Nucleosynthesis in the Innermost Ejecta of Neutrino-driven Supernova Explosions in Two Dimensions

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Abstract

We examine nucleosynthesis in the innermost neutrino-processed ejecta (a few $10^{-3} M_\odot$) of self-consistent two-dimensional explosion models of core-collapse supernovae (CCSNe) for six progenitor stars with different initial masses. Three models have initial masses near the low-mass end of the SN range of 8.8 $M_\odot$ e8.8; electron-capture SN), 9.6 $M_\odot$ (z9.6), and 8.1 $M_\odot$ (u8.1), with initial metallicities of 1, 0, and $10^{-4}$ times the solar metallicity, respectively. The other three are solar-metallicity models with initial masses of 11.2 $M_\odot$ (s11), 15 $M_\odot$ (s15), and 27 $M_\odot$ (s27). The low-mass models e8.8, z9.6, and u8.1 exhibit high production factors (nucleosynthetic abundances relative to the solar abundances) of 100–200 for light trans-Fe elements from Zn to Zr. This is associated with an appreciable ejection of neutron-rich matter in these models. Remarkably, the nucleosynthetic outcomes for the progenitors e8.8 and z9.6 are almost identical, including interesting productions of $^{48}$Ca and $^{60}$Fe, irrespective of their quite different (O–Ne–Mg and Fe) cores prior to collapse. In the more massive models s11, s15, and s27, several proton-rich isotopes of light trans-Fe elements including the p-isotope $^{92}$Mo (for s27) are made, up to production factors of ~30. Both electron-capture SNe and CCSNe near the low-mass end can therefore be dominant contributors to the Galactic inventory of light trans-Fe elements from Zn to Zr and probably $^{48}$Ca and live $^{60}$Fe. The innermost ejecta of more massive SNe may have only subdominant contributions to the chemical enrichment of the Galaxy except for $^{92}$Mo.

Key words: nuclear reactions, nucleosynthesis, abundances – stars: abundances – stars: neutron – supernovae: general

1. Introduction

Core-collapse supernovae (CCSNe), the deaths of stars with initial masses heavier than about 8 $M_\odot$, have long been suggested to be important astrophysical sources of trans-Fe species as well as of intermediate-mass and Fe-group elements. Traditionally, studies of CCSN nucleosynthesis were based on artificial one-dimensional (1D) explosion models with free parameters such as the mass-cut (the location of the ejecta-remnant interface), the electron fraction ($Y_e$; the number of protons per nucleon) of the ejecta, and the explosion energy (e.g., Woosley & Weaver 1995; Thielemann et al. 1996; Rauscher et al. 2002; Limongi & Chieffi 2006; Tomimaga et al. 2007; Heger & Woosley 2010; Nomoto et al. 2013). Nucleosynthetic abundances made near the mass-cut in the innermost ejecta are, however, highly dependent on such free parameters.

The approach of replacing the complex multidimensional dynamics of a neutrino-driven SN with a piston or thermal bomb as an artificial explosion engine is likely to be problematic when it comes to studying the innermost ejecta. These parameterized models face two serious problems in modeling the innermost ejecta: first and most important, they fail to reproduce the thermodynamic history of the neutrino-driven ejecta, i.e., of material that makes its way deep into the gain region, is then heated by neutrinos, and is ejected in buoyant high-entropy bubbles. In these ejecta, any memory of the initial composition is lost as charged-current neutrino interactions completely reset the electron fraction $Y_e$ and the entropy per nucleon $S$. The final $Y_e$ and $S$ in the ejecta are determined by the interplay of the neutrino reactions and the expansion of the ejecta (see Pruet et al. 2005, 2006; Fröhlich et al. 2006; Wanajo et al. 2011; and in particular Section 2.2.1 of Müller 2016). It is therefore impossible to reproduce this ejecta component if one assumes that the only effect of the SN engine on the ejecta is to provide shock heating as in simpler models of the engine. Second, even for the explosive burning of material that is immediately ejected after being shocked, the simple engine models are problematic for material close to the mass-cut: for ejecta that are shock-heated while the SN engine is still feeding energy into the incipient explosion (which can take seconds; Müller 2015; Bruenn et al. 2016; Müller et al. 2017), the shock heating is sensitive to the nature of the engine and is therefore not sufficient to merely tune an artificial engine to produce plausible final explosion energies and Ni masses to capture the nucleosynthesis in this ejecta component. This makes it difficult to predict the nucleosynthesis of some Fe-group and trans-Fe (Zn and heavier) elements based on such parameterized models in these two ejecta components, which we term the “innermost ejecta.” It is therefore imperative to complement extant studies of the detailed explosive and hydrostatic nucleosynthesis with nucleosynthesis calculations based on more consistent models of the SN engine.

Recent work by Sukhbold et al. (2016) has advanced the 1D approach by adopting parameterized neutrino-powered explosion models, which is a first step toward addressing these problems. In
their models the boundary conditions (e.g., proto-neutron star contraction and core luminosity) are set deep inside the proto-neutron star and the physical conditions of the innermost ejecta are obtained as a result of hydrodynamical computations (see also Fröhlich et al. 2006; Perego et al. 2015 for comparable approaches beyond the piston and thermal bomb models in 1D). Similar to previous works, they found the production of elements from B to Cu to be in reasonable agreement with the solar ratio (except for some elements that had other contributors such as low-mass stars and SNe Ia). However, they reported a severe deficiency of light trans-Fe elements from Zn to Zr, which had been explained in part by the weak s-process in previous studies (e.g., Woosley & Heger 2007). In addition, astrophysical sources of \( ^{48}\text{Ca} \) (a neutron-rich isotope of Ca), \( ^{64}\text{Zn} \) (the main isotope of Zn), and \( ^{92}\text{Mo} \) (a p-isotope) still remain unresolved, for which a rare class of SNe Ia (Meyer et al. 1996; Woosley 1997), hypernovae (Umeda & Nomoto 2002; Tominaga et al. 2007), and neutron-rich nuclear equilibrium (Hoffman et al. 1996; Wanajo 2006), respectively, have been proposed as possible explanations. Note, however, that Sukhbold et al. (2016) did not include neutrino-heated (just shock-heated) ejecta in their analysis.

One of the fundamental problems in such 1D models is obviously the limitation of dimensionality. Among other things, they do not account for the spatial variations in the electron fraction \( Y_e \), entropy \( S \), and ejection velocity that is seen in multidimensional CCSN models. Two-dimensional (2D) simulations with sophisticated neutrino transport therefore provide a much better basis for studying nucleosynthesis in the innermost CCSN ejecta. Puet et al. (2005, 2006) used the hydrodynamical trajectories of a 15 \( M_{\odot} \) CCSN (Buras et al. 2006) for nucleosynthesis and showed that interesting amounts of Sc, Zn, and light p-nuclei were formed in the innermost proton-rich ejecta (see also Fujimoto et al. 2011) while the contribution of neutron-rich ejecta appeared to be subdominant (Hoffman et al. 2007). This was an important step beyond previous CCSN nucleosynthesis studies, but the explosion, though neutrino-driven, was still induced artificially like the piston-induced or thermal-bomb explosions of the previous 1D approaches, i.e., the 2D explosion was not obtained in a fully self-consistent manner. Moreover, based on the superior treatment of the neutrino transport in our present models we can re-investigate the question of nucleosynthesis in the neutron-rich ejecta, in which light trans-Fe species could be abundantly produced (e.g., Hoffman et al. 1996).

To our knowledge, the only extant studies of nucleosynthesis based on a self-consistent multidimensional explosion model are those by Wanajo et al. (2011b, 2013a, 2013b), which are based on a 2D simulation of an 8.8 \( M_{\odot} \) electron-capture SN (EC SN, a subclass of CCSNe arising from collapsing O–Ne–Mg cores; Nomoto 1987; Janka et al. 2008). They found an appreciable production of trans-Fe elements from Zn to Zr, \( ^{48}\text{Ca} \), and \( ^{60}\text{Fe} \) (radioactive nuclei) in the innermost neutron-rich ejecta, a very different result from nucleosynthesis studies (Hoffman et al. 2008; Wanajo et al. 2009) based on the corresponding 1D models (Kitaura et al. 2006; Janka et al. 2008). This demonstrates the importance of self-consistent multidimensional models for reliable nucleosynthesis predictions.

However, 2D explosion models only provide a first glimpse at the role of multidimensional effects in SN nucleosynthesis. Recent three-dimensional (3D) core-collapse simulations have shown qualitative and quantitative differences to 2D models (see, e.g., Janka et al. 2016 for a recent review) concerning shock revival (with a trend toward delayed or missing explosions) and concerning the multidimensional flow dynamics and energetics during the first phase of the explosion (Melson et al. 2015; Müller 2015). Thus nucleosynthesis studies based on 3D models will eventually be needed but are of course computationally much more demanding.

This paper aims at extending the 2D studies of Wanajo et al. (2011b, 2013a, 2013b) to examine the nucleosynthesis in the innermost ejecta of CCSNe arising from Fe-core progenitors. This will be an important step toward future nucleosynthesis studies from self-consistent 3D explosion models. In addition to the ECSN model studied by Wanajo et al. (2011b, 2013a, 2013b), we consider five CCSN models with zero-age main-sequence progenitor masses of 9.6 \( M_{\odot} \) (an initial metallicity of \( Z = 0 Z_{\odot} \)), 8.1 \( M_{\odot} \) (10\(^{-4}\) \( Z_{\odot} \)), and 11.2 \( M_{\odot} \), 15 \( M_{\odot} \), and 27 \( M_{\odot} \) (1 \( Z_{\odot} \)) (Section 2). By including the 8.1 \( M_{\odot} \) and 9.6 \( M_{\odot} \) progenitors side by side with an ECSN model, our study retains a strong focus on the low-mass end of the progenitor spectrum. With the uncertainties surrounding the ECSN channel (Poelarends et al. 2008; Jones et al. 2013, 2014; Doherty et al. 2015; Jones et al. 2016), one particular question that we seek to answer is whether ECSN-like nucleosynthesis can also be obtained for slightly different progenitor channels close to the Fe-core formation limit.

