SYSTEMATIC FEATURES OF AXISYMMETRIC NEUTRINO-DRIVEN CORE-Collapse SUPERNOVA MODELS IN MULTIPLE PROGENITORS
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Abstract

We present an overview of two-dimensional (2D) core-collapse supernova simulations employing neutrino transport scheme by the isotropic diffusion source approximation. Studying 101 solar-metallicity progenitors covering zero-age main sequence mass from 10.8 $M_\odot$ to 75.0 $M_\odot$, we systematically investigate how the differences in the structures of these multiple progenitors impact the hydrodynamics evolution. By following a long-term evolution over 1.0 s after bounce, most of the computed models exhibit neutrino-driven revival of the stalled bounce shock at $\sim$ 200 - 800 ms postbounce, leading to the possibility of explosion. Pushing the boundaries of expectations in previous one-dimensional (1D) studies, our results confirm that the compactness parameter $\xi$ which characterizes the structure of the progenitors is a key to diagnose the properties of neutrino-driven explosions also in 2D. Models with high $\xi$ undergo high ram pressure from the accreting matter onto the stalled shock and it affects the shock expansion under the influence of neutrino-driven convection and the standing accretion-shock instability, as well as mass of protoneutron star left at the center. We show that the accretion luminosity becomes higher for models with high $\xi$, which makes the diagnostic energy higher and the synthesized nickel mass bigger. We find that these explosion characteristics tend to show a monotonic increase as a function of the compactness parameter $\xi$.

Subject headings: hydrodynamics --- neutrinos --- supernovae: general

1. INTRODUCTION

Core-collapse supernova (CCSN) mechanism is essentially an initial value problem. This has been widely agreed among CCSN theorists recently (e.g., O'Connor & Ott 2011, Ugliano et al. 2012, Couch & Ott 2013). For low-mass progenitors with O-Ne-Mg core, the neutrino mechanism works successfully to explode in one-dimensional (1D) simulations because of the tenuous envelope (Kitaura et al. 2006). For more massive progenitors with iron core, multidimensional (multi-D) effects such as neutrino-driven convection and the standing-accretion-shock-instability (SASI) are crucially important (e.g., Burrows et al. 1993, Janka & Muel"{e}d 1996, Foglizzo et al. 2000, Murphy et al. 2013, Fernandez et al. 2014 and references therein). Recently this has been confirmed by a number of self-consistent two-(2D) and three-dimensional (3D) simulations (e.g., Buras et al. 2008, Ott et al. 2008, Marek & Janka 2009, Bruenn et al. 2013, Suwa et al. 2010, 2014, Muller et al. 2012, 2013, Takiwaki et al. 2012, 2014, Hanke et al. 2013, Dolence et al. 2014). Up to now, the number of these state-of-the-art models amounts to $\sim$ 40 covering the zero-age main sequence (ZAMS) mass from 8.1 $M_\odot$ (Muller et al. 2012) to 27 $M_\odot$ (Hanke et al. 2013).

Based on stellar evolutionary calculations, on the other hand, hundreds of CCSN progenitors are available now, depending on a wide variety of the ZAMS mass, metallicity, rotation, and magnetic fields (e.g., Nomoto & Hashimoto 1988, Woosley & Weaver 1995, Woosley et al. 2002, Woosley & Heger 2007, Heger et al. 2003, 2005, Limongi & Chieffi 2006). These huge sets of CCSN progenitors, aided as well as by development of high-performance computers and numerical schemes, make systematic numerical study of CCSNe possible.

By performing general-relativistic (GR) 1.5D simulations for over 100 presupernova models using a leakage scheme, O'Connor & Ott (2013) were the first to point out that the postbounce dynamics and the progenitor-remnant connections are predictable basically by a single parameter, the compactness of the stellar core at bounce (see also O'Connor & Ott 2013). Along this line, Ugliano et al. (2012) performed 1D hydrodynamic simulations for 101 progenitors of Woosley et al. (2002). By replacing the proto-neutron star (PNS) interior with an inner boundary condition, they followed an unprecedentedly long-term evolution over hours to days after bounce in 1D. Their results also lent support to the finding by O'Connor & Ott (2011) that the compactness parameter is a good measure to diagnose the progenitor-explosion and the progenitor-remnant correlation.

