GRavitational wave signatures of hyperaccreting collapsar disks

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ABSTRACT

By performing two-dimensional special relativistic (SR) magnetohydrodynamic simulations, we study possible signatures of gravitational waves (GWs) in the context of the collapsar model for long-duration gamma-ray bursts. In our SR simulations, the central black hole is treated as an absorbing boundary. By doing so, we focus on the GWs generated by asphericities in neutrino emission and matter motions in the vicinity of the hyperaccreting disks. We compute nine models with adding initial angular momenta and magnetic fields parametrically to a precollapse core of a 35 M⊙ progenitor star. As for the microphysics, a realistic equation of state is employed and the neutrino cooling is taken into account via a multi-flavor neutrino leakage scheme. To accurately estimate GWs produced by anisotropic neutrino emission, we perform a ray-tracing analysis in general relativity by a post-processing procedure. By employing a stress formula that includes contributions from both magnetic fields and SR corrections, we also study the effects of magnetic fields on the gravitational waveforms. We find that the GW amplitudes from anisotropic neutrino emission show a monotonic increase with time, whose amplitudes are much larger than those from matter motions of the accreting material. We show that the increasing trend of the neutrino GWs stems from the excess of neutrino emission in the direction near parallel to the spin axis illuminated from the hyperaccreting disks. We point out that a recently proposed future space-based interferometer like Fabry–Perot-type DECIGO would permit the detection of these GW signals within ≈100 Mpc.

Key words: gravitational waves – neutrinos – radiative transfer – relativistic processes – supernovae: general

Online-only material: color figures

1. INTRODUCTION

Gamma-ray bursts (GRBs) are one of the most energetic phenomena in the universe. Thanks to Swift observations,4 it has now become evident that GRBs are basically categorized into two, namely, short-hard and long-soft bursts (e.g., Hjorth & Bloom 2011; Nakar 2007 for recent reviews). More surprisingly, GRBs with some mixed features of the two types have been reported (e.g., Gehrels et al. 2006; Gal-Yam et al. 2006), possibly necessitating a new classification (Lü et al. 2010). The mystery of their central engines seems to be thickening, which has long puzzled astrophysicists since their accidental discovery in the late 1960s (see Meszaros 2006 for review). Regarding the long-duration GRBs (LGRBs), robust associations of the underlying supernovae with a handful of LGRBs (e.g., Galama et al. 1998; Hjorth et al. 2003; Stanek et al. 2003; Malesani et al. 2004; Modjaz et al. 2006; Pian et al. 2006, and collective references in Woosley & Bloom 2006; Zhang 2011) and the fact that their host galaxies are typically irregular with intense star formation (Fruhcht et al. 2006) suggest that they are likely related to the deaths of massive stars, and the “collapsar” model has been widely recognized as the standard scenario for LGRBs (Woosley 1993; Paczynski 1998; MacFadyen & Woosley 1999).

In this scenario, the collapsed iron core of a massive star forms a temporary disk around a few M⊙ black hole (BH), whose gravitational binding energy is the driving source of the central engine, and accretes at a high rate (∼0.1–10 M⊙ s−1, e.g., Popham et al. 1999; Di Matteo et al. 2002; Kohri et al. 2005; Chen & Beloborodov 2007; Zalamea & Beloborodov 2011; and references therein). Paczynski (1990) and Meszaros & Rees (1992) pioneeringly proposed that pairs of neutrino and anti-neutrino illuminated from the hyperaccreting disks that annihilate into electron and positron (e.g., ν + ¯ν → e− + e+, hereafter “neutrino pair annihilation”) can supply sufficient energy to launch GRB outflows by heating material in the polar funnel regions. In addition, it is suggested that the strong magnetic fields in the cores of order 1015 G also play an active role both for driving the magneto-driven jets and for extracting a significant amount of energy from the central engine (e.g., Blandford & Znajek 1977; Mizuno et al. 2004a, 2004b; McKinney 2006; McKinney & Narayan 2007a, 2007b; Komissarov & Barkov 2007, 2009; Komissarov & McKinney 2007; Barkov & Komissarov 2008; Nagataki 2009, and references therein).

Although various possibilities including magnetar models (e.g., Dai & Lu 1998; Thompson et al. 2004; Uzdensky & MacFadyen 2007; Bucciantini et al. 2007) have been proposed so far, there has been no direct evidence to pin down the mechanism of the central engine. This is mainly because it is difficult to extract the information from conventional astronomy by electromagnetic waves, since high-energy photons are absorbed through interactions in the source and by the photon backgrounds. Alternatively, gravitational waves (GWs) are expected to be a primary observable to decipher the mechanism of the engine, because they imprint live information hidden deep inside the stellar core and they carry the information directly to us without being affected when propagating to Earth.

Currently, long-baseline laser interferometers such as LIGO (Abbott et al. 2005), VIRGO,5 GEO600,6 and TAMA300 (Ando

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4 http://www.swift.psu.edu/
5 http://www.ego-gw.it/
6 http://geo600.aei.mpg.de/
& the TAMA Collaboration 2005) are operational (see, e.g., Hough et al. 2005 for a recent review). For these detectors, core-collapse supernovae (CCSNe) have been proposed as one of the most plausible GW sources, therefore an extensive study of the GW predictions based on sophisticated numerical modeling has been carried out (see, for example, Ott 2009; Fryer & New 2011; Kotake 2011 for recent reviews). It is noted however that most of them have paid attention to CCSNe that leave behind neutron stars (NSs) after explosions. For a reliable prediction of GWs from CCSNe, one needs to perform multidimensional hydrodynamic simulations equipped with a precise neutrino transport scheme which follows the dynamics starting from stellar core-collapse, core-bounce, through shock-stall and subsequent growth of hydrodynamic instabilities, the neutrino-driven shock revival, to stellar explosion in a consistent manner. This is one of the most challenging subjects in computational astrophysics (e.g., Janka et al. 2007; Kotake et al. 2012a).

However, the numerical modeling to test the collapsar scenario could be much more demanding. One needs to trace a new path that bifurcates from the above story after bounce, namely, to the BH formation (phase 1), evolution of the surrounding accretion disk including energy deposition to the polar funnel region by neutrinos and/or magnetic fields (phase 2), to the launching of fireballs (phase 3). This apparently necessitates multidimensional magnetohydrodynamic (MHD) simulations not only with general relativity (GR) to handle BH formation, but also with multi-angle neutrino transfer for treating highly anisotropic neutrino radiation from the disks (e.g., Kotake et al. 2012a). In the business of CCSN simulations, the most up-to-date simulations8 can now follow multi-angle neutrino transport but are limited to a Newtonian case (Hubeny & Burrows 2007; Ott et al. 2008; Sumiyoshi & Yamada 2012), or handle GR with a sophisticated neutrino transport (Müller et al. 2010) but not applicable to a rapidly rotating case (due to the assumption of conformal flatness).