All these SN explosions were obtained self-consistently, i.e., with no free parameters, in 2D axisymmetric simulations adopting an elaborate neutrino transport scheme (Janka et al. 2012; Müller et al. 2012a, 2012b). Nucleosynthetic abundances are calculated by applying an up-to-date reaction network code to the hydrodynamic trajectories in a post-processing step (Section 3). Some details of the nucleosynthesis mechanisms operating in the innermost ejecta are described by taking the result of the metal-free 9.6 \( M_{\odot} \) star as representative of our models. The mass-integrated nucleosynthetic yields are compared to the solar abundances to test if the innermost ejecta of these SNe can be major sources of light trans-Fe elements and some other species in the Galaxy (Section 4). Our conclusions follow (Section 5).

2. SN Models

2.1. Numerical Methods and Progenitor Models

All the SN hydrodynamical trajectories have been computed from 2D general relativistic simulations with energy-dependent ray-by-ray-plus neutrino transport based on a variable Eddington factor technique as implemented in the SN code VERTEX (Rampp & Janka 2002; Buras et al. 2006; Müller et al. 2010). Except for the pseudorelativistic ECSN model of Wanajo et al. (2011b), general relativity is treated in the extended conformal flatness approximation (Cordero-Carrión et al. 2009). A modern set of neutrino interaction rates (the “full rates” set of Müller et al. 2012b) has been used for the simulations. The explosions were obtained self-consistently and thus the models contained no free parameters.

The initial pre-SN models were adopted from Nomoto (1987; an 8.8 \( M_{\odot} \) star with an O–Ne–Mg core with solar metallicity), A. Heger (unpublished);\(^8\) 9.6 \( M_{\odot} \) and 8.1 \( M_{\odot} \) models\(^9\) with

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\(^8\) Extension of Heger & Woosley (2010).

\(^9\) Both the 9.6 \( M_{\odot} \) of zero metallicity and the 8.1 \( M_{\odot} \) model of 10\(^{-4}\) \( Z_{\odot} \) were at the low-mass ends of CCSN progenitors for their respective metallicities in the calculations using the KEPLER code. See also Ibeling & Heger (2013) for the general metallicity dependence of the CCSN lower mass limit.
Table 1
Parameters of SN Explosion Models

| Model | $M_{\text{pre}}$ (M$_\odot$) | $Z_{\text{pre}}$ (Z$_\odot$) | $M_{\text{CO}}$ (M$_\odot$) | $M_{\text{He}}$ (M$_\odot$) | $\Delta\phi$ | $t_{\text{expl}}$ (ms) | $t_{\text{fin}}$ (ms) | $m_{\text{fin}}$ (M$_\odot$) | $E_{\text{shock}}$ ($10^{50}$ erg) |
|-------|-----------------|-----------------|-----------------|-----------------|--------|-----------------|-----------------|-----------------|-----------------|
| e8.8  | 8.8             | 1               | 1.377           | 1.377           | 1.4    | 92              | 362             | >1.377          | 0.9             |
| z9.6  | 9.6             | 0               | 1.371           | 1.69            | 1.4    | 125             | 1420            | 1.373           | 0.6             |
| s8.1  | 8.1             | $10^{-4}$       | 1.385           | 1.92            | 1.4    | 177             | 335             | 1.373           | 0.4             |
| s11$^f$ | 11.2           | 1               | 1.89            | 2.84            | 2.8    | 213             | 922             | 1.39            | 1.3             |
| s15n$^a$ | 15              | 1               | 2.5             | 4.5             | 2.8    | 569             | 779             | 1.45            | 1.3             |
| s27$^n$ | 27              | 1               | 7.2             | 9.2             | 1.4    | 209             | 790             | 1.74            | 1.9             |

Notes.

a Progenitor mass at the zero-age main sequence.
b Metallicity at the zero-age main sequence.
c CO core mass at collapse.
d He core mass at collapse.
e Angular resolution.
f Post-bounce time of explosion, defined as the point in time when the average shock radius ($r_{sh}$) reaches 400 km.
g Final post-bounce time reached in the simulation.
h Mass coordinate corresponding to the average shock radius at the end of the simulation.
i Diagnostic explosion energy at the end of simulation.

j Note that an artificial envelope with $\rho \propto r^{-3}$ was used in Wanajo et al. (2011b) outside the He core. It is therefore not appropriate to give a more precise mass coordinate for the final position of the shock in the hydrogen shell.

k By the end of the simulation, the shock has crossed the outer grid boundary at a mass coordinate of 1.373 M$_\odot$ in model z9.6.

l Same as model s11.2 in Woosley & Weaver (2002); a solar metallicity model, s11.2 in their paper; and s11 in their paper, respectively, Woosley & Weaver (1995; a 15 M$_\odot$ solar metallicity model, s15s7b2 in their paper), and Woosley et al. (2002; a 27 M$_\odot$ solar metallicity model, s27.0 in their paper). The corresponding explosion models are labeled as e8.8, z9.6, u8.1, s11, s15, and s27 hereafter. The relevant model parameters are summarized in Table 1. For more details on the individual explosion models, we refer the reader to the original publications on e8.8 (Wanajo et al. 2011b), z9.6 (Janka et al. 2012), u8.1 (Müller et al. 2012a), s11 and s15 (Müller et al. 2012b), and s27 (Müller et al. 2012a).

Note that although we include two models of subsolar metallicity, the limited number of models does not permit us to discuss the dependency on metallicity in this paper.

2.2. Explosion Dynamics

Snapshots of the electron fraction, $Y_e$, and the entropy per nucleon, $S$, in these simulations at late times are shown in Figure 1. One can see roughly spherical structures for e8.8 and z9.6 (top panels) and more strongly asymmetric features with a dipolar or quadrupolar geometry for the other models (middle and bottom panels). The different explosion dynamics reflect the core structures of pre-SN stars with steeper to shallower density gradients in the order of e8.8, z9.6, u8.1, s11, s15, and s27 (see Figure 8 in Janka et al. 2012). The ECSN progenitor for e8.8 is a super-asymptotic giant branch (SAGB) star with an O–Ne–Mg core surrounded by a very dilute H–He envelope, while progenitors of CCSNe have He cores embedded by dense O–Si shells. For e8.8, therefore, the explosion sets in very early at a postbounce time of $t_{\text{pb}}$ $\sim$ 80 ms before vigorous convection can develop. Overturn driven by the Rayleigh–Taylor instability only occurs when the explosion is underway, but the plumes do not have sufficient time to merge into large structures so no global asymmetry emerges. For the CCSN cases (except for z9.6), by contrast, the explosions gradually start after the multidimensional effects, i.e., convection or the standing-accretion-shock instability (SASI; Blondin et al. 2003), have reached the nonlinear regime (the onset of the explosion occurs at $t_{\text{expl}}$ $\sim$ 500 ms for s15 and $t_{\text{pb}}$ $\sim$ 150 – 200 ms for the others; see Table 1 and Figure 14 in Janka et al. 2012). Moreover, due to the presence of a relatively dense and massive O shell, the shock expansion is slow enough for a dominant unipolar or bipolar asymmetry to emerge after shock revival.

Models z9.6 and u8.1 stand apart from the more massive CCSN models s11, s15, and s27 since stars close to the Fe-core formation limit exhibit evolutionary and structural similarities to ECSN progenitors (Woosley & Heger 2010; Jones et al. 2013, 2014). Specifically, these progenitors have very thin O and C shells between the core and the low-density He and H envelope. Because of these peculiarities, model z9.6 (and to some degree u8.1) is a case on the borderline of ECSN-like explosion behavior (Müller 2016). Its progenitor exhibits the steepest core-density gradient near the core-envelope interface among the CCSN cases, resulting in a similarly rapid expansion of the neutrino-heated ejecta and hence in $Y_e$ and $S$ structures rather similar to those of e8.8. The progenitor of u8.1 has a slightly shallower core-density gradient than that of z9.6, so the propagation of the shock is slightly slower than for z9.6 but is still faster than for typical Fe-core progenitors. Since the width of the ECSN channel is subject to considerable uncertainties (see Poelarends et al. 2008; Jones et al. 2013, 2014; Doherty et al. 2015; Jones et al. 2016, and references therein), CCSNe from this mass range are particularly interesting as a possible alternative source for “ECSN-like” nucleosynthesis (i.e., the characteristic nucleosynthesis in neutron-rich neutrino-heated ejecta with moderate entropy as discussed in Wanajo et al. 2011b, 2013a, 2013b).

