GRavitational Wave SIGNAtures of MAGNETOHYDRODynamically Driven Core-Collapse Supernova Explosions

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ABSTRACT

By performing a series of two-dimensional, special relativistic magnetohydrodynamic (MHD) simulations, we study signatures of gravitational waves (GWs) in the MHD-driven core-collapse supernovae. In order to extract the gravitational waveforms, we present a stress formula including contributions both from magnetic fields and special relativistic corrections. By changing the precollapse magnetic fields and initial angular momentum distributions parametrically, we compute 12 models. As for the microphysics, a realistic equation of state is employed and the neutrino cooling is taken into account via a multiflavor neutrino leakage scheme. With these computations, we find that the total GW amplitudes show a monotonic increase after bounce for models with a strong precollapse magnetic field (10^{12} G) and with a rapid rotation imposed. We show that this trend stems both from the kinetic contribution of MHD outflows with large radial velocities and also from the magnetic contribution dominated by the toroidal magnetic fields that predominantly trigger MHD explosions. For models with weaker initial magnetic fields, the total GW amplitudes after bounce stay close to zero because the contribution from the magnetic fields cancels with the contribution from the hydrodynamic counterpart. These features can be clearly understood with a careful analysis of the explosion dynamics. We point out that the GW signals with the increasing trend, possibly visible to the next-generation detectors for a Galactic supernova, would be associated with MHD explosions with explosion energies exceeding 10^{51} erg.

Key words: gravitational waves – hydrodynamics – magnetohydrodynamics (MHD) – neutrinos – stars: magnetars – supernovae: general

Online-only material: color figures

1. INTRODUCTION

Successful detection of neutrinos from SN1987A paved the way for the neutrino astronomy (Hirata et al. 1987; Bionta et al. 1987) alternative to conventional astronomy by electromagnetic waves. Core-collapse supernovae are now expected to provide us an important clue to the solution of the supernova explosion mechanism. Since the GW signatures imprint live information of the asphericity at the moment of explosion, they are expected to offer us a potential method to understand the supernova mechanism.

So far, most of the theoretical predictions of GWs have focused on the bounce signals in the context of rotational core-collapse supernovae (e.g., Mönchmeyer et al. 1991; Zwerger & Mueller 1997; Katake et al. 2003b; Shibata & Sekiguchi 2004; Ott et al. 2004, 2007a, 2007b; Dimmelmeier et al. 2002, 2007, 2008; Scheidegger et al. 2008). For the bounce signals to have a strong and characteristic signature, the iron core must rotate rapidly enough. The waveforms are categorized into the following three types, namely types I, II, and III. Type II and III waveforms have been shown to be less likely to appear than type I waveforms because a combination of general relativity (GR) and electron capture near core bounce suppresses multiple bounces in the type II waveforms (Dimmelmeier et al. 2007; Ott et al. 2007a, 2007b). In general, a realistic nuclear equation of state (EOS) is stiff enough to prohibit type III waveforms. After bounce, asymmetries due to convection (Burrows & Hayes 1996; Müller & Janka 1997; Fryer 2004; Müller et al. 2004), SASI (Marek et al. 2009; Kotake et al. 2007, 2009a, 2009b; Murphy et al. 2009), and g-mode oscillations of protoneutron stars (Ott et al. 2006) are expected to account for sizable GW signals. In general, detection of these GW signals in the postbounce phase (except for the g-mode oscillation) is far more difficult than the bounce signals because they do not possess the same clear signature, but change stochastically with time as a result of chaotically growing convection and SASI in the nonlinear hydrodynamics therein). The explosion geometry is expected to be unipolar and bipolar for the former two cases, and bipolar for the MHD mechanism. Since the GW signatures imprint live information of the asphericity at the moment of explosion, they are expected to provide us an important clue to the solution of the supernova mechanism.
(Kotake et al. 2009b, 2011; Marek et al. 2009; Murphy et al. 2009).

Rapid rotation, necessary for strong bounce signals, is likely to be obtained in only \( \sim 1\% \) of the massive star population (e.g., Woosley & Bloom 2006). However, this can really be the case for progenitors of rapidly rotating metal-poor stars, which experience so-called chemically homogeneous evolution (Woosley & Heger 2006; Yoon & Langer 2005). The high angular momentum of the core and the strong precollapse magnetic field are preconditioned for the MHD mechanism because the MHD mechanism relies on the extraction of rotational free energy of the collapsing core via magnetic fields. The energetic MHD explosions have received great attention recently as to their possible relevance to magnetars and collapsars (e.g., Harikae et al. 2009, 2010 for collective references), which are presumably linked to the formation of long-duration gamma-ray bursts (GRBs; e.g., Meszaros 2006).

Among the previous studies mentioned above, only a small portion of studies has been spent on determining GW signals with the MHD mechanism (Yamada & Sawai 2004; Kotake et al. 2004b; Obergaulinger et al. 2006b; Cerdá-Durán et al. 2007; Shibata et al. 2006; Scheidegger et al. 2010). This may be because the MHD effects on the dynamics as well as their influence over the GW signals can be visible only for cores with precollapse magnetic fields over \( B_0 \gtrsim 10^{12} \mathrm{G} \) (Obergaulinger et al. 2006b; Kotake et al. 2004b). Considering that the typical magnetic field strength of GRB progenitors is at most \( \sim 10^{11}–10^{12} \mathrm{G} \) (Woosley & Heger 2006), this is already an extreme situation. Interestingly, in a more extreme case of \( B_0 \sim 10^{13} \mathrm{G} \), a secularly growing feature in the waveforms was observed (Obergaulinger et al. 2006a; Shibata et al. 2006; Scheidegger et al. 2010). Moreover, Obergaulinger et al. (2006a) identified as type IV a waveform in which quasi-periodic large-scale oscillations of GWs near bounce are replaced by higher frequency irregular oscillations. Some of these MHD simulations follow adiabatic core collapse in which a polytropic EOS is employed to mimic supernova microphysics. At this level of approximation, the bounce shock generally does not stall and a prompt explosion occurs within a few tens of milliseconds after bounce. Therefore, the main focus in these previous studies has been rather limited to the early postbounce phase (\( \lesssim 10 \) ms). However, for models with weaker precollapse magnetic fields akin to the current GRB progenitors, the prompt shocks first stall in the core like a conventional supernova model with more sophisticated neutrino treatment (e.g., Burrows et al. 2007a). In such a case, the onset of MHD explosions, depending on the initial rotation rates, can be delayed until \( \sim 100 \) ms after bounce (Burrows et al. 2007a; Takiwaki et al. 2009). There remains room to study GW signatures in such a case, which we hope to do in this work.