Various approximate approaches have therefore been undertaken in the business of collapsar simulations. In phase 1, GR simulations (Shibata et al. 2006) updated to implement neutrino cooling have been reported recently (Sekiguchi & Shibata 2011), in which the dynamics after the BH formation to the formation of the accretion disk was first consistently followed. The numerical studies of phase 2 are concerned with the subsequent evolution of the accretion disk and the outflow formation in the polar funnel region until the jets become mildly relativistic. The central BH has been traditionally treated by a fixed metric technique in GR simulations (e.g., Mizuno et al. 2004a; De Villiers et al. 2005; Hawley & Krolik 2006; McKinney & Narayan 2007b; Komissarov & McKinney 2007; Barkov & Komissarov 2008) or by an absorbing boundary in Newtonian simulations (e.g., MacFadyen & Woosley 1999; Proga et al. 2003; Fujimoto et al. 2006; Nagataki et al. 2007; López-Cámara et al. 2009, 2010) or simple relativistic (SR) simulations (e.g., Harikae et al. 2009). As for the microphysics in these simulations, except for Fujimoto et al. (2006), Nagataki et al. (2007), Harikae et al. (2009), and Shibata et al. (2007), a realistic nuclear equation of state (EOS) has been replaced by a very phenomenological one (like a gamma law or polytrope) and neutrino cooling (and heating) has often been neglected for simplicity. Numerical studies phase 3 are mainly concerned with the dynamics later on, namely, from the jet propagation to the breakout from the star, by assuming a manual energy input to the polar funnel region (see, e.g., Aloy et al. 2000; Zhang et al. 2003; Lazatti & Begelman 2009; Nagakura et al. 2011a, 2011b, and references therein). All of the studies mentioned above may be regarded as complementary in the sense that different epochs are focused on, with different initial conditions for the numerical modeling being undertaken.

To the best of our knowledge, Ott et al. (2011) were the only ones who extracted the GW signals based on their collapsar simulations (in phase 1). Based on their three-dimensional GR simulations of a 75 $M_\odot$ star using a polytropic EOS, they pointed out that the significant GW emission is associated with the moment of BH formation, which can be a promising target of the advanced LIGO for a Galactic source. In contrast to such a paucity of GW predictions based on numerical simulations of collapsars, a number of semianalytical estimates have been reported so far, which predict a significantly strong GW emission due to possible density inhomogeneities (Mineshige et al. 2002), bar or fragmentation instabilities in the collapsar’s accretion torii (e.g., van Putten 2001; Davies et al. 2002; Fryer et al. 2002; Kobayashi & Mészáros 2003; Piro & Pfahl 2007; Corsi & Mészáros 2009, and collective references in Fryer & New 2011), and the precession of the disks due to GR effects (Romero et al. 2010; Sun et al. 2012). The predicted GW amplitudes are typically high enough to be visible to advanced-LIGO class detectors for a 100 Mpc distance scale, which is about four orders of magnitudes larger than the numerical estimate at the black hole formation (Ott et al. 2011). In addition to these GWs produced by non-spherical matter motions, Hiramatsu et al. (2005) and Suwa & Murase (2009) pointed out that anisotropic neutrino emission from the accretion disk could be the source of GWs from collapsars, which was originally proposed as an equally important GW source to the matter GW in the context of CCSNe (Epstein 1978). Since these GWs from collapsars would be a smoking-gun signature of the central engine coinciding with the conventional electromagnetic messengers as well as neutrinos,9 it will be very important to put forward theoretical predictions of GW signals based on the collapsar simulations, as has been done in the matter of CCSNe.

In this work, we study possible GW signatures in the hyperaccreting collapsar disks10 by performing two-dimensional (2D) SR MHD simulations of accretion torii around a BH. We compute nine models by adding angular momenta and magnetic fields parametrically to a precollapse core of a 35 $M_\odot$ progenitor star (Woosley & Heger 2006). As for the microphysics, a realistic EOS is employed and neutrino cooling is taken into account via a multi-flavor neutrino leakage scheme. In our SR simulations, the central BH is treated as an absorbing boundary. By doing so, we focus on the GWs generated by asphericities in neutrino emission and matter motions in the vicinity of the accretion disks. To accurately estimate GWs produced by anisotropic neutrino emission, we perform a ray-tracing analysis in GR using a post-processing procedure. By employing a stress formula that includes contributions from both magnetic fields and SR corrections, we also study the effects of magnetic fields on the gravitational waveforms. Then we discuss their detectability by performing spectrum analysis.

The paper opens up with descriptions of the initial models and numerical methods (Section 2). The main results are given

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7 Here, for convenience we call each stage as phase 1, 2, etc.
8 Assisted by accelerating computer powers.
9 Especially for a nearby GRB source, e.g., Ando et al. (2005).
10 Which corresponds to evolution in phase 2.
in Section 3. We summarize our results and discuss their implications in Section 4.

2. NUMERICAL METHODS AND MODELS

2.1. Initial Models

Regarding our precollapse model, we take a 35 $M_\odot$ star (model 35OC in Woosley & Heger (2006) that is supposed to be one of the most promising GRB progenitor models. To study the effects of rotation systematically, we take the following rotational profiles

$$\Omega(r, \theta) = \frac{\Omega_0 X_0^4 + \alpha \Omega_{\text{iso}}(M(X)) X_0^4}{X_0^4 + X^4}, (1)$$

where $\alpha$, $\Omega_0$, and $X_0$ are model parameters. $M(X)$ is the mass coordinate at the cylindrical radius ($X = r \sin \theta$), and $\Omega_{\text{iso}}$ is given by $\Omega_{\text{iso}} = J_{\text{iso}}/X_0^2$, where $J_{\text{iso}}$ is the specific angular momentum in the last stable orbit of the Schwarzschild BH (e.g., Bardeen et al. 1972; Proga 2005; López-Cámara et al. 2009). By changing $\alpha$ in the range of $0.6 \leq \alpha \leq 1.2$, we compute seven non-magnetized models of $\alpha = 0.6, 0.7, 0.8, 0.9, 1.0, 1.1$, and 1.2, where they are labeled as model J0.6, J0.7, and so on (see Figure 1). To see the effects on the magnetic fields, we compute two more models: one in which the initial magnetic field ($B_0 = 10^{10}$ G or $10^{11}$ G) is added to model J0.8 (model J0.8B10 or model J0.8B11) and another in which the initial field is assumed to be purely poloidal and also assumed to be uniform and parallel to the spin axis.

