THE PROPAGATION OF NEUTRINO-DRIVEN JETS IN WOLF–RAYET STARS

HIROKI NAGAKURA
Yukawa Institute for Theoretical Physics, Kyoto University, Kyoto 606-8502, Japan; hiroki@heap.phys.waseda.ac.jp
AND
Advanced Research Institute for Science & Engineering, Waseda University, 3-4-1 Okubo, Shinjuku, Tokyo 169-8555, Japan

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ABSTRACT

We numerically investigate the jet propagation through a rotating collapsing Wolf–Rayet star with detailed central engine physics constructed based on the neutrino-driven collapsar model. The collapsing star determines the evolution of the mass accretion rate, black hole mass, and spin, all of which are important ingredients for determining the jet luminosity. We reveal that neutrino-driven jets in rapidly spinning Wolf–Rayet stars are capable of breaking out from the stellar envelope, while those propagating in slower rotating progenitors fail to break out due to insufficient kinetic power. For progenitor models with successful jet breakouts, the kinetic energy accumulated in the cocoon could be as large as $\sim 10^{51}$ erg and might significantly contribute to the luminosity of the afterglow emission or to the kinetic energy of the accompanying supernova if nickel production takes place. We further analyze the post-breakout phase using a simple analytical prescription and conclude that the relativistic jet component could produce events with an isotropic luminosity $L_{\nu,\text{iso}} \sim 10^{52}$ erg s$^{-1}$ and isotropic energy $E_{\nu,\text{iso}} \sim 10^{54}$ erg. Our findings support the idea of rapidly rotating Wolf–Rayet stars as plausible progenitors of GRBs, while slowly rotational ones could be responsible for low-luminosity or failed GRBs.

Key words: black hole physics – hydrodynamics

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1. INTRODUCTION

Long-duration gamma-ray bursts (GRBs) are thought to originate from the death of massive stars (Woosley & Bloom 2006). It is widely recognized that the study of GRBs provides important knowledge on the final evolutionary stage in the life of massive stars. Although the nature of GRBs remains elusive, one viable scenario to produce a GRB is the neutrino-driven collapsar model (Woosley 1993; MacFadyen & Woosley 1999). The gravitational collapse of the rapidly rotating core is believed to create a fast rotational Kerr black hole. The copious amounts of neutrinos and their antiparticles, which are emitted from the hot accretion disk, annihilate and create an electron–positron pair, and then produce an electron–positron pair, and then produce an electromagnetic cascade of secondary particles, including gamma rays. These gamma rays are emitted from the black hole and are detected as GRBs.

Several studies have been done (see, e.g., Aloy et al. 2000; Zhang et al. 2003; Kohri & Mineshige 2002; Di Matteo et al. 2002; Kohri et al. 2005; Gu et al. 2006; Chen & Beloborodov 2007; Liu et al. 2010; Zalamea & Beloborodov 2011; Liu et al. 2012) to understand the jet energetics in time-dependent simulations. In these studies, however, the jet luminosity was assumed to be constant or to directly follow the mass accretion rate (see, e.g., MacFadyen et al. 2001; Nagakura et al. 2012). In order to judge whether neutrino-driven jets are capable of successfully breaking out of their progenitors, and to explore the effects of rotation, one needs to take into account how the jet power scales with the neutrino energy deposition generated by the accompanying accretion.

In this paper, we present, for the first time, the propagation of neutrino-driven jets employing an accurate neutrino energy deposition rate as calculated by Zalamea & Beloborodov (2011). The evolution of mass accretion rate and black hole mass and spin, all of which are necessary to evaluate the energy deposition by neutrinos, are evaluated here using an inner boundary condition in the post-breakout phase using a simple analytical prescription and conclude that the relativistic jet component could produce events with an isotropic luminosity $L_{\nu,\text{iso}} \sim 10^{52}$ erg s$^{-1}$ and isotropic energy $E_{\nu,\text{iso}} \sim 10^{54}$ erg. Our findings support the idea of rapidly rotating Wolf–Rayet stars as plausible progenitors of GRBs, while slowly rotational ones could be responsible for low-luminosity or failed GRBs.
simulation, although the current calculation is still not self-consistent as it fails to resolve the accretion disk and assumes no feedback (e.g., a disk wind). The purpose of this study is (1) to clarify whether neutrino-driven jets can successfully break out from the stellar surface and (2) to determine the progenitor’s rotation rate necessary for successful jet breakout. As we will show in this paper, the final outcome of the explosion depends sensitively on the rate of rotation, and differences in rotation rate could be responsible for the observational differences seen between GRBs, low-luminosity GRBs (LLGRBs), and failed GRBs.

