newtonian MHD equations,

\[
\frac{D\rho}{Dt} + \rho \nabla \cdot \mathbf{v} = 0, \quad (2)
\]

\[
\frac{D\mathbf{v}}{Dt} = -\nabla P - \rho \nabla (\Phi + \Psi) + \frac{1}{4\pi} (\nabla \times \mathbf{B}) \times \mathbf{B}, \quad (3)
\]

\[
\frac{D}{Dt} \left( \frac{\rho}{\rho} \right) = -\nabla \cdot \mathbf{v} - L_v, \quad (4)
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}), \quad (5)
\]

\[
\Phi = 4\pi G\rho, \quad (6)
\]

where \( \rho, \mathbf{v}, P, \Phi, \Psi, \mathbf{B}, e, \) and \( L_v \) are the mass density, the fluid velocity, the pressure other than the magnetic pressure, the gravitational potential of the fluid, the gravitational potential of the central object, the magnetic field, the internal energy density, and the neutrino cooling rate, respectively. We denote the Lagrange derivative as \( D/Dt \).

The numerical code for the MHD calculation employed in this paper is based on the ZEUS-2D code (Stone & Norman 1992). We have extended the code to include a realistic EOS (Kotake et al. 2004) based on the relativistic mean field theory (Shen et al. 1998). For the lower density regime \((\rho < 10^5 \) g cm\(^{-3}\)), where no data are available in the EOS table with the Shen EOS, we use another EOS, which includes contributions from an ideal gas of nuclei, radiation, and electrons and positrons with arbitrary degrees of degeneracy (Blinnikov et al. 1996). We carefully connect the two EOSs at \( 10^5 \) g cm\(^{-3}\) so that physical quantities vary continuously in density at a given temperature.

We consider neutrino cooling through electron-positron pair capture on nuclei, electron-positron pair annihilation, and nucleon-nucleon bremsstrahlung. The total neutrino cooling rate is evaluated with a simplified neutrino transfer model based on the two-stream approximation (Di Matteo et al. 2002), with which we can treat the optically thin and thick regimes of neutrino reaction approximately. We ignore resistive heating, whose properties are highly uncertain, in contrast to the method of Proga et al. (2003). We note that the change in the electron fraction, \( Y_e \), is ignored in the MHD calculations (or \( DY_e/Dt = 0 \)).

Spherical coordinates, \((r, \theta, \phi)\), are used in our simulations, and the computational domain is extended over 50 km \( \leq r \leq 10,000 \) km and \( 0 \leq \theta \leq \pi/2 \) and covered with a 200\((r) \times 24(\theta)\) mesh, with which we were able to resolve the fastest growing mode of MRI for most models (see details in \S~4.1). The location of the inner boundary is at a greater radius than that in Proga et al. (2003; \( \sim 10 \) km). We discuss how the location of the inner
boundary affects our numerical results below (§ 4.2). We fix the ratios of the interval sizes of the mesh as $\Delta r_{k+1}/\Delta r_k = 1.02$ ($\Delta r \geq 1 \times 10^6$ cm) and $\Delta \theta_{k+1}/\Delta \theta_k = 1$. We assume that the fluid is axisymmetric and has mirror symmetry on the equatorial plane. We mimic strong gravity around the black hole in terms of the pseudo-Newtonian potential (Paczyński & Wiita 1980):

$$\Psi = -\frac{GM}{r - r_g},$$

where $G$ is the gravitational constant and $r_g = 2GM/c^2$ is the Schwarzschild radius.

2.2. Initial Conditions

We set the initial profiles of the density, temperature, internal energy density, and electron fraction to those of the spherical model of a 40 $M_\odot$ massive star before collapse (Hashimoto 1995). The radial and azimuthal velocities are set to be zero initially and increase due to the collapse induced by the central black hole and self-gravity of the star. The computational domain is extended from the iron core to an inner oxygen layer. A star of about 4 $M_\odot$ is enclosed with the computational domain. The boundaries of the silicon layers between the iron core and the oxygen layers are located at about 1800 km (1.88 $M_\odot$) and 3900 km (2.4 $M_\odot$), respectively. We adopt an analytical form of the angular velocity $\Omega(t)$ of the star before the collapse:

$$\Omega(t) = \Omega_0 = \frac{R_0^2}{r^2 + R_0^2},$$

as in the previous study of collapsars (Mizuno et al. 2004a, 2004b) and supernovae (SNe; Kotake et al. 2004; Yamada & Sawai 2004). Here $\Omega_0$ and $R_0$ are parameters of our model. We consider two sets of $(\Omega_0, R_0)$: $(10$ rad s$^{-1}, 1000$ km) (case with rapidly rotating core) and $(0.5$ rad s$^{-1}, 5000$ km) (case with slowly rotating core). For these sets of $\Omega_0$ and $R_0$, the maximum specific angular momentum is about $10^{17}$ cm$^2$ s$^{-1}$, which is comparable to that of the Keplerian motion around a black hole of 3 $M_\odot$ at 50 km. Therefore, the centrifugal force can be larger than the gravitational force of the central black hole and the formation of a disklke structure is expected near the black hole.