Figure 1 shows profiles of the initial angular momentum for some models. As shown, our models are taken to possess much more rapid rotation especially in the range of 1–4 $M_\odot$ in the mass coordinate compared to the original profile of model 35OC (black dashed line). We set such a rotational profile, otherwise the accretion disk cannot survive later than $\sim 2$ s after the onset of gravitational collapse in our simulation. Later on, the accretion disk is swallowed by the central object due to the neutrino cooling which removes the pressure support in the disk. In this case, one cannot account for the duration of long bursts, the activity of which is typically longer than $\sim 2$ s, and can last up to tens of minutes. So, we experimented by choosing to adjust the initial angular momentum as in Equation (1).

Although our models do have higher initial angular momentum, the deviation may not be so serious (compare the light blue line that is for the angular momentum amplified by a factor of two in the original progenitor) considering uncertainties in stellar evolution calculations (such as in the treatment of mass loss, weak interactions, and fluid instabilities (e.g., convection, semiconvection, rotation, and magnetic fields)). To mimic the original progenitor structure, we take $X_0$ to be the size of the Fe core ($\approx 3000$ km). By setting $\alpha = 0.8$, $\Omega_0 \approx 1$ rad s$^{-1}$, the above profile becomes closest to the original profile for mass coordinates larger than 5 $M_\odot$. As a side note, the most prevalent way is to tune the initial angular momentum to satisfy its local centrifugal force to be several percent levels of the local gravitational binding energy (e.g., MacFadyen & Woosley 1999). Our modeling may be more realistic in the sense that the assumed initial angular momentum profiles capture the basic trend obtained in the current progenitor models.

2.1.1. Hydrodynamics

As already mentioned in the introduction, a number of Newtonian and SR collapsar simulations in phase 2 have conventionally focused on the evolution of the collapsar disks by treating the central BH as an absorbing boundary. Following the tradition of setting the boundary from the beginning of the simulation, we also initially impose an absorbing inner-boundary condition at the radius of $\max(10 \text{ km}, 2r_g)$ with $r_g$ being the Schwarzschild radius that is estimated from the accumulating mass inside the inner boundary. Then we solve the dynamics outside the inner boundary up to the outer boundary of the computational domain (30,000 km in radius) using our SRMHD code and assuming axisymmetry and equatorial symmetry (see Harikae et al. 2009 for more details). In our 2D simulations, spherical coordinates are employed with logarithmic zoning in the radial direction ($r$) and regular zoning in the polar direction ($0 \leq \theta \leq \pi/2$). The computational domain is covered by 300($r$) 40(\theta) mesh points. Regarding the microphysics, a realistic nuclear EOS by Shen et al. (1998) is included and neutrino cooling is treated by a multi-flavor leakage scheme (e.g., Epstein & Pethick 1981; Rosswog & Liebendörfer 2003; Kotake et al. 2003a; see Takiwaki et al. 2009 for more details). The gravitational potential is estimated from the sum of the Paczynski–Witt type potential which mimics the gravitational pull from the central BH and the self-gravity of material outside the excited region that is determined by the Poisson equation (see Equations (5) and (6) in Harikae et al. 2009). In this paper, we examine numerical models without viscosity such as the alpha prescription (see, however, MacFadyen & Woosley 1999; Lindner et al. 2010; López-Cámara et al. 2010).

2.2. Extraction of Gravitational Waveforms

2.2.1. GWs from Matter Motions

To extract the gravitational waveforms from matter motions and magnetic fields, we employ the stress formulae derived in Takiwaki & Kotake (2011) and Kotake et al. (2004). For convenience, we briefly summarize them in the following.

In our axisymmetric case, the only non-vanishing quadrupole term is the plus mode ($h_+$) in the metric perturbation (e.g., Möhne et al. 1991), which can be written as

$$h(X, r) = \frac{1}{8} \sqrt{\frac{15}{\pi}} \sin^2 \theta A_{30}^{E2}(t - \frac{r}{c}) \frac{A_{30}^{E2}}{R}$$

(2)

Figure 1. Profiles of the specific angular momentum for model 35OC (Woosley & Heger 2006, labeled $j_{35OC}$ and the one amplified by a factor of two (labeled by $2j_{35OC}$) and for models J0.6, J0.8, J1.0, and, J1.2 (from bottom to top), respectively.

(A color version of this figure is available in the online journal.)
The quadrupole amplitude of matter GWs $A_{20}^{\text{matter}}$ consists of the following three parts

$$A_{20}^{\text{matter}} = A_{20}^{\text{hyd}} + A_{20}^{\text{grav}} + A_{20}^{\text{mag}}. \quad (3)$$

The first term in Equation (3) represents the contribution from non-spherical hydrodynamic motions, which is expressed as

$$A_{20}^{\text{hyd}} = \frac{G}{c^4} \frac{32\pi}{\sqrt{15}} \int_0^1 \int_0^{\infty} r^2 dr \rho \nu^2 (v_r^2 (3\mu^2 - 1) + v_\phi^2 (2 - 3\mu^2) - v_r^2 - 6v_r v_\phi \mu \sqrt{1 - \mu^2}). \quad (4)$$

in which $\rho_e$ is the effective density defined as

$$\rho_e = \rho + \frac{e + p + |b|^2}{c^2}. \quad (5)$$

Here, $\rho$, $e$, $p$, $c$, and $G$ denotes the baryon density, the internal energy, the pressure, the speed of light, and the gravitational constant, respectively, and $|b|^2 = b^\mu b_\mu$ is related to the energy density of the magnetic fields with $b_\mu$ representing the magnetic field in the laboratory frame (e.g., Takiwaki et al. 2009). $W = 1/\sqrt{1 - \nu^2} v_\nu$ is the Lorentz boost factor with $v_\nu$ denoting the spatial velocity in spherical coordinates ($i = r, \theta, \phi$). $\mu = \cos \theta$ is a directional cosine. The second term in Equation (3) represents the contribution from gravity as

$$A_{20}^{\text{grav}} = \frac{G}{c^4} \frac{32\pi}{\sqrt{15}} \int_0^1 \int_0^{\infty} r^2 dr \left[ \rho h(W^2 + (v_k/c)^2) + \frac{2}{c^2} \left( \frac{e + p + |b|^2}{c^2} \right) \right] \times \left[ - r \partial_r \Phi (3\mu^2 - 1) + 3 \partial_\theta \Phi \mu \sqrt{1 - \mu^2} \right]. \quad (6)$$

where $\Phi$ denotes the gravitational potential of the self-gravity. The last term in Equation (3) is the contribution from the magnetic fields as

$$A_{20}^{\text{mag}} = - \frac{G}{c^4} \frac{32\pi}{\sqrt{15}} \int_0^1 \int_0^{\infty} r^2 dr \left[ b_r^2 (3\mu^2 - 1) + b_\phi^2 (2 - 3\mu^2) - b_r^2 - 6b_r b_\phi \mu \sqrt{1 - \mu^2} \right] \quad (7)$$