As in Wanajo et al. (2011b) and similar to Wongwathanarat et al. (2017), we compute the trajectories for our nucleosynthesis
calculations from the 2D data files (with a time spacing of 0.25 ms) instead of co-evolving tracer particles during the simulation. In order not to follow the tracers during multiple convective overturns with the risk of accumulating discretization errors, we integrate the tracer trajectories backward in time from an appropriate point during the simulation (i.e., starting either

Figure 1. Late-time snapshots of models e8.8 (left-top, 266 ms after bounce), z9.6 (right-top, 317 ms), u8.1 (left-middle, 315 ms), s11 (right-middle, 922 ms), s15 (left-bottom, 776 ms), and s27 (right-bottom, 790 ms). Each panel shows the distribution of the electron fraction, $Y_e$ (left halves of panels; color bar at right-bottom), and the entropy per nucleon, $S$, in units of $k_B$/nucleon (right; color bar at right-top) at a time when the high-entropy plumes of neutrino-heated matter have reached a radius of roughly 3000 km. The black line in each panel indicates the shock front. The vertical and horizontal axes show the distance from the center. Note that the shock has already propagated well beyond 4000 km in e8.8 at this stage and is therefore no longer visible in the plot.
from the end of the simulation or from a time when the bulk of the ejecta have cooled sufficiently as in the case of z9.6). This procedure greatly reduces the number of required trajectories to adequately sample the ejecta: selecting the final locations of the tracer particles only within the region of interest reduces the total mass of ejecta that must be covered (less than 0.03 $M_\odot$) and makes it easy to use higher mass resolution where it is most needed, i.e., in the neutrino-processed ejecta. For a more general overview of the problems in extracting tracer trajectories from multidimensional simulations, we refer the reader to Harris et al. (2017).

### 2.3. Nucleosynthesis Conditions in the Ejecta

Some nucleosynthesis-relevant properties for all SN models are summarized in Table 2. It is important to note that we only calculate the nucleosynthesis for the “early ejecta”—i.e., we consider only the material that has been ejected in neutrino-driven outflows or undergone explosive burning in the shock by the end of the simulations. The contributions of the outer shells (i.e., the H–He envelopes of e8.8, z9.6, and u8.1 and material outside the middle of the O-burning shell in s11, s15, and s27) and later neutrino-driven ejecta to the total yields are neglected. For the massive CCSNe (s11, s15, s27), the nucleosynthetic contribution for the outer part of the O shell therefore remains uncertain, and long-term simulations to several seconds (Müller 2015; Müller et al. 2017) will be necessary to determine it. As the shock has only traversed part of the O shell in these models we also miss part of the contribution of hydrostatic and explosive burning to heavy element production from O/C/He shells, so our yields can only be understood as lower limits for some species (among them $^{26}$Al and $^{60}$Fe). We also disregard the possibility of late-time fallback of the initial ejecta, but this may not be a major uncertainty since the most advanced parameterized 1D models of neutrino-driven explosions suggest that this plays a minor role (Ertl et al. 2016) for the progenitor mass range considered here.

For models e8.8, z9.6, and u8.1, explosive burning is already complete at the end of the simulations: the outer shell (composed of H and He) will not add significant yields of heavy elements and no more explosive nucleosynthesis is taking place as the shock propagates to the stellar envelope after the end of the simulations. However, we still miss small amounts of ejected material from the neutrino-driven wind (Duncan et al. 1986) in these progenitors. This material may undergo weak r- or Yp-process nucleosynthesis, though the most recent calculations of wind nucleosynthesis for e8.8 suggest only a relatively unspectacular production of Fe-group elements and some Yp-process nucleides with small production factors (Pllumbi et al. 2015), which turn out to be insignificant compared to the contribution from the early ejecta that we investigate here. For e8.8, z9.6, and u8.1, the yields presented here thus cover essentially the complete nucleosynthesis of heavy elements in these progenitors.

The case is different for s11, s15, and s27, where the shock has progressed only to the middle of the O/Si shell. For these models, substantial amounts of ashes from O, Ne, C, and He burning thus remain to be ejected and will contribute significantly to the total production factors. The postshock temperatures at the end of our simulations also remain sufficiently high for some additional explosive O burning to take place. Moreover, strong accretion downflows still persist in these models. This keeps the neutrino luminosities high and allows the ejection of neutrino-heated matter to continue well beyond 1 s after the onset of the explosion (Müller 2015; Bruenn et al. 2016). Our nucleosynthesis calculations therefore place only a lower bound on the production of Fe-group and trans-Fe group elements in the SN core of these models. It is noteworthy that this may partly contribute to the unexpectedly small mass of $^{56}$Ni obtained for these models, which remains significantly smaller than expected for “ordinary” SNe (e.g., 0.07 $M_\odot$ for SN 1987A; Bouche et al. 1991).

Aside from the limitation of our analysis to the nucleosynthesis during the first few hundred milliseconds after shock revival, the nucleosynthesis in our models remains subject to other uncertainties: the simulations were conducted assuming axisymmetry (2D), and the explosion energies at the end of simulations are only $(0.3–1.5) \times 10^{50}$ erg (see Figure 15 in Janka et al. 2012), which are either less energetic than typical observed events (several $10^{50}$ erg; e.g., Kasen & Woosley 2009; Nomoto et al. 2013) or still in a phase of steep rise, i.e., not converged to their final values. We discuss the possible repercussions of this in Sections 4.8 and 5.

Figure 2 shows the $Y_e$ distributions in the ejecta for all models evaluated at $T_9 = 10$, where $T_9$ is the temperature in

### Table 2

Properties of Innermost Ejecta

| Model | Type  | $M_{\text{PNS}}$ | $M_\odot$ | $M_{\text{ej},n}$ | $Y_{e,\text{min}}$ | $Y_{e,\text{max}}$ | $S_{\text{min}}$ | $S_{\text{max}}$ |
|-------|------|-----------------|-----------|-------------------|-------------------|-------------------|-----------------|-----------------|
| e8.8  | ECSN | 1.36            | 11.4      | 5.83              | 0.398             | 0.555             | 9.80            | 383             |
| z9.6  | CCSN | 1.36            | 12.4      | 4.94              | 0.373             | 0.605             | 12.6            | 27.8            |
| u8.1  | CCSN | 1.36            | 7.69      | 0.592             | 0.464             | 0.598             | 6.78            | 36.7            |
| s11   | CCSN | 1.36            | 14.1      | 0.133             | 0.474             | 0.551             | 6.64            | 34.7            |
| s15   | CCSN | 1.58            | 15.9      | 0.592             | 0.464             | 0.598             | 6.78            | 36.7            |
| s27   | CCSN | 1.65            | 27.3      | 0.759             | 0.387             | 0.601             | 5.19            | 44.0            |

Notes.

a Baryonic mass of the proto-neutron star at the end of simulation.
b Total mass in the innermost ejecta.
c Ejecta mass with $Y_e < 0.4975$.
d Minimal $Y_e$ evaluated at $T_9 = 10$.
e Maximal $Y_e$ evaluated at $T_9 = 10$.
f Minimal asymptotic entropy (at the end of the simulation).
g Maximal asymptotic entropy (at the end of the simulation).
Figure 2. $Y_e$ histograms for the SN ejecta when the temperatures have decreased to $T < 10^9$. The ejecta masses ($\Delta M_{ej}$) are shown as functions of $Y_e$ with a bin size of $\Delta Y_e = 0.005$.

units of $10^9$ K. For a trajectory with a maximum temperature of $T_{\text{max}} < 10$, we show $Y_e$ at $T = T_{\text{max}}$ instead. $\Delta M_{ej}$ is the ejecta mass in each $Y_e$ bin with an interval of $\Delta Y_e = 0.005$. The minimum and maximum values of $Y_e$ for all models are given in Table 2 (8th and 9th columns, respectively). The $Y_e$ distributions shown here are similar to those inferred from the hydro data at late times except for small (up to a few %) shifts toward $Y_e \approx 0.5$ as a result of some numerical diffusion and mixing (clipping of extrema) that suppresses the tails of the $Y_e$ distribution in the hydro at late times.\(^{10}\) We find that the low-mass models (e8.8, z9.6, and u8.1) have appreciable amounts of neutron-rich ejecta (40%–50%; 7th column in Table 2). This is due to the faster growth of the shock radii for these low-mass cases: as a result of the higher ejecta speed, less time is available to increase the $Y_e$ by neutrino processing as material is ejected from the neutron-rich environment near the gain radius, where conditions close to equilibrium between neutrino absorption and electron (and to some extent positron) captures are maintained and lead to a low $Y_e$. By contrast, the bulk of the ejecta are proton-rich (96%–99%; Table 2) in the massive models (s11, s15, and s27). Here the ejecta expand more slowly, allowing neutrino absorption to reset the $Y_e$ to an equilibrium value determined by the competition of electron neutrino and antineutrino absorption (Qian & Woosley 1996; Fröhlich et al. 2006). With similar electron neutrino and antineutrino mean energies, the equilibrium value tends to lie above $Y_e = 0.5$ because of the proton-neutron mass difference.

Figure 3 illustrates the distributions of ejecta masses as functions of $Y_e$ and the asymptotic entropy per nucleon, $S$ (in units of Boltzmann’s constant $k_B$). All SN models show spikes of $Y_e \approx 0.50$ components as leftovers of the initial composition of shocked material that never undergoes strong neutrino processing, and they show a positive correlation of $S$ with $Y_e$ in the neutrino-processed ejecta. The latter fact is reasonable because neutrino heating raises both $Y_e$ and $S$. We find, however, a larger scatter of the entropies at a given $Y_e$ for the massive SN models (s11, s15, and s27). This reflects the more vigorous motions arising from convective instability and SASI for more massive cases as well as stronger temporal and spatial variations in the neutrino irradiation.