Figure 1 shows the specific angular momentum distribution on the equatorial plane for two cases and for the previous works (MacFadyen & Woosley 1999; Proga et al. 2003). The angular momenta for the two cases are comparable near the outer boundary of the computational domain. We find that the specific angular momentum adopted in MacFadyen & Woosley (1999) is similar to that for the model with $\Omega_0 = 10$ rad s$^{-1}$ and $R_0 = 1000$ km.

The initial magnetic field is assumed to be uniform, parallel to the rotational axis, and weak elsewhere ($\beta = 8\pi P/B^2 \gg 1$). We consider cases with the initial magnetic field $B_0 = 10^8, 10^{10},$ and $10^{12}$ G. It should be noted that the magnetic pressure is much smaller than the other pressure even if $B_0$ is equal to $10^{12}$ G. The models are labeled R8, R10, R12, S8, S10, and S12, in which the letter, R (rapidly rotating core) or S (slowly rotating core), indicates the set of $\Omega_0$ and $R_0$, and the numeral, 8, 10, or 12, equals log ($B_0/\mu G$).

3. NUMERICAL RESULTS

We summarize model parameters and features in Table 1. Columns (2)–(4) give model parameters, with which initial conditions are represented for each model. Column (5) gives the time at the end of each run, $t = t_f$, after the onset of collapse, at which $t$ is equal to 0. Column (10) gives the mass of the black hole at $t = t_f$. Simulations for models with $B_0 = 10^8, 10^{10},$ and $10^{12}$ G are performed for $3 \times 10^5, 2 \times 10^5,$ and $1 \times 10^5$ time steps, respectively. Columns (6) and (7) [(8) and (9)] give ratios of the rotational (magnetic) to the potential energy integrated over the computational domain at $t = 0$ and $t_f$, respectively.

As we see below, jets are produced near the black hole for R10, R12, S10, and S12. The properties of the jets are summarized in Table 2. Column (2) gives $t_{\text{jet}}$ when the jets have passed the

![Figure 1](image-url)
region the density drops below $10^7$ g cm$^{-3}$ finally drives a jet, whose expansion velocity reaches 0.1 km s$^{-1}$.

The disk is in a quasi–steady state and supported mainly by the centrifugal force, as well as the gas pressure, which becomes comparable to the magnetic pressure ($r > 5$) near the central remnant along the rotational axis. In the polar region, the radial and toroidal components of the magnetic field are amplified exponentially in time and the toroidal component is rapidly enhanced. We note that the MRI of the vertical magnetic field produces the $X$-component of the magnetic field and the wrapping of the poloidal field generates the toroidal field. We find that the amplification of the radial and azimuthal components of the field is mainly due to compression of the field initially and MRI in the later phase and that the toroidal field can be amplified by wrapping of the seed poloidal field. In fact, the radial and azimuthal fields are amplified as $t \sim 0.1$ s, as shown in Figure 4, and the growth timescale of MRI, $\tau_{MRI}$, is $\sim 0.1$ s. On the other hand, the characteristic timescale for the field wrapping, $\tau_{\text{wrap}}$, is calculated as

$$\tau_{\text{wrap}} = 4\pi \left\{ \frac{\partial B_i}{\partial t} \right\}^{-1} = 4\pi \left\{ \frac{B_X}{B_\phi} \frac{\partial \Omega}{\partial X} + \frac{B_Z}{B_\phi} \frac{\partial \Omega}{\partial Z} \right\}^{-1} \quad (9)$$

(Takiwaki et al. 2004). Here $B_X$ and $B_Z$ are the $X$ and $Z$ components of the magnetic field, respectively. The toroidal magnetic field is therefore rapidly enhanced through the field wrapping if the poloidal field dominates over the toroidal field and the gas is differentially rotating. This is the case in R10 at $t \sim 0.1$ s. We note that $\tau_{\text{wrap}}$ is shorter than $\tau_{MRI}$ during the rapid amplification phase of the toroidal field ($t \sim 0.1$ s) as in a magnetized core-collapse SN (Takiwaki et al. 2004). As the toroidal field increases, $\tau_{\text{wrap}}$ becomes longer (eq. [9]). After the toroidal field becomes comparable to the poloidal field ($t > 0.1$ s), the growth timescale of the toroidal magnetic field is comparable to that of the poloidal field. Note that $\tau_{MRI}$ $\sim 0.1$ s becomes comparable to or shorter than $\tau_{\text{wrap}}$ in this phase. Because our two-dimensional axisymmetric simulations cannot follow the amplification of the toroidal magnetic field due to MRI, the toroidal field may be saturated at a lower level. At the end of the simulation, the total magnetic, kinetic, and rotational energies are $0.57 \times 10^{51}$, $1.52 \times 10^{51}$, and $1.50 \times 10^{51}$ ergs, respectively.