In this study, we choose to take precollapse magnetic fields less than \( 10^{12} \mathrm{G} \) based on a recent GRB-oriented progenitor models. By this choice, it generally takes much longer after bounce compared with the adiabatic MHD models to amplify magnetic fields strong enough to overwhelm the ram pressure of the accreting matter, leading to MHD-driven explosions. Even if the speed of jets in MHD explosions is only mildly relativistic, Newtonian simulations are not numerically stable because the Alfvén velocity (\( \propto B/\sqrt{\rho} \)) could exceed the speed of light unphysically, especially when the strongly magnetized jets (i.e., large \( B \)) propagate to a stellar envelope with decreasing density (\( \rho \)). To follow a long-term postbounce evolution numerically stably, we perform special relativistic MHD (SRMHD) simulations (Takiwaki et al. 2009) in which a realistic EOS is employed and the neutrino cooling is taken into account via a multiflavor neutrino leakage scheme. Note in our previous study of GWs in magneto-rotational core collapse (Kotake et al. 2004b) that we were unable to study properties of the GWs in the postbounce phase because the employed Newtonian simulations quite often crashed, especially in the case of strong MHD explosions. To include GR effects in this study, we follow a prescription in Obergaulinger et al. (2006a), which is reported to capture basic features of full GR simulations quite well. By changing precollapse magnetic fields as well as initial angular momentum distributions parametrically, we compute 12 models. By doing so, we hope to study the properties of GWs in MHD explosions systematically and also address their detectability.

Section 2 presents descriptions of the initial models and numerical methods employed in this work. Formalism for calculating the gravitational waveforms in SRMHD is summarized in Section 3. The main results are given in Section 4. We summarize our results and discuss their implications in Section 5.

2. MODELS AND NUMERICAL METHODS

2.1. Initial Models

We make precollapse models by taking the profiles of density, internal energy, and electron fraction distribution from a rotating presupernova model of E25 in Heger & Langer (2000). This model has a mass of \( 25 M_\odot \) at the zero-age main sequence, however, it loses the hydrogen envelope and becomes a Wolf-Rayet (WR) star of \( 5.45 M_\odot \) before core collapse. Our computational domain involves the whole iron core of \( 1.69 M_\odot \). Note that this model is suggested as a candidate progenitor of long-duration GRBs because type Ib/c core-collapse supernovae that originated from WR stars have an observational association with long-duration GRBs (e.g., Woosley & Bloom 2006).

Since little is known about the spatial distributions of rotation and magnetic fields in evolved massive stars, we add the following profiles in a parametric manner to the non-rotating core mentioned above. For the rotation profile, we assume a cylindrical rotation of

\[
\Omega(X, Z) = \Omega_0 \frac{X^3 + X_0^3}{X^3 + X_0^3} Z_0^4, \tag{1}
\]

where \( \Omega \) is the angular velocity and \( X \) and \( Z \) denote the distance from the rotational axis and the equatorial plane, respectively. The parameter \( X_0 \) represents the degree of differential rotation, which we choose to change in the following three ways: \( 100 \) km (strongly differential rotation), \( 500 \) km (moderately differential rotation), and \( 2000 \) km (uniform rotation). The parameter \( Z_0 \) is fixed to \( 1000 \) km.

Regarding the precollapse magnetic field, we assume that the magnetic field is nearly uniform and parallel to the rotational axis in the core and dipolar outside. This can be modeled by the following effective vector potential:

\[
A_r = A_\theta = 0, \tag{2}
\]

\[
A_\phi = \frac{B_0}{2} \frac{r_0^3}{r^3 + r_0^3} r \sin \theta, \tag{3}
\]

where \( A_r, A_\theta, A_\phi \) are the vector potentials in the \( r, \theta, \) and \( \phi \) directions, respectively, \( r \) is radius, \( r_0 \) is radius of the core, and \( B_0 \) is a model constant (see Takiwaki et al. 2004 for detail).
In this study, $r_0$ is set to 2000 km, which is approximately the size of the precollapse iron core.

By changing the initial angular momentum, degree of differential rotation, and strength of magnetic fields, we compute 12 models. The model parameters are shown in Table 1. The models are named with the first part, B12 or B11, representing the strength of the initial magnetic field, the second part, X1, X5, or X20, indicating the degree of differential rotation ($X_0 = 100, 500, 2000$ km, respectively), and the third part, $\beta = 0.1$ or 1, showing the rotation parameter $\beta$. Here, $\beta$ represents the ratio of the rotational energy to the absolute value of the gravitational energy prior to core collapse. The original progenitor of model E25 in Heger & Langer (2000) has a uniform rotation profile in the iron core and the initial $\beta$ parameter is $\sim 0.15\%$. Thus the initial angular momentum in our models of B11X5 and B12X5.1 is similar to the original one. In a GRB-oriented progenitor of model 35OB (Woosley & Heger 2006), the precollapse magnetic fields reach $\sim 10^{11} - 10^{12}$ G and $\beta \sim 0.2\%$, which is not significantly different from the parameters chosen here. Although little is known about the degree of precollapse differential rotation, an extremely strong one—for example, our model series of X1—is unrealistic. These models, albeit rather academic, are examined in order to clearly see the effects of differential rotation.