2. METHODS AND MODELS

We perform two-dimensional, relativistic hydrodynamical axisymmetric (and also equatorial symmetric) simulations of the accretion and subsequent jet propagation. The numerical code employed in this paper is essentially the same as those used in previous papers (Nagakura et al. 2011, 2012). The initial stellar density distribution is fixed as in model 16TI in Woosley & Heger (2006). As is the case with previous studies, we cut the inner portions of the star from a certain radius. The self-gravity of matter in the active numerical regions is calculated by solving Poisson equations, and the monopole gravity is added as the point mass at the inner excised region. The mass accretion rate ($\dot{M}$) is estimated by the mass flows through the inner boundary (see Equation (1) in Nagakura et al. 2012). The mass and angular momentum in the excised region are assumed to be the same mass and angular momentum of the black hole. The time evolution of mass and spin of the black hole is calculated by integrating mass and angular momentum flux crossing the inner boundary. It should be noted that when the specific angular momentum (SAM) at the location of the inner boundary in the equatorial region becomes larger than SAM at the innermost stable circular orbit (ISCO; see cross marks in Figure 1), we alter our prescription when calculating the angular momentum to

$$f_a = M \times J_{\text{ISCO}},$$

where $f_a$ and $J_{\text{ISCO}}$ denote the angular momentum flux and SAM at the ISCO, respectively. This treatment comes from the fact that the angular momentum of matter in the disk is transported outward due to the turbulent viscosity or non-axisymmetric waves, and finally the matter falls into a black hole with $\sim J_{\text{ISCO}}$.

According to Zalamea & Beloborodov (2011), the jet luminosity by the neutrino process is determined by

$$L_j = 1.1 \times 10^{52} x_{\text{ms}}^{-4.8} \left( \frac{M_{\text{bh}}}{3 M_\odot} \right)^{-2} \times \begin{cases} 0 & (M < M_{\text{ign}}) \\ \dot{m}^{\frac{3}{4}} & (M_{\text{ign}} < M < M_{\text{trap}}) \\ \dot{m}_{\text{trap}}^{\frac{3}{4}} & (M > M_{\text{trap}}), \end{cases}$$

(2)

where $\dot{m} \equiv \dot{M}/(M_\odot/s)$ and $x_{\text{ms}} \equiv r_{\text{ms}}/(2GM_{\text{bh}}/c^2)$ ($r_{\text{ms}}$ denotes the marginally stable orbit). $G$ and $c$ denote the gravitational constant and the speed of light, respectively. The characteristic mass accretion rate $M_{\text{ign}}$ and $M_{\text{trap}}$ are given as a function of the viscous constant and the speed of light, respectively. The evolution of the jet accretion rate, which is sensitive to the location of the inner boundary (see Nagakura et al. 2012), can thus be better captured by our simulations. However, as a result, these simulations become rather computationally expensive, and we are only able to conduct them until the jet bow shock reaches the stellar surface or the black hole mass reaches $10^{10} M_\odot$. The evolution of the post-breakout phase is then analyzed by using a simple analytic formalism (see Equations (3)–(6)). The results of an extended numerical simulation will be nonetheless compare with the analytical approach for the reference model in order to confirm that the analytical approach qualitatively captures the evolution of the jet dynamics (see Section 3.2 and Figure 4).

We employ the gamma-law equation of state with $\gamma = 4/3$. The jet injection parameters such as the Lorentz factor and the specific internal energy are the same as those used in the standard model of Nagakura et al. (2012), where the initial Lorentz factor and specific internal energy are fixed to $\Gamma = 400$ and $\epsilon = 0.01$, respectively. It should be noted that for a fixed $\Gamma$ and $\epsilon$, the overall jet dynamics depend solely on $\theta_{\text{sp}}$ (see Nagakura et al. 2012). In this study, we assume $\theta_{\text{sp}} = 9^\circ$, which agrees well with the opening angles deduced for long GRBs (Goldstein et al. 2011). The dependence of our results on the $\theta_{\text{sp}}$ will be discussed in Section 3.