For higher initial magnetic field of the progenitor, the magnetic field is more rapidly amplified than R10. Figure 5 shows contours of density (left) and of the ratio of $P_{\text{mag}}$ to the pressure (right) for R12 at $t = 0.2378$ and 0.3560 s. A jet launches by the magnetic pressure (top right) just after the formation of a disk (top left) near the central remnant along the rotational axis. In the polar region, the radial and toroidal components of the magnetic field are comparable and reach $10^{15}$ G just before the launch of the jet. It should be emphasized that the density of the jet in R12

\begin{table}[h]
\centering
\caption{Jet Properties}
\begin{tabular}{ccccccc}
\hline
Model & $t_{\text{ft}}$ & $M_0$ & $(E_{\text{mag}})_0$ & $(E_{\text{k}})_0$ & $(E_{\text{kin}})_0$ & $(E_{\text{rot}})_0$ \\
(1) & (s) & (M$_\odot$) & (10^{50} ergs) & (10^{50} ergs) & (10^{50} ergs) & (10^{50} ergs) \\
\hline
R10 & 2.58 & 0.0010 & 2.89E-4 & 0.0274 & 0.0840 & \\
R12 & 0.20 & 0.883 & 0.997 & 4.58 & 2.79 & \\
S10 & 1.30 & 0.0053 & 0.045 & 0.138 & 0.606 & \\
S12 & 0.25 & 0.033 & 0.014 & 3.57 & 5.27 & \\
\hline
\end{tabular}
\end{table}

Notes.—Col. (2): jets have passed through 1000 km at time $t_{\text{ft}}$ from the central black hole. Col. (3): ejected mass via jets. Cols. (4), (5), and (6): magnetic, kinetic, and internal energies of the jets, respectively.
Fig. 2.—Contours of (left) the density and (right) the ratio of $P_{\text{mag}}$ to the pressure for R10 at $t = 1.6625$, 2.0193, and 2.6225 s (from top to bottom).
is much higher than that of R10. Figure 6 shows various time-scales in jets. Solid, dashed, and dotted lines denote the ejection time, the neutrino cooling time through pair capture, and that via pair annihilation, respectively. We find that the ejection time is much smaller than the neutrino cooling time and that jets cannot be cooled through neutrino processes.

For lower initial magnetic field, the hydrodynamic features are similar to R10. For R8, a disklike structure forms inside a shock surface as in R10. The disk is in a quasi-steady state and is supported by the centrifugal force and the pressure gradient. The radial velocity is lower than the rotational velocity, which has a Keplerian profile, by about 3 orders of magnitude. The magnetic pressure is lower than the gas pressure elsewhere, and the toroidal component of the field attains $10^{14}$ G.

3.1.2. Disk Properties

To examine the properties of the disk, we present physical quantities near the equatorial plane ($\theta = 88.1^\circ$) in Figure 7, for R10 at 1.6625 s, at which the weak shock propagates near 2000 km (Fig. 2, top left). The radial velocity is lower than the near-Keplerian rotational velocity by about 2 orders of magnitude. The disk is in equilibrium and mainly supported by the centrifugal...
force, which is greater than the pressure gradient near the black hole. Inside the shock surface, the flow is highly convective (Fig. 2, bottom left), which is clearly seen in Figure 8. Here contours of the ratio of $P_{\text{mag}}$ to the pressure are shown at an inner region (500 km × 500 km) for R10 at $t = 1.6625$ s (Fig. 2, top right). We find that inside the disk, the velocity field is highly complicated, so the magnetic field also has a rather complex structure.

Figure 9 shows contours of the logarithmic specific neutrino cooling rate through pair capture on nuclei, for R10 at 1.6625 s, defined as

$$q_{\text{cap}} = 9.2 \times 10^{33} \rho X_{\text{nuc}} \left( \frac{T}{10^{11} \text{ K}} \right)^6 \text{ergs s}^{-1} \text{cm}^{-3},$$

(10)

where $X_{\text{nuc}}$ is the mass fraction of nucleons and is approximated by

$$X_{\text{nuc}} = 8.2 \times 10^9 \frac{T_{\text{MeV}}^{9/8}}{\rho^{3/4}} \exp \left( -\frac{7.074}{T_{\text{MeV}}} \right),$$

(11)

or unity, whichever is smaller with the temperature in MeV, $T_{\text{MeV}}$ (Woosley & Baron 1992). Density contours of $10^8, 10^9, 10^{10}$, and $10^{11}$ g cm$^{-3}$ are plotted in Figure 9. We find that neutrino cooling is efficient inside the shock, where the temperature is higher than $10^{10}$ K via shock heating, and very efficient near the equatorial plane of the disk with high densities because of more efficient pair capture in denser regions. The neutrino luminosity stays at a constant level of $\sim 5 \times 10^{51}$ ergs s$^{-1}$, as we see below (Fig. 17). We find that neutrino cooling via $e^\pm$ capture dominates over those through electron-positron pair annihilation and nucleon-nucleon bremsstrahlung by 1 or 2 orders of magnitude inside the disk (Fig. 10, bottom).

At the inner region of the disks, the neutrino cooling time via $e^\pm$ capture, $\tau_{\text{cap}}$, which is $0.1–1$ s, is comparable to or smaller than the accretion time, $|r/v_r|$ (Fig. 10, bottom). Thus, the disks cool not by advection but by neutrino emission in the region. Indeed, we find that the density and temperature profiles are well described with those of not ADAF but NDAF, or $\rho \sim r^{-51/26}$ and $T \sim r^{-3/10}$ (Fig. 10, top). For ADAF, the profiles have different power laws, or $\rho \sim r^{-3/2}$ and $T \sim r^{-5/8}$ (e.g., Fujimoto et al. 2001). We note that the gas pressure is dominant over the radiation and degenerate pressure in our models other than in the outer part of the disks, where the radiation pressure is comparable to the gas pressure (Fig. 7, bottom right). The magnetic pressure is lower than the other pressures other than in the region near the rotational axis where the toroidal component of the field is amplified to $10^{15}$ G (Fig. 7, top right).