### 2.2. Special Relativistic Magnetohydrodynamics

Numerical results in this work are calculated by the SRMHD code developed in Takiwaki et al. (2009). In the following, we first mention the importance of SR and then briefly summarize the numerical schemes.

The Alfvén velocity of MHD jets propagating into the outer layer of the iron core can be estimated as $v_A \sim 10^{10} \text{cm s}^{-1} (B/10^{13} \text{G})(\rho/10^8 \text{g cm}^{-3})^{-1/2}$, where $\rho$ and $B$ are the typical density and the magnetic field along the rotational axis. It can be readily inferred that the Alfvén velocity could exceed the speed of light unphysically in Newtonian simulations. SR corrections are also helpful in correctly capturing the dynamics of infalling material in the vicinity of the protoneutron star because their free-fall velocities and rotational velocities become close to the speed of light. Such conditions are quite ubiquitous in MHD explosions. We have learned that even if the propagation speeds of the jets are only mildly relativistic (at the least), SR treatments are quite important for keeping the stable numerical calculations in good accuracy over a long-term postbounce evolution (e.g., Harikae et al. 2009).

The MHD part of our code is based on the formalism of De Villiers et al. (2003). The state of the relativistic fluid element at each point in the space time is described by its density, $\rho$, specific energy, $e$, velocity, $v_i$, and pressure, $p$, and the magnetic field in the laboratory frame is described by the four-vector $\sqrt{4\pi} b^a = *F^{\mu\nu} U_\nu$, where $*F^{\mu\nu}$ is the dual of the electromagnetic field strength tensor and $U_\nu$ is the four-velocity. The basic equations of the SRMHD code are written as

\begin{equation}
\frac{\partial D}{\partial t} + \frac{1}{\sqrt{\gamma}} \partial_i \sqrt{\gamma} D v^i = 0
\end{equation}

\begin{equation}
\frac{\partial E}{\partial t} + \frac{1}{\sqrt{\gamma}} \partial_i \sqrt{\gamma} E v^i = -p \frac{\partial W}{\partial t} - \frac{p}{\sqrt{\gamma}} \partial_i \sqrt{\gamma} W v^i - \mathcal{L}_v
\end{equation}

\begin{equation}
\frac{\partial (S_i - b^i b_i)}{\partial t} + \frac{1}{\sqrt{\gamma}} \partial_j \sqrt{\gamma} (S_j v^j - b_j b^j) = -\frac{1}{2} \rho h (W v_k)^2
\end{equation}

\begin{equation}
\partial_t (W b^i - W v^i b^j) + \partial_j (W v^i b^j - W v^j b^i) = 0
\end{equation}

\begin{equation}
\partial^i \partial_k \Phi = 4\pi \left[ \rho h (W^2 + (v_k)^2) + 2 \left( p + \frac{|b|^2}{2} \right) \right] - (h^2) + (b^2)
\end{equation}

where $W = 1/\sqrt{1 - v^2}$, $D = \rho W$, $E = eW$, and $S_i = \rho h W^2 v_i$ are the Lorentz boost factor, auxiliary variables corresponding to density, energy, and momentum, respectively. All of them are defined in the laboratory frame. Equations (4)–(6) represent the mass, energy, and momentum conservation, respectively. In Equation (6), note that the relativistic enthalpy, $h = (1 + e/\rho + p/\rho + |b|^2/\rho)$, includes the magnetic energy. Equation (7) is the induction equation in SR. In solving the equation, the method of characteristics is implemented to propagate accurately all modes of MHD waves (see Takiwaki et al. 2009 for more detail). Equation (8) is the Poisson equation for the gravitational potential of $\Phi$, which is solved by the modified incomplete Cholesky conjugate gradient method. For an approximate treatment of GR gravity, $\Phi_{\text{tot}}$ in Equation (6) includes a GR correction to $\Phi$ as in Buras et al. (2006). We employ a realistic EOS based on the relativistic mean field theory (Shen et al. 1998).

We approximate the neutrino cooling by a neutrino multiflavor leakage scheme (Epstein & Pethick 1981; Rosswog &
MHD jets pass through 1000 km along the polar direction. The explosion energy is evaluated when the neutrino (Liebendörfer 2003), in which three neutrino types, electron neutrino ($\nu_e$), electron antineutrino ($\bar{\nu}_e$), and heavy-lepton neutrinos ($\nu_\mu, \nu_\tau, \nu_\tau$, collectively referred to as $\nu_X$), are taken into account. The implemented neutrino reactions are electron capture on proton and free nuclei, positron capture on neutron, and photo-, pair, plasma processes (Fuller et al. 1985; Takahashi et al. 1978; Itoh et al. 1989, 1990). A transport equation for the lepton fractions (namely $Y_e - Y_\mu, Y_\tau, Y_\bar{\nu}_e$, and $Y_{\nu_X}$) is solved in an operator-splitting manner (see Equation (7) in Takiwaki et al. 2009 for more details). $L_\nu$ in Equation (5) represents the neutrino cooling rate summed over all the reactions, which can also be estimated by the leakage scheme (see Epstein & Pethick 1981; Rosswog & Liebendörfer 2003; Kotake et al. 2003a for more details).

