Massive star evolution and nucleosynthesis: Lower end of Fe-core-collapse supernova progenitors and remnant neutron star mass distribution

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In order to explore various aspects of stellar evolution, supernovae, gamma ray bursts, and nucleosynthesis, we have developed a new efficient stellar evolution code. In this paper, we describe this new code and compare the results with those calculated by the previous code. Specifically, we focus on the progenitor evolution of the lower end of Fe-core-collapse supernovae, and the mass distribution of remnant neutron stars. We describe how different assumptions will lead to different neutron star mass distributions. We also review some of the recent work by our research group.

1. Introduction

Massive stars end their lives as supernovae (SNe), leaving neutron stars behind, or by forming black holes without explosions if they are not rotating. The critical initial stellar mass for black-hole formation is usually considered to be about 20 to 25 \( M_\odot \), but this is uncertain [1]. This is because the explosion mechanism for supernovae is not yet definitely known [2–7], and there are also still uncertainties in the progenitor models.

The lower end for neutron-star-forming supernovae is also still uncertain (e.g., Ref. [8]). Here, the uncertainties in the stellar evolution theory may be even larger. It is well known that stars above approximately 10 \( M_\odot \) form an Fe core at the end of their evolution. The formation of the Fe core is relatively simple for an \( M > 13M_\odot \) star (Ref. [9] for a review). In such a star, the hottest region is the center, and this is mostly the case all the way through its evolution. Therefore, the heaviest element is synthesized around the center, forming an Fe core in the end.

On the other hand, the evolution of \( M \lesssim 13M_\odot \) stars is more complicated. Less massive stars tend to show temperature inversion between the center and outside regions. This phenomenon is well known for intermediate-mass stars with \( M < 8M_\odot \) after carbon burning, for which the carbon burning starts off-center (e.g., Ref. [10]). For an \( M \lesssim 11M_\odot \) star, off-center O and/or Si burning occurs. The off-center burning front propagates inward through complicated burning stages, eventually forming an Fe core. This off-center burning sometimes becomes very violent and could cause some mass ejection [11,12] (see also Sect. 3.1 in this paper). However, such calculations have not been done recently with updated input physics, so it is currently not clear if such mass ejection really occurs.

For less massive stars that do not ignite Ne even off-center, a cool degenerate O–Ne core is first formed. This O–Ne core grows in mass gradually when the star is in a super AGB phase. Depending on the mass-loss rate, this O–Ne core reaches the critical mass for core collapse as a result of
electron capture \cite{8,13}. It is considered that such a star explodes by the neutrino energy transport mechanism \cite{14}.

Calculations for the progenitors of such low-mass core-collapse SNe with updated input physics are interesting and important. However, their evolution can be quite different if the initial mass is only slightly different, say by 0.01–0.1 $M_\odot$ (e.g., Ref. \cite{8}; see also Sect. 3.1). To understand the whole story in this mass range, therefore, we need to calculate stellar evolution in very fine mass grids. This is quite time-consuming and, therefore, the currently well-used progenitor models in the literature (e.g., Refs. \cite{12,15–19}) do not fully deal with this mass range.

To tackle this problem, we have developed a stellar evolution code for efficient computation. In this paper, we describe this new code, the Yoshida–Umeda (YU) code \cite{20}, and compare the results with those calculated by one of the authors, H.U., using the Umeda–Nomoto (UN) code \cite{16,17,21,22}. We also briefly review some of the recent work using the YU code and other work in our research group.

This paper is organized as follows. In Sect. 2, we describe the new stellar evolution code and its differences from previous codes. In Sect. 3, massive star evolutions calculated with this code are compared with those calculated with a previous code. In this section we also describe the progenitor evolution of the lower end of Fe-core-collapse SNe in some detail. Section 4 describes nucleosynthetic aspects, including some reviews of our recent work. In Sect. 5, remnant neutron star masses are given as a function of progenitor mass. We describe how different assumptions will lead to different neutron star mass distributions. Section 6 reviews other recent work of our research group, and discussions and future prospects are given in Sect. 7.

2. Stellar evolution code

2.1. Umeda–Nomoto (UN) code

Before describing the YU code, we briefly describe the UN code because we will compare the results of these codes. This code is mostly based on the Nomoto–Hashimoto (NH) code \cite{12} and the Saio–Nomoto–Kato (SNK) code \cite{23}. In the UN code, the input physics such as the equation of state (EOS) are same as in the NH code, except for updates of radiative opacity, neutrino emissivity, and electron capture rates.

There are several differences between the NH and SNK codes:

i) The NH code is an He-star code, which means that it cannot solve stellar atmosphere and hydrogen-burning phases. Thus it cannot provide stellar radii correctly. On the other hand, the SNK code solves atmosphere.

ii) Since the SNK code is not designed for calculating the later evolutionary stages of massive stars, it does not deal with nuclear burning after the carbon-burning stages. Also, because of this, the SNK code does not include the acceleration term \cite{24}, and omits the inertial term in the equation of motion, while the NH code can include it. This acceleration term becomes important for constructing supernova progenitor models just before iron-core collapse.

iii) The treatment of convection is different. Both codes adopt the Schwarzschild criterion for convection but the method of mixing is different. The NH code assumes instantaneous mixing in convective regions, while, in the SNK code, matter is mixed diffusively using the formalism of Ref. \cite{25}, taking semi-convection effects into account. The energy transfer in a convective region is also different. In the NH code, time-dependent mixing length theory \cite{26,27} can be
included for convective energy transfer, though this effect is not so important for the massive star evolution forming an Fe core.

In the UN code, atmosphere is calculated as in the SNK code, the acceleration term is included, and the formalism of Ref. [25] is adopted for convective mixing. Convective energy transfer is treated as in the NH code.

The UN code differs from the NH and SNK codes in the calculations of nucleosynthesis and nuclear energy generation. The NH code assumes quasi-nuclear statistical equilibrium during Si burning while the UN code solves full nuclear reaction networks below \( \log_{10} T_C(K) \lesssim 9.6 \). Above that temperature both codes assume nuclear statistical equilibrium (NSE). In order to calculate nuclear energy generation rates, in the UN code nuclear reaction networks are solved simultaneously with Henyey relaxation, while in the NH and SNK codes the abundance is fixed during Henyey relaxation. In this sense the UN code solves abundance implicitly, while the NH and SNK codes solve it explicitly. Solving abundance implicitly is the best way to obtain consistency in the energy generation rates and the abundance evolution. However, this has a disadvantage in efficient calculations because solving large reaction networks involves time-consuming matrix inversion calculations.

2.2. Yoshida–Umeda (YU) code

The basic structure of this code is based on the SNK code [23]. As mentioned above, although the UN code has been successfully used for progenitor calculations, it has an disadvantage in the calculation time. To finish a calculation from ZAMS to Fe-core collapse typically takes a few months. Therefore this code is not suitable for a large-parameter search. Hence, we have developed a new, more efficient code (the YU code). The main differences between the YU code and the UN code are the calculations of nucleosynthesis and energy generation.

The YU code solves a full nuclear reaction network from hydrogen burning up to \( \log_{10} T_C(K) \sim 10 \). The nuclear reaction network consists of \( \sim 300 \) species of nuclei from \( n, p \) to Br. NSE is not assumed in the calculations. On the other hand, the nucleosynthesis is solved before Henyey relaxation in the YU code, like the SNK code. Using the abundance in the next step, the energy generation is calculated during the Henyey relaxation. Therefore, a smaller time step is required.

After the carbon-burning stage, the time interval is typically determined to satisfy the conditions \( \Delta \log_{10} T / \log_{10} T \leq 0.001 \) and \( \Delta \log_{10} \rho / \log_{10} \rho \leq 0.003 \) for each mass coordinate. This treatment requires more calculation steps but still saves calculation time, especially in late phases of the evolution. With this code, we can calculate one model in typically 1–2 weeks for massive star evolutions.

At present, this code does not yet implement the acceleration term. We will include this term in the near future.

