

Explodability matters: how realistic neutrino-driven explosions change explosive nucleosynthesis yields

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ABSTRACT

Context. Explosive nucleosynthesis is affected by many uncertainties, particularly regarding assumptions and prescriptions adopted during the evolution of the star. Moreover, simple explosion models are often used in the literature, which can introduce large errors in the assumed explosion energy and mass cut.

Aims. In this paper, our goal is to analyze the explosion properties and nucleosynthesis of a large range of progenitors from three different stellar evolution codes: FRANEC, KEPLER, and MESA. In particular, we will show the differences between the neutrino-driven explosions simulated in this work with the much simpler bomb and piston models that are typically widely used in the literature. We will then focus on the impact of different explodabilities and different explosion dynamics on the nucleosynthetic yields.

Methods. We adopt the neutrino-driven core-collapse supernova explosion code GR1D+, i.e. a spherically symmetric model with state-of-the-art microphysics and neutrino transport and a time-dependent mixing-length model for neutrino-driven convection. We carry out explosions up to several seconds after bounce, and then calculate the nucleosynthetic yields with the post-processing code SkyNet.

Results. We find that our 1D+ simulations yield explosion energies and remnant masses in agreement with observations of type II-P, IIb, and Ib supernovae, as well as with the most recent 3D simulations of the explosion. We provide a complete set of yields for all the stars simulated, including rotating, low-metallicity, and binary progenitors. Finally, we find that piston and bomb models, compared to more realistic neutrino-driven explosions, can artificially increase the production of Fe-peak elements, whereas the different explodability tends to cause discrepancies in the lighter elements.

1. Introduction

Core-collapse supernovae (CCSNe) are responsible for enriching the interstellar medium with almost half of the elements in the periodic table (Burbidge et al. 1957; Arnett 1987; Kobayashi et al. 2020). Broadly speaking, these elements can be synthesized either in the hydrostatic burning processes during the post-main-sequence phase of the star (Weaver et al. 1978; Aufderheide et al. 1991; Limongi & Chieffi 2003; Nomoto et al. 2013; Iwamoto & Thielemann 1999) and then ejected by the supernova, or during the explosion itself (Arnett 1969; Limongi & Chieffi 2003; Arcones & Thielemann 2023; Fröhlich et al. 2006a).

Both the pre-supernova and explosive nucleosynthesis yields are swept up by the supersonic bubble that propagates into the interstellar medium, and can therefore be visible in the ensuing supernova light curve (Gal-Yam 2017) as well as, at later times, in the supernova remnant (Vink 2020; Milisavljevic et al. 2024). Eventually, they will be completely mixed into the interstellar medium and will make up the material from which the next generation of stars will form (Diehl & Prantzos 2020; Kobayashi et al. 2020). Therefore, accurate predictions of both the pre-supernova and the explosive yields are crucial for explaining the observed supernova light curves and spectra, the observations of supernova remnants, and the derived photospheric spectral abundances of stars.

The uncertainties that dominate current predictions of CCSN yields are due both to nuclear physics (i.e., uncertainties in reac-

tion rates) and astrophysics (i.e., uncertainties in the astrophysical models). In this paper, we focus on the latter. The largest uncertainty comes from the stellar models. Despite the recent advancements in computational astrophysics, stellar evolution remains a problem that is intractable for multi-dimensional simulations, due to the extremely long timescales of stellar evolution compared to the dynamical timescales of convective processes. Therefore, we rely on spherically symmetric simulations that incorporate convection using mixing-length theory (Bohm-Vitense 1958) models. Other multidimensional effects, such as rotation and binary interactions, are typically ignored, although recently more and more models incorporate these effects in 1D (Limongi & Chieffi 2018; Laplace et al. 2021; Farmer et al. 2023). The other main source of uncertainty is the explosion. CCSNe are complex environments where extremely high densities are reached, and neutrino interactions play a crucial role (Mezzacappa et al. 2020; Janka 2025; Burrows & Vartanyan 2021; Müller 2020; Boccioli & Roberti 2024). Therefore, one has to consider uncertainties in the nuclear equation of state at high densities, as well as in transport processes of neutrinos (Fischer et al. 2017). Finally, multi-dimensional effects have been shown to be the key ingredient for the revival of the shock (Couch & Ott 2015; Abdikamalov et al. 2015; Radice et al. 2016, 2018) and, therefore, for the successful explosion of the star.

Most of the predictions of CCSN yields currently rely on relatively simple explosion models to avoid the complexities of CCSN simulations. Moreover, stellar evolution uncertainties are rarely addressed. The goal of this paper is to provide nucleosyn-

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thesis yields calculated using a sophisticated explosion model, with state-of-the-art neutrino transport and nuclear equation of state physics. Moreover, this is the first and only work where a self-consistent supernova model is used to explode progenitors calculated using different stellar evolution codes. For this purpose, we describe the pre-supernova progenitors, the supernova explosion model, and the explosive nucleosynthesis calculations in Section 2. We present various explosion properties in Section 3. We discuss some details of the nucleosynthesis, with particular focus on radioactive isotopes, in Section 4. We discuss explosion energies and remnant masses in Section 3.1. We present the yields and compare our results to previous work in Section 3. Finally, we summarize the discussion in Section 6.

2. Methods

All explosive nucleosynthesis calculations rely on two crucial ingredients: (1) the pre-supernova and (2) the explosion models. In this work, we explore the uncertainties related to both ingredients. To address the uncertainties related to the pre-supernova model, we adopt sets of progenitors calculated using three of the most popular stellar evolution codes available: KEPLER (Woosley & Heger 2007, hereafter WH07), FRANEC (Limongi & Chieffi 2018, hereafter LC18), and MESA (Farmer et al. 2023, hereafter F23). To address the uncertainties related to the explosion model, instead, we compare our results obtained with a realistic neutrino-driven explosion with previous works where the explosion of the same progenitors was induced (by means of, e.g., piston, thermal or kinetic bomb prescriptions) and, in one case, another realistic neutrino-driven explosion.

2.1. The KEPLER models (WH07)

The WH07 models are 32 single star progenitors at solar metallicity ($Z = 0.0149$, Lodders 2003), with zero-age main sequence (ZAMS) masses from $12 M_{\odot}$ to $33 M_{\odot}$ with a spacing of one solar mass, plus stars of 35, 40, 45, 50, 55, 60, 70, 80, 100, and $120 M_{\odot}$. They were evolved using a coupled 19-isotope network during main sequence and post-main sequence phases, which was replaced by a 128-isotope network in quasi-statistical equilibrium coupled to a small reaction network below magnesium at the end of oxygen burning. The models were then post-processed with a ~ 1500 isotopes network.