Note that the very high entropies at $Y_e \approx 0.5$ for e8.8 (up to $S_{\text{max}} = 383 k_B$; last column in Table 2) stem from the shock heating of the outgoing material colliding with the dilute SAGB envelope. However, a sizable postshock entropy, $S > 100 k_B$, is reached only in shells that never reach nuclear equilibrium ($T < 3$; Janka et al. 2008; Kuroda et al. 2008), so these shells do not contribute to the production of heavy elements (such as an r-process; suggested by Ning et al. 2007). Except for this component of shock-heated ejecta with low temperatures, e8.8 has a maximum entropy of nucleosynthesis-relevant material of $\approx 25 k_B$, and more massive models have larger values ($S_{\text{max}}$; last column in Table 2).

The core-collapse simulations were stopped at $t = t_{\text{fin}}$ for our models, where $t_{\text{fin}} = 0.423$ s, 0.605 s, 0.474 s, 0.767 s, 0.947 s, and 1.13 s after core bounce for e8.8, z9.6, u8.1, s11, s15, and s27, respectively. At these times, the temperatures in the ejecta were still high, in particular for massive models, so nucleosynthesis would be still active. The temperature ($T$) and density ($\rho$) thus need to be extrapolated for nucleosynthesis calculations. It is known that the late-time evolution of the density is well approximated by $\rho \propto t^{-2}$ (e.g., Arcones et al. 2007). We thus assume

$$\rho(t) = c_1 (t - t_1)^{-2} \quad (t > t_{\text{fin}}),$$

where the constants $c_1$ and $t_1$ are determined to get a smooth connection of the density at $t = t_{\text{fin}}$. The expansion of the

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\(^{10}\) In practice, we find this effect to be small up to the point that the ejecta reach a radius of around 1000 km and the radial computational grid becomes coarser.
ejecta is almost adiabatic at this stage (i.e., $T^3/\rho \approx$ constant),\(^{11}\) and thus the temperature is extrapolated as

$$T(t) = c_2(t - t_1)^{-2/3} \quad (t > t_{\text{fin}})$$ \(2\)

with $t_1$ from Equation (1) and $c_2$ determined so that the temperature matches the value at $t = t_{\text{fin}}$. The radius $r$ for $t > t_{\text{fin}}$, which is needed to calculate the rates of neutrino interactions, is obtained from Equation (1) with the assumption of steady-state conditions, i.e., $r^2\rho = \text{constant}$ (Panov \& Janka 2009; Wanajo et al. 2011a), where $v_r$ is the radial component of velocity at $t = t_{\text{fin}}$ and constant afterward. Selected temporal evolutions of density and temperature are shown in Figures 4 and 5, respectively, where the extrapolations are indicated by red curves. The evolutions for e8.8 and z9.6 appear to be nearly self-similar as expected from their almost spherical expansions (Figure 1). For more massive models, in particular s11, s15, and s27, a variety of evolutions can be seen because of highly asymmetric vigorous convective, and SASI motions.

3. Nucleosynthesis

The nucleosynthesis yields in each SN trajectory are computed in a postprocessing step by solving an extensive nuclear reaction network (Wanajo et al. 2001) with the temperature and density histories described in Section 2. The numbers of processed trajectories are 2343 (e8.8), 6310 (z9.6), 4672 (u8.1), 2739 (s11), 2565 (s15), and 2312 (s27). The up-to-date network consists of 7435 isotopes between the proton- and neutron-drip lines from single neutrons and protons up to isotopes with $Z = 110$. All the reaction rates are taken from REACLIB V2.0\(^{12}\) (Cyburt et al. 2010), making use of experimental data when available. As we see in Section 3.1, the nucleosynthetic abundances are mostly determined in nuclear equilibrium in the regions relatively close to $\beta$ stability where experimental masses are available. Uncertainties arising from nuclear data are thus expected to be small. Rates for electron capture (Langanke \& Martinez-Pinedo 2001) as well as for neutrino interactions on free nucleons (McLaughlin et al. 1996) and $\alpha$ particles (Woosley et al. 1990) are also included. The radiation field computed in the SN simulations is used as input for computing the neutrino interactions in our nucleosynthesis calculations. We retain the full dependence of the radiation field on radius, latitude, and time as computed in our ray-by-ray-plus approximation—i.e., we use the local neutrino energy density, number density, and energy moments of the distribution function at the current position of each tracer particle as input for the network calculations.

Each nucleosynthesis calculation is initiated when the temperature decreases to $T_{6} = 10$ with initial mass fractions of $Y_e = 1 - Y_p$ for free neutrons and $Y_p = Y_e$ for free protons, respectively.\(^{13}\) For the trajectories with $T_{6, \text{max}} < 10$, the initial compositions adopted from our hydrodynamical simulations are utilized.

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\(^{11}\) Entropy generation due to electron-positron annihilation after the freezeout from nuclear statistical equilibrium (NSE) is not taken into account for deriving Equation (2). However, we expect that its effect is minor because of the strong dependencies of nucleosynthetic abundances on $Y_e$ rather than on entropy or the expansion timescale.

\(^{12}\) https://groups.nscl.msu.edu/jina/reaclib/db/

\(^{13}\) At such a high temperature (and density), the matter immediately attains NSE and thus any initial composition with the total charge of $Y_e$ is available.
3.1. Types of Nucleosynthesis

We first analyze the nucleosynthesis in detail for z9.6, since this model covers the widest range in $Y_e$ ($Y_e = 0.373$–$0.603$; Figure 2). Figure 6 shows the neutron-, proton-, and $\alpha$-to-seed abundance ratios ($Y_n/Y_e$, $Y_p/Y_e$, and $Y/\alpha$) at a point when the temperatures decrease to $T_0 = 3$ for all the trajectories of model z9.6. We consider this temperature as it approximately corresponds to the end of nuclear equilibrium; this temperature is close to the conventional critical temperature for the termination of quasi-nuclear equilibrium (NSE; $T_0 = 4$; Meyer et al. 1998) to the beginning of a $\nu p$-process ($T_0 = 3$; Fröhlich et al. 2006), and the beginning of an $\nu$-process ($T_0 = 2.5$; Woosley et al. 1994). Here the seed abundance $Y_0$ is defined as the total abundance of all nuclei heavier than He. The other SN models show similar dependencies of $Y_n/Y_e$, $Y_p/Y_e$, and $Y/\alpha$ on $Y_e$ (not shown here), indicating that $Y_e$ (rather than entropy and the expansion timescale) is most crucial for the nucleosynthesis in our case. For this reason, we describe our nucleosynthetic results for model z9.6 in terms of $Y_e$. According to Figure 6, we can distinguish the following types of nucleosynthesis regimes in our SN models: among the (strongly) neutrino-processed ejecta, we find an NSE regime for $Y_e < 0.43$, a QSE regime for $0.43 \leq Y_e < 0.5$, and nucleosynthesis by charged-particle capture processes for $Y_e \geq 0.5$. No $\nu$-process is expected in our models because of neutron-to-seed ratios far below unity over the entire range of $Y_e$ (Figure 6). We note that in all of these regimes (and different from “classical” explosive nucleosynthesis), the neutrino-processed material is initially in NSE and the QSE/charged-particle capture regime is reached after an NSE phase as the temperatures decrease during the expansion of the ejecta.

In addition to these three regimes, there is a fourth regime of material with $Y_e \approx 0.5$ and lower entropy; this comprises the ejecta that immediately expand after being shocked and undergo classical explosive nucleosynthesis. The nucleosynthesis in this regime is not the primary subject of this paper; for models e8.8, z9.6, and u8.1 it is unimportant in terms of yields and for s11, s15, and s27, we merely cover the explosive burning of a part of the O shell, so there is no basis for an extensive analysis of this regime. Moreover, the physical principles of classical explosive nucleosynthesis are well known in the literature (Arnett 1996).

It is worth noting that although the regimes of NSE/QSE are also encountered in classical explosive nucleosynthesis, the conditions for the realization of these regimes are very different: for classical explosive burning, the nucleosynthesis is determined by the peak temperature (which depends on the explosion energy and the mass–radius profile of the progenitor) and in the NSE/QSE regime it is determined on the $Y_e$ in the progenitor. In the case of the neutrino-processed ejecta, the initial temperature is always high enough for NSE to obtain, and the interplay of multidimensional fluid flow and weak interactions of neutrinos, electrons, and positrons determines the conditions at freeze-out from NSE/QSE, i.e., $Y_e$, entropy, and the expansion timescale.

3.1.1. NSE

For $Y_e < 0.43$, all the ratios $Y_n/Y_e$, $Y_p/Y_e$, and $Y/\alpha$ are considerably smaller than unity ($<0.01$) at $T_0 = 3$ (Figure 6). The global distribution of final nucleosynthetic abundances is thus mostly determined in NSE at a high temperature ($T_0 > 5$). The subsequent QSE does not substantially change the abundance distribution because of the small amounts of free nucleons and $\alpha$ particles. This is due to relatively small entropies ($S \sim 14 k_B/\text{nuc}$; Figure 3) for these neutron-rich ejecta and the neutron-richness itself. Under such conditions the three-body process $\alpha(\alpha n, \gamma)$ $^6$Be (followed by $^7$Be $(\alpha, \gamma)^{12}$C), rather than triple-$\alpha$, is fast enough to form the NSE cluster.