Joining in these efforts but going beyond the previous 1D approach, we perform neutrino-radiation hydrodynamics simulations in 2D using the whole presupernova series (101 models) with solar metallicity of Woosley et al. (2002). Without the excision inside the PNS, we can self-consistently follow a long-term evolution starting from the onset of core-collapse, bounce, neutrino-driven shock-revival, until the revived shock comes out of the iron core. The goal of our 2D models is not to determine the very final fate of a massive star (which requires 3D-GR models with detailed trans-
Figure 1. Entropy distribution in unit of $k_B$ per baryon for selected nine models at $t_{pb} = 400$ ms. Shown are models s11.2 (a) to s24.0 (i), from top-left to bottom-right. Note that each model presents a different scale as shown in the panel.

Figure 2. Same as Figure 1 but for models s25.0 (a) to s75.0 (i), from top-left to bottom-right.

2. NUMERICAL SETUP

The employed numerical methods are essentially the same as those in [Takiwaki et al. 2014]. Our 2D models are computed on a spherical polar grid of 384 non-equidistant radial zones from the center up to 5000 km and 128 equidistant angular zones covering $0 < \theta < \pi$. We employ the equation of state by Lattimer & Swesty (1991) with a compressibility modulus of $K = 220$ MeV. For the calculations presented here, self-gravity is computed by a Newtonian monopole approximation and our code is updated, from the ZEUS-MP [Hayes et al. 2006], to use high-resolution shock capturing scheme with an approximate Riemann solver of Einfeldt (1988). As described in Nakamura et al. (2014), we take into account explosive nucleosynthesis and the energy feedback into hydrodynamics by solving a 13 $\alpha$-nuclei network including $^4$He, $^{12}$C, $^{16}$O, $^{20}$Ne, $^{24}$Mg, $^{28}$Si, $^{32}$S, $^{36}$Ar, $^{40}$Ca, $^{44}$Ti, $^{48}$Cr, $^{52}$Fe, and $^{56}$Ni. The nuclear energy compensates for energy loss via endothermic decomposition of iron-like NSE nuclei to lighter elements (see Appendix of Nakamura et al. 2014 for more details).

To solve spectral transport of electron- ($\nu_e$) and antielectron neutrinos ($\bar{\nu}_e$), we employ the isotropic diffusion source approximation (IDSA, Liebendörfer et al. 2009). We take a ray-by-ray approach, in which the neutrino transport is solved along a given radial ray assuming that the hydrodynamic medium for the direction is spherically symmetric (e.g., Buras et al. 2006). Although one needs to deal with the lateral transport more appropriately (e.g., Sumiyoshi et al. 2014; Dolence et al. 2014), this approximation is useful because of the high computational efficiency in parallelization, which allows us to explore the more systematic progenitor survey based on the radiation-hydrodynamics models than ever in this study. Regarding heavy-lepton neutrinos ($\nu_\tau = \nu_\mu, \nu_\tau, \bar{\nu}_\mu, \bar{\nu}_\tau$), we employ a leakage scheme to include the $\nu_\tau$ cooling via pair, photo and plasma processes (see Takiwaki et al. 2014 for more details).

The investigated progenitors with iron cores (Woosley et al. 2002) are given in 0.2 $M_\odot$ steps between 10.8 $M_\odot$ (s10.8) and 28.2 $M_\odot$ (s28.2) and further from 30 $M_\odot$ (s30) up to 75.0 $M_\odot$ (s75) in 1.0 $M_\odot$ steps. The structure of these stars, such as density profiles and the pre-collapse masses have been already described in Ugliano et al. (2012). To induce non-spherical instability, we add initial seed perturbations by zone-to-zone random density variations with an amplitude less than 1%.