In our two-dimensional simulations, the spherical coordinates are employed with $300(r) \times 60(\theta)$ grid points to cover the computational domain assuming axial and equatorial symmetry. The radial grid is nonuniform, extending from 0 to 4000 km with a finer grid near the center. The finest grid for the radial direction is set to 1 km. The polar grid uniformly covers from $\theta = 0$ to $\pi/2$. The finest grid for the polar direction is 25 m. The numerical tests and convergence with this choice of the numerical grid points are given in Section 5 of Takiwaki et al. (2009).

To measure the strength of explosion, we define the explosion energy as:

$$E_{\text{exp}} = \int_D dV \rho W (W - 1),$$

(9)

Here $\rho W$ is the energy density, $\rho$ is the density, and $W$ is the specific internal energy. The domain $D$ in Equation (9) represents the domain where the local energy ($\rho W$) is positive, indicating that the matter is gravitationally unbound. The explosion energy is evaluated when the MHD jets pass through 1000 km along the polar direction.

3. FORMULAE FOR GRAVITATIONAL WAVES IN SRMHD

To extract the gravitational waveforms in MHD explosions, we present the stress formula in SRMHD for later convenience. As shown below, this can be done straightforwardly by extending the Newtonian MHD formulation presented in Kotake et al. (2004b).

From the Einstein equation, one obtains the following formula as a primary expression for the leading part of the gravitational quadrupole field emitted by the motion of a fluid in the post-Newtonian approximation (e.g., Mönchmeyer et al. 1991; Finn & Evans 1990; Blanchet et al. 1990),

$$h_{ij}^{TT}(X, t) = \frac{4G}{c^4 R} P_{ijk}(N) \int d^3 x T_{kl},$$

(14)

where $G$ and $c$ are the gravitational constant and the velocity of light, respectively, $T_{ij}$ is the energy momentum tensor of the source, and $R = |\delta_{ij} X^{k} X^{j}|^{1/2}$ is the distance between the observer and the source. $P_{ijkl}$, with $N = X/R$, denotes the transverse-traceless (TT) projection operator on the plane orthogonal to the outgoing wave direction $N$ (e.g., Mönchmeyer et al. 1991).

$T_{ij}$ consists of the three parts, namely perfect fluid, electromagnetic field, and gravitational potential as follows:

$$T_{ij} = T_{ij}(\text{hyd}) + T_{ij}(\text{mag}) + T_{ij}(\text{grav}).$$

(15)

The first term, which we refer as the hydrodynamic part, is explicitly written as

$$T_{ij}(\text{hyd}) = \rho_s W^2 v_i v_j + p \delta_{ij},$$

(16)

where $\rho_s$ is effective density defined as

$$\rho_s = \rho + \frac{E}{c^2} + \frac{|b|^2}{c^2}.$$  

(17)

The second term in Equation (15) represents the contribution from the magnetic fields as

$$T_{ij}(\text{mag}) = -b_i b_j.$$  

(18)

The last term in Equation (15) is the contribution from the gravitational potential,

$$T_{ij}(\text{grav}) = \frac{1}{4\pi G} \left( \Phi_j \Phi_i - \frac{1}{2} \delta_{ij} \Phi_m \Phi^m \right),$$

(19)

where $\Phi$ corresponds to the self-gravity in Equation (8).

In our axisymmetric case, only one non-vanishing quadrupole term remains in the metric perturbation, namely $\ell = 2, m = 0$ in terms of the pure-spin tensor harmonics, as

$$h_{ij}^{TT}(X, t) \equiv \int_{-\infty}^{\infty} \frac{1}{R} A_{20}^{E2} (t - \frac{R}{c}) T_{ij}^{E2,20} (\theta, \phi),$$

(20)

where $T_{ij}^{E2,20} (\theta, \phi)$ is

$$T_{ij}^{E2,20} (\theta, \phi) = \int_{0}^{\pi} \frac{15}{\pi} \sin^2 \theta \, d \theta.$$  

(21)

(e.g., Thorne 1980). The projection operator in Equation (14) acts on $T_{ij}$ as

$$P_{ijkl} T^{kl} = 2 T^{zz} - T^{zx} - T^{zy}.$$  

(22)

Transforming Equation (14) to the spherical coordinates, and expressing $b_i$ and $v_i$ in terms of unit vectors in the $r, \theta, \phi$ directions, we obtain for $A_{20}^{E2}$ the expression

$$A_{20}^{E2} = A_{20}^{E2}(\text{hyd}) + A_{20}^{E2}(\text{mag}) + A_{20}^{E2}(\text{grav}),$$

(23)

where

$$A_{20}^{E2}(\text{hyd}) = \frac{G}{c^4} \frac{32 \pi^{3/2}}{\sqrt{15}} \int_0^1 \int_0^\infty dr f_{20}^{E2}(\text{hyd}),$$

(24)

$$f_{20}^{E2}(\text{hyd}) = \rho_s W^2 (v_\phi^2 (\mu^2 - 1) + v_\theta^2 (2 - 3 \mu^2) - v_r^2 - 6v_r v_\phi \mu \sqrt{1 - \mu^2}),$$

(25)
Figure 1. Gravitational waveforms with the increasing (left) or the cancellation trend (right; see the text for more details). At the right bottom in each panel, the model names are given such as B12X1β0.1 (top left, for example). The total wave amplitudes are shown by the red line, while the contributions from the magnetic fields and from the sum of the hydrodynamic and gravitational parts are shown by blue and green lines, respectively (e.g., Equation (29) and Equations (25) and (27)). (A color version of this figure is available in the online journal.)