In order to examine whether the disk is opaque to neutrinos or not, we calculate the height $h_\nu$ at which the neutrino optical depth $\tau_\nu(R)$ is equal to $1/2$,

$$\tau_\nu(R) = \int_{h_\nu(R)}^{h_{\text{max}}(R)} \rho(R, z) \kappa_\nu(R, z) dz = \frac{2}{3},$$

(12)

as in Surman & McLaughlin (2004). Here $h_{\text{max}}$ and $\kappa_\nu$ are the height of the computational domain from the equatorial plane and the neutrino opacity at the distance from the rotational axis, $R$. We note that the optical depth is simply set to be the sum of the depths for $e^\pm, \mu^\pm$, and $\tau$ neutrinos. The neutrino decoupling surface is found to be torus-like and very limited to the inner region ($\lesssim 100$ km) of the disk near the black hole if it exists. For $t \geq 0.3$ s, the surface is located at the region below 30 km from the equatorial plane near the black hole. We find that almost the entire disk is therefore optically thin to neutrinos and that neutrinos are hardly trapped inside the disk.

Figure 11 shows the time evolution of the accretion rate through the inner boundary at 50 km, defined as

$$\dot{M} = 4\pi r^2 \int_{0}^{\pi/2} \rho v_r \sin \theta \, d\theta,$$

(13)

where $\rho$ and $v_r$ are estimated at the boundary. The rates through parts of the boundary, $M_{\text{pol}}(\theta \leq 20^\circ)$ and $M_{\text{disk}}(\theta \geq 50^\circ)$, are also presented in Figure 11. At the beginning of the collapse, material with low angular momentum falls through the inner boundary at very high accretion rates ($\geq 10 M_\odot$ s$^{-1}$) for the free-fall time of the core ($\sim 0.1$ s). As gas accretes onto the black hole with high angular momentum ($t \geq 0.1$ s), the polar infall dominates over the disk accretion: $M_{\text{disk}}$ is much lower than $\dot{M}$, which is comparable to $M_{\text{pol}}$ by 1 order of magnitude. During quasi-steady collapse ($t \geq 1.0$ s), $\dot{M}$ is $\sim 0.2–0.5 M_\odot$ s$^{-1}$ with fast fluctuations and gradual decrease, while $M_{\text{disk}}$ varies between 0.02 and 0.07 $M_\odot$ s$^{-1}$ in this phase and sharply decreases after the ejection of the jet.

3.2. Collapse of a Star with a Slowly Rotating Core

Now we move on to the models with slowly rotating cores, S8, S10, and S12. As we can see in Figure 1, the initial angular momentum for these models is smaller than $10^{17}$ cm$^2$ s$^{-1}$ except in the outer part of the computational domain. Most of the inner region of the star therefore collapses to a black hole. Material with high angular momentum is expected to form a disklike structure.

Figure 12 shows the contours of density (left) and of the ratio of $P_{\text{mag}}$ to the pressure (right) for S10 at $t = 1.0207$ and 1.3134 s. The core for S10 rotates more slowly than that for R10. The star collapses with a nearly spherical configuration for a longer time. When the rapidly rotating gas falls near the black hole ($t \approx 0.45$ s), the magnetic field has been amplified due to field compression and wrapping initially and MRI later. Eventually, a magnetically driven jet is found to be generated from a central region near the black hole (Fig. 12, bottom), although the core rotates slowly. We find that the jet is less collimated than that in R10 (Figs. 2 and 12, bottom left).

A disklike structure forms as in R10, but magnetically driven winds arise near the equatorial plane (Fig. 12, top left), which disappear in R10. Figure 13 shows physical quantities near the
equatorial plane ($\theta = 88.1^\circ$) for S10 at $t = 0.5676$ s. When gas is infalling to the remnant with high angular momentum, the magnetic field has been amplified so much that the magnetic pressure is comparable to the gas pressure near the central remnant. Consequently, the winds are driven via magnetic and gas pressure for S10, in contrast to R10. These magnetically driven winds have high velocity ($\sim 0.1c$; Fig. 13, bottom left) and relatively high entropy of $s \sim 20$ (Fig. 13, top left), where $s$ is the entropy per baryon in units of the Boltzmann constant. Distributions of the density, temperature, and pressure cannot be described with the use of simple power laws in $r$, in contrast to R10, due to the winds.

As in cases with a rapidly rotating core, the toroidal magnetic field is more rapidly amplified in models with a larger initial magnetic field. We find that S12 also launches a jet from a central region near the black hole along the rotational axis. Figure 14 shows contours of density (left) and of the ratio of $P_{\text{mag}}$ to the pressure (right) for S12 at $t = 0.2812$ s. The jet is collimated in a wider region than that in R12 (Fig. 14, bottom) because of slow core rotation in S12. It is noted that the jet has a density lower than that in R12.