2.2. The FRANEC models (LC18)

The LC18 models are 108 single star progenitors with ZAMS masses of 13, 15, 20, 25, 30, 40, 60, 80, and $120 M_{\odot}$, initial rotation velocity of 0, 150, and 300 km/s, and metallicities of $[\text{Fe}/\text{H}] = 0, -1, -2, -3$ ¹. The adopted solar chemical composition is the one provided by Asplund et al. (2009) ($Z = 1.345 \times 10^{-2}$), while in the case of sub-solar metallicities, an enhancement of α -element abundances is considered in the scaled solar chemical composition. As a result, the metal fraction corresponding to $[\text{Fe}/\text{H}] = -N$ is $Z = 3.236 \times 10^{-(N+2)}$ (see also Roberti et al. 2024). The nuclear network includes 335 isotopes (from n to ²⁰⁹Bi), whose temporal evolution is solved simultaneously with the stellar structure and mixing equations.

¹ We remind here that $[A/B] = \log_{10}(X^A/X^B) - \log_{10}(X_{\odot}^A/X_{\odot}^B)$, where X^A and X^B are the abundances of the species A and B respectively, and $X_{\odot}^A$ and $X_{\odot}^B$ are their solar abundances.

2.3. The MESA models (F23)

The F23 models are 35 single and 31 binary-stripped stars at solar metallicity ($Z = 0.0142$, Grevesse & Sauval 1998), with zero-age main sequence (ZAMS) masses from $11 M_{\odot}$ to $45 M_{\odot}$ with a spacing of one solar mass. Four binary-stripped models with ZAMS masses of 24, 25, 28, and $29 M_{\odot}$ did not reach core-collapse due to numerical issues, and were therefore not included in the study. The binary-stripped stars were all case B systems, where stable Roche-lobe overflow occurs during the Hertzsprung gap. The nuclear network includes 162 isotopes, from n to ⁶⁴Zn.

2.4. Explosion simulations

The explosion was simulated using the open-source, spherically symmetric code GRID² (O'Connor & Ott 2010), with neutrino transport following an M1 moments scheme and neutrino opacities from NuLib O'Connor (2015), and a time-dependent mixing-length theory (MLT) model for neutrino driven convection based on Couch et al. (2020) and Boccioli et al. (2021).

The spatial grid has 850 radial zones, with a linear spacing of 0.3 km until 20 km, and then logarithmically increasing up to the radius at which the density of the progenitor falls below 500 g/cm^3 , which, depending on the progenitor, is typically around 10^{10} cm . The nuclear equation of state adopted is the SFHo (Steiner et al. 2013), based on the relativistic mean-field model of Hempel & Schaffner-Bielich (2010).

The neutrino energy grid adopted has 18 energy groups logarithmically spaced from $\sim 1 - 280 \text{ MeV}$, and muon and tauon neutrinos and antineutrinos are treated as a single “heavy-lepton” specie ν_x , since they only interact via neutral-current interactions. The opacities included are emission and absorption of ν_e and $\bar{\nu}_e$ on nucleons and a representative heavy nucleus, pair production, elastic scattering on baryons, inelastic scattering on electrons, and nucleon-nucleon Bremsstrahlung for production of ν_x adopting the formalism of Bruenn (1985) and Burrows et al. (2006). Weak magnetism and virial corrections are included following Horowitz (2002) and Horowitz et al. (2017), respectively.

The MLT parameter was fixed to be $\alpha_{\text{MLT}} = 1.51$, based on the calibration performed against 3D simulations by Boccioli et al. (2021), Boccioli et al. (2022), and Boccioli et al. (2023), as well as the most recent bounds on the fraction of failed supernovae (Adams et al. 2017; Neustadt et al. 2021).

2.5. Nucleosynthesis calculations

The nucleosynthesis calculations were performed in post-process using the open-source code SKYNET (Lippuner & Roberts 2017), with a nuclear network of 1500 isotopes up to nuclei with a mass number of 100 (i.e. the same network used for the test-case hydrostatic burning of carbon and oxygen in Lippuner & Roberts (2017)). The trajectories were extracted from the explosion simulations, and continued up to $t_{\text{fin}} = 30$ seconds by extrapolating the temperature and density following a power-law (Arcones et al. 2007; Ning et al. 2007):

$$T(t) = T(t_{\text{fin}})(1 + (t - t_{\text{fin}}))^{-2/3} \quad (1)$$

$$\rho(t) = \rho(t_{\text{fin}})(1 + (t - t_{\text{fin}}))^{-2}. \quad (2)$$

It should be noted that the extrapolation method can indeed change the yields by up to 10 – 20% (Wang & Burrows 2024a),

² <https://github.com/evanconnor/GRID>

and is therefore an active source of uncertainty. The time post-bounce when the simulation stops varies between $\sim 3 - 7$ seconds, depending on the progenitor structure. In all cases, the peak temperature reached by the shock at the end of the simulation is below 1.5 GK.

Neutrino interactions are ignored in this work since negligible differences were observed by Wang & Burrows (2024b). However, we expect significant ν -process (which is however not included in SKYNET) and, to a lesser extent, νp -process to occur (Fröhlich et al. 2006b; Pruet et al. 2006; Arcones & Thielemann 2013; Fischer et al. 2024; Friedland et al. 2025). We impose a temperature threshold for the network to switch to nuclear statistical equilibrium (NSE) of 7 GK. Notice that the choice of the threshold can indirectly affect the electron fraction Y_e of each trajectory, as discussed in Section 4.1. The innermost tracer was chosen to be the one with the smallest mass coordinate that has an increasing radius at the end of the simulation, and that spent at least 20 ms below NSE temperatures.