2.3. Mass-loss rates

In the UN and YU codes, the same mass-loss rates are used for the solar metallicity before the Wolf–Rayet star phases shown in this paper. In OB stars, where the surface temperature is larger than \( 1.2 \times 10^4 \) K and the surface H mass fraction is larger than 0.4, the mass-loss rate from Ref. [28] is used. In yellow supergiants and red-giant branches, the mass-loss rate in Ref. [29] is used. In the UN code, the metallicity-dependent factor of \((Z/0.02)^{0.5}\) is multiplied to the rate for metal-poor stars as in Ref. [30]. The YU code uses the metallicity dependence in a different manner. The metallicity dependence in the main-sequence stage is taken from Ref. [28]. In the YU code presented in this
8.5
9
9.5
10

\( \log T_C \) [K]

\( \log \rho_c \) [g/cm\(^3\)]

12M\(_\odot\)
15M\(_\odot\)
20M\(_\odot\)
25M\(_\odot\)

Fig. 1. Evolution tracks of \( M = 12, 15, 20, \) and \( 25M_\odot \) stars in terms of central temperature and density. The mass fraction distributions of a \( 15M_\odot \) star denoted by the points \( a-e \) are shown in Fig. 2.

paper, the case A mass-loss rate in Ref. [20] is used. In yellow supergiants and red-giant branches, the metallicity-dependent factor \((Z/0.02)^{0.64}\) is multiplied to the mass-loss rate. The power index is the same as that of B supergiants in Ref. [28]. See Ref. [20] for details.

3. Progenitor evolution

The basic properties of massive star evolution have been discussed in many papers (e.g., Refs. [12,16–19]), so we do not repeat the details here, except for the \( M \simeq 10–13M_\odot \) models in the next subsection. We briefly describe the evolution in the advanced stages of an \( M = 15M_\odot \) star as an example. Figure 1 shows the evolution tracks of \( M = 12, 15, 20, \) and \( 25M_\odot \) stars in terms of the central density and temperature. We also show the mass fraction distributions of a \( 15M_\odot \) star at five stages during its evolution in Fig. 2. Figure 2(a) indicates the mass fraction distribution after core helium burning. There is a CO core of \( 2.2M_\odot \) in the central region. An He layer of \( 2M_\odot \) and an H-rich envelope surround the CO core. The He layer consists of ashes from the hydrogen burning. After core helium burning, shell helium burning produces \( ^{12}\text{C} \) in the He layer. Then, the carbon ignites at the center when the central temperature becomes \( \log_{10} T_C(K) \sim 8.8 \). The core carbon burning converts \( ^{12}\text{C} \) into \( ^{20}\text{Ne} \) and \( ^{23}\text{Na} \) to form an O/Ne core [Fig. 2(b)]. Some Mg–Si is also produced in the core.

Shell carbon burning is followed by core carbon burning. This converts the O/C layer into an O/Ne layer. However, \( ^{12}\text{C} \) of 0.09 by mass fraction remains even after shell carbon burning. When the central temperature becomes \( \log_{10} T_C(K) \sim 9.2 \), core neon burning starts and an O/Si core forms [Fig. 2(c)]. After Ne burning, oxygen ignites at the center with a central temperature of \( \log_{10} T_C(K) \sim 9.3 \). Core oxygen burning produces a \( \sim 1M_\odot \) Si core [Fig. 2(d)]. The main components of the Si core are \(^{28}\text{Si}, ^{32}\text{S}, \) and \(^{34}\text{S} \), and the electron fraction decreases by electron capture. The Si core extends through the subsequent shell oxygen burnings and the central temperature increases by contraction. When the central temperature is \( \log_{10} T_C(K) \sim 9.55 \), Si burning starts and an Fe core forms [Fig. 2(e)]. The composition in the O layer also gradually changes through shell burnings. After Fe core formation, the Fe core extends with shell Si burning and finally collapses to explode as a supernova.

Here we stress that one of the most important factors in massive star evolution is the carbon abundance after core helium burning, \( X_C(^{12}\text{C}) \). In general, a smaller carbon abundance leads to a larger
Fig. 2. Mass fraction distributions of an $M = 15M_\odot$ star after He burning (a), C burning (b), Ne burning (c), O burning (d), and Si burning (e). The final mass is $M_f = 13.6M_\odot$. “Si” and “Fe” indicate the sum of the mass fractions of the elements Si–Sc and Ti–Br.

iron core at core collapse. This is because, when the carbon abundance is smaller, shell carbon burning is weaker after helium burning. Then the core is less supported by the convective carbon shell burning and the time between the central helium burning to the oxygen burning is shorter. As a result, for a smaller carbon abundance a larger core with higher entropy is formed because neutrino cooling is less effective \[11,12\].

The carbon abundance also affects nucleosynthesis substantially. When the abundance is too low, carbon-burning products such as Ne, Na, Mg, and Al are underproduced compared with the solar abundance pattern, while they are overproduced when the abundance is too large \[31\]. They also found that the abundances of S, Ar, and Ca are anti-correlated with Ne, Na, Mg, and Al.

Despite these important factors, the carbon abundance is very uncertain because it sensitively depends on the $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ rate and the treatment of convection. Reference \[31\] discussed the
Table 1. Carbon mass fraction after core He burning for UN and YU models.

| UN model | | | | | | |
|---|---|---|---|---|---|---|
| $M (M_{\odot})$ | 11.5 | 12 | 13 | 15 | 20 | 22 | 25 |
| $X_C^{(12C)}$ | 0.36 | 0.36 | 0.33 | 0.36 | 0.34 | 0.28 | 0.26 |

| YU model | | | | | | |
|---|---|---|---|---|---|---|
| $M (M_{\odot})$ | 11 | 12 | 13 | 15 | 18 | 20 | 25 |
| $X_C^{(12C)}$ | 0.26 | 0.25 | 0.31 | 0.24 | 0.23 | 0.23 | 0.20 |

Table 2. Final, core, and remnant masses for UN, YU, and NH88 models.

| UN model | | | | | | |
|---|---|---|---|---|---|---|
| $M (M_{\odot})$ | 11.5 | 12 | 13 | 15 | 20 | 22 | 25 |
| $M_f (M_{\odot})$ | 11.2 | 11.6 | 12.7 | 13.9 | 17.5 | 17.6 | 17.7 |
| $M_{\text{He}} (M_{\odot})$ | 2.8 | 3.0 | 3.3 | 4.0 | 5.3 | 6.7 | 7.9 |
| $M_{\text{CO}} (M_{\odot})$ | 1.62 | 1.72 | 1.95 | 2.51 | 3.91 | 5.42 | 6.54 |
| $M_{\text{rem}} (M_{\odot})$ | 1.40 | 1.43 | 1.45 | 1.51 | 1.57 | 1.66 | 1.70 |
| $M_g (M_{\odot})$ | 1.26 | 1.29 | 1.30 | 1.35 | 1.40 | 1.46 | 1.50 |

| YU model | | | | | | |
|---|---|---|---|---|---|---|
| $M (M_{\odot})$ | 10 | 11 | 12 | 13 | 15 | 18 | 20 |
| $M_f (M_{\odot})$ | 9.5 | 10.5 | 11.4 | 12.1 | 13.6 | 16.2 | 17.6 |
| $M_{\text{He}} (M_{\odot})$ | 2.6 | 2.8 | 3.1 | 3.2 | 4.2 | 5.4 | 6.2 |
| $M_{\text{CO}} (M_{\odot})$ | 1.47 | 1.61 | 1.81 | 1.87 | 2.64 | 3.58 | 4.34 |
| $M_{\text{rem}} (M_{\odot})$ | 1.29 | 1.32 | 1.44 | 1.45 | 1.64 | 1.77 | 1.96 |
| $M_g (M_{\odot})$ | 1.18 | 1.19 | 1.29 | 1.30 | 1.45 | 1.55 | 1.69 |

| NH88 model | | | | | | |
|---|---|---|---|---|---|---|
| $M (M_{\odot})$ | ∼13 | ∼15 | ∼20 | ∼25 |
| $M_{\text{He}} (M_{\odot})$ | 3.3 | 4 | 6 | 8 |
| $M_{\text{rem}} (M_{\odot})$ | 1.27 | 1.33 | 1.61 | 1.77 |
| $M_g (M_{\odot})$ | 1.15 | 1.20 | 1.43 | 1.55 |

The fact that, due to the nucleosynthetic constraint described above, in their model the $^{12}$C($\alpha, \gamma$)$^{16}$O rate needs to be $1.7 \pm 0.5$ times the Caughlan and Fowler (1988) (CF88 hereafter) rate [32]. In the NH88 model [12], the rate of Ref. [33] was taken, which corresponds to 2.3–2.4 times the CF88 rate. In NH88, $X_C^{(12C)} = (0.25, 0.22, 0.19)$ for the helium star model for $M_{\text{He}} = (3.3, 6, 8) M_{\odot}$, which roughly corresponds to the ZAMS $M = (13, 20, 25) M_{\odot}$, respectively. In the UN model the $^{12}$C($\alpha, \gamma$)$^{16}$O rate was chosen to be 1.3 times larger than the CF88 rate, so that the abundance patterns of EMP and VMP stars are reproduced [22,34].