For S8, as in R8 and R10, a disklike structure forms inside a shock surface located at 800 km on the equator at the end of the computation, and the disk is well described by a quasi-steady
NDAF rotating with a Keplerian profile. The magnetic pressure, whose toroidal component attains $10^{15}$ G, is lower than the gas pressure elsewhere. Magnetocentrifugal winds do not appear in S8, in contrast to S10, and a jet cannot be produced in S8, as in R8.

In brief, the collapsing star can originate well-collimated jets if the star has a sufficiently large magnetic field and rapidly rotating core. On the other hand, when the star has not a rapidly rotating core but an envelope with sufficiently high angular momentum, the star can drive less collimated jets. The jet could be generated from collapsing stars with a nonrotating core and a sufficiently high angular momentum envelope and large magnetic field.

### 3.3. Time Evolution of Collapsing Star

We present the time evolution of various quantities, such as the magnetic energy, the mass accretion rate, and the neutrino luminosity of our models to compare the properties of each model.

Figure 15 shows the time evolution of the magnetic energy integrated over the computational domain, $E_m = \int \frac{B^2}{8\pi} dV$, for all models. We note that the toroidal field dominates over the other components of the field, as in R10 (Fig. 4). The fields are amplified earlier as the initial magnetic field increases and the core rotation becomes faster. This is because the magnetic field is mainly amplified by the wrapping of the field. Rapid differential rotation is therefore required from the amplification.

Figure 16 shows the time evolution of $\dot{M}$ for rapidly rotating models (R models; top) and slowly rotating models (S models; bottom). At the beginning of the simulations ($t < 0.1$ s), $\dot{M}$ are roughly comparable for all models because the collapse is almost spherical and the effects of rotation and magnetic field can be neglected in this phase. After the infall of higher angular momentum material, $\dot{M}$ rapidly drops due to the centrifugal barrier. This is because $\dot{M}$-values in R models are smaller than those of S models during the initial phase ($t \leq 0.5$ s). Once a disklike structure forms (for R8, R10, S8, and S10), $\dot{M}$ stays 0.2–0.3 $M_\odot$ s$^{-1}$ with fast fluctuations. We note that $\dot{M}_{\text{disk}}$ is much lower than $\dot{M}$ by about 1 order of magnitude for all models, as in R10 (Fig. 11).

**Fig. 6.**—Timescale inside jets for R12 at 0.1997 s showing the ejection timescale (solid line) and the neutrino cooling timescale through pair capture (dashed line) and through pair annihilation (dotted line).

**Fig. 7.**—Physical quantities near the equatorial plane ($\theta = 88.1'$) for R10 at 1.6625 s. Top left: Logarithmic density in grams per cubic centimeter (solid line), logarithmic temperature in kelvins (dashed line), and entropy per baryon in the Boltzmann constant (dotted line). Top right: Logarithmic magnetic fields. Solid, dashed, and dotted lines represent the radial, azimuthal, and toroidal components of the magnetic field in units of gauss, respectively. Bottom left: Logarithmic velocities in centimeters per second: radial velocity (solid line), angular velocity (dashed line), Alfvén speed (dotted line), sound speed (double-dashed line), and Keplerian velocity around a 2 $M_\odot$ black hole (dot-dashed line). Bottom right: Logarithmic pressure in ergs per cubic centimeter.
Figure 8.—Contours of the ratio of $P_{\text{mag}}$ to the pressure at an inner region (500 km × 500 km) of an accretion disk for R10 at $t = 1.6625$ s.

Figure 17 shows the time evolution of the neutrino flux integrated over the computational domain, or the neutrino luminosity, $L_{\nu}$, for all models. Neutrinos are initially emitted from a quasi-spherically collapsing dense core near the black hole at a maximum rate of $\sim 10^{52}$ erg s$^{-1}$ around $\sim 0.1$ s. For R8 and R10, material with high angular momentum reaches the black hole before the infall of the entire dense core to the black hole. Neutrinos continue to be emitted from a disk formed near the equator ($\sim 500$ km) of an accretion disk for R10 at $t = 1.6625$ s. Density contours of $10^{8}$, $10^{9}$, $10^{10}$, and $10^{11}$ g cm$^{-3}$ are also shown.

However, the core is stable against MRI in this phase because $\Omega$ is nearly constant (see the initial distribution, eq. (8)).

After the formation of a disklike structure, the disk becomes unstable to MRI. The growth timescale of MRI, $\tau_{\text{MRI}}$, is comparable to that of the wrapping of the field $\tau_{\text{wrap}}$ ($\sim 0.1$ s). However, the relation (14) holds at a later time ($\sim 0.2$ s for R10) and our numerical grid can resolve, albeit marginally, the fastest growing mode of MRI, by which the field is amplified mainly by the field wrapping. As the field increases, $v_A$ becomes greater. We find that the relation (14) holds at a later time ($\sim 0.2$ s for R10) and our numerical grid can resolve, albeit marginally, the fastest growing mode of MRI, by which the field is amplified to a saturated value. For R8, however, the relation (14) cannot hold, so the magnetic energy is possibly saturated at a lower amplitude. We need finer numerical grids to resolve MRI for R8.

### 4.2. The Location of the Inner Boundary and the Numerical Grids

In this study, the inner boundary of the computational domain is set to be 50 km, which is larger than that in Proga et al. (2003; $\sim 10$ km). In order to examine the dependence of disk properties on the location of the inner boundary, $r_{\text{in}}$, we have performed a MHD simulation of R10 with a smaller inner boundary of 10 km and with 220 finer radial mesh intervals until $t \sim 0.8$ s.