Following O'Connor & Ott (2011), we estimate the compactness parameter as the ratio of mass $M$ and the enclosed radius $R(M)$,

$$\xi_M \equiv \frac{M/M_\odot}{R(M)/1000\text{km}}. \quad (1)$$

The previous studies used $M = 2.5 \ M_\odot$ (O’Connor & Ott 2011; Ugliano et al. 2012) or 1.75 $M_\odot$ (O’Connor & Ott 2013) and estimated $\xi$ at the time of core bounce. On the other hand, the outer radius of our computational domain (5000 km) is too small to contain 2.5 $M_\odot$ for all models and even 1.75 $M_\odot$ for some less massive models. In this Letter, we estimate $\xi$ at $M = 2.5 \ M_\odot$ ($\xi = \xi_{2.5}$) directly from pre-collapse data. It should be noted that our definition of $\xi$ gives almost the same compactness estimated at bounce, because the radius $R$ enclosing 2.5 $M_\odot$ is far from the center and...
the radial velocity $v_R$ there is very small (e.g., for $s15.0$ model $R = 1.7 \times 10^9$ cm and $v_R = -6.8 \times 10^6$ cm s$^{-1}$). Actually $\xi_{2.5}$ of $s15.0$ model in our definition is 0.145, which is very close to the value (0.150) estimated by O’Connor & Ott (2011) at bounce.

### 3. RESULTS

For all the employed 101 models, the bounce shock stalls in a spherically symmetric manner and only after that, we observe a clear diversity of the multi-D hydrodynamics evolution in the postbounce (pb) phase. Figures 1 and 2 show a snapshot of entropy distribution for selected 18 models at $t_{ob} = 400$ ms. For some less massive progenitors (e.g., model $s11.2$ in Figure 1a), the shock is reaching close to the outer boundary of the computational domain with developing pronounced unipolar and dipolar shock deformations. At this time, the shock of the most massive progenitor ($s75.0$ in Figure 2i)) is reaching an average radius of $\langle r \rangle \sim 1000$ km, whereas the shock of $s24.0$ (in Figure 1i) still wobbles around at $\langle r \rangle \sim 200$ km. This demonstrates that the ZAMS mass is not a good criterion to diagnose the possibility of explosion.

This is more clearly visualized in the top left panel of Figure 3. Taking six models as an example in the mass range between 19.2 $M_\odot$ and 24.0 $M_\odot$, the shock revival is shown to occur earlier for $s20.0$ (red line) and $s22.0$ (blue line) compared to the lighter progenitor $s19.2$ (green line). Comparing with the bottom left panel of Figure 3 it can be seen that the compactness parameter (Equation (1)) is smaller for $s20.0$ (labeled by red arrow) and $s22.0$ (by blue arrow) in the chosen mass range above 3. The smaller compactness is translated into smaller mass accretion rate onto the stalled bounce shock. For model $s20.0$ (red line in the top panel), the relatively earlier shock revival ($\sim 100$ ms postbounce) coincides with the sharp decline of the accretion rate (dashed red line). After that, the accretion rate gradually decreases to $\sim 0.1$ $M_\odot$ s$^{-1}$ till $t_{400} = 420$ ms at this time the revived shock has expanded to an average radius of $\langle r \rangle = 400$ km. Here $t_{400}$ is a useful measure to qualify the vigor of the shock revival (e.g., Hanke et al. 2012). On the other hand, high compactness (model $s21.0$, black arrow in the bottom panel) leads to the high accretion rate (black dashed line in the top panel) and it takes $\sim 500$ ms for the sloshing shock (black line in the top panel) to gradually turn into a pronounced expansion later on ($t_{700} = 700$ ms). The growing diagnostic explosion energy (top right panel) and the almost converged mass of PNS (bottom right) are also shown as a function of post-bounce time. The PNS mass has a clear correlation with $\xi_{2.5}$. Regarding the diagnostic energy, the increasing rate ($E_{\text{dia}}/t_{\ln} - t_{400}$) in unit of $10^{51}$ erg s$^{-1}$) tends to become higher for models with high $\xi_{2.5}$ (0.754, 1.12, and 1.44, for $s22.0$, $s23.0$ and $s24.0$, respectively). This also indicates that the diagnostic energy of the high-$\xi$ models might become higher later on. Note that in previous 1D studies with simplified neutrino heating and cooling scheme (O’Connor & Ott 2011) or with the excision in-
side the PNS (Ugliano et al. 2012), the relation between the compactness and these explosion properties could not be determined in a self-consistent manner.