\footnote{In this paper, we use the term NSE even for the case of $\alpha$-deficient QSE (Meyer et al. 1996; Wanajo et al. 2013a), in which NSE serves as a reasonable guideline for abundance determinations.}
by assembling free nucleons and α particles. Such neutron-rich conditions also disfavor α emission (i.e., to avoid being more neutron-rich). In NSE, the resulting abundance distribution is independent of specific reactions. What determines the abundance distribution are the binding energies per nucleon (\(B/A\); shown in Figure 7, left). As the \(Y_e\) of nuclides (indicated by white lines) decreases from \(\sim 0.50\) to \(\sim 0.40\), the nuclides with the maximal \(B/A\) shift from \(^{56}\text{Ni}\) \((Z = N = 28)\) to \(^{48}\text{Ca}\) \((Z = 20\) and \(N = 28)\) and \(^{84}\text{Se}\) \((Z = 34\) and \(N = 50)\); see Hartmann et al. 1985).

We find such NSE-like nucleosynthesis features in Figure 8, which shows the abundance distribution in selected trajectories of model z9.6 when the temperatures decrease to \(T_9 = 3\). The top three panels correspond to the trajectories with NSE-like conditions of \(Y_e < 0.43\) and relatively low entropies (\(\sim 14\, k_B/\text{nuc}\)). Major abundance peaks are already formed at \(T_9 = 5\) (red lines). The abundance patterns are almost frozen when the temperature decreases to \(T_9 = 4\) (defined as the end of NSE) and do not change significantly during the subsequent evolution.

### 3.1.2. QSE

For \(0.43 \leq Y_e < 0.50\), the α concentration becomes important \((Y_e/Y_b \sim 0.01–10)\) but \(Y_e/Y_b\) is less than unity at \(T_9 = 3\). In this case, the final abundances are mainly determined in QSE, a subsequent stage after the α-rich freeze-out from NSE (Woosley & Hoffman 1992; Meyer et al. 1998). At the end of NSE \((T_9 \sim 5)\), the single NSE cluster splits into two QSE clusters with one consisting of free nucleons and α particles and the other consisting of heavy nuclei. In the latter QSE cluster, the heavy nuclei are absorbed in the “α-bath” and its distribution is determined by the α separation energies (shown in the right panel of Figure 7), independent of specific nuclear reactions. With a modest neutron-richness of \(Y_e \sim 0.43–0.49\) (values are indicated by white lines in Figure 7), the nuclides near \(N = 28\) and \(50\) such as \(^{64}\text{Zn},^{88}\text{Sr},^{89}\text{Y},^{90}\text{Zr},\) and \(^{92}\text{Mo}\) are preferentially formed in QSE owing to their greater α separation energies (see also Hoffman et al. 1996; Wanajo 2006).

We find in the left-middle and middle panels of Figure 8 \((Y_e = 0.450\) and 0.475, respectively) that the abundance distributions substantially change during QSE \((T_9 \sim 5\) to 4\). The abundance patterns are determined roughly when the temperature decreases to \(T_9 = 3\), and further evolutions are unimportant (although an enhancement of nuclei with \(A = 10–50\) can be seen as a result of α-particle capture).

### 3.1.3. Charged-particle Capture Process

For \(Y_e \geq 0.50\), \(^{56}\text{Ni}\) dominates in the heavy QSE cluster, and the α-rich freeze-out from QSE \((T_9 \sim 4)\) leads to an α-process.
(Woosley & Hoffman 1992). This greatly enhances the abundances of \(\alpha\)-elements with \(A \sim 12-40\) (with multiples of four) as can be seen in the right-middle panel of Figure 8 (for \(Y_e = 0.500\)). For \(Y_e > 0.51\), both \(Y_n/Y_h\) and \(Y_p/Y_h\) are greater than unity at \(T_\text{fl} = 3\) (Figure 6). The freeze-out is thus followed by \(\alpha\)-capture and proton-capture processes. We find in the bottom three panels of Figure 8 (\(Y_e = 0.525, 0.550,\) and \(0.603\) that nuclei in a wide range of \(A \sim 10-70\), including odd-\(Z\) elements, are substantially enhanced by these charged-particle capture processes after the temperature drops below \(T_\text{fl} = 3\).

Figure 9 compares the final abundances with (blue) and without (cyan) neutrino reactions for the trajectory that has the highest \(Y_e = 0.603\) in model z9.6 (same as that in the right-bottom panel of Figure 8). This indicates that the enhancement of nuclei with \(A = 60-70\) is due to a \(\nu\nu\)-process in which the faster \((n, p)\) and \((n, \gamma)\) reactions with the free neutrons supplied by \(\nu_e\) capture on free protons replace the slower \(\beta^+\)-decays (Fröhlich et al. 2006; Pruett et al. 2006; Wanajo 2006). Note that the \(\nu\nu\)-process in our result is very weak and only the nuclei up to \(A \sim 70\) are produced despite its substantial proton-richness (up to \(Y_e = 0.603\)). By contrast, Wanajo et al. (2011a) have shown that the \(\nu\nu\)-process in neutrino-driven wind with similar proton-richness can produce nuclei of up to \(A \sim 120\) (see also Pruett et al. 2006; Arcones et al. 2012). This discrepancy is due to the lower entropies and longer expansion timescales in the early dynamical ejecta, which reduce \(Y_n/Y_h\) at the beginning of a \(\nu\nu\)-process (\(T_\text{fl} \sim 3\), than those in the late-time neutrino-driven wind.

### 3.2. Dependencies of Isotope Productions on \(Y_e\)

In Section 3.1 we found that the iron-group (from Ca to Cu) and light trans-iron (from Zn to Mo) species can be produced in the innermost ejecta of z9.6, the model with the widest range in \(Y_e\) in the ejecta. No heavier elements are produced in any of our six models. In Figures 10 and 11, the final mass fractions of stable isotopes from K to Mo are presented as functions of \(Y_e\) for all the trajectories of model z9.6. The mass fractions of selected radioactive isotopes (before decay) are also shown in the right-bottom panel of Figure 11. We find from these figures that few isotopes exhibit maximum abundances near \(Y_e = 0.5\), in particular the light trans-iron species (from Zn to Mo).

Overall, Ni and light trans-iron species are predominantly formed in neutron-rich ejecta, although the weak \(\nu\nu\)-process in proton-rich ejecta plays a subdominant role for those up to Ge. Several isotopes such as \(^{44}\text{Ca}, ^{50}\text{Ti}, ^{54}\text{Cr},\) and the radioactive nuclide \(^{56}\text{Fe}\) are made only in very neutron-rich ejecta with \(Y_e \sim 0.40-0.43\) (Wanajo et al. 2013a, 2013b). This sensitivity clearly demonstrates the importance of nucleosynthesis studies based on multidimensional SN simulations with detailed multigroup neutrino transport, which are indispensable for accurately determining the \(Y_e\) distribution in the ejecta.

### 4. Possible Contribution to the Solar Abundance Distribution

For each SN model, we calculate the mass-integrated abundances from the nucleosynthetic outcomes of all trajectories. In the following subsections, we discuss the possible contribution of products from the innermost ejecta of SNe with different explosion dynamics during the first second after collapse to the Galaxy by comparing our nucleosynthesis yields to the solar values. In particular, we focus on the production of several key species such as \(^{44}\text{Ca}, Zn\) isotopes, light trans-iron elements, \(p\)-nuclide \(^{52}\text{Mo},\) and the radioactive isotopes \(^{56}\text{Ni}\) and \(^{56}\text{Fe}.\) It is important to stress that our goal here is not to predict the full CCSN nucleosynthesis across the whole range of progenitors, which would require both longer simulations and a larger grid of models at different metallicities. However, even with our small set of models, we can already establish whether the innermost ejecta can provide an important contribution to the solar inventory for plausible rates of the different explosion behaviors exemplified by our six models. It is particularly noteworthy that focusing on the neutrino-processed ejecta allows us to neglect the influence of the progenitor metallicity on the yields, since any memory of the progenitor composition is erased in this ejecta component; i.e., the production of heavy elements in the neutrino-heated ejecta is always a primary process.
4.1. Comparison with the Elemental Solar Abundance

For all SN models, the mass-integrated yields are compared to the solar abundance \cite{Lodders2003}. Figure 12 shows the elemental mass fractions in the total ejecta with respect to their solar values (not to the initial compositions of progenitors), i.e., “production factors,” for these models. The total ejecta mass from each SN is taken to be the sum of the ejected mass from the core and the outer envelope—that is, \( M_{\text{prog}} - M_{\text{PNS}} \) in Tables 1 and 2. Here the outer envelope is assumed to be metal free (i.e., H and He only), which is reasonable for the models near the low-mass end of the SN range, e8.8, z9.6, and u8.1: the part of the H and He shells that were not included in our nucleosynthesis calculations would not contribute substantially to the yields for the range of nuclei considered here \( Z \geq 19 \) regardless of progenitor metallicity. For the more massive models s11, s15, and s27, one should bear in mind that there is an important additional contribution to heavy elements (those heavier than He) from the outer envelope \cite{Woosley2002, Woosley2007}.

In each panel of Figure 12, we show a normalization band in yellow that covers the range within 1 dex of the maximum production factor and one tenth of that. The elements that reside in this band can (at least in part) originate from SNe represented by each model, provided that their production factors are greater than \( \sim 10 \) \cite{Woosley2007}.

As an order-of-magnitude estimate, the production factors of \( \sim 40 \) for light trans-iron elements from Zn to Zr in our representative

Figure 8. Nucleosynthetic abundances for model z9.6 when the temperatures decrease to \( T_9 = 5, 4, \) and 3 as well as those at the end of the calculations (“final” in panels). Selected trajectories are those with initial \( Y_e = 0.373 \) (left-top), 0.400 (middle-top), 0.425 (right-top), 0.450 (left-middle), 0.475 (middle), 0.500 (right-middle), 0.525 (left-bottom), 0.550 (middle-bottom), and 0.603 (right-bottom). The asymptotic entropy \( S \) is also shown in the legend of each panel.