Figure 18 shows radial profiles of the density (solid lines) and temperature (dashed lines) near the equatorial plane ($\theta = 88.1^\circ$) in R10 at 0.5 s for $r_{\text{in}} = 10$ and 50 km. Thick and thin lines correspond to the profiles for $r_{\text{in}} = 10$ and 50 km, respectively. We also plot profiles of the density (dotted line) and temperature (double-dashed lines) of NDAF, as shown in Figure 10. We find that the profiles are independent of the location of the inner boundary. Radial profiles of other physical quantities inside the disk are also independent of the location. We also find that the disk is neutrino cooling dominated even if we adopt the smaller inner boundary of 10 km. The neutrino luminosity slightly increases to $1.20 \times 10^{52}$ from $0.81 \times 10^{52}$ erg s$^{-1}$ for $r_{\text{in}} = 50$ km.
The luminosity has a trend of gradual decrease after the formation of the disklike structure for $r_{\text{in}} = 50$ km (Fig. 17). For $r_{\text{in}} = 10$ km, the luminosity decreases to $< 1 \times 10^{52}$ ergs s$^{-1}$ after $t = 0.55$ s and stays roughly constant at the rate of $(7-8) \times 10^{51}$ ergs s$^{-1}$ for $t = 0.6-0.8$ s. Furthermore, the disk is found to be optically thin to neutrinos for $t > 0.5$ s, as in the case with $r_{\text{in}} = 50$ km.

We impose mirror symmetry on the equatorial plane, which is not assumed in Proga et al. (2003). If we calculate without the symmetry, the convective motion near the equatorial plane may be weakened and jets may become less collimated. We need calculations in a computational domain with a grid to examine such a possibility. Moreover, in light of the collapsar model, it may be important to investigate the standing accretion shock instability, which does not appear in calculations with mirror symmetry and whose importance has been recognized in SN explosions (Blondin et al. 2003).

4.3. Comparison between Our Models and Previous Works

In MHD simulations of a 25 $M_\odot$ collapsar (Proga et al. 2003) and a 15 $M_\odot$ collapsar (Mizuno et al. 2004a, 2004b), a jet has been shown to be produced magnetically near the black hole,
as in our results. We compare our results and those obtained by other groups (Proga et al. 2003; Mizuno et al. 2004a, 2004b).

We first compare the disk properties of our model R10 (Fig. 7) and those in Proga et al. (2003; their Fig. 3). We note, however, that Figure 7 shows the properties at $t = 1.6625$ s while Figure 3 in Proga et al. (2003) shows time-averaged properties during $t = 0.2629 - 0.2818$ s. In Proga et al. (2003), the $25 M_\odot$ collapsar is simulated with a distribution of angular momentum and magnetic field similar to that in our models after the formation of a black hole of $1.7 M_\odot$, which is the mass of an entire iron core. They assume nonrotating inner Si and O layers and set the initial distributions of velocity and density to be those of one-dimensional free-falling gas. They have taken into account resistive heating and neutrino cooling in an optically thin regime and ignored the self-gravity of the star.

On the other hand, we have simulated the collapse of a $40 M_\odot$ star, whose core is assumed to collapse to a black hole promptly. The initial distributions of physical quantities are set to be those of a presupernova, and collapse to the black hole is mimicked with the absorption of material through the inner boundary. We have ignored resistive heating but considered neutrino cooling with the two-stream approximation (Di Matteo et al. 2002), with which we can treat the optically thin and thick regimes for neutrino reactions approximately.
The radial velocity for R10 is much slower than that in Proga et al. (2003; their Fig. 3, top right). The difference is attributed to the existence of a weak shock, where the radial velocity sharply drops (Fig. 7, bottom left). The cause of the weak shock is possibly the centrifugal barrier as seen in § 3.1.1. Note that a similar accretion shock is also revealed in hydrodynamic calculations of the collapse of a $35M_{\odot}$ star (MacFadyen & Woosley 1999). The disappearance of the weak shock in Proga et al. (2003) is probably a result of their initial distributions of the inner layers mentioned above. We note that the radial velocity inside the disk is much smaller than that in Proga et al. (2003), even if the calculation is performed for $r_{\text{in}} = 10$ km.

The slower radial velocity leads to a longer accretion time. Thus, the cooling time via neutrino processes is shorter than the
accretion time in our study (Fig. 10), in contrast with Proga et al. (2003). Hence, the disk is neutrino cooling–dominated in our models while advection-dominated in their model except for a small region inside the torus where the density reaches maximum. We note that the disk structure obtained from the MHD calculation in our study is different from that through the hydrodynamic calculation with α-viscosity (MacFadyen & Woosley 1999), in which the structure is roughly coincident with that of ADAF.