Although there have been improvements in the estimation of the $^{12}$C($\alpha, \gamma$)$^{16}$O rate [35,36], uncertainty in the rate is still about a factor of two. Therefore, the nucleosynthetic method is still the best way to constrain the carbon abundance.

We show in Table 1 the central carbon abundance after core helium burning for the UN and YU models. In the YU model, the $^{12}$C($\alpha, \gamma$)$^{16}$O rate is chosen to be 1.5 times as large as the CF88 rate so that the carbon abundance $X_C^{(12C)} = 0.20$ for the $M = 25 M_{\odot}$ model. As shown in the table, mainly because of the difference in the $^{12}$C($\alpha, \gamma$)$^{16}$O rate, the YU models presented in this paper have systematically smaller carbon abundances than the UN models. In Table 2, we also show the final, $M_f$, He core, $M_{\text{He}}$, and CO core masses, $M_{\text{CO}}$. Here the latter two masses are defined where the mass fractions of H and He are less than 0.001, respectively.
Table 3. Ignition properties of 10–13 $M_\odot$ stars.

| Mass   | 10 $M_\odot$ | 11 $M_\odot$ | 12 $M_\odot$ | 13 $M_\odot$ |
|--------|--------------|--------------|--------------|--------------|
| C ignition at the center or not central composition | C–O | C–O | C–O | C–O |
| O ignition at the center or not central composition | off-C | off-C | C | C |
| Ne ignition at the center or not central composition | O–Ne | O–Ne | O–Ne | O–Ne |
| Si ignition at the center or not central composition | off-C | off-C | off-C | off-C |

Fig. 3. Evolution tracks of $M = 9.82, 9.89, 10, 11, \text{ and } 12 M_\odot$ stars in terms of central temperature and density.

3.1. Evolution of 10–13 $M_\odot$ stars

Here we describe the evolution of 10–13 $M_\odot$ stars calculated by the YU code, since these stars correspond roughly to the lightest Fe core collapse model, and their evolution after oxygen burning is quite different from more massive $M > 13 M_\odot$ stars. Table 3 shows the differences in the ways the stars commence nuclear burning of Ne, O, and Si. In the table ‘C’ represents the case that the ignition occurs at the center of the core and ‘off-C’ shows that the ignition occurs off-center. In the central composition line, the letters represent the main nuclei at the center. The differences also appear in Fig. 3, in which evolutional tracks are shown in terms of the central temperature and density. In this mass range, a star forms a CO core around the Chandrasekhar limiting mass. The degeneracy of the core is the important property and affects its evolution.

We will present a more detailed description in Takahashi et al. (manuscript in preparation, see also Ref. [37]), in which stellar evolutions are calculated in very fine grids. So far, we have found that the $M = 9.89 M_\odot$ model ends up with Fe core formation, though the $M = 9.82 M_\odot$ model does not form an Fe core. As described below, the off-center Si burning is more and more violent for less massive stars. In this mass range, only a 0.01 $M_\odot$ difference in the zero-age main-sequence mass would result in quite different evolution. However, the global evolutionary properties and final density structure of an Fe core are similar to the 10 $M_\odot$ model presented here as far as Fe core formation is concerned.
Fig. 4. Mass fraction distribution for a 10 $M_\odot$ star during a shell Ne+O burning, when $\log_{10} \rho_C = 7.9358$.

$10 M_\odot$ star: For a 10 $M_\odot$ star, the evolution from the zero-age main-sequence to the C-burning phase proceeds roughly in the same way as more massive stars that end up with Fe core formation. However, after O–Ne core formation through C burning, different evolutionary aspects come about.

Because the mass of the C–O core of 1.47 $M_\odot$ is only slightly larger than the critical mass for Ne ignition, the O–Ne core first contracts without nuclear burning. Eventually, the importance of the pressure of degenerate electron increases. Such semi-degenerate cores often show inverse temperature distributions; temperature does not take its highest value at the center. This is due to the property of degenerate electrons. Since neutrino cooling is more effective at higher density as well as at higher temperature, the O–Ne core loses its heat by the cooling. As heat escapes, the core contracts and compressional heating supplies energy to the gas of ions, electrons, and radiation. If the core is supported by pressure of ions, the gas temperature increases as a result of compressional heating. On the other hand, if electron pressure supports the core, it is hard to increase the gas temperature because the internal energy of degenerate electrons does not depend on temperature but on density. What the high-density O–Ne core supports mainly is the pressure of degenerate electrons, so the denser the region is, the slower the temperature increase is. Owing to this temperature inversion, a shell Ne burning commences at an off-center region with a higher temperature than the center. Since Ne and O have close ignition temperatures, once the shell Ne burning commences, heat generation increases the temperature and the shell O-ignition succeeds. Then a degenerate O–Ne core is surrounded by a hot Si-shell (Fig. 4). These evolutionary properties are fully consistent with previous work (e.g., Refs. [8,12]).

Some of the nuclei of the Si cluster produced by the O burning can capture electrons in such a high density O–Ne core. Since the core is mainly supported by the pressure of degenerate electrons, a decrease in the number of electrons causes core contraction. Heat production by the contraction increases the temperature. When the temperature at the base of the Si shell becomes high enough to ignite Si, the shell Si burning commences and transforms the Si shell into an Fe shell (Fig. 5). The energy production rate of the burning from Ne to Fe is very high. So at every ignition, the core repeats the expansional cooling and compressional heating almost adiabatically, and the burning front approaches the center little by little. After the burning front reaches the center of the core, an Fe core forms. Then the star collapses. Some vertical ascents in Fig. 3 show moments when a base

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1 This is seen in spikes around $\log_{10} \rho_C \sim 8$ in Fig. 3.
Fig. 5. Mass fraction distribution for a 10 $M_\odot$ star during a shell Ne+O+Si burning, when $\log_{10} \rho_C = 8.0958$.

Fig. 6. Mass fraction distribution for an 11 $M_\odot$ star during a shell Ne+O burning, when $\log_{10} \rho_C = 7.6181$.

of an off-center burning reaches the center of the core. For a 10 $M_\odot$ star, the central temperature suddenly increases at $\log_{10} \rho_C = 8.4319$, when the shell Ne+O+Si burning reaches the center.

3.2. 11 $M_\odot$ star

For an 11 $M_\odot$ star with a larger C–O core of 1.61 $M_\odot$ compared to a 10 $M_\odot$ star, the electron degeneracy of the core is less. This causes a difference in evolution after Ne and O ignite off-centrally.

Because there is less degeneracy at the center, the difference between the maximum and central temperatures is less, so a shell O+Ne ignition commences at a deeper region than in the 10 $M_\odot$ star (Fig. 6). The temperature at the base of the shell increases by compressional heating in the same way as the 10 $M_\odot$ star during O+Ne-burning front propagation. However, the burning front arrives at the center before the base temperature reaches Si ignition and, thus, an Si core forms. This is the main difference between 10 and 11 $M_\odot$ stars.

Later, as the Si core contracts, the outer O+Ne shell burning continues moving its base outward, while neutron-rich nuclei such as $^{36}\text{S}$ and $^{50}\text{Ti}$ are produced in the inner region. The degeneracy of electrons also increases in the contracting core and the above-mentioned temperature inversion appears. When the central density reaches a value of $\log_{10} \rho_C = 8.9867$, the first off-center Si ignition occurs at $M_r = 0.5933M_\odot$. This ignition causes the core to expand adiabatically. After the burning terminates, the core contracts adiabatically again. When the central density becomes $\log_{10} \rho_C = 9.0253$, the second off-center Si ignition occurs at $M_r = 0.1103M_\odot$ (Fig. 7). After the second off-center Si ignition, the burning front reaches the center. The core transforms into an Fe core, then collapses.
Fig. 7. Mass fraction distribution for an $11 \, M_\odot$ star during a shell Si burning, when $\log_{10} \rho_c = 8.0481$.

Fig. 8. Mass fraction distribution for a $12 \, M_\odot$ star during a shell Si burning, when $\log_{10} \rho_c = 7.9694$.

3.3. $12 \, M_\odot$ star

A $12 \, M_\odot$ star evolves in a similar way to an $11 \, M_\odot$ star but the properties of Ne and O ignitions are different. Because it has a larger C–O core of $1.85 \, M_\odot$, the temperature takes its highest value at the center of the O–Ne core. Therefore, Ne and O ignite at the center. Note that Table 3 shows that when O ignites at the center, the central composition is O–Si, and not O–Ne. This is just because Ne has burned ahead of O, producing an O–Si core.