Here, we write the total gravitational amplitude as follows for later convenience.

$$h_{\text{TT}}^{(\text{matter})} = h_{\text{TT}}^{(\text{hyd})} + h_{\text{TT}}^{(\text{mag})} + h_{\text{TT}}^{(\text{grav})}. \quad (8)$$

The second term in Equation (3) represents the contribution from gravity as

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non-zero matter GWs from axisymmetrically but dynamically rotating collapsar disks are generated. In the following computations, we assume that the observer is located in the equatorial plane ($\theta = \pi/2$ in Equation (2)), and also that the distance to the GW source is comparable to nearby GRB-associated CCSNe ($R = 100$ Mpc) unless stated otherwise.

### 2.2.2. GWs from Anisotropic Neutrino Emission

To compute the gravitational waveforms from anisotropic neutrino emission, we follow the formalism pioneeredly proposed by Epstein (1978) and Müller & Janka (1997). In the case of our 2D axisymmetric case, the only non-vanishing component is the plus mode for the equatorial observer,

$$h_\nu = \frac{4G}{c^4 R} \int_0^\pi d\theta' \Phi(\theta') \frac{dl_\nu(\theta', t')}{d\Omega}, \quad (9)$$

where $\Phi(\theta')$ depends on the angle measured from the symmetry axis ($\theta'$)

$$\Phi(\theta') = \sin \theta' (\phi - 2|\cos \theta'|) $$. (10)

As given in Figure 1 of Kotake et al. (2007), this function has positive values in the north polar cap for $0^\circ \leq \theta' \leq 60^\circ$ and in the south polar cap for $120^\circ \leq \theta' \leq 180^\circ$, but becomes negative between $60^\circ < \theta' < 120^\circ$. To determine the anisotropy in neutrino emission (i.e., $dl_\nu/d\Omega$ in Equation (9)), we perform a ray-tracing analysis, which was initially proposed to be applicable in Newtonian gravity (Kotake et al. 2009a, 2009b, 2011) and later improved to be applicable to SR and GR (Harikae et al. 2010a, 2010b).

Applying the formalism in Lindquist (1966), the Boltzmann equation for the neutrino occupation probability $f_\nu(\epsilon_\nu, \Omega)$ for a given neutrino energy ($\epsilon_\nu$) along a specified direction of $\Omega$ can be expressed as

$$\frac{df_\nu(\epsilon_\nu, \Omega)}{d\lambda} = \frac{n_\nu}{Q_\nu(1 - f_\nu) - \kappa f_\nu} = \frac{n_\nu}{Q_\nu - \kappa^* f_\nu}, \quad (11)$$

where $n_\nu$ is the proper number density of the external medium with which neutrinos interact and thus measured in its own local rest frame, $\epsilon_\nu$ is the neutrino energy measured in the local proper frame, $\lambda$ denotes an affine parameter along the geodesics, $Q_\nu$ and $\kappa$ represents neutrino emissivity and absorptivity, and $(1 - f)$ represents the Pauli blocking term (e.g., see Harikae et al. 2010b for more detailed information to derive the equation). Here, we consider only $v_r$ and $\tilde{v}_r$ for simplicity, and the energy and momentum transfer via neutrino scattering are neglected, which is not only difficult to treat with the ray-tracing technique but is also a major undertaking in the radiative transport problem in general. In the final expression of Equation (11), $\kappa^*$ is defined to represent the effective absorptivity. As for the opacity sources of neutrinos, electron capture on proton and nuclei, positron capture on neutron, and neutrino scattering with nucleon and nuclei are included (Fuller et al. 1985; Takahashi et al. 1978; Bruehn et al. 2010). Here, $\kappa^*$ is estimated as

$$\kappa^* = \Sigma \tilde{n}_{\text{target}} \cdot \sigma(\epsilon_\nu) \text{ with } \tilde{n}_{\text{target}} \text{ and } \sigma(\epsilon_\nu) \text{ being the target number density of each reaction and the corresponding cross-sections, respectively.}$$

According to Zink (2008), the formal solution of Equation (11) can be given as

$$f_\nu(\epsilon_\nu, \Omega) = \int_{\lambda_{\text{out}}}^{\lambda_{\text{in}}} n(\lambda') Q_\nu(\lambda', \epsilon_\nu) \times \exp \left[ - \int_{\lambda'}^{\lambda_{\text{out}}} n(\lambda') \kappa^*(\lambda') d\lambda' \right] d\lambda', \quad (12)$$

where the quantities of the right hand side of the equation are defined by combining Equations (2) and (3) with Equations (4), (6), and (7). Note that by dropping the $O(v/c)$ terms, the above formulae reduce to the conventional quadrupole formula employed in Newtonian simulations (e.g., Mönchmeyer et al. 1991). In the collapsar disk that we consider in this work, the condition of $v_\nu \gg v_\phi, v_r$ in Equation (4) is generally satisfied inside the disk. If the disk is in a perfectly stationary state, which means that the centrifugal force ($\sim \rho v_\phi^2$) balances the gravitational forces (e.g., Equation (6)), no GWs can be emitted. As will be explained later, the disk attains mass continuously due to mass accretion, the specific angular momentum of which increases outward (e.g., Figure 1). This is the primary reason why...
which is referred to as the rendering equation of the radiation transport problem. We perform a line integral along the geodesics from every point on the surface of neutrinospheres (λ_out) from which neutrinos can escape freely, up to the outermost boundary of the computational domain (λ_out). On the neutrinospheres, the neutrino distribution function is assumed to take a Fermi–Dirac type distribution \( f^{FD} = \frac{1}{\epsilon_c/k_B T + 1} \) with a vanishing chemical potential.\(^{11}\) For the neutrino energy bins (\( \epsilon_c \)), we use 16 logarithmically spaced energy bins reaching from 3 to 300 MeV.

The path integral in Equation (12) is done explicitly along the geodesics. In doing so, we determine each integration step from 3 to 300 MeV.