We investigate the postbounce shock evolution in more details for two specific models, s11.2 and s15.0. The density profile of model s11.2 falls rapidly off with radius and the compactness parameter $\xi_{2.5}=0.005$ is one of the smallest values among the 101 progenitors, whereas model s15.0 has moderate compactness ($\xi_{2.5}=0.145$) casting more extended envelope out to the iron core. Color-coded panels in Figure 1 show the differences in the postbounce evolution for s11.2 (left) and s15.0 (right), respectively. Model s11.2 explodes rather early ($t_{400}=150$ ms, top panels) and convective activity as well as the oscillations of the shock is moderate before the onset of the explosion (see the bottom left panel). Note in the bottom panels that the anisotropic velocity $v_{\text{aniso}}$ (upper), and the normalized pressure perturbation $\Delta p$ (lower), is defined respectively as $v_{\text{aniso}} = \sqrt{\langle (\rho v_r - \langle v_r \rangle)^2 + v_{\theta}^2 \rangle / \langle \rho \rangle}$ and $\Delta p = \sqrt{\langle p^2 \rangle - \langle p \rangle^2 / \langle p \rangle}$, where $\langle A \rangle$ represents the angle average of quantity $A$ (e.g., Takiwaki et al. 2012, Kuroda et al. 2012).

For model s15.0 with higher compactness parameter, sloshing motions of the shock are more clearly visible (right panels of Figure 4). These features regarding the dominance of the SASI over neutrino-driven convection for high-ξ models are quantitatively consistent with those observed in Müller et al. (2012). Near at the bottom of the gain region, the accreting flows receive an abrupt deceleration, below which the regions are convectively stable. A strong pressure perturbation forms there (seen, in the lower part of the bottom-right panel in Figure 4 as a boundary between the regions colored by white and deep blue at a radius of $r \lesssim 100$ km after $t_{\text{pb}} \sim 100$ ms). Subsequently the pressure perturbations propagate outward before they hit the shock, which is seen as a up-going stripe in the $\Delta p$ plot. This leads to the formation of the next vortices (e.g., yellow regions behind the shock in the $v_{\text{aniso}}$ plot). These features, as previously identified, are natural outcomes of the SASI and neutrino-driven convection (e.g., Fernández et al. 2014 and references therein). When the residency timescale becomes enough long, due to these multi-D effects, compared to the neutrino-heating timescale in the gain region, the runaway shock expansion initiates at $t_{\text{pb}} \sim 500$ ms ($t_{400} \sim 600$ ms) for this model. Such features have been also presented in previous 2D simulations with more detailed transport scheme (e.g., Müller et al. 2012, Bruenn et al. 2013). Here we emphasize that the use of the leakage scheme, together with the omission of inelastic neutrino scattering on electrons and GR effects in the present scheme, is likely to facilitate artificially easier explosions (see discussions in Takiwaki et al. 2014 for more details). Discrepancies between our Newtonian models and GR models with detailed transport (e.g., Müller et al. 2012) become remarkable especially for progenitors with high compactness (e.g., compare model s27.0 in Figure 2 with the bottom right panel of Figure 1 in Müller et al. 2013). For models with moderate compactness parameter, on the other hand, we like to note that some key hydrodynamic features in the post-bounce phase (such as the onset time of shock revival $t_{100}$) are rather similar (surprisingly) between ours versus the Garching models ($t_{400} = 640$ vs. 580 ms for model s18.4, 320 vs. 400 ms for s19.6, 540 vs. 560 ms for s21.6, 460 vs. 460 ms for s22.4, 420 vs. 440 ms for s27.0, respectively, e.g., Janka et al., TAUP Conference, 2013).

In Figure 5, we plot various quantities to summarize our results for the 101 simulations as a function of $\xi_{2.5}$. In the left panels, the mass accretion rate $\dot{M}$ at 500 km, electron-type neutrino luminosity $L_{\nu_e}$, and the shock revival time $t_{100}$ are shown. Three panels in the right column show a diagnostic explosion energy $E_{\text{dia}}$, a mass of central remnants $M_{\text{PNS}}$, and mass of nickel in the ejected material $M_{\text{Ni}}$, at the final time $t = t_{\text{fin}}$ of our simulations. Here the PNS is defined by the region where the density $\rho > 10^{11}$ g cm$^{-3}$. For most of our models $t_{\text{fin}}$ is defined as the time when the maximum shock radius arrived at the outer boundary.