Figure 9. Same as the final abundance pattern in the right-bottom panel of Figure 8 but neglecting neutrino interactions below \( T_9 = 10 \) (cyan).
model z9.6 suggest that such low-mass CCSNe can be the major sources of these elements if these events account for a few 10% of all CCSNe and would still contribute sizable amounts if they made up \( \sim 10\% \) of all CCSNe. This suggests that SNe at the low-mass end of the progenitor spectrum as represented by models e8.8, z9.6, and u8.1 could be an important source of light trans-iron elements (Section 4.6).

Note that there is a remarkable agreement of the nucleosynthesis of z9.6 with that of e8.8, which is a consequence of the similarity of their pre-SN core structures with a steep density gradient outside the core and a very dilute outer envelope. In both cases, the resulting fast explosion leads to the ejection of appreciable amounts of neutron-rich material. However, model u8.1, with a core structure very similar to (but with a slightly shallower density gradient than) that of z9.6, results in a different nucleosynthetic trend with deficiencies of several elements between Zn and Zr. This indicates that only a slight difference of pre-SN core-density structures can lead to substantially different nucleosynthesis outcomes.

For more massive models, the production factors are smaller than \( \sim 10 \) for all elements, despite their greater ejecta masses compared to low-mass models \( (M_{ej} \text{ in Table 2}) \). The maximum production factor of \( \sim 1 \) for s11 provides no element responsible for the Galactic chemical evolution. For s15 and s27, only Zr with the production factor \( \sim 10 \) is the element that can be originated from these types of CCSNe. However, these SNe are possible contributors of some isotopes as discussed in Section 4.2. Moreover, the non-neutrino processed “outer ejecta” from these stars, which are not fully included in our study, contribute significantly to the chemical enrichment of the Galaxy as discussed in Section 4.1.

4.2. Comparison with the Isotopic Solar Abundance

Figure 13 compares the isotopic abundances with the solar values for all SN models. The maximum isotopic production factor for each model is generally greater than that of elements (Section 4.1) and therefore places tighter constraints on the contribution of relevant SNe to the Galaxy. In Table 3, we list the five largest production factors for each model. Note that most of the isotopes listed here are made in nuclear equilibrium, so uncertainties in individual nuclear reaction rates are irrelevant. Models e8.8 and z9.6 exhibit appreciable production factors of 354 (\(^{86}\text{Kr}\)) and 168 (\(^{82}\text{Se}\)), respectively. This implies that such low-mass SNe account for \( \sim 10\% \) of all CCSN events (according to the reason described in Section 4.1), supposing that these trans-iron elements originate
solely from this class of events. A similar contribution of u8.1-like events can be expected with its largest production factor of 149 (74Se). The maximum production factor of 1.48 for s11 indicates that no species are dominantly produced in the innermost ejecta of such SNe. For s15 and s27, the largest production factors are 24.8 and 43.1 (74Se), respectively, which are sizably greater than those of elements (~10; Figure 13). This indicates that such intermediate-mass and massive CCSNe can be important sources of several species listed in Table 3.

To be more quantitative, we consider 82Se, with the largest production factor for model z9.6, as representative. By assuming $f_{9.6}$ to be the fraction of z9.6-like events to all CCSNe, we have (Wanajo et al. 2011b)

$$\frac{f_{9.6}}{1 - f_{9.6}} = \frac{X_{\odot}(82\text{Se}) / X_{\odot}(16\text{O})}{M_{9.6}(82\text{Se}) / (M(16\text{O}))} = 0.164, \quad (3)$$

where $X_{\odot}(82\text{Se}) = 1.38 \times 10^{-8}$ and $X_{\odot}(16\text{O}) = 6.60 \times 10^{-3}$ are the mass fractions in the solar system (Lodders 2003), $M_{9.6}(82\text{Se}) = 1.91 \times 10^{-5} M_{\odot}$ is the ejecta mass of 82Se for z9.6, and $(M(16\text{O})) = 1.5 M_{\odot}$ is the production of 16O by massive CCSNe averaged over the stellar initial mass function (IMF) between $13 M_{\odot}$ and $40 M_{\odot}$ (see Wanajo et al. 2009). Equation (3) gives $f_{9.6} = 0.14$, which corresponds to a mass window of $\Delta M_{\text{prog}} \sim 1 M_{\odot}$ near the low-mass end of the SN progenitor spectrum, $M_{\text{prog}} \sim 9 M_{\odot}$. It is currently uncertain whether this mass window for z9.6-like progenitors is reasonable or not, as only a slight difference of core-density structure leads to a very different nucleosynthetic result of u8.1. If we apply Equation (3) to model e8.8 by replacing 82Se with 86Kr (Table 3), then we get $f_{8.8} = 0.085$. This corresponds to a mass window of $\Delta M_{\text{prog}} \sim 0.5 M_{\odot}$, which is in reasonable agreement with the prediction from synthetic SAGB models at solar metallicity (Poelarends et al. 2008). If other production channels were to contribute significantly to these elements, then the allowed rate of ECSNe or a z9.6-like explosion would, of course, be lower.

It is thus conceivable that a combination (or either) of ECSNe and low-mass CCSNe accounts for the production of light trans-iron species in the Galaxy. For the more massive models s11, s15, and s27, a lack of self-consistent nucleosynthesis yields in the outer envelopes precludes a quantitative estimate of their contributions.
4.3. Neutron-rich versus Proton-rich Ejecta

Our results described in Sections 4.1 and 4.2 shows that SNe near the low-mass end produce appreciable amounts of trans-iron species despite their small ejecta masses (compared to those of massive models; $M_{ej}$ in Table 2). The reason can be found in Figure 14, which shows the fraction of the ejecta that originates in neutron-rich matter ($Y_e < 0.4975$, $M_{ej,n}$ in Table 2) for a given species. For low-mass models e8.8, z9.6, and u8.1, the light trans-iron isotopes of $A = 64–90$ are almost exclusively produced in neutron-rich ejecta, whereas proton-rich matter plays a minor role (see also Figures 10 and 11). The subdominant roles of SNe from massive progenitors (represented by s11, s15, and s27) to the production of these species can be understood as a result of the small mass of neutron-rich ejecta for these stars ($M_{ej,n}$ in Table 2 and Figure 2). Among the massive models, s27 has a relatively larger amount of neutron-rich ejecta compared to the two others, with the $Y_e$ down to $\sim 0.4$. This is a consequence of the fact that model s27 exhibits an earlier explosion with a rapidly increasing shock radius because of prominent SASI activity that is absent in other models. These neutron-rich ejecta lead to a relatively flat trend of production factors in this model, similar to what we found for e8.8 and z9.6 (Figures 12 and 13). However, the bottom panel of Figure 13 indicates that about half of the species between $A = 64$ and 90 originate from neutron-rich and proton-rich ejecta. In fact, 43% of $^{64}$Zn in s27 comes from the proton-rich ejecta. As described in Section 3.1.3, a weak $\nu p$-process is responsible for the production of these isotopes in the proton-rich ejecta. Greater entropies of the ejecta (Figure 3), slower expansion compared to that for low-mass models, and higher neutrino luminosities and mean energies also enhance the efficiency of the $\nu p$-process.

4.4. $^{48}$Ca

Models e8.8 and z9.6 exhibit appreciable production factors of $^{48}$Ca (18.9 and 18.6, respectively), which is made in neutron-rich NSE (or $\alpha$-deficient QSE; Meyer et al. 1996; Wanajo et al. 2013a) as can be seen in Figures 8 and 10 (middle-top panels, $Y_e \sim 0.4$). The amounts are still not large enough to regard these low-mass SNe as the main contributors of $^{48}$Ca, although a reduction of the entropies by about 30% would lift the values to a satisfactory level (Wanajo et al. 2013a). A rare class of high-density SNe Ia, in which similar physical conditions are expected, has also been suggested as a possible source of $^{48}$Ca (Woosley 1997).

4.5. Zn Isotopes

One of the outstanding features of our nucleosynthesis result is the production of all the stable isotopes of Zn ($A = 64, 66, 67, 68$, and 70), an element whose origin remains a mystery. Among our explored models, the low-mass models e8.8 and z9.6 exhibit nearly flat production factors over $A = 64–70$. This fact implies that the element Zn, or all its stable isotopes, could originate from ECSNe or low-mass CCSNe as represented by e8.8 and z9.6, respectively. As can be seen in the right-bottom panel of Figure 10, Zn isotopes are predominantly made in neutron-rich ejecta with $Y_e \sim 0.4–0.5$ and models e8.8 and z9.6 produce a sufficient amount of ejecta in this range to contribute most of the solar inventory of Zn. There is still a discrepancy of a factor of 10 between the largest ($^{66}$Zn) and smallest ($^{67}$Zn) production factors, which might be cured if the $Y_e$ distributions for these models were slightly modified. Alternatively, some of the neutron-rich Zn isotopes may originate from hydrostatic neutron-capture processes in massive stars (Woosley & Weaver 1995; Nomoto et al. 2013). CCSNe
represented by u8.1 cannot be the single source of all Zn isotopes because of the descending trend of production factors (Figure 15). Model s27, which shows a flat trend of production factors, cannot be representative of Zn contributors either because of the small production factors (1.2–5.2).