With \( f(\epsilon_c, \Omega) \) at the outermost boundary, which is obtained using the above procedure, the emergent neutrino energy fluxes along the specified direction of \( \Omega \) can be estimated as

\[
\frac{dl_v(\Omega, \epsilon_c)}{d\Omega dS} = \int f(\epsilon_c, \Omega) \cdot (c\epsilon_c) \cdot \frac{\epsilon_c^2 d\epsilon_c}{(2\pi h c)^3}.
\]

By summing up the energy fluxes with the weight of the area in the plane perpendicular to the rays (\( dS \)), we can find \( dl_v/d\Omega \) along the specified direction \( \Omega \),

\[
\frac{dl_v(\Omega)}{d\Omega} = \int dl_v(\Omega) d\Omega dS dS.
\]

Repeating the above procedures, \( dl_v(\Omega)/d\Omega \) can be estimated for all directions and by which we can find the amplitudes of the GWs from neutrinos through Equation (9).

\(^{11}\) \( k_B \) is the Boltzmann constant.

3. RESULTS

First, we pay attention to the properties of GWs as well as the hydrodynamic features in models without magnetic fields.

Figure 2 shows the total GW amplitudes for all the computed models without magnetic fields (see also Table 1 for a model summary). As seen, the total GW amplitudes generally show a positively growing feature with time. In addition, a sudden disappearance in the signals can be seen for some slowly rotating models (models J0.6 (red line) and J0.7 (green line)). Figure 3 shows the gravitational waveforms contributed only by anisotropic neutrino emission (left panel) and only by matter motions (right panel), respectively. The increasing trend of the wave amplitude with positive sign is shown to come from the neutrino contribution (left), which is much larger than the matter contribution (right). To understand these properties, we first briefly summarize a hydrodynamic evolution of our 2D non-magnetized models in Section 3.1. Then, in Section 3.2, we move on to analyze the reason for the positive growth in the neutrino GWs by performing ray-tracing analysis. After that, we analyze the matter GWs, paying particular attention to the magnetic effects in Section 3.3. We then discuss their detectability in Section 3.4.

3.1. Hydrodynamic Features

To capture hydrodynamic features in our models, Figure 4 shows the evolution of neutrino luminosities for some selected models. In \( t \sim 1.0 \) s after we start our simulations (\( t \equiv 0 \) s), all the models experience rapid infall and the subsequent shock formation in the center, which leads to a drastic increase and the subsequent decrease in the neutrino luminosities.\(^{12}\)

Here, we take model J0.8 as a reference, because the precolapse angular momentum is adjusted to be closest to the original progenitor. The top left panel in Figure 5 shows a snapshot at \( t = 0.42 \) s near the shock formation. Note that the accretion mass in this epoch is typically greater than 2–3 \( M_\odot \) in the center. The maximum mass of the NS of the Shen EOS is in the same

\(^{12}\) Note that the shock formation is not because the central density exceeds the nuclear density, but because the matter pressure becomes so high due to compression that it pushes back the ram pressure of accreting material in the vicinity of the inner boundary. In contrast to the state-of-the-art simulations of BH-forming CCSNe (e.g., Sumiyoshi et al. 2006; Fischer et al. 2009; Ott et al. 2011), the inability to capture dynamics correctly especially before the BH formation is one of major drawbacks of collapsar simulations in general.
Figure 3. Similar to Figure 2 but for the gravitational waveforms only from anisotropic neutrino emission (left panel, e.g., Equation (9)) and only from matter motions (right panel, e.g., Equation (8)). Note that only the selected models are drawn in this figure so as not to make it messy.

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

Figure 4. Time evolution of neutrino luminosities (the sum of all the neutrino species) for models J0.6, J0.75, J0.8, J0.85, J1.0, and J1.2, respectively. Note that the time is measured from the epoch when the simulations are started (i.e., the onset of gravitational collapse).

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

mass range (Shen et al. 1998, see also O’Connor & Ott 2011; Kiuchi & Kotake 2008). Recent full GR simulations show that the mass of the central object just before collapsing to a BH is typically \( \lesssim 2.3 \, M_\odot \) (e.g., Ott et al. 2011) with typical radius of several km (as inferred from their Figure 3). These evidences might support our very crude assumption of the prompt BH formation that is modeled by setting the BH initially in the center, although such assumption can only be tested by full GR simulations using the same progenitor model. Later on, the density configuration (compare the left half in each panel) deforms to become more oblate with time due to accretion of material with higher angular momentum outside (e.g., Figure 1). As will be explained in the later section, the luminous accretion disk is the primary source of GWs from anisotropic neutrino emission.

The top right panel in Figure 5 shows a snapshot at \( t = 1.89 \, s \). The bluish regions near the spin axis of the accretion disk (left half, density) correspond to the so-called polar funnel regions. The entropy becomes highest near the surface of the accretion disk due to the shock heating when the accreting material hits the wall of the disk (right panel). Comparing Figure 3 (left panel) to Figure 4, the GWs from neutrinos typically deviate from zero later than \( t \gtrsim 2–3 \, s \), when the (total) neutrino luminosities become as high as \( \sim 10^{52} \, \text{erg s}^{-1} \). Later on, the neutrino luminosities show a gradual increase with time, reflecting the increase in the mass accretion to the newly formed accretion disk.

In Figure 4, the neutrino luminosity for models J0.6 (red line) and J0.7 (green line) is shown to steeply decrease at \( \sim 3.3 \, s \) and 7.5 s, respectively. This is because the accretion disk is swallowed to the center (bottom panel in Figure 6) mainly because the pressure support in the accretion disk is reduced by the neutrino cooling. The disappearance of the accretion disks is also the reason for the sudden decrease in the GW signals observed both in the neutrino and matter sectors (e.g., Figures 2 and 3). Before the disappearance, the accretion disk is observed to show rapid expansion and contraction (as indicated in the top panels of Figure 6). At the same time, the mass flux on the Lagrange point around the disk changes violently, which may be related to the so-called runaway instability (e.g., Abramowicz et al. 1983; Font & Daigne 2002). Except for these slowing rotating models, the neutrino luminosities gradually settle to become nearly constant typically later than \( t \sim 6 \, s \) (Figure 4). The saturation of the neutrino luminosity is because the disk is already depletonized and the neutrino emission there is suppressed (e.g., Harikae et al. 2009 for more details). The neutrino luminosities at this epoch become as high as \( 10^{52}–10^{53} \, \text{erg s}^{-1} \), which touches the level of \( \sim 10\% \) of the accretion luminosities (see also Chen & Beloborodov 2007; Sekiguchi & Shibata 2011). Apparently the accretion disk is neutrino-cooling dominated.