After Si core formation, the evolution proceeds in the same way as an $11 \, M_\odot$ star. Neutron-rich nuclei are produced around the center after several shell Ne$^+$O burnings in the outer regions of the contracting core. Then off-center Si burning occurs (Fig. 8), and after the front reaches the center an Fe core forms.

3.4. Final mass and metallicity

In Fig. 9 we show the metallicity dependence of the final and core masses of the YU models [20]. We see a clear metallicity dependence in the final mass among the stars with $M \gtrsim 20M_\odot$. In $Z = 0.02$, an $M = 30M_\odot$ star has the maximum final mass $M_f = 21.3M_\odot$. $M \gtrsim 40M_\odot$ stars indicate a roughly constant final mass of $\sim 10M_\odot$. These stars lose all their H and He layers and become Wolf–Rayet stars. The effect of mass loss in $Z = 0.01$ stars is less than that in $Z = 0.02$ stars. Increasing the main-sequence mass, the final mass also increases in $M \leq 30M_\odot$ and $40 \leq M \leq 80M_\odot$ but it decreases in $30 \leq M \leq 40M_\odot$ and $M \geq 80M_\odot$. Metal-poorer stars have a similar $M$ dependence of the final mass. In this figure, the zero-age main-sequence mass, where the mass loss becomes effective, and
the final mass for a given main-sequence mass become larger for metal-poorer stars. We do not see a clear mass-loss effect among $M \leq 80 M_\odot$ stars with $Z = 10^{-4}$.

The metallicity dependence of the He core is small [Fig. 9(b)]. On the other hand, most of the He layer has been lost during the evolution in $M \geq 40, 50$, and $100 M_\odot$ stars for $Z = 0.02, 0.01$, and $0.004$. These stars lose their He/C/O envelope and their surface He abundance decreases during the He burning. They become C- and O-enriched Wolf–Rayet stars (WO stars) with a surface He mass fraction of $\sim 0.2$. The mass of the He-rich shell is $1.4 - 5.5 M_\odot$ in $Z = 10^{-4}$ stars, in which the mass-loss effect is small. Since the fraction of the He layer is small, mass loss brings about the removal of the He layer rather than a reduction in the He core mass.

Fig. 9. The final mass (panel (a)), the He core mass (panel (b)), and the CO core mass (panel (c)) with the relation to the main-sequence mass. The adopted metallicities are $Z = 10^{-4}$, 0.001, 0.004, 0.01, and 0.02. In panel (b), stars more massive than the attached number of the main-sequence mass ($M_\odot$) evolve to WO stars. The surface He mass fraction of the stars is about 0.2.
Fig. 10. Mass fraction distribution of the supernova progenitors evolved from 15 (a) and 20 $M_\odot$ (b) stars with the metallicity of $Z = 0.02$. The corresponding final masses are 13.6 and 17.6 $M_\odot$. The mass fraction distribution in the H-rich envelope outside the range of each figure is the same as that at the outer edge of the figure.

The metallicity dependence of the CO core mass is seen in $Z > 0.004$ and $M \gtrsim 40 - 50M_\odot$ in Fig. 9(c). The CO core mass is roughly constant with $\sim 10M_\odot$ in $Z = 0.02$ and $M \geq 40M_\odot$ stars. In $Z = 0.01$, the maximum mass of the CO core is $22M_\odot$ of an $M = 80M_\odot$ star. On the other hand, the CO core mass can be more massive in metal-poorer stars. The CO core mass is $35 - 40M_\odot$ in $Z \leq 0.004$ and $M = 100M_\odot$.

4. Nucleosynthesis

Nucleosynthesis in these massive stars occurs mainly in two stages. One is before core collapse and the other is during supernova explosion. It is well known that, inside these massive stars, nuclear fusion takes place up to Fe synthesis. Before gravitational core collapse, Fe core is produced at the center and lighter elements form onion-like structures from inside to the surface (Fig. 10).

Nucleosynthesis during a supernova explosion is usually called explosive nucleosynthesis. During the explosion, a shock wave propagates out of the Fe core to the stellar surface. Behind the shock wave, matter heats up and nuclear burning takes place. Explosive nucleosynthesis is important for the synthesis of Si and heavier elements (e.g., Ref. [16]).

In our previous work (e.g., Refs. [16,17,22,38]), we use a simple 1D model for supernova explosions to calculate explosive nucleosynthesis. We inject thermal or kinetic energy below the mass cut or just above the Fe core to initiate supernova shock. We may call this procedure ‘instant energy injection’.
We use a 1D PPM code for the hydrodynamical calculations and solve small alpha-networks together to calculate nuclear energy generation. Then we calculate the detailed nucleosynthesis by post-processing by solving a large nuclear reaction network.

Explosive nucleosynthesis should depend on how the star explodes, but we do not yet know how gravitational collapse leads to the explosion. Therefore, there are still several proposals for successful supernova explosions. Fortunately, the properties of supernova shock outside the Fe core are almost independent of the explosion mechanism, because in most models the supernova shock gains energy inside the Fe core. Then, as far as spherical symmetry is assumed, explosion energy $E$ is the only parameter that determines the properties of the supernova shock. For example, the kinetic to thermal energy ratio does not particularly affect the propagation of the shock outside the Fe core, because it quickly converges into the same solution. In this case, the ‘instant energy injection’ method is sufficient to calculate nucleosynthesis. For most cases, nucleosynthesis of Fe peak elements including Zn and lighter elements can be safely calculated with this method. In Sect. 7, we will discuss the limitations of this simple procedure.

Behind the shock front, temperature distribution is almost homogeneous, and the radiation energy is strongly dominant. Then the radiation temperature $T$ can be related to the explosion energy, $E$, by $E = (4\pi/3)a r^3 T^4$, where $a$ is the radiation constant and $r$ is the shock radius. Thus the maximum temperature after shock passage, $T_s$, for the mass elements located at radius $R$ in the progenitor is approximately given as

$$T_s = \left(\frac{3E}{4\pi a R^3}\right)^{1/4}. \quad (4.1)$$

For a given progenitor model and explosion energy $E$, the propagation of shock and thus the time evolution of density distribution is also determined uniquely. Behind the shock front, pressure and entropy are also radiation-dominated and can be written as $P_\gamma = aT^4/3$ and $S_\gamma = (4/3)aT^3/\rho$. Since the shocked region cools roughly adiabatically at first, the radiation entropy is often convenient to specify the explosion.

Figure 11 shows the mass fraction distribution of the innermost region of the SN ejecta evolved from 15 $M_\odot$ and 20 $M_\odot$ stars with a metallicity of $Z = 0.02$. In the innermost region of supernova ejecta where $T_s$ is higher than $5 \times 10^9$ (K), Si burns completely. In such a region, Si mostly decomposes to alpha-particles at first; this is an endothermic reaction. Then, as the temperature decreases, the alpha-particles start to recombine and alpha-rich freezeout nucleosynthesis takes place. This phase is exothermic. The energy changes due to these processes are typically about 10%, and thus the equation above (4.1) still provides a useful approximation.

In this complete Si-burning region, $^{56}$Ni is dominantly produced but Co, Ni, and Zn are also mostly produced here. The final abundance depends on entropy because the mass fraction of alpha-particles during the alpha-rich freezeout phase increases with entropy [39].

For example, one of the authors has shown that high entropy during the hypernova explosion explains well the larger ratio of [Zn/Fe] and [Co/Fe] towards lower metallicity, [Fe/H], in extremely metal-poor (EMP) stars [16,22,34,40]. They also discussed the fact that a simple 1D model cannot explain the abundance of EMP stars by the hypernova models, because Fe or $^{56}$Ni is over-ejected to satisfy the large [Zn/Fe] ratio. This problem and its solution are described in Subsect. 4.3.

When $T_s$ is between $4 - 5 \times 10^9$ (K), Si burns incompletely. This region is called the incomplete Si-burning region. The main products here are Cr, Mn, and $^{56}$Ni. For $T_s \simeq (3-4) \times 10^9$ (K), oxygen burns partially and Si, S, Ar, Ca, Ti, and V are the main products. In these regions, nuclear burning
Fig. 11. Mass fraction distribution of the innermost region of the supernova ejecta evolved from 15 (a) and 20 $M_\odot$ (b) stars with a metallicity of $Z = 0.02$. The explosion energy is set at $1 \times 10^{51}$ erg.

Fig. 12. Mass fraction distribution of Li, B, F, Sc, V, and Mn produced through the $\nu$-process in the supernova ejecta evolved from a 15 $M_\odot$ star with a metallicity of $Z = 0.02$. Solid and dotted lines indicate the mass fractions of the elements with and without the $\nu$-process, respectively.

also produces some amount of energy. In summary, previous work has shown that explosive nucleosynthesis up to Zn can be described well by instant energy injection models, though hypernova models over-produce Fe in simple spherically symmetric models.