3. Results: Explosion properties

3.1. Explosion Energies and Remnant Masses

The energy of the explosion is one of the most important observables that can constrain the models and, despite being subject to several modeling assumptions which should not be underestimated (Popov 1993; Arnett 1982; Blinnikov et al. 1998), can be directly inferred from the light curve. Typically, the simulations presented in this paper have been run until 2.5-seconds after bounce, depending on the model. Therefore, they are very far from shock breakout, when a true explosion energy can be defined. However, it is standard practice in the field to use the so-called diagnostic explosion energy (Buras et al. 2006; Marek & Janka 2009; Müller et al. 2012; Bruenn et al. 2016; Müller et al. 2017). We use the following definition of the diagnostic energy from Müller et al. (2017):

$$E_{\text{diag}} = \int_{e_{\text{tot}} > 0} \rho e_{\text{tot}} dV, \quad (3)$$

where dV is the volume element (which includes general relativistic corrections), ρ is the density and e_{tot} is the total energy per mass. As discussed in Müller et al. (2017), defining the gravitational energy in general relativity (GR) is not obvious, and we direct the reader to that paper for a more thorough discussion. Essentially, in GR, the most intuitive definition of e_{tot} leads to a double-counting of the gravitational potential. This can be circumvented by subtracting from e_{tot} the Newtonian potential $\Phi_{\text{grav, out}}$ of the shells outside a given radius r . This leads to:

$$e_{\text{tot}} = \alpha \left[(c^2 + \epsilon + P/\rho) W^2 - P/\rho \right] - Wc^2 - \Phi_{\text{grav, out}}(r), \quad (4)$$

with

$$\Phi_{\text{grav, out}}(r) = \int_r^\infty \frac{G\rho dV}{r'}. \quad (5)$$

In the above equations α is the lapse function, ϵ is the sum of the internal and turbulent energy of the matter, P is the pressure, W is the Lorentz factor, c the speed of light and G the gravitational constant. The use of the Newtonian potential for the outer shells is justified because the matter contributing to the total energy is located in the exploding region at radii $r \gtrsim 100$ km, where GR corrections are negligible.

The definition of the internal energy is another crucial detail that can significantly change the value of the diagnostic energy.

Progenitor	$\xi_{2.0}$	t_{fin} (s)	$M_{\text{PNS}}^{\text{fin}}$ ($M_\odot$)	$M_{\text{NS}}^{\text{cold}}$ ($M_\odot$)	$R_{\text{shock}}^{\text{max}}$ (km)	E_{expl} (B)
060b300	0.65	0.92	2.45	2.07	7449	2.06
080b000	0.58	0.35	2.54	> 2.08	629	0.00
080b300	0.60	0.59	2.49	> 2.08	2762	0.18
080c000	0.58	0.64	2.49	> 2.08	2586	2.22
080d000	0.57	0.53	2.50	> 2.08	2357	2.21

Table 1. List of properties of selected LC18 progenitors that lead to successful explosions and also form a black hole. In the first column, the name of the progenitor indicates its ZAMS mass in $M_\odot$ (first three characters), metallicity (a, b, c, and d stands for $z_\odot$, $10^{-1}z_\odot$, $10^{-2}z_\odot$, and $10^{-3}z_\odot$, respectively), and initial rotation in km/s (last three characters). The remaining columns are the compactness (defined in Eq. (7)), the final time post bounce t_{fin} when the simulation stops (i.e. very close to BH formation), the final baryonic mass of the PNS $M_{\text{PNS}}^{\text{fin}}$, the gravitational mass of the cold neutron star $M_{\text{NS}}^{\text{cold}}$ corresponding to $M_{\text{PNS}}^{\text{fin}}$, the maximum shock radius $R_{\text{shock}}^{\text{max}}$ reached during the simulation, and finally the explosion energy E_{expl} at the end of the simulation defined in Eq.(6).

What needs to go into the expression of ϵ is the thermal energy, as pointed out by Bruenn et al. (2016). However, the reference point for the internal energy in the equations of state for nuclear matter adopted in supernova simulations typically accounts for the binding energy of nuclei. Therefore, to calculate the thermal energy, this zero point should be properly readjusted. To do that, we shift the internal energy in our EOS table so that it matches the internal energy of a pure ^{16}O gas at a temperature of 0.01 MeV and a density 10^6 g/cm 3 , calculated using a Helmholtz EOS (Timmes & Swesty 2000).

The explosion energy is then obtained by subtracting from E_{diag} the energy of the overburden (Bruenn et al. 2016), defined as the binding energy of the material outside the shock with negative e_{tot} :

$$E_{\text{expl}} = E_{\text{diag}} - E_{\text{ov}} \quad (6)$$

The explosion energies and ejected ^{56}Ni for all of the simulated progenitors are shown in Figure 1. We compare progenitors that have retained a fraction of their hydrogen envelope at the pre-SN stage to Type II-P SNe. Some outliers in the LC18 set are not shown in Figure 1. Their main properties are listed in Table 1. These are progenitors with high compactnesses $\xi_{2.0} > 0.5$, where

$$\xi_M = \frac{M/M_\odot}{R(M)/1000 \text{ km}}, \quad (7)$$

is the compactness as defined by O’Connor & Ott (2010). Notice that in the original definition, they calculate compactness at bounce, whereas we calculate it using the pre-collapse profile, but the two yield similar values.

What is peculiar about these progenitors is that in the simulations, the shock radius is successfully revived, but the PNS is so massive that, after losing thermal pressure due to neutrino emission, it leads to black hole formation.

As shown in Table 1, the baryonic mass of the PNS corresponds to a cold neutron star with a gravitational mass above the maximum mass allowed by the EOS ($\sim 2.08 M_\odot$)³. However, a hot PNS can sustain a mass slightly above the maximum mass

³ Notice that the only exception is the 60 $M_\odot$ progenitor in the first row of Table 1, which gets to central densities above the upper limit of the table, but we expect it to otherwise be very similar to all the others.

of a cold neutron star, and therefore, for a few hundred milliseconds, the added thermal pressure from neutrinos is able to sustain this very large PNS until it cools enough to lead to BH formation.