The 1000 non-uniform radial grids cover the entire computational region, while the meridian section is covered by 60 uniform grids. The three-level adaptive mesh refinement technique, similar to that used in Nagakura et al. (2011, 2012), is also employed in order to decrease computational cost. We set up the stellar rotation in a similar manner as used in López-Cáceres et al. (2010). The SAM distribution is separated into radial and polar components as $J(r, \theta) = j(r)\Theta(\theta)$, where $r$ and $\theta$ are the spherical radius and the polar angle, respectively. In the reference model (Mref), $j(r)$ is given by model 16TI. For models M150, M70, and M50, $j(r)$ is multiplied by 1.5, 0.7, and 0.5, respectively (see Figure 1 for the SAM distribution of...
Figure 2. Density contour (log scale) in the meridian section for each model at $t = 17.6$ s (left), $t = 27.8$ s (middle), and $t = 46.5$ s (right). Each time corresponds to the time of jet breakout for M150, Mref, and M70, respectively (see also $t_{br}$ in Table 1). The spatial size of each box is $-5 \times 10^9$ cm $< x < 5 \times 10^9$ cm and $0 < z < 1 \times 10^{10}$ cm. The upper and lower panels correspond to models M150, Mref, M70, and M50. Note that some panels in M150 and Mref are lacking since we stop the simulation at the time of breakout.

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

Table 1

| Model | Breakout | $t_{ij}$ (s) | $t_{br}$ (s) | $L_p$ ($10^{50}$ erg s$^{-1}$) | $E_{i\beta}$ ($10^{51}$ erg) | $E_j$ ($10^{51}$ erg) | $E_{j\rightarrow L_{50}}$ ($10^{51}$ erg) | $E_{j\rightarrow L_{40.5}}$ ($10^{51}$ erg) | $T_{j\rightarrow L_{50}}$ (s) | $T_{j\rightarrow L_{40.5}}$ (s) |
|-------|----------|-------------|-------------|-----------------------------|------------------|--------------|-----------------------------|-----------------------------|----------------|----------------|
| Mref  | Yes      | 10.9        | 27.8        | 1.9                         | 1.4              | 7.4          | 4.4                         | 5.6                         | 27.8           | 46.2           |
| M150  | Yes      | 2.2         | 17.6        | 3.2                         | 1.6              | 11.7         | 8.8                         | 9.8                         | 40.0           | 55.0           |
| M70   | Yes      | 15.5        | 45.5        | 1.0                         | 0.9              | 3.6          | ...                         | ...                         | ...             | ...             |
| M50   | No       | 21.6        | ...         | 0.6                         | ...              | ...          | ...                         | ...                         | ...             | ...             |

Note: Our models assume rigid-body rotation on shells, i.e., $\Theta(\theta) = \sin^2 \theta$.

It is important to highlight that the simulations in this study do not cover the black hole accretion disk system. Even if the simulations cover the full computational domain, our numerical calculations cannot treat the disk evolution appropriately, since the general relativistic effect and microphysics, which are important in determining the disk evolution, are not incorporated. However, the analytical formula proposed by Zalamea & Beloborodov (2011) allows us to estimate neutrino luminosity without resolving the black hole accretion disk system. Owing to these prescriptions, we can study the jet penetration phase as determined by the neutrino-driven energy injection. The study of the coupling between the black hole and disk accretion system is beyond the scope of this paper.

We also note that the injection of the jet is delayed in the simulation as a result of the inner core not possessing enough angular momentum to create a disk (Woosley & Heger 2006; Lee & Ramirez-Ruiz 2006; Zalamea & Beloborodov 2011). Based on the standard neutrino-driven collapsar model (and assumptions used in Chen & Beloborodov 2007; Zalamea & Beloborodov 2011), the central engine starts to operate after the accretion disk is formed around a black hole. Therefore, in our simulations, the jet is injected only when the SAM of matter at the inner boundary in the equatorial plane exceeds $J_{ISCO}$.

3. RESULTS

3.1. Basic Features and Jet Penetrability

The overall evolution of the collapse of the progenitor seen in our simulations is not surprisingly similar to that found in Nagakura et al. (2011, 2012). The infall of the stellar envelope generates a rarefaction wave, which propagates outward. The envelope contraction is almost identical among all models and follows a rather spherical contraction, since the centrifugal force plays a minor role. Note that we find that the density distribution at the inner boundary is slightly oblate, but this does not affect the subsequent jet propagation, although it might have consequences for the jet production (which is not properly simulated here).