From Figure 4, the neutrino luminosity of model J0.8 is shown to be the highest (e.g., blue line). Models with higher initial angular momentum have more extended disks due to larger centrifugal forces (compare the bottom panels in Figures 5 and 7), leading to lower density and temperature in the disks. This is the reason why the neutrino luminosities for models J1.0 (light blue line) and J1.2 (yellow line in Figure 4) are lower in this order. Reflecting this, the GWs from neutrinos become highest for model J0.8 (left panel of Figure 3; see also \( |h_c| \) in Table 1). Regarding the matter GWs, it should be noted that their maximum absolute amplitudes are obtained not for the most rapidly rotating model (model J1.2), but for model J0.8 (right panel of Figure 3). This situation may be akin to the GW signals emitted near core-bounce in CCSNe (see Kotake et al. 2006, 2011; Ott 2009 for recent reviews). Too much initial angular momentum works to suppress the matter compression.
3.2. GWs from Anisotropic Neutrino Emission

In this section, we move on to look more into the detail of the properties of gravitational waveforms mentioned in the previous section. By taking model J0.8 as a reference, we first focus on the neutrino GWs.

Figure 8 shows the local neutrino energy fluxes \( \frac{d \nu}{d \Omega dS} \) (Equation (13)) at \( t = 9.0 \) s (e.g., bottom panel of Figure 5) seen from the polar (left) and the equatorial direction (right). Note that the polar direction is taken to be parallel to the spin axis of the accretion disk. The top panel is for the Minkowski geometry, which corresponds to a purely Newtonian case. Reflecting the shape of the accretion disk in axisymmetry, the contours of the neutrino flux are deformed oblately when seen from the equator.
shows the neutrino luminosity per solid angle \(dL_ν/dΩ\) we place 4\(M_☉\) in the central region, which mimics the event horizon of the BH. The effects of the BH spin are examined by setting the Kerr parameter by hand to be \(a = 0\) and \(a = 0.999\) for the Schwarzschild and the extreme Kerr geometry, respectively. The bottom panels of Figure 8 are for the extreme Kerr geometry, which, however, looks very similar to the middle panels.

For a more detailed comparison, the left panel of Figure 9 shows the neutrino luminosity per solid angle \(dL_ν/dΩ\) for the Minkowski (indicated by “SR”), Schwarzschild (“GR(\(a = 0\))”), and extreme Kerr geometry (“GR(\(a = 0.999\)).” The most important message in this panel is that in every case, the neutrino luminosity seen from the direction near the parallel to the spin axis (\(θ = 0\)) becomes higher than the one seen from the equatorial direction (\(θ = π/2\)). This is because the cross-section of the pancake-like accretion disk seen from the spin axis becomes larger compared to the one seen from the horizontal direction (Figure 8). Remembering again that \(Φ(\theta')\) in Equation (10) is positive near the north and south polar caps, the dominance of the polar neutrino luminosities makes the neutrino GWs positive in the polar cap regions (right panel in Figure 9). This is the reason why the neutrino GWs increase monotonically with time (e.g., left panel of Figure 3). Here, it is worth mentioning that from neutrino luminosity (\(L_ν\), in Figure 4), their typical duration (\(Δt_ν\)), anisotropy in the neutrino radiation (\(α_ν\)), read from the left panel of Figure 9), the GW amplitudes from neutrinos can be estimated from Equation (9) to be

\[
h_ν \approx 10^{-22} (\frac{\alpha_ν}{0.2}) (\frac{L_ν}{10^{53} \text{ erg s}^{-1}}) (\frac{Δt_ν}{10^5 \text{ s}}) (\frac{D}{100 \text{ Mpc}})^{-1},
\]

which is in good agreement with the numerical results (Figure 2).

Comparing the SR case with the GR case in the left panel of Figure 9, the neutrino luminosity becomes smaller by \(\sim 20\%-30\%\) in the GR case due to the GR redshift and also due to the bending effects. Comparing the green line with the blue line, the frame-dragging effect due to the BH spin barely affects the emergent neutrino luminosity. This is probably because the GR effect on the neutrino luminosity becomes important only in the central regions very close to the BH (e.g., Figure 15, Harikae et al. 2010b). Our results indicate that a ray-tracing calculation in the Schwarzschild geometry is at least needed to accurately estimate the neutrino GWs in the collapsar’s environment.

### 3.3. GWs from Matter Motions and Magnetic Fields

Now we briefly analyze the matter GWs in model J0.8. As summarized in Section 2.2, the matter GWs are estimated by the sum of the hydrodynamic and gravity parts in Equations (4) and (6) for non-magnetized models. Among the kinetic energy terms (\(c_ρ \nu_i \nu_j\)) in Equation (4), the largest contribution comes from the rotational energy (i.e., \(-ρ_ν W^2 ν_i ν_j\)), which is shown in Figure 10 (green line). In contrast to the negative contribution by this term, the gravity part (blue line labeled “Gravity” in Figure 10) is shown to make a positive contribution. If a star rotates perfectly stationary, the centrifugal forces balance out the gravitational force, leading to no GWs. In the collapsar disk studied in this work, the disk attains mass continuously due to mass accretion with its specific angular momentum.
Figure 8. Neutrino energy fluxes of $d\nu / (d\Omega dS)$ (Equation(13)) seen from either polar (left) or equatorial directions (right panels) for model J0.8 at $t = 9.1$ s. The top and middle panels are calculated for the Minkowski geometry without or with special relativistic corrections, and the bottom panels are for the extreme Kerr geometry ($a = 0.999$; see the text for more details). Note that the Z-axis in the right panels coincides with the spin direction of the accretion disk.

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increasing outward. This is the reason why the disk is not perfectly stationary, leading to a non-zero GW emission from matter motions (see pink line labeled by “Matter” in Figure 10). However, their GW amplitudes are much smaller compared to those from anisotropic neutrino emission (compare the pink line with the light blue line (neutrino GWs) in Figure 10).

The left panel of Figure 11 shows the normalized contribution of each term in $\Delta I^{(2)}$ at $t = 9$ s for model J0.8, which is estimated by the volume integral of $\Delta I^{(2)}$ within a given sphere enclosed by a certain radius. The region within a radius between 50 km and 160 km is shown to contribute to the production of GWs, which corresponds to a high-density region in the accretion disk and the timescale, which we may take to be the rotational period, respectively, and $\epsilon$ is the degree of the non-sphericity, which we take optimistically to be 10% in the case of the accretion disk.