So far we have not mentioned the effects of the neutrino process ($\nu$-process), but the process is important for some elements such as Li, B, F, and Mn [38,41–50]. Of course, neutrino emission depends on the explosion model and thus we have to assume a model to include the effects (see Ref. [38] for details). Figure 12 shows the mass fraction distribution of Li, B, F, Sc, V, and Mn of
the supernova ejecta evolved from a $15 \, M_\odot$ star with $Z = 0.02$. In this calculation, we assume that the neutrino luminosity decreases exponentially with a decay time of 3 s. The total neutrino energy is set to be $3 \times 10^{53}$ erg. The neutrino spectra are assumed to obey a Fermi distribution with zero chemical potential. The temperatures of $\nu_\mu$, $\bar{\nu}_\mu$, $\nu_e$, $\bar{\nu}_e$ are set to be $6 \, \text{MeV}/k$ and $4 \, \text{MeV}/k$, respectively, as in Ref. [38]. Almost all Li and B are produced through the $\nu$-process in the O/Ne-layer and He-layer. When the $\nu$-process is not taken into account, the mass fractions of Li and B are smaller than the lowest value in this figure. Large F production through the $\nu$-process is obtained in the O/Ne layer (see also Refs. [49,50]). Some F is produced in the explosive He burning from $^{15}\text{N}$. Although the F production from $^{15}\text{N}$ strongly depends on metallicity, the F production through the $\nu$-process depends very little on metallicity and, thus, it is important throughout Galactic chemical evolution. Some Sc, V, and Mn are produced though the $\nu$-process, especially in the complete and incomplete Si-burning regions.

4.1. $X_C^{(12}\text{C})$ and abundance pattern of Ne to Ca

In the previous section, we mentioned that the carbon abundance after helium burning is important for the abundance of Ne to Ca. Here we look at the results of the UN and YU models more closely. In Fig. 13 we show $[X_i/O]$ vs. atomic number. Here $X_i$ represents Salpeter’s initial mass function (IMF) weighted yields of Ne to Ca isotopes, and $[A/B] = \log_{10}(Y_A/Y_B) - \log_{10}(Y_A/Y_B)_\odot$, where $Y_A$ and $Y_B$ are the abundances of elements A and B. If any point is far from the solar value, $[X_i/O] = 0$, the model fails to explain the present-day abundance, though a chemical evolution model should be applied for a detailed discussion. This figure show that, for both models, each point is roughly in the $\pm 0.3$ range, thus these models would yield a roughly solar abundance pattern.

Since the UN model has a larger $X_C^{(12}\text{C})$ after helium burning, this model has larger Ne/O and Na/O ratios than the YU model, as expected. A larger $X_C^{(12}\text{C})$ means a smaller $X_C^{(16}\text{O})$ and thus explosive oxygen burning products, such as S, Ar, Ca, would be less abundant. This is also seen in the figure.

The UN yields calculated with the same parameter choices were applied to the Galactic chemical evolution model in Ref. [51]. It was shown that the model could reproduce the observed abundance pattern reasonably well. If we look at the results closely, however, $[\text{Na/Fe}]$ is slightly larger than zero at $[\text{Fe/H}] = 0$, and $[\text{Ca/Fe}]$ and $[\text{Ar/Fe}]$ are slightly smaller than zero at $[\text{Fe/H}] = 0$. This suggests that a slightly larger $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ rate may reproduce the solar abundance better at $[\text{Fe/H}] = 0$. 

Fig. 13. The IMF weighted yield ratios $[X_i/O]$ as a function of atomic weight. Red and blue lines indicate the yield ratios in the YU model and UN model, respectively.
Using the YU codes we will investigate the best choice of the $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ rate, or $X_{C}(^{12}\text{C})$, in a forthcoming paper.

### 4.2. Zn abundance in EMP stars and asphericity of hypernovae

Observations of EMP stars, which are defined as halo stars with $[\text{Fe}/\text{H}] < -4$, revealed that their [Zn/Fe] and [Co/Fe] are larger than higher-metallicity stars (e.g., Ref. [52]). Reference [16] proposed that these large ratios are explained by the ejecta of complete Si burning of hypernovae. This idea explains also why such EMP stars have low [Fe/H]. In the early universe, where the metal abundance in interstellar matter (ISM) was low, the Fe to H ratio was determined by the amount of Fe produced by a supernova over that of hydrogen swept by the supernova shock. Since the swept hydrogen mass is roughly proportional to the explosion energy, the ejecta of a hypernova tends to have lower [Fe/H] than that of a supernova.

As mentioned above, however, there is one problem with this idea. Zn and Co are produced mainly in the complete Si-burning region while Fe or $^{56}\text{Ni}$ is produced in both the complete and incomplete Si-burning regions. Therefore, [Zn, Co/Fe] becomes larger when relatively more material is ejected from the complete Si-burning region. Under the assumption of a spherical symmetric explosion, this means that the mass-cut has to be taken sufficiently deep, then the ejected mass of Fe or $^{56}\text{Ni}$ also increases. On the other hand, if too much Fe is ejected, the O/Fe ratio, for example, becomes too small compared with observations. To quantify this problem, Ref. [16] introduced the ‘mixing-fallback model’. In the model, yields are calculated as follows. First, the innermost matter is mixed between the mass coordinates $M_{\text{in}}$ and $M_{\text{out}}$. For hypernova models, $M_{\text{out}}$ is set to the upper boundary of the incomplete Si-burning region, and $M_{\text{in}}$ is chosen sufficiently deep to eject Zn. Then it is assumed that only a fraction $f$ of the matter is ejected, or the fraction $1 - f$ fallbacks, from this region. To explain the abundance pattern of EMP stars $f \sim 0.1$ is required.

For supernova models with $M < 25 M_{\odot}$, $f = 1$ is fine to reproduce the abundance pattern of very metal-poor stars with $-3 < [\text{Fe}/\text{H}] < -2$ [34]. This suggests that the explosion mechanism for supernovae and hypernovae is different.

Interestingly, almost the same yield as the $f \sim 0.1$ hypernova yield could be obtained when we calculate jet-like hypernova models in 2D [53]. In such models, it is assumed that an unknown central engine ejects jets along the polar directions. These jets can blow up the entire star above the Si layer, but the explosive Si burning mainly takes place along the jet directions. As a result, the mass fraction of complete Si-burning products becomes smaller than the spherically symmetric models. In Sect. 6.3, we show such 2D calculations applied to an initially $M = 110 M_{\odot}$ star. Since observed hypernovae show some asphericity [54], this jet-like explosion model is certainly interesting, though we need to know the mechanism for the central engine.

It is sometimes claimed that Zn and Co in EMP stars may be explained by the innermost matter from the ‘hot bubble’ region of normal supernovae (e.g., Ref. [55]). Reference [56] explored this possibility and discussed that this explanation is not likely because fine tuning of $Y_{e} (= 0.500 - 0.501)$ is required for a 0.06 $M_{\odot}$ ejecta to produce sufficient amount of Zn and Co. These conditions are difficult to satisfy for the hot bubble matter of supernovae, but relatively easy for the innermost ejecta of hypernovae.

### 4.3. Nucleosynthesis of (weak-) r-process elements

It is quite likely that, during explosion, supernovae produce elements heavier than Zn, because some EMP stars have r- and weak-r-process elements. For example, CS22892-052 shows an r-process
abundance pattern that is roughly same as the present-day r-process pattern \[57\]. This suggests that r-process nucleosynthesis is almost universal. There are other classes of stars that show a greater abundance of weak r-process elements such as Sr, Y, Zr than the universal r-process pattern. There are also some stars with enhancements of Mo and Ru as well as Sr, Y, Zr stars \[58\].

As mentioned above, spherically symmetric instant energy injection models cannot synthesize such elements because, to synthesize such elements, a much larger entropy and/or more neutron-rich environment are required. On the other hand, Fe-peak elements including Zn can be synthesized from matter with \(Y_e \simeq 0.50\). The most recent 1D core collapse supernovae simulations have shown that the innermost matter in the ejecta has \(Y_e \simeq 0.50\) or even greater than 0.5 because of the interaction with neutrinos (e.g., Refs. \[59–61\]).