To date, only hypernova models (Umeda & Nomoto 2002; Tominaga et al. 2007; Nomoto et al. 2013) have been proposed as sources of $^{64}$Zn (except for an ECSN model in Wanajo et al. 2011b), the dominant isotope of Zn in the solar system. In their hypernova models high entropies of ejecta lead to a strong $\alpha$-rich freeze-out from nuclear equilibrium, resulting in an appreciable production of $^{64}$Zn. However, this mechanism does not coproduce the other isotopes and requires additional contributors such as hydrostatic neutron captures in massive stars. As described in Section 3.1.3, the $vp$-process also produces $^{64}$Zn only and therefore cannot be the main mechanism responsible for making Zn.

4.6. Light Trans-iron Elements

The nearly flat production factors over the wide range of $A = 64–90$ (except for Ga and As; Figure 13) in e8.8 and z9.6 also suggest that such low-mass SNe are the dominant sources of light trans-iron elements from Zn to Zr as suggested by Wanajo et al. (2011b) if their mass window is $\Delta M_{\text{prog}} \sim 0.5–1$. This is a consequence of the fact that the abundant neutron-rich ejecta for these models (Figure 2) cover the range of $Y_e$ (down to $\sim 0.40$) where the production of these species becomes maximal (Figures 10 and 11). By contrast, the results for model u8.1 (with a pre-SN core structure similar to that of z9.6) imply that such SNe can only be a source of the proton-rich isotopes of these trans-iron elements. This is due to the small amount of neutron-rich ejecta with $Y_e \sim 0.40–0.45$.

In previous studies, the production of such light trans-iron species has often been attributed to the weak $s$-process (e.g., Käppeler et al. 2011). Woosley & Heger (2007) showed that
the $s$-process in massive stars produces appreciable amounts of light trans-iron isotopes but only between $A = 65$–$85$ with a descending trend of production factors. It also has been claimed that significant production of these elements can occur in stars above $25 \ M_\odot$ via the weak $s$-process (Pignatari et al. 2010). However, Sukhbold et al. (2016) pointed out that such light trans-iron species (in particular those above As) were sizably underproduced in their work because a large number of stars formed black holes instead of exploding and therefore did not contribute to weak $s$-process nuclei. A major contribution of the weak $s$-process from stars above $25 \ M_\odot$ is also unlikely considering observational evidence against the successful explosions of stars in this mass range (Smartt 2015). Note that Chieffi & Limongi (2013) find that sufficient rotation-induced mixing in massive stars appreciably increases the amount of weak $s$-process elements. We speculate that the weak $s$-process in massive stars may still be responsible for several trans-iron elements, in particular such as Ga and As, which are underproduced in our models e8.8 and z9.6 (Figures 12 and 13).

4.7. $^{92}$Mo

It is interesting to note that model s27 gives abundant $^{92}$Mo (with the third largest production factor, 30.7, in Table 3), an important $p$-nuclide whose astrophysical origin has been a long-lasting problem. In our case $^{92}$Mo is made in slightly neutron-rich ejecta ($Y_p \sim 0.47$; see middle panel of Figure 8 and middle-bottom panel of Figure 11) during nuclear equilibrium (as suggested by Hoffman et al. 1996; Wanajo 2006). With a production factor of $\sim 30$, such a type of CCSNe—i.e., relatively early explosions with dense outer envelopes—$s$27-like explosions would need to account for $\sim 30\%$ of all CCSNe to explain the solar abundance of $^{92}$Mo. A potential problem is that the other $p$-isotope $^{94}$Mo cannot be made in neutron-rich nuclear equilibrium (Hoffman et al. 1996; Wanajo 2006). However, the $\nu p$-process in the neutrino-driven wind may add these isotopes with a high $^{94}$Mo/$^{92}$Mo ratio (Pruet et al. 2006; Wanajo 2006; Wanajo et al. 2011a). Note that we do not find large production factors for the $p$-isotopes of Ru and Pd as in Pruets et al. (2006). This is because these isotopes were made by the $\nu p$-process (Section 3.1.3) in their late-time ($\gtrsim 1$ s) wind ejecta with high entropies ($S \sim 70 \ k_B$), while the hydrodynamical simulations of models s15 and s27 stopped at $\sim 0.8$ s after core bounce with entropies still below $S \sim 50 \ k_B$.

4.8. $^{56}$Ni

The masses of $^{56}$Ni ejected from all models are listed in Table 4 (7th column), along with those of $^{57}$Ni (last column). These values, in the range $\sim 0.002$–$0.006 \ M_\odot$, are about one order of magnitude smaller than typical observed values (e.g., 0.07 $M_\odot$ for SN 1987A; Bouchet et al. 1991). However, the $^{56}$Ni masses in our study should be taken as lower limits for the massive models s11, s15, and s27, for which we omit the outer envelopes including large parts of the Si/O layer. We therefore likely underestimate the amount of $^{56}$Ni produced by explosive nucleosynthesis in the SN shock. Moreover, 2D models face a generic difficulty in determining the amount of $^{56}$Ni made by explosive burning in the shock, since the shocked material is funneled around the neutrino-driven outflows onto the proto-neutron star with little mixing by the Kelvin–Helmholtz instability into the neutrino-driven ejecta (Müller 2015). The production of $^{56}$Ni by explosive burning also depends on the postshock temperatures and could be related to the slow rise of the explosion energy to only $(0.3$–$1.5) \times 10^{50}$ erg at the end of the simulations.

In addition, the core-collapse simulations of s15 and s27 stopped at a time when accretion was still ongoing and the mass ejection rate in the neutrino-driven outflows was still high so that we may miss some late-time contribution to $^{56}$Ni in the neutrino-driven ejecta (see, e.g., Wongwathanarat et al. 2017). In fact, the late-time ejecta of s15 and s27 are mostly proton rich and are dominated by $^{56}$Ni and $\alpha$ particles (Figure 6 and the bottom right panel of Figure 11).

For the low-mass models e8.8, z9.6, and u8.1, the listed values can be regarded as the final ones; the $^{56}$Ni mass is definitely very small ($\sim 0.002$–$0.003 \ M_\odot$) for these cases. It has already been argued before for e8.8 (Kitaura et al. 2006; Wanajo et al. 2011b; Wanajo 2013) that the small $^{56}$Ni may be consistent with the value estimated for some observed low-luminosity SNe (e.g., Hendry et al. 2005; Pastorello et al. 2007) with small ejecta masses. The $^{56}$Ni mass and the other explosion properties of the three low-mass models are also consistent with the remnant composition and reconstructed light curve of the Crab supernova SN 1054 (Smith 2013; Tomimaga et al. 2013; Moriya et al. 2014). From the nucleosynthetic point of view, low-mass iron-core SNe thus appear to be an equally viable explanation for these events compared to ECSNe.

4.9. $^{60}$Fe and Other Radioactive Isotopes

The low-mass models e8.8 and z9.6 produce appreciable amounts of $^{60}$Fe, $3.61 \times 10^{-5} \ M_\odot$, and $3.14 \times 10^{-5} \ M_\odot$, respectively (6th column in Table 4), which is comparable to the IMF-averaged ejection mass from CCSNe, (2.70–3.20) $\times 10^{-5} \ M_\odot$ in Sukhbold et al. (2016). Note that in massive stars, $^{60}$Fe is produced by successive neutron captures from Fe isotopes (Timmes et al. 1995; Limongi & Chieffi 2006) in the outer envelope, which is not included in our analysis of the massive models s11, s15, and s27. In e8.8 and z9.6, $^{60}$Fe forms in neutron-rich NSE ($Y_p \sim 0.42$–0.43; right-bottom panel of Figure 11) as suggested by Wanajo et al. (2013b). This indicates that ECSNe or low-mass CCSNe...
account for at least $\sim 10\%$ (see $f_{e8.8}$ and $f_{z9.6}$ in Section 4.2) of live $^{60}\text{Fe}$ in the Galaxy. Models other than e8.8 and z9.6 produce little $^{60}\text{Fe}$ in their innermost ejecta because of their small masses of neutron-rich ejecta.

Table 4 also lists the masses of other important radioactive isotopes, $^{26}\text{Al}$, $^{41}\text{Ca}$, $^{44}\text{Ti}$, and $^{53}\text{Mn}$, which are, however, negligibly small compared to the contribution from the outer envelopes of massive stars (and from neutrino-driven wind for $^{44}\text{Ti}$; Wongwathanarat et al. 2017). Recent work by Sukhbold et al. (2016) suggests the IMF-averaged amount of $^{26}\text{Al}$ from massive stars to be $(2.80 - 3.63) \times 10^{-3} M_\odot$. Their resultant mass ratio of $^{60}\text{Fe}$ to $^{26}\text{Al}$, $\sim 1$ (although a factor of two smaller than previous estimates, e.g., Woosley & Heger 2007), conflicts with the mass ratio $\sim 0.34$ inferred from gamma-ray observation (flux ratio 0.148 $\pm$ 0.06; Wang et al. 2007). Taking the observational value as a constraint, this would suggest that Sukhbold et al. (2016) overestimated the $^{60}\text{Fe}$ mass by about a factor of three. This can be attributed to uncertainties in the relevant reaction rates (Woosley & Heger 2007; Tur et al. 2010) while for our cases $^{60}\text{Fe}$ forms in nuclear equilibrium and thus individual reactions are irrelevant. If this is the case, then ECSNe or low-mass SNe can contribute to live $^{60}\text{Fe}$ up to $\sim 30\%$ of the amount in the Galaxy.