We now discuss the effects of magnetic fields on the waveforms by taking model J0.8B11, which has the strongest initial magnetic field in our models. In Figure 12, the pink line represents the GWs from magnetic fields (Equation 7). But, first of all, let us briefly comment on the burst-like feature of the quadrupole moment of $I_{ij}$, and $M, R,$ and $T$ represents the typical mass and radius of the accretion disk and the timescale, which we may take to be the rotational period, respectively, and $\epsilon$ is the degree of the non-sphericity, which we take optimistically to be 10% in the case of the accretion disk. This estimate, roughly consistent with the numerical results (e.g., right panel of Figure 3), also shows that anisotropic neutrino emission, producing the GW amplitude on the order of $10^{-22}$ for a source of 100 Mpc, is the primary source in the long-term evolution of collapsars ($\sim 10$ s; see Equation (15)).

Putting these numbers to the standard GW stress formula (e.g., Shapiro & Teukolsky 1983), an upper bound of the matter GW amplitudes may be estimated to be

$$h_{\text{matter}} = \frac{2G}{c^3D} I_{ij} \sim \frac{2G}{c^3D} \frac{MR^2}{T^2} \lesssim 10^{-23} \left( \frac{\epsilon}{0.1} \right) \left( \frac{100 \text{ Mpc}}{D} \right) \times \left( \frac{M}{4M_{\odot}} \right) \left( \frac{R}{100 \text{ km}} \right)^2 \left( \frac{T}{4 \text{ ms}} \right)^{-2},$$

where $D$ is the distance to the source, $I_{ij}$ is the second time derivative of the quadrupole moment of $I_{ij}$, and $M, R,$ and $T$ represents the typical mass and radius of the accretion disk and the timescale, which we may take to be the rotational period, respectively, and $\epsilon$ is the degree of the non-sphericity, which we take optimistically to be 10% in the case of the accretion disk. This estimate, roughly consistent with the numerical results (e.g., right panel of Figure 3), also shows that anisotropic neutrino emission, producing the GW amplitude on the order of $10^{-22}$ for a source of 100 Mpc, is the primary source in the long-term evolution of collapsars ($\sim 10$ s; see Equation (15)).

Now, we discuss the consequences of magnetic fields on the waveforms by taking model J0.8B11, which has the strongest initial magnetic field in our models. In Figure 12, the pink line represents the GWs from magnetic fields (Equation 7). But, first of all, let us briefly comment on the burst-like feature of the green line (contributed from matter motions) at $t \sim 0.24$ s in the figure. This may look similar to the burst of GWs emitted at the moment of BH formation possibly followed by a damped sinusoidal oscillation (Seidel 1990, 1991), but this can only be captured in full GR simulations (Baiotti et al. 2007; Ott et al. 2011; Sekiguchi & Shibata 2011; Kuroda et al. 2012). As already seen in the luminosity plot of Figure 4, the above burst simply corresponds to the shock formation at the center in our SR models. In our main point, we analyze the increasing trend of the GWs from magnetic fields in the following.

The left panel of Figure 13 depicts a snapshot of density (left half), entropy (right top), and plasma $\beta$ (right bottom, the ratio of magnetic to the matter pressure) at the final simulation time ($t = 541$ ms) for model J0.8B11. The high entropy regions in the slightly off-axis region (seen as reddish in the
Figure 11. Left panel shows a normalized contribution of each term in $A^{E2}_{20}$ (e.g., Equation (3) as similar to Figure 10) as a function of radius for model J0.8 at $t = 9$ s. “Gravity ($r/\theta$)” indicates contributions from radial and lateral derivative of the gravitational potential (corresponding to the last two terms in Equation (6)). Right panel shows the radial profiles of the rotational period (red line) and the enclosed mass normalized by $M_\odot$ (green line).

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

Figure 12. Similar to Figure 2 but for model J0.8B11. The pink line represents the GWs from magnetic fields ($h_{(\text{mag})}$ in Equation (8)).

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

top right panel) correspond to the magnetohydrodynamically driven outflows that are pushed outward by the twisted toroidal magnetic fields (also seen as reddish in the bottom right panel, indicated “Toroidal”).

The right panel of Figure 13 shows contributions to the total GW amplitudes (Equation (9), in which the left-hand-side panels are for the sum of the hydrodynamic and gravitational parts (indicated “Matter”), namely, $\log(\pm[A^{E2}_{20}\text{(hyd)} + A^{E2}_{20}\text{(grav)}])$ (left top(+) /bottom(−) (Equations (4) and (6)), and the right-hand-side panels are for the magnetic part, namely, $\log(\pm A^{E2}_{20}\text{(mag)})$ (right top(+) /bottom(−)) (e.g., Equation (7)). By comparing the top two panels, it can be seen that the positive contribution overlaps with the regions where the MHD outflows exist. The major positive contribution is from the kinetic term of the MHD outflows with large radial velocities (e.g., $+\rho_0W^2v_r^2$ in Equation (4)). The magnetic part also contributes to the positive trend (see top right half in the right panel (labeled by mag(+))). This comes from the toroidal magnetic fields (e.g., $+b_0^2$ in Equation (7)), which dominantly contribute to drive MHD explosions.

Unfortunately, our MHD code becomes numerically unstable when the strong MHD jets propagate to a stellar mantle with decreasing density, which prevents us from studying the resulting GWs for a longer period. The neutrino GWs are much smaller than the other GW sources (e.g., Figure 12) simply due to the shorter simulation time. We expect that the increasing trends of the magnetic fields remain as the MHD shocks propagate farther out (e.g., Takiwaki & Kotake 2011). To confirm it, we need to implement a numerical technique especially developed to solve the force-free fields, which is a major undertaking (e.g., McKinney 2006).

3.4. Detectability

Finally, Figure 14 depicts the GW spectra for models J0.8 (left) and J1.0 (right). The GW spectra in the frequency domain between 1 and 100 Hz becomes slightly larger for model J0.8 than for model J1.0. This reflects a more efficient release of the gravitational binding energy for the moderately rotating case (model J0.8) as already mentioned in Sections 3.2 and 3.4. For the cosmological distance scale of GRBs ($\sim$100 Mpc), these low-frequency GW signals are unfortunately very hard to detect even by advanced detectors (like the advanced LIGO or KAGRA/LCGT) whose sensitivity is severely limited by seismic noises (Ando & the TAMA Collaboration 2005; Abbott et al. 2005; Weinstein 2002; Kuroda & LCGT Collaboration 2010). The good news is that these signals could be detectable by a recently proposed future space-based interferometer-like Fabry–Perot-type DECIGO (Kawamura et al. 2006; black line in Figure 14). Two low-luminosity LGRBs were already observed at distances of $\sim$40 Mpc (980425, Galama et al. 1998) and $\sim$130 Mpc (060218, Ferrero et al. 2006). Their local rate, being much higher than that of normal bursts, is expected to be as large $\approx 0.1 D_{100}$ yr$^{-1}$ with $D_{100}$ representing the distance normalized by 100 Mpc (see discussions in Corsi & Mészáros 2009). Our results suggest that the GW astronomy of collapsars could become a reality with the DECIGO-class GW detectors, hopefully in the near future.
4. SUMMARY AND DISCUSSION