Ugliano et al. (2012) showed that these quantities are not a monotonic function of ZAMS mass. We confirm it in our 2D simulations and find that these values are nearly in the linear correlation with the compactness parameter $\xi_{2.5}$. This can be interpreted as follows: the core of high-ξ models is surrounded by high-density Si/O layers and the mass accretion rate therefore remains high long after the stalled shock has formed (Figures 5(a) and 3). This makes the PNS mass of the high-ξ models heavier (Figure 5(c)). Due to the high accretion rate, the accretion neutrino luminosities become higher for models with high ξ (Figure 5(b)), see also O’Connor & Ott (2013). As a result, we obtain the high diagnostic explosion energy (Figure 5(d)), as well as outgoing unbound material rich in nickel (Figure 5(f)) for high-ξ models. On the other hand, the time of shock revival ($t_{400}$, Figure 5(c)) shows a large scatter and weaker correlation with $\xi_{2.5}$. This may come from the stochasticity of the initial random perturbations that should affect the onset of an explosion. It should be noted that the diagnostic energy of almost all models, which ranges in $\sim 0.2 - 0.7 \times 10^{51}$ erg, is still increasing at the final time of our simulation.

4. CONCLUSIONS AND DISCUSSION

We have conducted a systematic study of neutrino-driven CCSNe in 2D to explore the difference that the structure of progenitors makes on the explosion characteristics, such as mass of central remnants. We find, as Ugliano et al. (2012) did in 1D, that these characteristics exhibit a non-monotonic variability as a function of the progenitor ZAMS mass, which is a consequence of the non-monotonocities of the progenitor structure. Thus, we adopted the compactness parameter ξ introduced by O’Connor & Ott (2011) to characterize the progenitor structure. We find that the accretion luminosity becomes higher for models with high ξ, which makes the diagnostic energy higher and the synthesized nickel mass bigger. We point out that these explosion characteristics tend to show a monotonic increase as a function of the compactness parameter ξ.

Our simulations are limited in space ($r < 5000$ km) and time ($t \leq 1.5$ s). The simulations are terminated before the diagnostic energies are saturated. Later on neutrino energy deposition would get smaller with time as the neutrino luminosity as well as post-shock density becomes smaller. Further global simulation, taking account of gravitational energy of an envelope and nuclear energy
released via recombination process behind the shock, is necessary to determine the final explosion energy and the quantities shown in the right panels of Figure 5. Moreover, the finding of this study should be reexamined by 3D models ([Hanke et al. 2012; Dolence et al. 2013; Couch 2013; Couch & O’Connor 2014; Nakamura et al. 2014], in which neutrino transport is appropriately solved (see, e.g., [Hanke et al. 2013; Takiwaki et al. 2014; Nagakura et al. 2014; Mezzacappa et al. 2014] for different levels of sophistication). To get a more accurate amount of the synthesized nickel and other isotopes, a post-process calculation with a larger nuclear network is needed. After all, long-term 3D full-scale simulations are needed to unambiguously clarify the critical $\xi$ parameter, below or above which neutron stars or black holes will be left behind.

In this work we have reported our results of progenitors from one modeling group. Currently we are conducting the same sort of CCSN simulations for sets of progenitors from the other groups including a variety of metallicity, rotation, and magnetic fields, which will be presented in the forthcoming work. We hope that the compilation of our results would be helpful to many researchers to study nucleosynthetic yields from CCSNe and its impacts on chemical and dynamical evolution of galaxies, and also detectability of neutrino and gravitational-wave signals by next-generation detectors.