\[
A_{20}^{E2 (grav)} = \frac{G}{c^4} \frac{32\pi^{3/2}}{\sqrt{15}} \int_0^1 d\mu \int_0^\infty r^2 dr f_{20}^{E2 (grav)}, \tag{26}
\]

\[
f_{20}^{E2 (grav)} = \left[ \rho h (W^2 + (v_k/c)^2) + \frac{2}{c^2} \left( p + \frac{|b|^2}{2} \right) - \frac{1}{c^2} (b^{(0)})^2 + (b_k)^2 \right] \left[ -r \delta^i \Phi (3\mu^2 - 1) + 3\partial_\mu \Phi \sqrt{1 - \mu^2} \right]; \tag{27}
\]

and

\[
A_{20}^{E2 (mag)} = -\frac{G}{c^4} \frac{32\pi^{3/2}}{\sqrt{15}} \int_0^1 d\mu \int_0^\infty r^2 dr f_{20}^{E2 (mag)}, \tag{28}
\]

\[
f_{20}^{E2 (mag)} = \left[ b_1^2 (3\mu^2 - 1) + b_2^2 (2 - 3\mu^2) - b_0^2 - 6b_r b_0 \sqrt{1 - \mu^2} \right]. \tag{29}
\]

where \( \mu = \cos \theta \). For later convenience, we write the total GW amplitude as

\[
h_{TT} = h_{TT}^{(hyd)} + h_{TT}^{(mag)} + h_{TT}^{(grav)}, \tag{30}
\]
where the quantities of the right hand are defined by combining Equations (20) and (23) with Equations (25), (27), and (29).

By dropping $O(v/c)$ terms, the above formulae reduce to the conventional Newtonian stress formula (e.g., Mönchmeyer et al. 1991). In the following computations, the observer is assumed to be located in the equatorial plane ($\theta = \pi/2$ in Equation (21)), and the source is assumed to be located near our Galactic center ($R = 10$ kpc).

4. RESULTS

The gravitational waveforms obtained in this work can be categorized into two types, which we call the increasing type and the cancellation type for convenience. Note that the latter type does not mean a new waveform, as will be explained later in this section. The former type was presented in previous literature (Obergaulinger et al. 2006a; Shibata et al. 2006; Scheidegger et al. 2010), however, its properties have not yet been clearly understood. In Section 4.1, we first overview their characters, which are peculiar in the case of MHD explosions. In Section 4.2, we analyze their properties by carefully comparing each contribution in Equation (30) to the total GW amplitudes. Then in Section 4.3, we perform the spectra analysis and discuss their detectability.

4.1. Properties of Waveforms in the MHD Exploding Models

Figure 1 shows examples of the two categories, which we refer to as the increasing type (left panels) and the cancellation type (right panels), respectively. In the increasing type, the total wave amplitudes (red line) have a monotonically increasing trend after bounce ($t - t_b = 0$ in the figures) while in the cancellation type (right panels), the total amplitudes after bounce stay almost zero. This is because the contribution from the magnetic fields (blue line, Equation (29)) cancels with the contribution from the sum of the hydrodynamic and gravitational parts (green line, Equations (25) and (27)). Regardless of the difference in the two types, it is common that the magnetic contribution (blue line) increases almost monotonically with time. Not surprisingly, the bounce GW signals ($t - t_b \lesssim 20$ ms) are categorized into the so-called type I or II waveforms. Note here that the MHD simulations are terminated at around 100 ms after bounce for all the computed models. This is simply because the GW amplitudes in a later phase decrease because the MHD shock comes out of the computational domain and the enclosed mass in the domain becomes smaller.

Figure 2 depicts a classification of the computed models, in which “C” and “I” indicate the cancellation and increasing types, respectively. I* indicates the mixture of the two types, which we call the intermediate type. The figure shows that the bifurcation of the two types is predominantly determined by the precollapse magnetic fields ($B_0 \times 10^{12}$ G) are basically classified as the increasing type. For models shown in bold, which are slow rotators with uniform rotation in our models, the field amplification works less efficiently than for models with stronger differential rotation (such as $X_0 = 100, 500$ km; see also Cerdá-Durán et al. 2008). This suppresses the increase in the wave amplitudes due to magnetic fields, which gives rise to the intermediate state between the two types.

Figure 3 is a summary showing intervals measured from the stall of the bounce shock to the MHD-driven revival of the stalled shock (top panel) and the one measured from core bounce (bottom panel). The computed models are classified according to whether the launch of the MHD jets occurs relatively promptly after bounce (models colored by red, with the intervals typically being shorter than $\sim 30$ ms) or rather later (models colored by green), which we refer to as prompt or delayed MHD explosion for convenience in this work (see text for more details). Note that these timescales are estimated just by looking at the velocity evolutions along the polar axis.

Figure 4 shows a summary regarding the explosion energy (defined in Equation (9)). Compared with Figures 2–3, the explosion energies for the models with the increasing trend generally exceed $10^{51}$ erg (models shown in yellow).
Having summarized the waveform classification together with the explosion dynamics, we move on to look in more detail at what causes the difference between the two types in the next section.

4.2. Analysis of Waveforms

4.2.1. Increasing Type

By taking model B12X1β0.1 as a reference, we first focus on the increasing type waveform. The left panel of Figure 5 depicts distributions of entropy (left half) and plasma $\beta$ (right half, the ratio of matter to magnetic pressure) at 100 ms after bounce. It can be seen that the outgoing jets indicated by the velocity fields (arrows in the left half) are driven by the magnetic pressure behind the shock (see bluish region, i.e., low plasma $\beta$ in the right half of the panel).