During the jet propagation phase, on the other hand, the results of jet evolution are very different among each model (see Table 1
out and almost stagnates around the inner boundary despite the onset of the central engine. We also show the summary of our results in Figure 3. (A color version of this figure is available in the online journal.)

\[ \eta = \frac{\dot{M}}{L_j} \]

\[ \dot{M} = \frac{L_j}{\beta c^2} \]

\[ L_j = \eta \dot{M} \]

The most important result in this study is that the jet succeeds to break out from the star except for M50, which corresponds to the model with the slowest rotation rate. The neutrino-driven jet with \( \theta_{\text{op}} = 9^\circ \) produced in a rapidly rotating compact Wolf–Rayet star can potentially give rise to GRBs. We also find that the time-averaged accretion-to-jet conversion efficiency among successful jet breakout models is roughly \( \eta \sim 10^{-3} \), while \( \eta \) for M50 cannot reach \( 10^{-3} \) and never accomplishes the jet breakout. This result is roughly consistent with our previous work (Nagakura et al. 2012). The analytical criteria in Nagakura et al. (2012) also show the opening angle dependence for the successful jet breakout, which is the threshold \( \eta \) increases to \( (\theta_{\text{op}})^2 \). Therefore, according to this result, we would like to give an important caution that jets with wider opening angle have more difficulty penetrating the star than in the present results, which indicates that the threshold progenitor rotation rate differs in accordance with the jet opening angle.

### 3.2. Post-breakout Phase: Analytical Formula

As we have already mentioned, our numerical simulations are terminated at the time of jet breakout. However, it is interesting to extend the result of our numerical simulations to post-breakout stage, which allows us to estimate the expected observational differences among the computed models (see Section 3.3). We employ the following analytic approximations in order to see the subsequent evolution after the jet breakout:

\[ t = t_{(b)} + \int_{r_{(b)}}^r \frac{dM_b}{d\tilde{r}}(\tilde{r}) d\tilde{r} \]

\[ M_{bh} = M_{bh(b)} + \int_{r_{(b)}}^r \frac{dM_b}{d\tilde{r}}(\tilde{r}) d\tilde{r} \]

where

\[ t_{(b)} = \beta \sqrt{\frac{GM_f}{p^3}} \]

\[ \eta \approx \frac{L_j}{\dot{M} c^2} \] as a function of time from the onset of the collapse. The chief cause for the different jet propagation behavior is the sensitive dependence of the neutrino luminosity on the Kerr parameter (see Equation (2)). For the fast rotational model, the angular momentum of the black hole is very large and increases with time (see the fourth panel in Figure 3), which produces a powerful jet as a result of the large neutrino deposition energy rates. It is also important to note that the onset timing of the central engine \( (t_{i}) \) also greatly affects the outcome of the explosion. As shown in Table 1 and illustrated in Figure 2, the jet production is significantly delayed for a slower rotational model. This is due to the neutrino luminosity being weaker for both smaller accretion rates and larger black hole masses (see Equation (2) for the dependence of \( \dot{m} \) and \( M_{bh} \)). In fact, for M50, although the Kerr parameter reaches \( \gtrsim 0.9 \) at the end of our simulation, the neutrino luminosity is not large enough to move out of the forward shock wave.

According to these results, we infer that neutrino-driven jets may not penetrate progenitors with extended envelopes, since significant large mass might be able to accrete to the black hole before jet breakout. The weak neutrino luminosity resulting for a sizable increase in black hole mass could result in the jet becoming non-relativistic ejecta. Therefore, a compact progenitor is an inevitable requirement for a successful neutrino-driven jet breakout. For the massive envelope progenitors such as PoP III or red (blue) supergiants, a different process might be required to powered GRBs (see also discussions in Suwa & Ioka 2011).

Figure 3 shows the evolution of hemispherical neutrino luminosity, mass accretion rate, black hole mass, spin parameter, and conversion efficiency from accretion energy to neutrino luminosity \( (\eta \equiv L_j / \dot{M} c^2) \) as a function of time from the onset and Figure 2). We also show the summary of our results in Table 1. For model M50, the forward shock wave does not move out and almost stagnates around the inner boundary despite the successful operation of the central engine \( (t_i \text{ in Table 1 denotes the time of initiation of the central engine}) \). In fact, no collimated feature can be seen for M50 in the lowest panels in Figure 2. This is attributed to the fact that the jet power does not exceed the ram pressure of the inflowing material, and the forward shock wave stagnates or is advected inward. For models with successful jet breakout, on the other hand, the jet also cannot move forward quickly after the initiation of the central engine. However, due to an increase in the Kerr parameter of the black hole over time, the jet power eventually exceeds the ram pressure of the inflowing material (Figure 3). Once the forward shock wave is able to move out, the jet interacts with the stellar mantle and gives rise to a cocoon. The hot cocoon helps jet confinement and helps to preserve the jet’s strong outgoing momentum and energy flux, eventually leading to a successful breakout.