5. Conclusions

We have examined the nucleosynthesis in the innermost ejecta ($0.01-0.03 M_\odot$) of CCSNe, including ECSNe, by adopting thermodynamic trajectories obtained from self-consistent (general-relativistic with one exception) 2D CCSN explosion models with multigroup neutrino transport. We explored the six models e8.8 ($8.8 M_\odot$ ECSVN, 1 Z$_\odot$), z9.6 ($9.6 M_\odot$, 0 Z$_\odot$), u8.1 ($8.1 M_\odot$, 10$^{-4}$ Z$_\odot$), s11 (11.2 M$_\odot$, 1 Z$_\odot$), s15 (15.0 M$_\odot$, 1 Z$_\odot$), and s27 (27.0 M$_\odot$, 1 Z$_\odot$), with a progressively shallower core-density gradient in that order (Figure 8 in Janka et al. 2012). In this paper, we focused on the

Figure 14. Fraction of isotopic yields contributed by neutron-rich matter ($Y_r < 0.4975$, $M_{nj,n}$ in Table 2). $X_{\text{neutron-rich}}/X = 1.0$ and 0.0 mean that a given isotope originates exclusively from neutron-rich and proton-rich ejecta, respectively.
effects of these differences in the pre-SN core-density structure on nucleosynthesis and did not attempt to address the exact dependence of nucleosynthesis on mass or metallicity (owing to the limited number of available models). However, our results indicate that low-mass progenitors close to the Fe-core formation limit (which are characterized by the off-center ignition of O burning) stand apart from massive stars due to shared structural and evolutionary features: the low-mass models e8.8, z9.6, and u8.1 (near the low-mass end of SN progenitors) have abundant neutron-rich matter (40%–50%) in their ejecta because of fast shock expansion after the onset of the explosion. By contrast, proton-rich matter dominates the ejecta of the more massive models s11, s15, and s27 by far because slower shock expansion allows neutrino interactions to efficiently raise $Y_e$.

For this reason, our nucleosynthesis calculations in the postprocessing steps resulted in remarkably different outcomes between the low-mass (e8.8, z9.6, and u8.1) and massive (s11, s15, and s27) models. We determined that low-mass SNe, in particular those represented by e8.8 (ECSN) and z9.6, could be the dominant source of light trans-Fe elements from Zn to Zr (except for Ga and As), if these events account for $\sim 8%$–$14\%$ (or a mass window of $\Delta M_{\text{prog}} = 0.5–1 M_\odot$) of all CCSNe (including ECSNe). These species are made predominantly in neutron-rich ejecta during nuclear equilibrium phases. Good agreement of the nucleosynthetic yields with the solar abundance pattern over a wide range of isotopes with $A = 64–90$ strongly supports our conclusion, which sets this candidate site apart from other previously suggested astrophysical sites, e.g., hypernovae as the origin of Zn (Umeda & Nomoto 2002; Tominaga et al. 2007; Nomoto et al. 2013) and the weak $s$-process in massive stars (as the origin of light trans-Fe elements; Woosley & Heger 2007; Käppeler et al. 2011). These SNe (e8.8 and z9.6) could also be important sources of the neutron-rich isotope $^{48}$Ca and the live radioactive species $^{60}$Fe in the Galaxy and supplement the chemo-Galactic contribution of other sites such as rare SNe Ia (Woosley 1997) with high ignition densities and massive stars (e.g., Sukhbold et al. 2016). However, model u8.1 was found to have contributions of proton-rich isotopes of light trans-Fe elements only, despite its similar (but slightly shallower) core-density gradient compared to that of z9.6.

We found that the innermost ejecta of massive SNe make little contribution to the chemical inventory of the Galaxy, except for proton-rich isotopes of light trans-Fe elements (in s15 and s27). The most massive model, s27, exhibited an interesting production of $p$-nucleus $^{92}$Mo, which could explain the solar amount of $^{92}$Mo if such s27-like events accounted for $\sim 30\%$ of all CCSNe. This was traced back to a sizable amount of slightly neutron-rich ejecta ($Y_e \sim 0.47$) with moderately high entropies ($\sim 30 k_B/N_{\text{nuc}}$), in which $^{92}$Mo is made in nuclear equilibrium (Hoffman et al. 1996; Wanajo 2006). Note that our calculations did not include the outer envelopes, from which Fe-group and lighter elements would be ejected in these massive SNe. Instead, we focused mostly on the neutrino-processed ejecta during the first few hundreds of milliseconds after shock revival. This implies that our results on the production of Fe-group and trans-Fe elements in the more massive models (s11, s15, s27) are only indicative of the overall nucleosynthesis and may miss important contributions. For the ECSN progenitor e8.8 and the ECSN-like models z9.6 and u8.1 the nucleosynthesis in this range is essentially complete, barring minor contributions from the neutrino-driven wind that follows after our simulations were terminated.

Aside from the fact that the nucleosynthesis in s11, s15, and s27 is not yet complete, some further caveats should be added to our conclusions. Although nucleosynthesis calculations based on self-consistent 2D CCSN models represent a fundamental improvement over previous 1D studies (which mostly excluded neutrino-processed ejecta), the explosion dynamics in 3D has emerged as noticeably different from 2D (see Janka et al. 2016; Müller 2016 for recent reviews). The interaction between buoyant neutrino-heated ejecta and accretion downflows tends to brake both outflows and downflows (Melson et al. 2015; Müller 2015) and could therefore modify the final $S$ and $Y_e$ in the neutrino-driven ejecta. So far, model z9.6 offers the only opportunity for a direct comparison of the nucleosynthesis conditions. For this progenitor, 3D effects may moderately affect the explosion dynamics (Melson et al. 2015) and may only slightly raise the minimum $Y_e$ in the ejecta without fundamentally changing the neutron-rich nucleosynthesis in this progenitor (Müller 2016). The situation for other low-mass Fe CCSNe and ECSNe is likely similar.

For more massive progenitors, larger systematic effects on the nucleosynthesis conditions in 3D cannot be excluded. The results of Müller (2015) for an 11.2 $M_\odot$ model suggest that the outflow velocities (which affect the $Y_e$) and the terminal entropies in 3D can be considerably lower, but these results have yet to be borne out for a broader range of progenitors. The mixing of shocked material into the neutrino-heated ejecta is likely more efficient in 3D and will affect the contribution of explosive nucleosynthesis in the shock to the total ejecta because this matter is swept outward instead of being accreted onto the newborn neutron star. Along with the low explosion energies of the massive 2D models, this may explain the unusually small mass of $^{56}$Ni found in our calculations.

Effects that may change the final $Y_e$ in the neutrino-processed ejecta by altering the neutrino emission also warrant further exploration. Aside from uncertainties in the neutrino microphysics, the impact of the genuinely 3D lepton-number emission self-sustained asymmetry (LESA) instability (Tamborra et al. 2014; Janka et al. 2016) on the nucleosynthesis conditions deserves investigation. As LESA produces a global asymmetry in the flux difference between electron neutrinos and antineutrinos it will
contribute to a spread in the Y_e distribution in the ejecta. A detailed analysis of the LESA effect in z9.6 is currently under way; the neutron-rich bubbles in the early ejecta remain a robust feature for this progenitor also in the presence of LESA (Melson et al., in preparation).

In summary, while our conclusions for low-mass SNe may be robust, those for massive models should only be taken as suggestive. Eventually, successful explosion models of such massive SNe in 3D are required for drawing firmer conclusions. However, it is also important to scan the low-mass end of the progenitor range with a finer resolution in the progenitor mass. While Wanajo et al. (2011b, 2013a, 2013b) suggested that ECSNe could be important contributors of light trans-Fe elements (including Zn), ^{44}Ca, and ^{56}Fe, our results show that CCSNe near the low-mass end with Fe cores (z9.6) have almost the same nucleosynthetic outcomes. However, a model (u8.1) with an only slightly shallower core-density gradient resulted in a substantially different result. Stellar evolution models and CCSN simulations therefore need to better address over which mass range the progenitors possess core-density profiles that are sufficiently steep to produce ECSN-like nucleosynthesis near the low-mass end of the progenitor-mass spectrum.

Note also that we neglect late-time ejecta in the neutrino-driven wind in the present study. These could also contribute to the enrichment of trans-Fe species to the Galaxy by a weak r-process (e.g., Wanajo 2013a) and a vp-process (Fröhlich et al. 2006; Pruet et al. 2006; Wanajo 2006). Finally, our result should be tested by a study of Galactic chemical evolution to reproduce the observational signatures of light trans-Fe elements in the Galaxy.

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### Table 4
Masses of Radioactive Isotopes in the Innermost Ejecta (M_⊙)

| Model  | ^{26}Al | ^{44}Ca | ^{44}Ti | ^{53}Mn | ^{56}Fe | ^{56}Ni | ^{57}Ni |
|--------|--------|--------|--------|--------|--------|--------|--------|
| e8.8   | 4.39E−08| 1.96E−07| 2.06E−06| 1.11E−06| 3.61E−05| 2.93E−03| 1.01E−04|
| z9.6   | 1.29E−07| 1.16E−07| 2.41E−06| 1.61E−06| 3.14E−05| 2.51E−03| 9.15E−05|
| u8.1   | 4.77E−08| 5.77E−08| 1.97E−06| 1.45E−06| 6.65E−07| 1.60E−03| 7.33E−05|
| s11^a  | 1.22E−09| 1.05E−09| 1.19E−06| 2.53E−06| 1.96E−02| 3.86E−03| 8.75E−05|
| s15^a  | 1.07E−08| 2.25E−07| 1.71E−06| 4.86E−06| 2.44E−18| 6.05E−03| 1.16E−04|
| s27^a  | 2.53E−08| 4.56E−07| 3.31E−06| 5.89E−06| 1.14E−06| 5.73E−07| 1.49E−04|

Note: ^a Values for these models should be taken as lower limits (see text).