The right panel of Figure 5 shows contributions to the total GW amplitudes (Equation (27)), in which the left-hand panels are for the sum of the hydrodynamic and gravitational part, namely $\log(\pm f^{E2}_{20}^{\text{hyd}} + f^{E2}_{20}^{\text{grav}})$ (left top(+) / bottom(−)) (Equations (25) and (27)), and the right-hand panels are for the magnetic part, namely $\log(\pm f^{E2}_{20}^{\text{mag}})$ (right top(+) / bottom(−)) (e.g., Equation (29)). By comparing the top two panels in Figure 5, it can be seen that the positive contribution overlaps 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., $+\rho_0 W^2 v_r^2$ in Equation (25)). The magnetic part also contributes to the positive trend (see the top right half in the right panel; labeled by mag(+)). This comes from the toroidal magnetic field (e.g., $+b_\phi^2$ in Equation (29)), which dominantly contributes to driving MHD explosions. The magnetic contribution was already mentioned in Kotake et al. (2004b). This study furthermore adds that the kinetic energy of MHD outflows more importantly contributes to the positive trend.

Figure 6 shows a normalized cumulative contribution of each term in $A_{20}^{E2}$, which is estimated by the volume integral of $A_{20}^{E2}$ within a given sphere enclosed by a certain radius. It can be seen that the respective contributions of the hydrodynamic and gravity parts (indicated by “hyd and grav”) are prominent for a radius outside ~1000 km, which stems from exploding regions with large kinetic energy as mentioned above. Note that the secular drift observed in the increasing waveform may come from ambiguities in estimating the gravitational potential in the stress formula as pointed out by Dimmelmeier et al. (2002). To exclude such potential issues, we plot the waveform calculated

![Figure 5](image-url)
by the first moment of momentum–density formalism (e.g., Finn & Evans 1990). As is shown in Figure 7, the increasing trend is also seen in the waveform estimated by the first-moment formalism (green line). As a side note, a smoother curve is obtained for the stress formula (red line), about which Mönchmeyer et al. (1991) pioneeringly mentioned that the numerical evaluation of the time derivative sometimes makes the waveform noisy.6

4.2.2. Cancellation Type

Now we proceed to focus on the cancellation-type waveform by taking model B11X1β1 as an example. The top panel in Figure 8 shows density distributions (left half) with the plasma β (right half). Note that the jet head of MHD outflows is located at \( \sim 400 \) km along the pole which sticks out of the plots in Figure 8 (compare the difference in scales of Figure 5). In fact, the bluish regions (low β) around the pole have a lower density because material there has already been blown up due to the passage of MHD-driven shocks. The middle panel of

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6 This may be the reason why the stress formula has often been employed in supernova studies so far.
contribution from magnetic fields (bluish region in the bottom right half of the middle panel) comes from the regions that are off-axis from the propagation of the MHD jets. In these regions, the poloidal components of magnetic fields are stronger than the toroidal ones, which makes the negative contribution mainly through $-b_1^2(3\mu_1^2 - 1) + b_0^2(2 - 3\mu_0^2)$ in Equation (29).

Looking at the sum of the hydrodynamic and gravitational part (left in the middle panel), a large negative contribution comes from regions near the rotational axis (in red, bottom left). The bottom panel of Figure 8 further shows a contribution from the hydrodynamic (left) and gravitational parts (right), separately. Regarding the hydrodynamic part, the negative contribution is highest in the vicinity of the equatorial plane which closely coincides with the oblately deformed protoneutron star (red in the bottom left in the bottom panel). This is because the negative contribution comes from the centrifugal forces (e.g., the term related to the rotational energy, $-\rho v^2_\phi$ in Equation (25)). For the gravitational part (right side in the bottom panel), a large negative contribution comes from regions in the vicinity of the rotational axis (bottom right). This comes from the term of $-r\partial_r \Phi(3\mu_1^2 - 1)$ in Equation (25), which is the radial gradient of the gravitational potential. Remembering that $\partial_r \Phi > 0$ is generally satisfied in self-gravitating objects, the sign of this term is determined by $\mu = \cos \theta$, which is a directional cosine measured from the rotational axis. As a result, the gravitational part makes a negatively contribution in the vicinity of the rotational axis ($\mu \sim 1$). These two factors coming both from the hydrodynamic and gravity part make the negative contribution to the GW amplitudes. In fact, the negative contribution is highest near the surface of the protoneutron star ($\sim 10$ km in Figure 9). Outside it, the magnetic part becomes almost comparable to the sum of the hydrodynamic and gravitational parts, which creates the cancellation type shown in the right panels of Figure 1.

### 4.3. Spectrum Analysis

Now we move on to performing a spectral analysis (Figure 10). For the two types (left panel, increasing type; right panel, cancellation type), the peak amplitudes in the spectra are around 1 kHz, which comes from the GWs near at bounce. It can be seen that the spectra for lower frequency domains (below $\sim 100$ Hz) are much larger for the increasing type (left panels) compared with the cancellation type (right panels). This reflects a slower temporal variation of the secular drift inherent to the increase-type waveforms (e.g., Figure 1).

As a measure for characterizing the dominance in the lower frequency domains, we define $\tilde{h}_{\text{low}}$, which represents average amplitudes below 100 Hz (see Figure 11). Although the peak amplitudes, $\tilde{h}_{\text{peak}}$, in the spectra have no clear correlation with the two types, we point out that the final GW amplitudes (the first column) and the $\tilde{h}_{\text{low}}$ (the third column) are much larger for the increasing type (in yellow) compared with the cancellation type (in light blue). In Figure 10, the peak amplitudes near 1 kHz are, irrespective of the two types, marginally within the detection limits of the currently running detector of the first LIGO and the advanced LIGO for a Galactic supernova. It is true that the GWs in the low frequency domains mentioned above are relatively difficult to detect because of seismic noises, but a recently proposed future space interferometer like the Fabry–Perot type DECIGO is designed to be sensitive in those frequency regimes (Kawamura et al. 2006; Kudoh et al. 2006). The sensitivity curve of the detector is plotted with the black line in Figure 10. Our results suggest that these low-frequency signals, if observed, could be one important messenger of the increasing type waveforms that are likely to be associated with MHD explosions exceeding $10^{51}$ erg.