Based on the discussion above, these progenitors clearly represent edge cases, since the baryonic PNS mass needs to be just above the maximum allowed mass, although not too much, otherwise the BH is formed very quickly and the shock is never revived. However, in a more realistic multi-dimensional simulation, this could be a much more frequent scenario, due to the continued accretion after shock revival. One would therefore still have downflows that would contribute to increasing the mass of the PNS. These downflows would be even more pronounced for high-compactness progenitors, since they tend to develop a larger degree of asymmetry in the explosion (Burrows et al. 2024a). Therefore, as seen for the progenitors listed in Table 1, one expects the PNS to collapse to a black hole after shock revival, which would halt the neutrino emission and lead to a much weaker explosion and ejection of a much smaller amount of ^{56}Ni . Therefore, the estimated explosion energies and Nickel masses obtained for these progenitors are likely overestimated. These could instead be progenitors that could give rise to the so-called Black-Hole Supernovae (BHSNe), where a successful explosion occurs accompanied by the formation of a black hole that, depending on how soon after bounce it occurs, can produce very weak to moderately strong explosion of $0.01 - 2.5 \times 10^{51}$ erg (Chan et al. 2018; Burrows et al. 2023; Sykes & Müller 2025; Eggenberger Andersen et al. 2025). As further confirmation that this could indeed be the case for our outliers, the ZAMS masses of these progenitors are all above 30 solar masses, and their compactnesses $\xi_{2.0}$ above 0.5, which is exactly the properties found by the recent simulations of BHSNe mentioned above.

For all other progenitors, there is a known offset between the explosion models and the estimated explosion energies and ejected ^{56}Ni masses based on observations of type II-P SNe. This is shown in Figure 1, where our simulations are compared to the observational fits by Pejcha & Prieto (2015) and Fang et al. (2025). It is important to stress that both explosion energies and ejected masses of ^{56}Ni are derived quantities based on semi-analytical and/or radiation-hydrodynamic models used to reproduce the observed bolometric luminosities. Therefore, they will not only be affected by observational uncertainties but also (and, in some cases, mostly) by modeling uncertainties in the progenitor structure, explosion dynamics, lack of multi-dimensional effects, etc... In particular, Fang et al. (2025) have shown how wind model assumptions can skew the inference of explosion energies and ejected ^{56}Ni masses, and how by avoiding fits that fix the wind model one can reduce the tension between multi-dimensional simulations (magenta line in Figure 1) and observations (cyan line in Figure 1). This type of analysis goes beyond the scope of this paper. However, it is important to highlight that the 1D+ models analyzed in this work are in excellent agreement with more sophisticated 2D and 3D models (Burrows et al. 2024a; Bruenn et al. 2016; Nakamura et al. 2015; Janka 2025), as confirmed by the magenta line in Figure 1.

It is well known that the explosion energy directly correlates with the mass of the iron core and, therefore, with the progenitor's compactness (Ebinger et al. 2019; Janka 2017; Burrows et al. 2024a). The reason behind this is rather straightforward. More massive iron cores will lead to higher mass accretion rates during the stalled-shock phase, since the average densities are higher, which in turn causes larger neutrino luminosities and energies and, therefore, larger neutrino heating (Boccioli et al. 2025b) that will hence increase the explosion energy. Since higher explosion energies are achieved because of higher mass

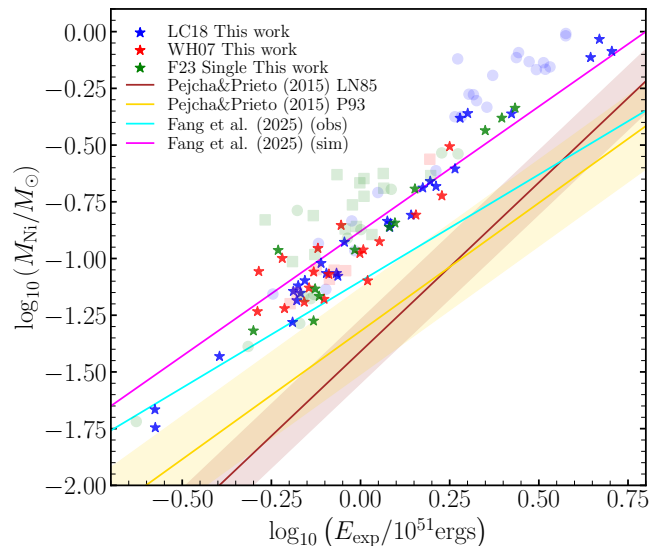


Fig. 1. Comparison of explosion energies and ejected ^{56}Ni masses obtained from the explosion simulations in this work (indicated by stars) with estimates based on hydrodynamical or semi-analytical modeling of observed light curves. The shaded circles and squares represent simulations of stripped (i.e., with less than $0.01 M_{\odot}$ of hydrogen in the envelope) and ultra-stripped (i.e., with less than $0.01 M_{\odot}$ of helium in the envelope) progenitors, respectively. The brown and yellow bands are the semi-analytical fits based on the light-curve models of Pejcha & Prieto (2015) using the calibrations of Litvinova & Nadezhin (1985) and Popov (1993), respectively. The magenta line is a fit presented by Fang et al. (2025) to the 2D simulations from Bruenn et al. (2016) and to the 3D simulations of Burrows et al. (2024b), which is in good agreement with our 1D+ simulations. The cyan line is a fit derived by Fang et al. (2025) based on observations of 32 type II SNe, where instead of assuming a fixed wind model, they used a model grid with different hydrogen envelope masses, which lessens the tension between simulations and observations.

accretion rates, a natural consequence of this is that larger explosion energies also correlate with the gravitational mass of the cold neutron star left behind by the explosion, which is larger for larger mass accretion rates and compactnesses (Boccioli & Fragione 2024). To calculate the gravitational mass of the cold neutron star, we first take the baryonic mass of the remnant to be the total mass at densities larger than 10^{12} g/cm 3 . Then we convert it to the gravitational mass of the cold neutron star by solving the TOV equations. As an alternative definition, we also selected the baryonic mass of the remnant to be the mass coordinate of the innermost tracer, and saw differences below 0.1 %.

To summarize, the mass of the iron core, the compactness, the explosion energy, and the mass of the cold neutron star are all correlated. This is what is shown in Figure 2, which can be compared to Figure 11 from Burrows et al. (2024a) to highlight again the excellent agreement of the explosion energies obtained by our 1D+ simulations and by state-of-the-art 3D simulations.