The accretion rate through the inner boundary is estimated to be about $5 \times 10^{12} \, g \, s^{-1} (=0.25 \, M_\odot \, s^{-1})$ with fast fluctuations (Proga et al. 2003), which is comparable to the rates in R8, R10, S8, and S10 after the formation of the quasi-steady disk (see Fig. 11). The radial profiles of the magnetic fields are similar to those of Proga et al. (2003), but the magnitudes are higher. The magnetic energy in Proga et al. (2003) is smaller than that in our models by up to 1 order of magnitude at the end of the calculation, in spite of greater initial magnetic energy; $E_{m,r}$ is about $2 \times 10^{47}$ ergs s$^{-1}$ (in their Fig. 1). A jet is magnetically driven from the inner region of the disk in our models and in those of Proga et al. (2003), and the jet velocity is comparable.

Moreover, we compare disk properties of our model R10 for $r_{in} = 10 \, km$ at $t = 0.5 \, s$ (Fig. 18) and those in Proga et al. (2003; their Fig. 3). We find that the radial profiles of the density and temperature in our calculation are similar to those in Proga et al. (2003). The density has a maximum of $10^{12} \, g \, cm^{-3}$ around

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**Fig. 15.**—Time evolution of the magnetic energy integrated over the computational domain, $E_m = \int (B^2/8\pi) \, dV$, for all models. The thick solid, thin solid, thick dashed, thin dashed, thick dotted, and thin dotted lines represent the magnetic energy for R8, S8, R10, S10, R12, and S12, respectively.

**Fig. 16.**—Time evolution of the accretion rate through the inner boundary at 50 km, $\dot{M}$, for all models. The red, green, and blue lines represent the rates for R8, R10, and R12 (top), respectively, and the rates for S8, S10, and S12 (bottom), respectively.

**Fig. 17.**—Time evolution of the neutrino luminosity integrated over the computational domain, for all models. Line types are as in Fig. 15.

**Fig. 18.**—Radial profiles of the density (solid lines) and temperature (dashed lines) near the equatorial plane ($\theta = 88.1^\circ$) in R10 at 0.5 s for $r_{in} = 10$ and 50 km. Thick and thin lines correspond to the profiles for $r_{in} = 10$ and 50 km, respectively. We also plot profiles of the density (dotted line) and temperature (double-dashed lines) of ADAF, as shown in Fig. 10. The profiles are independent of the location of the inner boundary.
~30–40 km. The radial velocity rapidly increases near the inner boundary because of strong gravity of the black hole. The density in our run is greater than that in Proga et al. (2003) due to smaller radial velocity, while the temperature is lower. In Proga et al. (2003), the accretion time is shorter than that in our run due to greater radial velocity. The liberated energy is therefore advected into the black hole with the accreting gas, which cools via neutrino processes inefficiently and becomes hotter than that in our model. The resistive heating, which is considered in their model and not in our models, could also heat the accreting gas to higher temperature. Consequently, the neutrino luminosity, which is very sensitive to the temperature (eq. [10]), is greater than that in our run.

Next we compare our results with those of scale-free, general relativistic MHD simulations (Mizuno et al. 2004a). For a $3 M_\odot$ black hole and a density scale of $10^{10}$ g cm$^{-3}$, Mizuno et al. (2004a) have simulated the collapse of a magnetized, rotating massive star after the formation of the black hole for a very short duration ~5 ms using a simplified (gamma-law) EOS and a very small computational domain of 360 km $\times$ 360 km. The initial profiles of the angular velocity and the magnetic field are similar to those in our models, but their model A2 corresponds to $\Omega_0 = 378$ s$^{-1}$ and $B_0 = 1.5 \times 10^{13}$ G, which is much faster rotation and greater initial field than R12. A magnetically driven jet is shown to emit near the black hole and propagate along the rotational axis with a speed of up to 0.2c at 5 ms, much earlier than in R12, because of faster rotation and greater initial magnetic field.

Our simulations are performed with a Newtonian MHD code, in which the pseudo-Newtonian potential is used to mimic the general relativistic effects of a black hole. If we examine the situation during the collapsing phase to the black hole or the more inner region near the black hole, we need a general relativistic MHD code as in Mizuno et al. (2004a, 2004b) and Sekiguchi & Shibata (2004, 2005). However, our numerical grid covers the region with $r \geq 50$ km, or 5.56$r_g$ of a 3 $M_\odot$ black hole. General relativistic effects are therefore too small for our models.

4.4. Implications for GRB Progenitors

Taking into account angular momentum transport through magnetic torque (Spruit 2002), Heger et al. (2004, 2005) and Maeder & Meynet (2004) have performed 1.5-dimensional simulations on the evolution of rotating massive stars. They showed that the iron core of the stars has specific angular momentum up to $\sim 10^{15}$ cm$^2$ s$^{-1}$ before core collapse. This is much smaller than the specific angular momentum of material in Keplerian motion around a $2 M_\odot$ black hole at the last stable circular orbit, $2 \times 10^{16}$ cm$^2$ s$^{-1}$. Consequently, single massive stars are claimed not to be a progenitor of GRBs if the magnetic braking mechanism of Spruit (2002) operates in the stars (Petrovich et al. 2005; Fryer & Heger 2005). Moreover, even in a binary companion star (Petrovich et al. 2005), a rapidly rotating core with the specific angular momentum of $2 \times 10^{16}$ cm$^2$ s$^{-1}$ is unlikely to be realized, although such a core can be produced in the merger of two helium cores during the common-envelope inspiral phase of a binary system (Fryer & Heger 2005). It should be emphasized, however, that the angular momentum transport through magnetic torque in a massive star is still uncertain.