References \[40\] and \[56\] explored the conditions for synthesizing weak-r-process elements, Sr, Y, and Zr. Typically, the innermost layers of normal supernova explosions have an entropy per baryon of \(s/k \sim 5\) and those of hypernovae have \(s/k \sim 15\). With such entropy, weak-r-process elements are not produced if \(Y_e \simeq 0.50\). Thus they relaxed the constraint \(Y_e \simeq 0.50\) and considered \(Y_e \) as low as 0.45. Such low \(Y_e\) matter cannot be ejected as long as we consider an exactly spherically symmetric explosion. However, as shown in Ref. \[62\], in 2D calculations, inner materials are mixed by convection during explosion and a small amount of low \(Y_e \sim 0.45\) and higher entropy \(s/k \sim 40–50\) matter may be ejected. Reference \[40\] showed that the \(s/k \gtrsim 15\) model could produce Sr, Y, Zr explosively. Therefore, such low \(Y_e\) matter ejected due to multi-dimensional effects can explain the origin of weak-r-process elements.

We note that \(s/k \lesssim 50\) models could produce up to Zr but not heavier elements. Because of the lack of observations, it is not yet clear whether weak-r stars with enhancements of Sr, Y, Zr always have Nb–Mo enhancements as well. If this is the case, the innermost matter of hypernovae or the hot bubble of normal supernovae considered in Ref. \[40\] may not be the main site for weak-r-process synthesis. To produce Nb–Mo, a much larger entropy \((s/k \sim 150)\) is required \[56\]. It is currently not clear which astronomical site has such an environment.

5. Remnant neutron star mass

In this section we present remnant neutron star masses in both the UN and YU models and discuss their implications. There are several factors to determine the mass, such as the CO-core mass, explosion energy and EOS. Among them, the most important factor is the CO-core mass, which is determined by the stellar evolution before core collapse. A larger CO core leads to a larger remnant mass because it typically leads to a larger Fe core and, more importantly, it increases the amount of mass above the Fe core.

Once the pre-explosion density structure is given, the explosion energy determines the mass-cut or remnant neutron star mass \(M_{\text{rem}}\). For a given progenitor model, a larger explosion energy blows up more material above the Fe core, leading to a smaller mass-cut. In this paper, however, as well as in our previous work, we do not determine the mass-cut in this dynamical way but determine it by the amount of ejected \(^{56}\)Ni, \(M(^{56}\text{Ni})\). This is because \(M(^{56}\text{Ni})\) is rather sensitive to explosion energy when we determine the mass-cut dynamically. In reality, each SN may eject a somewhat different amount of \(^{56}\)Ni. However, when we apply supernova yields, e.g., to the Galactic chemical evolution, it is more useful to provide averaged yields. For this purpose, it is better to determine the mass-cut by \(M(^{56}\text{Ni})\).

We set the explosion energy as \(E_{\text{exp}} = 10^{51}\) erg, and the ejected \(^{56}\)Ni amount as 0.07 \(M_\odot\) to determine \(M_{\text{rem}}\). \(M(^{56}\text{Ni}) = 0.07\) \(M_\odot\) is the value for the SN1987A \[63\] and is considered the typical
value for normal core-collapse SNe. Strictly speaking, to eject the same amount of $^{56}$Ni, a more massive core requires a larger $E_{\text{exp}}$. However, the resultant $M_{\text{rem}}$ is not very different, so we fix the $E_{\text{exp}}$ for simplicity.

In Table 2, we show the baryon mass of the remnant neutron star mass by $M_{\text{rem}}$. It correlates with the CO-core mass, which in turn correlates with the He-core mass. The UN model leaves a smaller $M_{\text{CO}}$ than the YU model for the same initial mass, $M$, mainly because $X_C(^{12}\text{C})$ after He burning is larger and thus $M_{\text{CO}}$ is smaller, as described in Sect. 3. As described, the choice of $M(^{56}\text{Ni}) = 0.07 \, M_\odot$ is a typical value, but this is, of course, not the only value (see, e.g., the review in Ref. [64]). Therefore, we should remember this fact when comparing with observations. For example, when $M(^{56}\text{Ni}) = 0.01 \, M_\odot$ in the YU model, $M_{\text{rem}} = (1.38, 1.41, 1.53, 1.53, 1.72, 1.86, 2.04) \, M_\odot$ for $M = (10, 11, 12, 13, 15, 18, 20) \, M_\odot$, respectively.

$M_{\text{rem}}$ is the baryon mass and is not the observable neutron star mass. The observed mass is about 10% smaller due to the general relativistic effect and is called the gravitational mass (see, e.g., Ref. [65]). The conversion from baryon to gravitational mass depends on the EOS of nuclear matter. In Table 2 we show the gravitational mass, $M_g$, for the UU model in Ref. [66] using the UV14+UVII EOS [67] for nuclear matter.

### 5.1. Comparison with the observed mass

Here we compare the $M_g$ of our models with observed neutron star (NS) masses. NS masses are most accurately determined when they are in double NS systems. In this case if two or more post-Keplerian parameters are obtained, NS masses are precisely determined (e.g., Ref. [68,69]). In Table 4 we show the NS masses and errors for these stars. Interestingly, the mass distribution shows a peak at 1.33 $M_\odot$ with a small dispersion at 0.06 $M_\odot$. These stars may be considered to keep their birth mass [70], so they are suitable to compare with our results. NSs in binaries with high-mass companions are also considered to roughly keep their birth mass. Reference [70] discussed the idea that the most likely values of the central mass and dispersions for these NSs are 1.28 $M_\odot$ and 0.24 $M_\odot$, respectively.

First, we note that the lowest NS mass in the table is $M_g = 1.16-1.20 M_\odot$, and this is consistent with the lowest $M_g$ model in the YU model $M = 10 M_\odot$. As described above, this model may roughly correspond to the lightest Fe-core-collapse SNe. For a smaller mass star such as $M = 9.5 \, M_\odot$ or less, the star forms an O–Ne degenerate core and its collapse may lead to a weak SN explosion [14].

In the calculation in Ref. [14], the remnant neutron star has $M_{\text{rem}} = 1.36 M_\odot$ or $M_g \simeq 1.23 M_\odot$. This seems to be larger than the observed minimum mass NS, though these O–Ne supernovae have

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**Table 4.** NS masses in double NS systems and errors. See Ref. [70] and references therein.

| Name          | Mass ($M_\odot$) | Error ($M_\odot$) |
|---------------|------------------|-------------------|
| J0737-3039    | 1.3381           | 0.0007            |
| pulsar B      | 1.2489           | 0.0007            |
| B1534+12      | 1.3332           | 0.0010            |
| companion     | 1.3452           | 0.0010            |
| J1756-2251    | 1.40             | 0.02              |
| companion     | 1.18             | 0.02              |
| J1906+0746    | 1.248            | 0.018             |
| companion     | 1.365            | 0.018             |
| B1913+16      | 1.4398           | 0.002             |
| companion     | 1.3886           | 0.002             |
| B2127+11C     | 1.358            | 0.010             |
| companion     | 1.354            | 0.010             |
not been well studied and currently only one progenitor model exists [71,72]; therefore, we do not know the general properties of these SNe yet.

For the observed double NS, the typical mass range is $M_g = 1.27-1.39M_\odot$, as mentioned above. This corresponds to $M \sim 11.5-20M_\odot$ for the UN model and $M \sim 12-14M_\odot$ for the YU model.

The largest NS mass in the table is $M_g = 1.44M_\odot$; this corresponds to $M \sim 22M_\odot$ for the UN model and $M \simeq 15M_\odot$ for the YU model. It is usually said that the boundary between NS and black-hole formation is at around $M = 20-25M_\odot$. Therefore it appears that the $M_g$ range of the UN model is consistent with the double NS mass range, though we have not yet confirmed if the UN model with $M \lesssim 11M_\odot$ would result in $M_g \sim 1.2M_\odot$ NSs.

If we adopt the YU model, the mass range of the double NS corresponds to a rather narrow range, $M \sim 12-14M_\odot$, and it is difficult to understand the reason for this. However, at this moment, we cannot say that the UN model represents reality better, because the number of double NS systems are still limited. In Ref. [70], using the model of Ref. [73], they also discussed the idea that the narrow mass range of $M_g$ in double NS systems is difficult to understand, and suggested a particular and rare formation channel for the systems. On the other hand, the NS mass in binaries with high-mass companion ranges $M_g = 1.04-1.52M_\odot$ [70]. Although this estimate is more uncertain than that in double NS systems, this wider range would be more easily understood in the YU model.