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\[ M_{bh} = M_{bh(b)} + \int_{r_{(b)}}^r \frac{dM_b}{d\tilde{r}}(\tilde{r}) d\tilde{r} \]

where

\[ t_{(b)} = \beta \sqrt{\frac{GM_f}{p^3}} \]
\[ M_{\Pi}(r) = \frac{dM_\Pi}{dr} \]
\[ = \frac{1}{\beta^2} \frac{8\pi G M r t_{\Pi} \rho}{3M_* - 4\pi r^3 \rho}. \]

The above analytical estimation is essentially similar to the approach presented in Nagakura et al. (2012) and Suwa & Ioka (2011). The free-fall time \( t_\Pi(r) \) can be determined based on the assumption of spherical symmetric envelope contraction. Here, \( t_\Pi \) denotes the time at jet breakout. Note also that the functions of \( M_\Pi(r) \) and \( \rho(r) \) are extracted from the table of model 16T1. \( r_b \) is determined by the assumption of \( M_{\Pi b}(t_b) = M_\Pi(r_b) \). The non-dimensional parameter \( \beta \) is determined to ensure that \( M(t_b) \) is equal to \( M_{\Pi}(r_b) \). According to this procedure, the neutrino luminosity and mass accretion rate can be smoothly connected from the results of numerical simulations. The evolution of the Kerr parameter is determined by Equation (1). Note that, since model M50 does not succeed in jet breakout, there is no post-breakout phase for this model. Results of these analytical extensions are also described in Figure 3.

Figure 4 shows the comparison between the results of extended numerical simulations and analytical estimation for the reference model. The extended numerical simulations are performed until \( t = 40 \) s, which corresponds to about 10 s after the jet breakout. Note that we do not broaden the computational region, since the stellar contraction is not affected by the jet dynamics in the outer parts of the star. As shown in this figure, the time evolution of black hole mass (\( M_{\rm bh} \)), Kerr parameter \( (a) \), and conversion efficiency \( (\eta) \) are almost identical between results of the extended numerical simulation and the analytical calculation. For luminosity \( (L_j) \) and mass accretion rate \( (\dot{M}) \), on the other hand, the analytical calculations are slightly larger than the results of our numerical simulations. This may be attributed to the fact that the analytical approach neglects stellar rotation, which increases the mass accretion rate and consequently overestimates the neutrino luminosity. However, these differences are within 10%. Therefore, we confirm that the above analytical approach qualitatively well describes the time evolution of the dynamics of the jet. In the following subsection, we discuss the observational consequence with the aid of the analytical approach in the post-breakout phase.

3.3. Observational Consequences

We first divide the energetics of the jet into two parts, which are the relativistic jet component \( (E_j) \) and the cocoon component \( (E_{\rm co}) \) (see Suwa & Ioka (2011); Matzner (2003)). The \( E_j \) is calculated based on the assumption that all the injected energy after the jet break out goes into the relativistic component, i.e., \( E_j \) is given by \( \int_{t_b}^\infty (L_j/2) dt \). Note that 1/2 factor comes from the assumption of equatorial symmetry.