3.2. Explodability and explosion dynamics

The advantage of having a full GR simulation with state-of-the-art neutrino transport and physically motivated model for neutrino-driven convection is that, compared to the legacy methods to carry out explosive nucleosynthesis calculations (i.e. bomb and pistons models), the explosion only depends on the

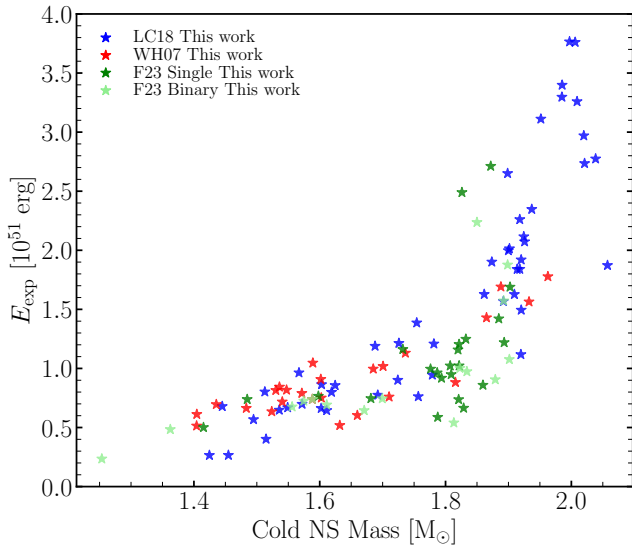


Fig. 2. This figure shows a clear correlation between the explosion energy and the gravitational mass of a cold neutron star at the end of the simulation, where different colors refer to the different sets of simulations described in Section 2. A similar correlation is also found in 3D simulations (Burrows et al. 2024a).

value of the mixing length, that can be calibrated against multi-D simulations as well as other observables (Boccioli et al. 2021, 2022). Therefore, the explosion arises self-consistently and, as discussed in Sect. 3.1, so do the mass cut and the explosion energy. Other methods, such as a bomb or a piston, with which the vast majority of explosive yield tables are constructed, require instead that the explosion energy, mass cut, and explosion dynamics are arbitrarily set. Usually, the explosion energy and mass cut are set to reproduce the explosion energy and ejected ^{56}Ni mass of SN1987A (Sonneborn et al. 1987). However, as shown in the previous section, the range of ^{56}Ni masses and explosion energies can be quite large, and therefore a single data point is not enough to represent the full population of supernovae. More recently, more accurate models (see Boccioli et al. 2025b, for a complete list of them, and the comparison to the present model) have been developed, and some of them have been, with degrees of physical approximations, adopted for nucleosynthesis calculations (Curtis et al. 2019; Sukhbold et al. 2016).

The differences among different one-dimensional explosive nucleosynthesis calculations can therefore be grouped into two main effects: explodability and explosion dynamics. Explodability changes which stars will enrich the interstellar medium (ISM). If a star explodes, its contribution to the galactic chemical evolution (GCE) will be through its winds (i.e. the mass lost during the hydrostatic evolution of the star), as well as all of the mass ejected by the explosion, which is enriched by the products of the explosive nucleosynthesis as well as by the products of pre-SN nucleosynthesis. If a star does not explode, however, only its winds will pollute the surrounding environment. This is, of course, a simplified picture that does not account for mass ejected in failed explosions (Antoni & Quataert 2023), the occurrence of BHSNe, or any multi-dimensional effects. Nonetheless, despite this apparent simplicity, its impact is often entirely neglected by assuming that all stars explode, or it is grossly approximated by setting an arbitrary cutoff of 20–25 $M_{\odot}$ that separates successful explosions (for lower mass stars) from failed

explosions (for higher mass stars). With explosion dynamics, we instead define anything that has to do with how the shock wave travels through the star, which depends on the mass cut, the final explosion energy, and how much energy is injected in the model. In a simplified picture, one can freely change two of these three quantities, and this is what is typically done in simple bomb and piston models, as we will briefly highlight in Sections 5.1, 5.2, and 5.3.)

The explodability obtained in this work is shown in Figures 4, 5, 6 for the WH07, LC18, and F23 sets, respectively. To illustrate how different the explodability can be depending on how the explosion was simulated, we add, when available, explodabilities from other 1D studies of the same sets. In Figure 4, alongside our explodability, the one from Curtis et al. (2019) (hereafter C19) is also shown. Similarly, in Figure 6, we also show the explodability derived from the criterion of Ertl et al. (2016), which was adopted by Farmer et al. (2023) as their benchmark. More details on what causes the differences between our explodability and those from C19 and Ertl et al. (2016) can be found in Boccioli et al. (2025b). In general, in contrast with C19 and E16, our model predicts all high compactness stars to explode in agreement with the most recent state-of-the-art 3D simulations (Burrows et al. 2024a), whereas it predicts lower compactness stars to explode or not depending on how pronounced the density jump at the Si/O interface is (Boccioli et al. 2023; Wang et al. 2022).

The explosion dynamics is instead less straightforward to compare, given that it is affected by explosion energy, mass cut, and shock velocities, all of which vary from one progenitor set to the other and from one explosion simulation to the other. The mass cut in WH07 and in F23 was chosen as the point in the star where the specific entropy goes above 4 (see sections 5.1 and 5.3 for a more complete discussion), which is essentially a proxy for the Si/O interface, and therefore is for the most part in agreement with the mass cut that emerges from our simulations since the accretion of Si/O interface through the shock often coincides with the shock revival (Boccioli et al. 2023, 2025b). It is important to stress that what matters for the explosion is a drastic change in density, which typically occurs when a pocket of Oxygen appears inside the Silicon shell, and therefore one should more correctly discuss the presence of a Si-Si/O interface. However, for simplicity and for consistency with previous literature, we will refer to this simply as Si/O interface.

The choice adopted by LC18 is instead to set the mass cut as the point in the star that ensures that 0.07 $M_{\odot}$ of ^{56}Ni is ejected by an explosion that starts at the edge of the iron core. This is motivated by the observational constraints on the ejected Nickel mass in SN1987A. However, as mentioned above in the context of explosion energies, the range of Nickel masses ejected in a supernova explosion is quite large. Moreover, since LC18 start the explosion at the edge of the iron core, a much higher explosion energy of a few Bethe is needed in order to eject all the material above it, which will change the shock velocity and, more relevant for the nucleosynthesis, the shock temperature as it travels through the outer mantle. The explosion in their case is simulated using a bomb, consisting of dumping a certain amount of kinetic energy in the layers just above the mass cut, which is also the explosion method adopted by F23. WH07 use instead a piston model, where an inner boundary is forced to move outward following a specific time-dependent velocity, as outlined in Woosley & Weaver (1995).