By performing axisymmetric SRMHD simulations, we investigated possible signatures of GWs in the context of the collapsar model of LGRBs. By cutting out the central BH, we focused on the GWs generated by asphericities in the neutrino emission and matter motions in the vicinity of the hyperaccreting disks. Nine models were computed by changing the initial angular momenta and magnetic fields parametrically in the precollapse core of a $35 \, M_\odot$ progenitor star. As for the microphysics, a realistic EOS was employed and neutrino cooling was taken into account via a multi-flavor neutrino leakage scheme. To accurately estimate GWs from neutrinos, we performed a ray-tracing analysis in GR via a post-processing procedure. We also studied the effects of magnetic fields on the gravitational waveforms by employing a stress formula that includes contributions from both magnetic fields and SR corrections. We found that the GW amplitudes from anisotropic neutrino emission shows a monotonic increase with time, and that these amplitudes are much larger than those from matter motions of the accreting material. We showed that the increasing trend of the neutrino GWs stems from the excess of neutrino emission in the direction near the parallel to the spin axis illuminated by the hyperaccreting disks. We pointed out that a recently proposed future space interferometer-like Fabry–Perot-type DECIGO would permit the detection of these signals within $\approx 100$ Mpc.

Here, it should be noted that our 2D simulations cannot capture any non-axisymmetric instabilities proposed thus far to provide a strong GW emission in the semianalytical (e.g., van Putten 2001; Davies et al. 2002; Fryer et al. 2002; Kobayashi & Mészáros 2003; Piro & Pfahl 2007; Corsi & Mészáros 2009) and in full GR simulations (e.g., Shibata & Sekiguchi 2005; Manca
et al. 2007). Therefore the present results might be regarded as a lower limit for the possible GW emission in collapsars. To go up the ladder beyond the 2D simulations is very numerically challenging; however, we need to handle it not only to test the outcomes of the proposed ideas about the non-axisymmetric instabilities but also to obtain more accurate waveforms from collapsars.

When studying the formation of BHs and the associated GW signals, the use of a pseudo-Newtonian potential can lead to significant errors. This is the reason why we had to limit our discussion only to the asphericities and the resulting GW emission in the vicinity of the accretion disk which is far away from the central object. In studying the dynamics from core-collapse to NS, a conformally flat condition (CFC) approximation has often been employed to solve the GR equations (Dimmelmeier et al. 2002, 2007, 2008). It has been tested that such a treatment is very good in capturing the results of full-GR results (e.g., Shibata & Sekiguchi 2003) for a wide variety of supernova progenitors with NS formations, but further investigation maybe needed as to whether the approximation is still valid in applications to collapsars, which are highly aspherical disk–BH systems with very dilute polar funnel regions along the rotational axis. Full GR simulations (e.g., Baiotti & Rezzolla 2006; Ott et al. 2011; Sekiguchi & Shibata 2011; Rezzolla et al. 2011; Kuroda et al. 2012, and references therein) are indeed one of the most important topics pointing to the final frontiers of stellar core-collapse simulations; however, it is generally computationally too expensive at present to follow the late-time collapsar evolution up to ~10 s, which is a typical duration of LGRBs. Moreover, the inclusion of microphysics such as neutrino heating is a major undertaking for the full GR simulations. The Cowling approximation (or the fixed metric approach, e.g., McKinney & Narayan 2007b; Komissarov & McKinney 2007; Barkov & Komissarov 2008) has been often used in collapsar simulations thus far, but how to treat the self-gravity of the accretion disk still remains a non-trivial issue. Finally, the pseudo-Newtonian approach we take in this work cannot unambiguously capture accurate properties of the flows in the vicinity of the BH as well as the associated GW signals. Sacrificing the central regions, such a simplified method is currently only one possible way to follow long-term evolution (especially for the disk evolution) including an appropriate treatment of microphysics. Needless to say, these four approaches (full GR, CFC, Cowling, and post-Newtonian) may be regarded as complementary to each other in studying the different epochs in the collapsar evolution. To bridge the gaps between them, it may be a good idea to employ the end results of fully GR simulations (e.g., Ott et al. 2011) in our simulation, which we consider to be the most urgent task to investigate in a sequel to this study.

In a data analysis to extract the true GW signals from the confusing detector noises, it is of primary importance to take a coincident analysis with the conventional electromagnetic observation as well as neutrinos. What could be the photon and neutrino signatures in our collapsar models? To answer this question, we need to perform a long-term simulation that continuously bridges phases 2 and 3 as mentioned in the Introduction. To generate neutrino-driven or MHD-driven outflows in numerical simulations of collapsars, GRMHD simulations that include the effects of neutrino heating are needed, the formulation of which is in steady progress (e.g., Shibata et al. 2011; Müller et al. 2010; Kuroda et al. 2012). Assisted by increasing computational power, these advanced simulations would be hopefully practical by utilizing the next-generation supercomputers. These updates should move forward not only our understanding of the dynamics of the collapsar engines, but also the theoretical predictions of observable multi-messengers (e.g., Ando et al. 2012 for a recent review13), including neutrino emission (e.g., Abbasi et al. 2011) and nucleosynthetic yields (e.g., Fujimoto et al. 2007 and references therein). The neutrino signals emitted from our collapsar models (e.g., from Figure 2 with its typical emergent neutrino energy of ~15–20 MeV) are visible up to the Local Group (~1 Mpc, S. Kawagoe et al. 2012, in preparation) by future megaton-class detectors (e.g., Hyper-Kamiokande, Memphys, and LBNE), large-scale scintillators (e.g., HALO (Engel et al. 2003), and by liquid-Argon detectors (GLACIER; see Scholberg 2010 for collective references therein). While this work might raise many more questions than it can answer, it definitely makes clear that our understanding of the GWs in collapsars is still in its infancy and that collapsars and BH-forming supernovae are “gold mines” in which a number of unsettled and fascinating research themes are hidden. We hope that our exploratory results at least give momentum to theorists to make GW predictions based on a more sophisticated numerical modeling of collapsars.

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13 See Kotake et al. (2012b) for a review about multi-messenger perspectives on core-collapse supernovae.