Moreover, even for stars with a slowly rotating iron core, an accretion disk around a black hole forms during the collapse of the stars and a jet can be produced from the disk near the black hole if the outer layers of the stars have a sufficiently high angular momentum, as shown in § 3.2. Therefore, progenitors of GRBs may not always require a rapidly rotating iron core; rather, layers with a high angular momentum and a high magnetic field could be sufficient.

5. CONCLUDING REMARKS

We have performed two-dimensional, axisymmetric MHD simulations of the collapse of a $40 M_\odot$, rapidly rotating star, whose core is assumed to collapse to a black hole promptly and whose angular momentum attains $10^{17}$ cm$^2$ s$^{-1}$, to examine the formation of an accretion disk around the black hole and the jet production near the black hole in light of the collapsar model of GRBs. Considering two distributions of the angular velocity and the uniform magnetic field, whose magnitude is $10^8$, $10^{10}$, and $10^{12}$ G, parallel to the rotational axis, inside the star before collapse, we investigate how angular momentum and magnetic field distributions inside the star affect the jet production and the disk properties. We summarize our conclusions as follows:

1. After material reaches the black hole with high angular momentum of about $10^{17}$ cm$^2$ s$^{-1}$, a disk forms inside a surface of weak shock, which appears near the black hole due to the centrifugal force and propagates outward slowly. The disks reach a quasi–steady state for models with initial magnetic field of less than $10^{10}$ G.

2. We find that the jet can be driven by the tangled up magnetic fields even if the central core does not rotate as rapidly as previously assumed as long as the outer layers of the star have sufficiently high angular momentum.

3. The jet is driven by the magnetic field, which is dominated by the toroidal component and is amplified due to the wrapping of the field, as long as the initial magnetic field is greater than $10^8$ G. The fields are chiefly amplified inside the disks and propagate to the polar region along the inner boundary near the black hole through the Alfvén wave. The jets cannot be cooled through neutrino processes.

4. The quasi-steady disk is not an advection-dominated accretion flow but a neutrino cooling–dominated disk. The accretion time is larger than the cooling time via neutrino processes. At an inner region (<100 km) of the disk, the profiles of density and temperature are similar to those in Proga et al. (2003).

5. The radial profiles of the density and the temperature of the quasi-steady disk are well described by those of a neutrino cooling–dominated disk, or $\rho \sim r^{-5/2}$ and $T \sim r^{-3/10}$, in which the gas pressure is dominant over the other pressure. These profiles are rather different from those of ADAF, or $\rho \sim r^{-3/2}$ and $T \sim r^{-5/8}$.

6. Mass accretion rates in the quasi-steady disks are greater than 0.01 $M_\odot$ s$^{-1}$ with large fluctuations. A small fraction of the rest-mass energy is liberated through neutrinos emitted from the dense part of the quasi-steady disks, which hardly trap neutrinos. The neutrino luminosity stays at a constant level of $<8 \times 10^{51}$ ergs s$^{-1}$, which is much smaller than those from a SN and in previous work (Proga et al. 2003).

The mass and total energy of the jets are 0.0018–0.037 $M_\odot$ and $5 \times 10^{48}$ to $5 \times 10^{50}$ ergs, respectively (Table 2). The jets are too heavy and weak to produce a relativistic fireball, or GRB, whose baryon mass and isotopic energy are required to be $\leq 10^{-4}$ to $10^{-5}$ $M_\odot$ and $10^{53}$ ergs, respectively (e.g., Hurley et al. 2006). The jets are baryon-rich and cannot be accelerated to a relativistic velocity, or failed GRB, as in the previous works (Proga et al. 2003; Mizuno et al. 2004a, 2004b). However, after the generation of the jets, the polar region near the black hole becomes baryon-poor and has a large magnetic field. If we continue to perform simulations, a baryon-poor outflow may be produced from the region.
We need much longer simulations and thus a much wider computational domain to examine the ejection of these multiple jets. Although no jet appears in R8 and S8 during the computation, a jet could originate near the black hole if the computation is performed for a longer time. However, at the end of the computation, the density sharply drops near the outer boundary because material inside the computational domain begins to be consumed to infall to the black hole, whose mass attains 2.64 \( M_\odot \) for R8. In order to perform the calculation for a longer time to investigate whether a jet can be produced or not, we need a larger computational domain. Our numerical grid cannot resolve the fastest growing mode of MRI for R8 and S8. Therefore, we need a longer calculation with a larger computational domain and finer numerical grids to examine the jet generation in R8 and S8. This is our future task.

The density and temperature of jets can become so high that material in the jets is in nuclear statistical equilibrium. For relatively low density jets, the composition is protons and neutrons near the black hole. As jets propagate along the rotational axis to decrease the density and temperature, \( \alpha \)-rich freezeout operates inside the jets, which are mainly composed of \(^4\)He and \(^{56}\)Ni. On the other hand, for dense jets, such as that in R12, the jets can be neutron-rich due to electron capture on protons, as suggested by one-dimensional calculations for a collapsar disk and jets (Fujimoto et al. 2004, 2005). Consequently, the r-process can operate in such dense jets and neutron-rich heavy nuclei, and up to third-peak elements can be produced inside the jets (Fujimoto et al. 2006).

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