In general, both for the bomb and for the piston cases, the explosion tends to be more energetic compared to the one that arises from our simulations, as shown in Figure 3 for selected

progenitors. More detailed studies have recently compared piston, bomb, and neutrino-driven models (Imasheva et al. 2023, 2025), which have generally concluded that both bomb and pistons struggle to reproduce some of the key properties of neutrino-driven explosions, especially regarding the correlations of explosion energy with the mass of the resulting cold neutron star and ejected Nickel mass shown in Figures 1 and 2. As we will show in the remainder of this section, we generally agree with those findings.

The bomb model of LC18, for the non-rotating $20M_{\odot}$ progenitor with metallicity $Z \simeq 10^{-1} Z_{\odot}$, shown in the left panel of Figure 3, has an explosion energy of ~ 2.5 B, to be compared to the 0.66 B obtained with GR1D. The reason for such a large explosion energy in the LC18 model is that the bomb is placed at the edge of the iron core, whereas in GR1D and in the piston model of WH07 (and, as discussed above, also in the bomb model of F23), the explosion sets in around the Si/O interface. Therefore, in LC18, the shock has to eject one to a few solar masses more than in the other cases, hence leading to larger explosion energies. The mass cut is then placed by hand to force $0.07 M_{\odot}$ of ^{56}Ni to be ejected, as mentioned above. The larger explosion energy, therefore, will allow the shock to maintain high temperatures for larger and more extended regions of the pre-SN structure, as clearly shown by the final abundances shown in the left panel of Figure 3. The ^{56}Ni is relatively similar, with LC18 having the peak of production at slightly inner masses and a more extended production at larger masses. A very similar situation occurs for the WH07 progenitor shown in the right panel, showing that indeed, as one would intuitively assume, the larger explosion energy causes much stronger explosions. Notice that in the piston model of WH07 the mass cut is smaller, and this is because the $30M_{\odot}$ has relatively large compactness and, therefore, in our simulations it explodes later, after the Si/O interface has been accreted through the shock, which as explained above roughly correspond to the $s = 4$ layer adopted by WH07 as their mass cut. It is perhaps even more instructive to analyze the progenitor on the rightmost panel of Figure 3, for which the explosion energy in our model is the closest to the piston model of WH07. The final abundances are much closer in the two models, as one expects. However, the piston still gives slightly larger peak temperatures, showing that even when the explosion energy is the same (and actually our model has a very slightly larger explosion energy), the explosion dynamics tends to be a bit stronger, and therefore the peak temperature decreases less steeply compared to our simulations.

In summary, there are two main shortcomings that both bomb and piston models experience. The first one is that they arbitrarily set the mass cut, explosion energy, and explosion dynamics. This inevitably causes them to not be able to reproduce robust trends seen in multi-D simulations among explosion energy, Nickel masses, and neutron star masses, which instead neutrino-driven explosions obtain. The second one is that they arbitrarily set the explodability, contrary to neutrino-driven explosions. However, it should be highlighted that explodability is still an actively debated issue, both in 3D as well as 1D neutrino-driven simulations, as can be clearly seen by the significant differences in Figures 4 and 6 between this work, C19, and E16. More details on comparisons of 1D simulations can be found in Baccioli et al. (2025b), who have also shown that the 1D+ model adopted in this work is in excellent agreement with both FORNAX (Vartanyan & Burrows 2023; Burrows et al. 2024a) and FLASH (Eggenberger Andersen et al. 2021, 2025; Li 2024; Andresen et al. 2024) multi-dimensional simulations, whereas all of the other 1D models lack validation by 3D simulations.

4. Results: Temperature, Y_e , and Radioactive Nuclei

4.1. The role of Y_e

An important quantity that can drastically change the nucleosynthesis in core-collapse supernovae is the electron fraction Y_e . Due to neutrino interactions occurring close to the PNS, Y_e could, in principle, be quite different from 0.5. Values of $Y_e < 0.5$ would favor weak r -process nucleosynthesis (Meyer et al. 1992; Woosley et al. 1994; Wanajo 2013), allowing for elements up to the first or, in rare cases, second r -process peak, to be produced. Values of $Y_e > 0.5$ would instead favor production of neutron-deficient nuclei through, for example, the rp -process (Schatz et al. 1998; Wanajo 2006) or the νp -process (Fröhlich et al. 2006b; Pruet et al. 2006). In recent multi-dimensional simulations (Wang & Burrows 2024b, 2023), it has been shown that a large range of Y_e can be achieved, due to neutrino interaction and, crucially, asymmetric explosions allowing for more matter to be exposed to neutrino fluxes near the PNS as well as, for low-mass progenitors, early explosions that prevent neutrinos to increase Y_e in the early ν -driven wind, allowing for a significant weak r -process.

Because of the lack of these key multi-dimensional effects, spherically symmetric simulations cannot reproduce these same conditions. In our models, we do not find any significant weak r -process, since Y_e never drops significantly below 0.48-0.49. Since, however, we consistently evolve neutrino interactions (contrary to the other studies WH07, LC18, F23, and E16), our inner mass shells can achieve Y_e of up to 0.55 in some cases, especially for low-mass models, which is consistent with other 3D studies (Melson et al. 2015; Müller et al. 2019; Stockinger et al. 2020; Sandoval et al. 2021; Wang & Burrows 2024c). Therefore, the same (or, at least, very similar) nucleosynthesis observed in proton-rich ejecta in 3D is also observed in our models. The only difference is what is usually referred to as re-heating due to secondary shocks at late times (Sieverding et al. 2023), which can sometimes re-heat matter from ~ 1.5 GK to temperatures above ~ 2.5 GK, hence causing α -chain reactions that can produce ^{44}Ti . The key difference still remains the amount of proton-rich ejecta, which in 1D can be a factor of 5-10 lower compared to 3D. The Y_e distributions of all of the simulated models are shown in Figure 7.