### 5. SUMMARY AND DISCUSSION

By performing a series of two-dimensional SRMHD simulations, we studied signatures of GWs in the MHD-driven core-collapse supernovae. In order to extract the gravitational waveforms, we presented a stress formula including contributions both from magnetic fields and special relativistic corrections. By changing the precollapse magnetic fields and initial angular momentum distributions parametrically, we computed 12 models. As for the microphysics, a realistic EOS was employed and the neutrino cooling was taken into account via a multitype neutrino leakage scheme. With these computations, we found that the total GW amplitudes show a monotonic increase after bounce for models with a strong precollapse magnetic field ($10^{12}$ G) that also have a rapid rotation imposed. We pointed out that this trend stems both from the kinetic contribution of MHD outflows with large radial velocities and from the magnetic contribution dominated by the toroidal magnetic fields that predominantly trigger MHD explosions. For models with weaker initial magnetic fields, the total GW amplitudes after bounce stay almost zero because the contribution from the magnetic fields cancels with the contribution from the hydrodynamic counterpart. These features can be clearly understood with a careful analysis of the explosion dynamics. It was pointed out that the GW signals with an increasing trend, possibly visible to the next-generation detectors for a Galactic supernova, would be associated with MHD explosions exceeding $10^{51}$ erg.

Although the presented simulations have utilized the leakage scheme to approximate the deleptonization, it would be more accurate (especially before bounce) to employ a formula developed by Liebendorfer (2005), which was designed to fit one-dimensional Boltzmann results. Figure 12 shows snapshots at around 33 ms after bounce for model B12X20/80.1 in which the leakage scheme (left panel) or the $Y_e$ prescription (right panel) is employed, respectively. Note that “G15” is taken in our simulation until bounce among the parameter sets in.
Figure 10. Gravitational-wave spectrum for representative models with the expected detection limits of the first LIGO (Abbott et al. 2005), the advanced LIGO (Weinstein 2002), LCGT (Kuroda & the LCGT Collaboration 2006), and Fabry–Perot-type DECIGO (Kawamura et al. 2006; Kudoh et al. 2006). It is noted that $h_{\text{char}}$ is the characteristic gravitational wave strain defined in Flanagan & Hughes (1998). The supernova is assumed to be located at a distance of 10 kpc.

Figure 11. Summary of GW amplitudes for all models. The colors used are the same as in Figure 2. The first column represents the GW amplitudes when we terminated the simulation (100 ms after bounce). The second column, $h_{\text{100ms}}$, is the peak GW amplitudes in the spectra. The third column $h_{\text{low}}$, is the average amplitude below 100 Hz. The supernova is assumed to be located at a distance of 10 kpc.

Now we mention a comparison between the obtained results and relevant MHD simulations. Model R4E1CF in Scheidegger et al. (2010), whose precollapse rotational parameter is $\beta = 0.5\%$ with a uniform rotation imposed and whose initial poloidal magnetic field is set to $10^{12}$ G, is close to our model B12X20/0.1. From their Figure 23, the jet propagates to $\sim 300$ km along the rotational axis at around $\sim 18$ ms after bounce. In our counterpart model, the MHD-driven shock revives after around 30 ms after bounce, and it reaches 300 km at around $\sim 10$ ms, which is equivalent to $\sim 40$ ms after bounce. Considering that our model ($\beta = 0.1\%$) is a slower rotator than model R4E1CF ($\beta = 0.5\%$), the delay of the shock revival for our model seems reasonable. Model s20A1B5-D3M12 is quite similar.

(Liebendörfer 2005) and is switched to the leakage scheme at the postbounce phase. As is shown, shock revival also occurs for the model with the $Y_e$ prescription (right panel). Figure 13 depicts the magnetic pressure (red line) versus ram pressure (blue line) along the polar axis (left panel) or the equatorial plane (middle panel) near the rebirth of the stalled shock for the model with the $Y_e$ prescription (top panels) or the leakage scheme (bottom panels). For the equator, the magnetic pressure is much less than the ram pressure (middle panel), while the magnetic pressure amplified by the field wrapping along the pole becomes as high as the ram pressure of the infalling material at the shock front, leading to the MHD-driven shock formation (see right panels). Regardless of the two different deleptonization schemes, these important features associated with the MHD explosions are shown to be quite similar.

Now we mention a comparison between the obtained results and relevant MHD simulations. Model R4E1CF in Scheidegger et al. (2010), whose precollapse rotational parameter is $\beta = 0.5\%$ with a uniform rotation imposed and whose initial poloidal magnetic field is set to $10^{12}$ G, is close to our model B12X20/0.1. From their Figure 23, the jet propagates to $\sim 300$ km along the rotational axis at around $\sim 18$ ms after bounce. In our counterpart model, the MHD-driven shock revives after around 30 ms after bounce, and it reaches 300 km at around $\sim 10$ ms, which is equivalent to $\sim 40$ ms after bounce. Considering that our model ($\beta = 0.1\%$) is a slower rotator than model R4E1CF ($\beta = 0.5\%$), the delay of the shock revival for our model seems reasonable. Model s20A1B5-D3M12 in