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REFERENCES

Bruenn, S. W., Mezzacappa, A., Hix, W. R., et al. 2013, ApJ, 767, L6
Buras, R., Rampp, M., Janka, H.-T., & Kifonidis, K. 2006, A&A, 447, 1049
Burrows, A., Hayes, J., & Fryxell, B. A. 1995, ApJ, 450, 830
Couch, S. M. 2013, ApJ, 775, 35
Couch, S. M., & Ott, C. D. 2013, ApJ, 778, L7
Couch, S. M., & O’Connor, E. P. 2014, ApJ, 785, 123
Dolence, J. C., Burrows, A., Murphy, J. W., & Nordhaus, J. 2013, ApJ, 765, 110
Dolence, J. C., Burrows, A., & Zhang, W. 2014, arXiv:1403.6115
Einfeldt, B. 1988, SIAM Journal on Numerical Analysis, 25, 294
Foglizzo, T., Scheck, L., & Janka, H.-T. 2006, ApJ, 652, 1436
Hanke, F., Marek, A., Müller, B., & Janka, H.-T. 2012, ApJ, 755, 138
Hanke, F., Müller, B., Wongwathanarat, A., Marek, A., & Janka, H.-T. 2013, ApJ, 770, 66
Hayes, J. C., Norman, M. L., Fiedler, R. A., et al. 2006, ApJS, 165, 188
Heger, A., Langer, N., & Woosley, S. E. 2000, ApJ, 528, 308
Heger, A., Woosley, S. E., & Spruit, H. C. 2005, ApJ, 626, 350
Janka, H.-T., & Mueller, E. 1996, A&A, 306, 167
—. 2012. Annual Review of Nuclear and Particle Science, 62, 407
Kitaura, F. S., Janka, H.-T., & Hillebrandt, W. 2006, A&A, 450, 345
Kuroda, T., Kotake, K., & Takiwaki, T. 2012, ApJ, 755, 11
Lattimer, J. M., & Swesty, F. D. 1991, Nuclear Physics A, 535, 331
Liebendörfer, M., Whitehouse, S. C., & Fischer, T. 2009, ApJ, 698, 1174
Figure 5. Resultant supernova properties from our 101 simulations as a function of compactness parameter $\xi_{2.5}$. Left: Mass accretion rate (a), and electron neutrino luminosity (b), estimated at time of shock revival $t_{400}$ (c). Right: diagnostic energy (d), mass of proto-neutron star (e), and mass of nickel in outgoing unbound material (f), at $t = t_{\text{fin}}$. Failed models which cannot carry the shock to the outer boundary during our simulation time are excluded from these panels.

Limongi, M., & Chieffi, A. 2006, ApJ, 647, 483
Marek, A. & Janka, H.-T. 2009, Astrophys. J., 694, 664
Müller, B., Janka, H.-T., & Heger, A. 2012, ApJ, 761, 72
Müller, B., Janka, H.-T., & Marek, A. 2013, ApJ, 766, 43
Murphy, J. W., Dolence, J. C., & Burrows, A. 2013, ApJ, 771, 52
Mezzacappa, A., Bruenn, S. W., Lentz, E. J., et al. 2014, arXiv:1406.7075
Nagakura, H., Sumiyoshi, K., & Yamada, S. 2014, arXiv:1407.5632
Nakamura, K., Takiwaki, T., Kotake, K., & Nishimura, N. 2014, ApJ, 782, 91
Nakamura, K., Kuroda, T., Takiwaki, T., & Kotake, K. 2014, arXiv:1403.7290
Nomoto, K., & Hashimoto, M. 1988, Phys. Rep., 163, 13
Ott, C. D., Burrows, A., Dessart, L., & Livne, E. 2008, ApJ, 685, 1069
O’Connor, E., & Ott, C. D. 2011, ApJ, 730, 70
O’Connor, E., & Ott, C. D. 2013, ApJ, 762, 126
Fernández, R., Müller, B., Foglizzo, T., & Janka, H.-T. 2014, MNRAS, 440, 2763
Sukhbold, T., & Woosley, S. E. 2014, ApJ, 783, 10
Sumiyoshi, K., Takiwaki, T., Matsufuru, H., & Yamada, S. 2014, arXiv:1403.4470
Suwa, Y., Kotake, K., Takiwaki, T., et al. 2010, PASJ, 62, L49
Suwa, Y., Yamada, S., Takiwaki, T., & Kotake, K. 2014, arXiv:1406.6414
Takiwaki, T., Kotake, K., & Suwa, Y. 2012, ApJ, 749, 98
Takiwaki, T., Kotake, K., & Suwa, Y. 2014, ApJ, 786, 83
Ugliano, M., Janka, H.-T., Marek, A., & Arcones, A. 2012, ApJ, 757, 69
Woosley, S. E., & Heger, A. 2007, Phys. Rep., 442, 269
Woosley, S. E., Heger, A., & Weaver, T. A. 2002, Reviews of Modern Physics, 74, 1015
Woosley, S. E., & Weaver, T. A. 1995, ApJS, 101, 181