Another key aspect of explosive nucleosynthesis is how to set the temperature threshold for considering matter to be in nuclear statistical equilibrium (NSE). In the literature, typical values have been chosen to be anywhere between 5 and 10 GK, mostly because nuclear reaction rates are not properly defined above 10 GK. Therefore solving the simplified NSE equations introduces a much smaller uncertainty than attempting an extrapolation of the reaction rates. For stellar nucleosynthesis calculations, this is only relevant in the last Si-burning stages, and one typically would expect that higher temperature thresholds would be more realistic, since a large enough network should be able to naturally reproduce the NSE results even at very high temperatures. For explosive nucleosynthesis, however, this is not true, although it is often assumed to be. The key difference is the role of neutrinos. If one were to use temperature thresholds that are too large, this would imply that, for tracers that were never exposed to high temperatures, one would perform the nucleosynthesis calculations starting from the pre-collapse composition. For example, by choosing 10 GK as the NSE threshold, any tracer that is not heated to those temperatures, would be post-processed starting from the initial composition, i.e. $Y_e \lesssim 0.5$. This would completely ignore all of the neutrino interactions that the tracer will have undergone, and therefore significantly changed its Y_e . Even

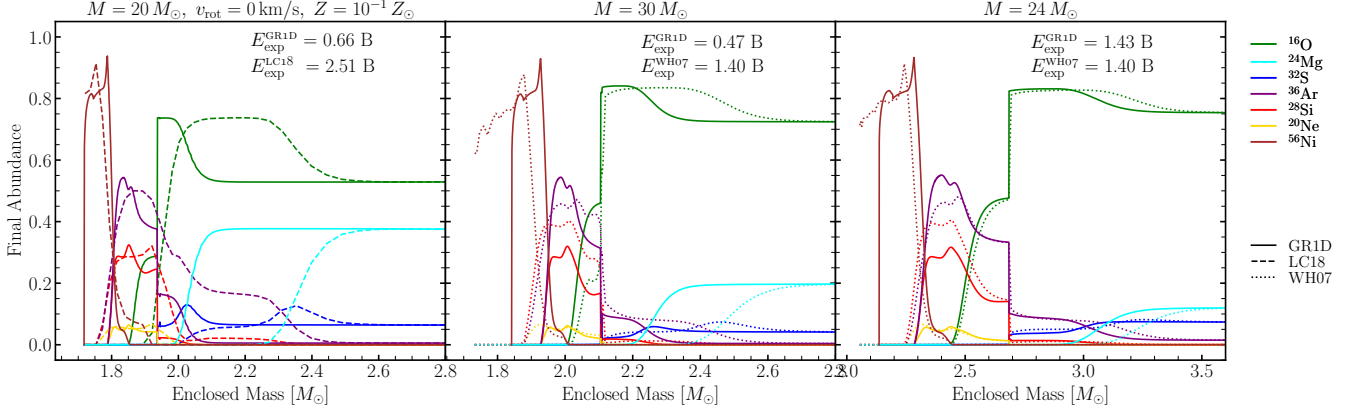


Fig. 3. The abundances for selected isotopes after the explosive nucleosynthesis are shown for three selected progenitors from LC18 (left) and WH07 (middle and right). The solid lines show the results from the explosion simulated in this work with GR1D, the dashed lines show the bomb model of LC18, and the dotted lines show the piston model of WH07. The progenitor on the left was chosen to be the one for which the mass cut derived with GR1D is closest to the one in the original work. The progenitor in the middle was chosen instead to be the one for which the explosion energy derived with GR1D is the most different from the one in the original work of 1.4 B. The progenitor on the right was chosen instead to be the one for which the explosion energy derived with GR1D is closest to the one in the original work of 1.4 B. As one can see, the left and middle panels are quite similar, showing that what causes the main difference in explosion dynamics is the explosion energy, regardless of whether a piston or a bomb is used. When the explosion energies are similar (i.e. the right panel), the differences are smaller.

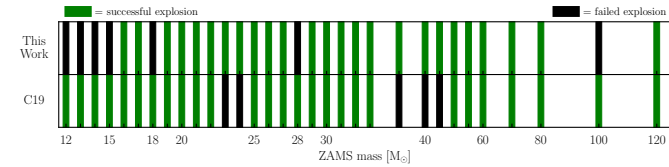


Fig. 4. Explodability of progenitor stars from Woosley & Heger (2007) as obtained in this work (upper panel) and by Curtis et al. (2019) (bottom panel). Successful explosions, defined as simulations for which the shock is successfully revived and crosses 500 km, are shown as green bands. Failed explosions are shown as black bands.

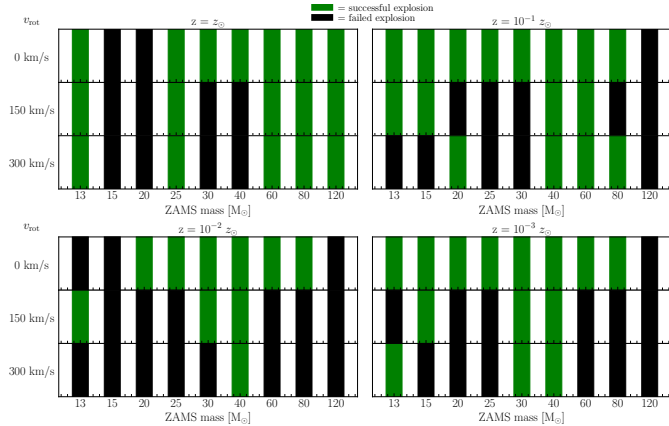


Fig. 5. Explodability of progenitor stars from Limongi & Chieffi (2018) as obtained in this work, shown as a function of initial rotational velocity at the beginning of the Main Sequence, and metallicity. Successful explosions, defined as simulations for which the shock is successfully revived and crosses 500 km, are shown as green bands. Failed explosions are shown as black bands.

if neutrino interactions are included in the network calculations, as they are in many nucleosynthesis codes (Hix & Thielemann 1999; Fröhlich et al. 2006a; Lippuner & Roberts 2017; Reichert et al. 2023), they are much less sophisticated than explicitly solv-

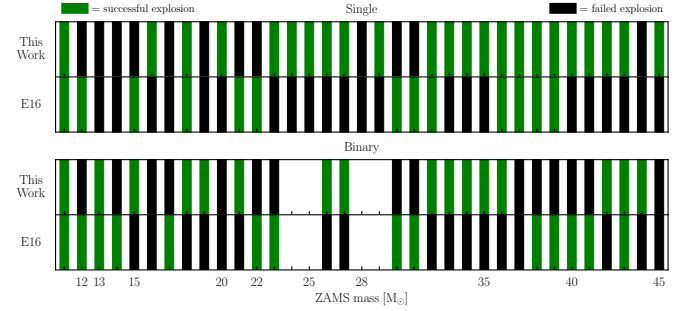


Fig. 6. Explodability of single (upper two panels) and binary-stripped (bottom two panels) progenitor stars from F23 as obtained in this work and as obtained by F23 using the explosion criterion from Ertl et al. (2016), indicated by the label “E16”. Successful explosions, defined as simulations for which the shock is successfully revived and crosses 500 km, are shown as green bands. Failed explosions are shown as black bands. Notice that the binary-stripped progenitors of 24, 25, 28, and 29 $M_{\odot}$ are missing because, as mentioned in F23, they failed to converge and reach the collapse of the iron core.

ing the Boltzmann equation as it is done in neutrino transport and, therefore, they can lead to different values of Y_e compared to the simulations.