In summary, NS mass observations would certainly provide interesting and important information to constrain the progenitor model, $X_C(^{12}C)$, and $^{12}C(\alpha, \gamma)^{16}O$ rate. Interestingly, all the observed NSs mentioned above have relatively small masses compared with ‘recycled’ NSs. NSs with white-dwarf companions, millisecond pulsars, and in low-mass X-ray binaries are known as ‘recycled’ and are currently undergoing accretion. It is now known that these NSs can surely be as massive as $M_g = 2.0M_\odot$ as the case for J1614-2230 [74]. Observations and theory are consistent in the sense that such massive NSs are rarely formed by normal supernova explosions. Theoretically, when we fix the explosion energy, $M_{\text{rem}}$ increases rapidly around $M \sim 25M_\odot$. This explains why massive NSs are rarely formed. It will be very interesting to constrain observationally the method by which massive NSs can be formed at birth.

6. Other recent work by our group

In this paper we have described our recent work on stellar evolution with some new results. Here we briefly review some other work by our group that is not mentioned above.

6.1. (Magneto-)hydrodynamical simulations

Kuroda and Umeda have developed a 3D magneto-hydrodynamical general relativistic code with adaptive mesh refinements [75]. This code was used to follow gravitational Fe core collapse and to calculate the spectra of gravitational waves. We will apply this code to various progenitor models to explore the explosion mechanism and nucleosynthesis.

6.2. Special relativistic hydrodynamical simulations for GRB jets

Okita and Umeda have developed a 2D special relativistic hydrodynamical code. Using this code, Okita and Umeda (manuscript submitted for publication) explored the conditions for successful ejection of ultra-relativistic jets in the collapsar model of GRBs. This code is also applied to explore various aspects of explosions and nucleosynthesis in SNe and GRBs. One example is given in the next subsection.
6.3. Evolution and nucleosynthesis in very massive stars which end up as Fe core-collapse SNe

Motivated by the discoveries of very luminous SNe, such as SN 1999as (SN Ic) [76] and SN 2006gy (SN IIn) [77], Umeda and Nomoto [17] calculated the evolution of metal-poor massive stars in the mass range $M = 20-100 M_\odot$ with metallicity $Z = 10^{-4}$ to study how much $^{56}$Ni is produced in core collapse SNe (CCSNe). They found that $^{56}$Ni of $\sim 13 M_\odot$ can be produced for a $100 M_\odot$ star. This amount is large enough to explain SN 1999as if the SN shines by $^{56}$Ni decay.

Since the actual bright SN Ic usually appears in a host galaxy with metallicity larger than $Z = 10^{-4}$, the mass-loss effect becomes important. Therefore, Yoshida and Umeda [20] studied the uncertainties in mass loss in detail for $Z = 0.004$, which corresponds to the host galaxy of SN 2007bi. SN 2007bi was a very bright SN Ic and $3.5-7.4 M_\odot$ of $^{56}$Ni is required if it shines by radioactive $^{56}$Ni. First, Gal-Yam et al. [78] discussed the idea that the SN was a pair-instability SN (PISN), but Moriya et al. [79] showed that it could be a CCSN if a $\sim 43 M_\odot$ CO star explodes with $E_{\text{exp}} = 3 \times 10^{52}$ erg.

Since the PISN model is much more massive than the CCSN model, these models predict quite different light curves. Unfortunately, without the light-curve data well before the maximum light, one cannot distinguish these two models. Yoshida and Umeda [20] found that for $Z = 0.004$, the required ZAMS ranges are $M = 110-280 M_\odot$ and $M = 515-575 M_\odot$ for the CCSN and PISN models, respectively. They showed that, if the progenitor was a single star and assuming Salpeter’s IMF, the ratio of the probabilities of CCSN to those of PISN appropriate for SN 2007bi is 42. We should remember, though, that for the CCSN model we are assuming that such a large CO star can explode energetically. So far, no one has shown such an explosion from first-principles calculations.

Under the assumption that it can explode, Okita, Umeda, and Yoshida simulated spherically symmetric and axisymmetric jet-like core-collapse supernova explosions of an $M_I = 43.2 M_\odot$ WO star in Refs. [20,80,81]. This progenitor is evolved from an $M = 110 M_\odot$ star with a metallicity of $Z = 0.004$. The CO core mass is $M_{\text{CO}} = 38.2 M_\odot$. For spherically symmetric calculations, the same method as described in Sect. 4 is used. For a jet-like explosion, the code mentioned in Sect. 6.2 was used. After the explosion simulations, nucleosynthesis was calculated post-processingly. First, we investigated the explosion-energy dependence of the $^{56}$Ni amount. We found that $^{56}$Ni larger than 3 $M_\odot$, enough to reproduce the light curve of SN 2007bi, is produced in the supernova ejecta if the explosion energy is larger than $2 \times 10^{52}$ erg. An investigation with jet-like explosions with different opening angles $\theta_{\text{op}}$ with $E_{\text{exp}} = 3 \times 10^{52}$ erg indicated that the ejected $^{56}$Ni yield strongly depends on the opening angle. Figure 14 shows the relation between the amount of ejected $^{56}$Ni and the opening angle. Although the spherical explosion releases the amount of $4 M_\odot$ $^{56}$Ni in the ejecta, the ejected $^{56}$Ni is much smaller when the opening angle is smaller than 78$^\circ$. The ejected amount of $^{56}$Ni is smaller than 2.3 $M_\odot$ for opening angles smaller than 68$^\circ$. This suggests that, if SN 2007bi was a CCSN, it exploded with a large opening angle or the explosion was more energetic than $3 \times 10^{52}$ erg.

6.4. Pop III very massive stars: Evolution, dark stars, stability

The formation of first-generation or Population (Pop) III stars in the Universe is considered to be quite different from that of later-generation stars. This is because, when the first-generation stars were formed, there was no metal, thus radiative feedback from the proto-star was weak and might not have been able to stop gas accretion. Also, the gas accretion rate itself was larger than that in the present universe. As a result, it was expected that typical masses of Pop III stars were 100 $M_\odot$ or more (e.g., Ref. [82]).
Fig. 14. The relation between the ejected $^{56}$Ni amount and opening angle $\theta_{\text{op}}$ in the axisymmetric supernova model. The progenitor is a 43.2 $M_{\odot}$ WO star. The explosion energy is set to be $3 \times 10^{52}$ erg.

In Ohkubo et al. [83], they calculated evolution and nucleosynthesis in very massive 500 and 1000 $M_{\odot}$ stars assuming that they explode energetically. Next, by taking account of a realistic mass accretion rate, Ref. [84] calculated Pop III star evolution with mass accretion. They found that typically 500−1000$M_{\odot}$ stars were formed. It is notable that this mass range is larger than the often-quoted mass range for a Pop III star, $M = 140−300 M_{\odot}$. If the mass range is in the 140−300$M_{\odot}$ range, most first stars should have exploded as PISNe. However, the abundance patterns of EMP and HMP stars with $[\text{Fe/H}] < -4$ are not consistent with the yields of PISNe [16].

These works, however, neglected the effects on the accretion disk likely formed around the first stars. References [85] and [86] investigated such effects and showed that the accretion disk was evaporated by radiative feedback. Then accretion onto a first star could stop before its mass grows more than 100 $M_{\odot}$, preventing the formation of a PISN [86].

Even if this is true, there is still a possibility for the formation of 1000 $M_{\odot}$ stars. First stars are formed around the center of a dark halo where the density of dark matter is much higher than in other places. References [87] and [88] showed that if the dark matter is self-annihilating WIMPs, the self-annihilation energy is sufficiently large to sustain the first stars. Such stars are called ‘dark stars’. These stars typically have much larger radii, thus surface temperatures are lower and radiative feedback is expected to be weak. Hirano, Umeda, and Yoshida [89] calculated the dark-star model with mass accretion. We confirmed that the stars keep their large radii until they reach the main sequence stage at around $M \sim 1000 M_{\odot}$.

Umeda et al. [90] also considered the accreting dark star model including the effects of captured dark matter annihilation, which becomes important after the main sequence stage. They showed that if the baryon–dark matter scattering cross section is as large as $\sigma = 10^{-38}$ cm$^{-2}$, the dark stars could grow more: they could be $10^4$ to $10^5$ $M_{\odot}$, depending on the mass accretion rates.

It has been known that such huge stars are vibrationally unstable against the epsilon mechanism (e.g., Ref. [91]). As long as the star is sustained by dark-matter annihilation and not nuclear burning, the epsilon mechanism does not operate. However, after the main sequence stage, nuclear burning soon dominates unless the captured dark matter effect is important. Sonoi and Umeda [92] explored the stability of such very massive stars against the epsilon mechanism. They found that the amount of mass loss is less than 10% of the whole stellar mass.