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|}
\hline
Model & $h_{\text{100ms}}[10^{-21}]$ & $\tilde{h}_{\text{peak}}[10^{-21}]$ & $\tilde{h}_{\text{low}}[10^{-21}]$ \\
\hline
B11X1/30.1 & 2.2 & 10.3 & 0.3 \\
B11X5/30.1 & 1.8 & 13.9 & 0.3 \\
B11X20/30.1 & 2.0 & 3.0 & 0.3 \\
B11X1/31 & 4.0 & 52.5 & 1.0 \\
B11X5/31 & 0.8 & 7.5 & 0.2 \\
B11X20/31 & 1.1 & 17.7 & 0.2 \\
B12X1/30.1 & 22.0 & 18.1 & 5.2 \\
B12X5/30.1 & 20.2 & 14.8 & 3.3 \\
B12X20/30.1 & 3.5 & 8.2 & 0.9 \\
B12X1/31 & 25.7 & 35.3 & 5.6 \\
B12X5/31 & 26.3 & 7.4 & 4.8 \\
B12X20/31 & 12.3 & 17.6 & 2.1 \\
\hline
\end{tabular}
\end{table}
Figure 12. Snapshots at 33 ms after bounce for model B12X20/0.1 in which the leakage scheme (left panel) or the $Y_e$ prescription (right panel) are employed. In each panel, density (logarithmic, left half) and entropy (right half) distributions are shown. The side length of each plot is $600 \times 600$ km.

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

Figure 13. Left and middle panels show magnetic pressure (red line) vs. ram pressure (blue line) for model B12X20/0.1 around 32 ms after bounce with the $Y_e$ prescription (top panels) or the leakage scheme (bottom panels) along the polar axis (left panel) or the equatorial plane (middle panel). Matter pressure is shown by a green line as a reference. The right panels show velocity profiles along the pole near after the stall of the bounce shock (red lines). The MHD-driven explosions are indeed obtained in both of the two different deleptonization schemes.

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

Cerdá-Durán et al. (2008), whose precollapse angular velocity is $4 \text{ rad s}^{-1}$ (the rotational parameter should be close to $\beta = 0.1\%$) with uniform rotation and whose initial poloidal magnetic field is $10^{12}$ G, is close to our model B12X20/0.1. From their Figure 10, the MHD jet propagates to ~240 km at 51 ms after bounce. The employed EOS is the same as ours (the Shen EOS), while the deleptonization scheme taken in their study was the $Y_e$ formula (Liebendörfer 2005). As mentioned above, the dynamics are rather close to our corresponding model. Among the computed models in Burrows et al. (2007a), model M15B12P2A1H, which has a precollapse angular velocity of $\pi \text{ rad s}^{-1}$ (the rotational parameter should be close to $\beta = 0.06\%$) with an initial dipolar magnetic field of $10^{12}$ G, is close to our model B12X20/0.1. The interval before the launch of the MHD shock for their model is 80 ms after bounce (e.g., their Table 1), which is much later than it is for our model (28 ms after bounce). This may be due to the larger initial angular momentum ($\beta = 0.1\%$) assumed in our study. Model A3B3G5-D3M13 in Obergaulinger et al. (2006b), which has a rotational parameter of $\beta = 0.9\%$ with a differential rotation imposed (the radial cutoff is 500 km) and the initial poloidal magnetic field is $10^{13}$ G, is closer to our model B12X1/0.1. The MHD jet reaches to 500 km at around 7 ms after bounce, which is also the case of our counterpart model. As discussed above, our results are compatible to the ones obtained in the relevant preceding results.

The major limitation of this study is the assumption of axisymmetry. Recently, it was reported in three-dimensional (3D) MHD core-collapse simulations (Scheidegger et al. 2008, 2010) that the fast growth of the spiral SASI hinders the efficient amplification of the toroidal fields, which could suppress the formation of jets that are rather easily realized in two-dimensional simulations. As a sequel of this study, we plan to investigate the 3D effects in SRMHD. Regarding a resolution dependence of our results, Figure 14 indicates that our standard resolution is adequate to follow the evolution of the computed models.
However, it is not at all sufficient to capture the magneto-rotational instability (MRI; e.g., Balbus & Hawley 1998). At least 10–100 times finer mesh points are required for resolving the MRI (Obergaulinger et al. 2009), which may require some adaptive mesh refinement treatment, a very important component that remains to be improved. If the MRI could be resolved, the increasing type waveform could also emerge for models with weak initial magnetic fields because a more efficient field amplification could be captured. Although the GR effects were treated only by rough approximation, we think that the GR should not drastically change our results qualitatively because the central protoneutron stars will not collapse to a black hole during our simulation time as inferred from a simple argument that the total amplitudes become larger when we include the neutrino heating may play a minor role in the waveforms since the timescales before the neutrino-driven explosions set in (e.g., Marek & Janka 2009; Bruenn et al. 2009; Suwa et al. 2010, and references therein) are much longer than the MHD explosions.

As one possible extension of this study, we think it interesting to study GWs from anisotropic neutrino radiation in the MHD models. Extrapolating the outcome obtained in previous studies (e.g., MÜLLER et al. 2004; KOTAKE et al. 2009a), we anticipate that the total amplitudes become larger when we include the neutrino GWs. This is because the neutrino GWs may make a positive contribution since the neutrino emissions from the oblately deformed protoneutron stars become stronger toward the polar direction (JANKA & MOENCHMEYER 1989; KOTAKE et al. 2003a; WALDER et al. 2005; OTT et al. 2008). If this is really the case, the total GW amplitudes especially for the increasing type should be much larger, possibly making its detection more promising. Furthermore, the neutrino signals from MHD explosions may have a sharp directional dependence through neutrino oscillations, reflecting the aspherical propagation of the shock to the stellar envelope (KAWAGOE et al. 2009). Taking a correlation analysis between the neutrino and GW signals could help to reveal the hidden nature of the central engines. In fact, several observational proposals have been made in this direction recently (van Eleywck et al. 2009; ASO et al. 2008; LEONOR et al. 2010). The MHD-driven core-collapse supernovae, albeit rather minor among typical type II supernovae explosions, seem to still contain a number of rich research subjects.

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