As shown in Table 1, \( E_j \) increases with an increase of stellar rotation. For the purpose of studying the outcome of explosion in more detail, we further divide the energy of the relativistic jet into \( E_{j>\nu L_{50}} \) and \( E_{j>L_{50}} \). \( E_{j>\nu L_{50}} \) is calculated by the same manner as \( E_j \), except the integration is carried out when \( L_j/2 > 10^{50} \) erg s\(^{-1} \), while \( E_{j>L_{50}} \) is calculated with a condition \( L_j/2 > 5 \times 10^{50} \) erg s\(^{-1} \). In addition, we also list the corresponding time duration of each component as \( T_{j>\nu L_{50}} \) and \( T_{j>L_{50}} \) in Table 1. In all cases, these timescales are several tens of seconds, which are comparable with the typical timescale of the prompt phase of GRBs. It should be noted, however, that, according to Bromberg et al. (2012), the duration of the prompt phase of GRBs may be modified by the duration of jet penetration, i.e., \( t_p = t_e - t_b \) (see also Equation (2) in Bromberg et al. 2012), where \( t_p \), \( t_e \), and \( t_b \) denote observed duration of the prompt phase, the central engine working time, and the duration of the jet penetration phase, respectively. Therefore, the actual observed duration will be smaller than the \( T_j \), and it is substantially modified especially for slower rotation models (e.g., M70, since \( t_b \) is larger with slower rotation). It is also interesting to note that model M70, which is the slowest model among successful jet breakout models, produces the weak explosion and does not have \( E_{j>L_{50}} \) due to the low luminosity of the jet. Note also that this luminosity is the upper limit for observed luminosity since we neglect the conversion process from hydrodynamical energy to gamma rays. According to these facts, we infer that the neutrino-driven jet from the compact Wolf–Rayet star, whose rotation rate is between models M70 and M50, may produce very low luminous type bursts, possibly LLGRBs. These results are qualitatively consistent with the previous studies (see, e.g., López-Cámara et al. 2009; Lindner et al. 2010; Milosavljević et al. 2012).
we find that our analysis is meaningful as it defines the neutrino luminosity observed light curve in reality. It should be noted, however, that $E_j$ is typically $t_{\text{prompt}}$ burst ($\sim 10$ s). This may be attributed to the fact that the jet with a slowly rotating core tends to be weaker and takes a longer time to penetrate the star. Therefore, despite its low jet luminosity, a large fraction of jet energy has been consumed for sweeping aside the stellar mantle and accumulated as the cocoon energy, which eventually reaches $\sim 10^{51}$ erg. The cocoon material is expected to contribute the subsequent explosive event after the prompt phase (Tominaga et al. 2007; Lazzati et al. 2012) and at the afterglow phase (Ramirez-Ruiz et al. 2002). The results of neutrino deposition presented in this paper are not able to discern whether sufficient large nickel production might take place to explain hypernova explosions (see also discussions of cocoon propagation in Mészáros & Rees 2001; Ramirez-Ruiz et al. 2002; Matzner 2003). If nickel is not effectively produced at the jet interaction region alternative pathways such as the disk wind by viscous heating or magnetic-driven wind from the central engine would be required to explain the link between the GRBs and hypernovae.

We further calculate the isotropic energy ($E_{j(\text{iso})}$) for $E_j$, $E_{j>L_{\text{iso}}}$, and $E_{j>L_{\text{iso}}}$, and also isotropic peak luminosity $L_{p(\text{iso})}$, which are shown in Figure 5. In this calculation, the jet opening angle is assumed to be $\theta_{\text{jet}}=9^\circ$, which is the same as the root of the injected jet in our simulations. Note again that we neglect the conversion efficiency from hydrodynamic energy to radiation, so our results are still at the qualitative level and give only the observational consequences with the aide of analytic extrapolation in the post-breakout phase. We show that every model except for M50 succeeds in breaking out of the star. Especially, Mref and M150, which correspond to models with sufficient rapidly rotation, have the relativistic outflow component as $L_{p(\text{iso})} \sim 10^{52}$ erg s$^{-1}$ and $E_{j(\text{iso})} \sim 10^{52}$ erg, which are sufficiently large to explain GRBs. On the other hand, the energy in the cocoon component $E_{d}$ is $\sim 10^{51}$ erg for models with successful jet breakouts, although it remains an open question whether the jet or the cocoon expansion could give rise to enough nickel production to explain the GRB-hypernova connection. One of the other important results in this study is that model M50, which corresponds to the model with slow rotation, cannot succeed in jet breakout (failed GRB). Therefore, there is the threshold SAM distribution between models M70 and M50 for the success of jet penetration. It should be noted, however, that the threshold rotation, no doubt, strongly depends on the progenitor rotation.

Finally, we would like to note that the results presented in this paper are optimistic. As one of the large uncertainties, the actual mass accretion rate would be smaller than the results obtained in this paper, since some fraction of the mass is expected to escape from the disk rather than being accreted to the black hole due to neutrino winds or viscous heating (MacFadyen & Woosley 1999). In addition, the increase rate of the Kerr parameter would be slower than the current result, since the SAM of the infall mater is smaller than the ISCO due to the effect of the pressure gradient. Note also that the disk wind also extracts the angular momentum of accretion matter. Since the neutrino deposition rate depends sensitively on the mass accretion rate and Kerr parameter, the jet dynamics would be affected by these effects. Note also that the neutrino-driven jet cannot explain the extremely long duration of bursts; other populations are necessary to explain these peculiar events. The other factors such as viewing angle may also cause the observational difference among the GRB population (see, e.g., Yamazaki et al. 2002, 2003; Granot et al. 2005). More quantitative discussions will be conducted in our forthcoming paper.

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