6.5. Dust formation in supernovae

H.U. has also been working on dust formation in supernovae with T. Nozawa, T. Kozasa and collaborators. For example, Nozawa et al. [93] calculated the dust yield for Pop III supernovae including
CCSNe and PISNe, and showed that a large amount of dust grains would be produced in the early universe by these SNe. Reference [94] calculated the dust destruction by reverse shock in a supernova remnant. The theory of dust formation was also applied to actual supernovae. References [95] and [96] for SN 2006jc and Cas A supernova remnants, respectively, showed that the observed data are reasonably reproduced by the theory. It was also applied to SNe Ia [97]. These dust grains in supernovae play important roles in star and galaxy formation, and chemical evolution in galaxies.

6.6. Presolar grains from supernovae

Presolar grains are recovered from primitive meteorites or interplanetary dust and are identified as grains with very large isotopic anomalies compared with solar-system materials (see, e.g., reviews [98,99]). The observed isotopic ratios of the grains are considered to indicate traces of nucleosynthesis in stars at their birth or Galactic chemical evolution. Small amounts of presolar grains are considered to originate from supernovae. They mainly indicate excesses of $^{12}$C, $^{15}$N, and $^{28}$Si. Some supernova grains show evidence for the original presence of radioactive $^{44}$Ti in Ca isotopic ratios, which strongly supports their origin. Isotopic ratios of heavy elements such as Mo and Ba have also been observed. However, it is still difficult to reproduce the observed isotopic ratios by supernova models. The bulk composition of supernova ejecta of supernova models does not reproduce the observed isotopic ratios of supernova grains. [100]. In order to reproduce the observed isotopic ratios, inhomogeneous mixing is required.

Yoshida and Hashimoto [101] and Yoshida [102] investigated supernova mixtures, reproducing several isotopic ratios of supernova-originating SiC and graphite grains. They divided supernova ejecta into seven different layers and investigated the mixing ratios, reproducing C, N, O, Al, Si, and Ti isotopic ratios as well as possible for individual grains. The mixing ratios of the mixtures strongly depend on the reproduced isotopic ratios. The main components of the mixtures are the innermost Ni layer and the outer He/C and He/N layers, so that inhomogeneous mixing of supernova ejecta is necessary. Yoshida, Umeda, and Nomoto [103] investigated the supernova mixtures of different stellar masses, reproducing Si isotopic ratios of supernova grains. They found that less massive supernovae with $M \lesssim 15M_\odot$ and hypernovae are preferable for reproducing Si isotopic ratios of supernova grains.

7. Discussions and future prospects

7.1. Massive star evolution

In this paper we have explained the differences between the newly developed efficient YU code and the previous UN code, and have shown that the YU code yields reasonable results, as shown in Sect. 3 in some detail. We need such an efficient code because the study of massive star evolution still requires a heavy amount of calculations, as described in the following.

As shown in Tables 1 and 2, the results of the UN and YU models presented in this paper are different mainly because $X_C(^{12}C)$ is different. The values of $X_C(^{12}C)$ sensitively depend on the $^{12}C(\alpha, \gamma)^{16}O$ rate and treatment of convection. Since it is currently not possible to determine the $^{12}C(\alpha, \gamma)^{16}O$ rate, we need to calculate for various choices of the rate to find a best set to fit observations. Traditionally, the nucleosynthetic argument shown in Sect. 4.1 has given the most stringent constraints on $X_C(^{12}C)$ and the $^{12}C(\alpha, \gamma)^{16}O$ rate. However, as shown in Fig. 13, it is not easy to distinguish even the UN and YU models because abundance data are only given for the IMF weighted yields.
We propose here that NS mass distribution will give alternative and independent constraints on the $X_{\text{C}}(^{12}\text{C})$ and $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ rate. As shown in Sect. 5 and Table 2, the UN and YU models predict different mass distributions for the remnant NSs. Though the differences are not so large, the observed NS masses are given with very fine precision for double NS systems (Table 4). Therefore, by calculating stellar evolution in binary systems, we may constrain the $X_{\text{C}}(^{12}\text{C})$ and $^{12}\text{C}(\alpha, \gamma)^{16}\text{O}$ rate more precisely, though many calculations are required for the various possible binary pairs.

One of our purposes in developing the efficient code was to tackle the evolutions just above and below Fe-core-forming SNe. As mentioned in Sect. 3.1, such stars experience violent shell flashes near the end of their evolution. It is not clear yet if such events cause mass loss and shock interactions to be observable. In order to follow these stages precisely, it is important for a code to include the acceleration term. As mentioned in Sect. 2, we will include the term in the YU code in the near future.

The stars just below the critical mass for Fe-core formation are very interesting for another reason. Such a star may become an electron capture supernova (ECSN). We would like to construct another progenitor model of ECSNe in addition to the one in Ref. [72] to see if the successful explosion in Ref. [14] is model independent or not. Since these calculations require high numerical resolution in both time and space [8], we need an efficient code to perform the computation.

We have also been working to develop a new code including the rotation effects as in Refs. [104–109]. This code is based on the YU code, and also aims to achieve efficient computation. Although rotating stellar models in the 1D formalism have been already calculated by these authors, angular momentum transfer is still quite uncertain. Therefore, our understanding of the evolution of a rotating star is still far from complete. Since rotation is inevitable for constructing realistic progenitor models for hypernovae and GRBs, and possibly even for normal CCSNe, we plan to calculate rotating progenitor models as well.

The uncertainties in the mass-loss rate are another reason why one needs to compute several cases to find a better set. Since knowledge on the rate is still very limited, one needs to constrain the rate by using various pieces of information, including the properties of massive stars, supernovae, compact remnants, ISM abundances, and so on. Unfortunately, this may not be easy, because other uncertain factors may be involved, such as rotation, binarity, and magnetic field effects. Supernovae may provide a key to understanding the rate. For example, as discussed in Ref. [20], if one can identify the SN 2007bi-like event as a CCSN or PISN, one can severely constrain the mass-loss rates.

### 7.2. Nucleosynthesis

As for our nucleosynthetic work, in Sect. 4 we described the success and limitations of the simple ‘instant energy injection’ model. This model reasonably reproduces supernova yields up to Zn. Although spherically symmetric hypernova models over-produce the Fe to lighter elements ratios, such as Fe/O, this problem may be avoided if one considers jet-like explosions in 2D for hypernovae even under the instant energy injection model. This suggests that nucleosynthesis up to Zn does not greatly depend on the details of the central engine.

On the other hand, nucleosynthesis of r- and weak-r-process elements depends on the details of the explosion model. This in turn implies that successful supernova and/or hypernova and/or GRB models must produce these elements.

As described in Ref. [40], there are observational indications that normal CCSNe produce lighter weak-r-process elements, Sr, Y, and Zr. These elements can be produced and ejected from the hot
bubble of normal SNe. However, to produce heavier weak-r-process elements, Mo and Ru, a much larger entropy and a neutron-rich environment is required [56]. It is currently not certain if such an environment can be realized in a CCSN. This is because long-term simulations of a CCSN have shown that high-entropy matter ejected from a CCSN is always almost neutral or proton-rich [61]. As mentioned in Sect. 4, it is not clear if all Sr–Zr-rich EMP stars are also Mo–Ru-rich. If this is the case, a CCSN model has to produce Mo–Ru somehow. One possibility is the $\nu p$-process (Refs. [110–113]) not considered in Ref. [56]. With this process, such weak-r-process elements may be synthesized in proton-rich matter if the entropy is sufficiently high. Thus it is critically important to observationally clarify if normal SNe produce Mo–Ru or not. We should note that the results in Ref. [61] are not obtained by first-principles calculations, but the explosion is driven by artificially enhancing the effective neutrino luminosity. The main reason that the matter becomes proton-rich is because of the interaction between matter and neutrinos. If the explosion is driven with the assistance of something else, such as rotation energy, neutron-rich matter may be ejected.

The arguments above also apply to the r-process elements. Currently, there is no clear evidence in EMP stars that r-process elements are produced in normal CCSNe and hypernovae. For example, Zn-rich EMP stars, which we consider to be polluted by a hypernova, show no enhancement of r-process elements. From the results in Ref. [61], it seems hopeless to produce r-process elements in normal CCSNe, though unknown massive stars should have polluted metal-poor r-rich stars, such as CS22892-052. Some authors consider NS–NS or NS–black-hole merger systems as r-process sites (e.g., Refs. [114–119]), though it is not yet clear if such sites are sufficient to explain the whole r-process abundance in the universe. Therefore, it is still interesting to explore various SN explosion models to examine whether they can produce r-process elements or not.

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