In practice, this can lead to moderate to severe differences in selected isotopes, particularly those produced in complete silicon burning ($T \geq 5$ GK) close to the PNS. A clear example of this is ^{44}Ti . To test how sensitive our results are to the choice of the NSE threshold, we decided to also run, for the WH07 progenitors, nucleosynthesis calculations with an NSE threshold set at 10 GK. Unfortunately, SkyNet is not able to handle NSE thresholds higher than 7-8 GK (Lippuner & Roberts 2017). Therefore, just as a proof of principle, we took the trajectories with peak temperatures from 7 to 10 GK, and we set a ceiling of 7 GK and initialized the abundances to the pre-SN values. Although this is not equivalent to doing a proper nucleosynthesis calculation between 7 and 10 GK, it still highlights the effect of ignoring the Y_e evolution for these tracers. More accurate methods should be

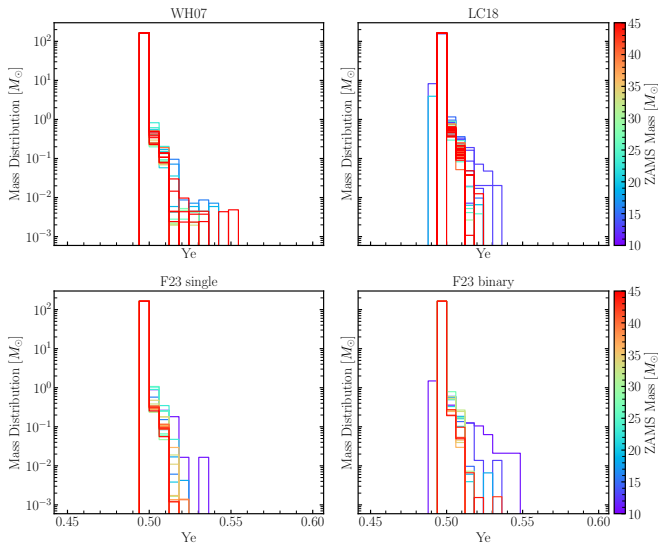


Fig. 7. Ye distributions of the ejecta color-coded by the ZAMS mass of the progenitor star for each of the three sets of simulations. Notice that we separated the single and binary stars of F23 in the bottom two panels. Lower mass stars reach higher Ye, as is also seen in 3D simulations (Melson et al. 2015; Müller et al. 2019; Stockinger et al. 2020; Sandoval et al. 2021; Wang & Burrows 2024c)

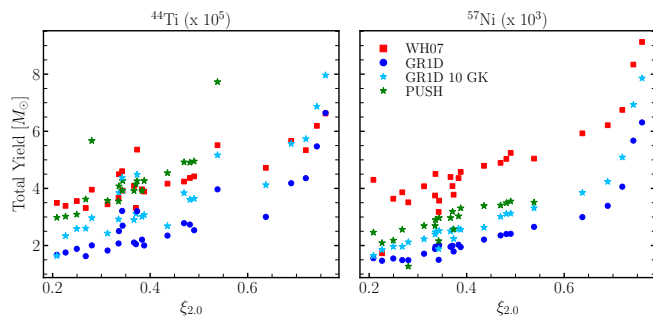


Fig. 8. Total Yield of ^{44}Ti (left) and ^{57}Ni (right) multiplied by a factor of 10^5 and 10^3 , respectively. The cyan symbols refer to a nucleosynthesis calculation that mimics one where 10 GK is used as the NSE threshold (although see text for a detailed explanation of how that was actually done). The increase in the yield of both isotopes compared to the standard run (blue symbols) shows that part of the discrepancy with the C19 results lies in the choice of the threshold.

adopted for a thorough study, which is however beyond the scope of this paper.

The comparison of our standard nucleosynthesis results with this approximate 10 GK threshold is shown in Figure 8 alongside the results from WH07 and C19. In the latter study, an actual NSE threshold of ~ 10 GK was used, and as one can clearly see, our ^{44}Ti and ^{57}Ni yields with the approximate threshold at 10 GK are much closer to the results of C19 compared to our standard calculation. It is important to highlight that the yields from WH07 are even larger, which is a consequence of the larger explosion energies adopted in their work, as explained in previous sections.

4.2. Radioactive Nuclei

In this section, we briefly discuss the two most important radioactive nuclei ejected during the explosion: ^{56}Ni and ^{44}Ti . The

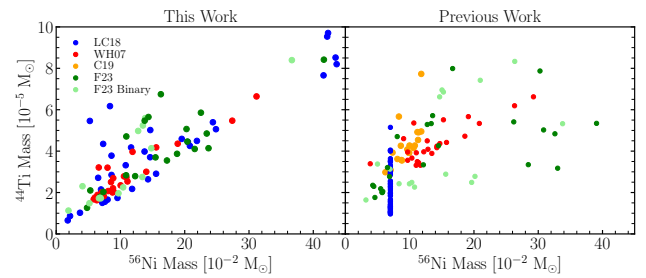


Fig. 9. The left panel shows the total yields of ^{44}Ti and ^{56}Ni for the nucleosynthesis calculations carried out in this work. Different colors refer to the different sets adopted in this paper. The right panel shows the same quantities, but the nucleosynthesis calculations are the ones carried out in the original LC18, F23, and WH07 papers. In addition to that, we also show the results from C19, who exploded the WH07 progenitors. Notice that, in their bomb explosion model, LC18 fix the amount of ^{56}Ni ejected $0.07 M_{\odot}$.

former decays to ^{56}Co with a half-life of ~ 6 days, which subsequently decays to ^{56}Fe with a half-life of ~ 77 days, whereas the latter decays to ^{44}Sc with a half-life of ~ 60 years, which subsequently decays to ^{44}Ca with a half-life of ~ 4 hours. This means that the former is responsible for powering the early light-curve of core-collapse supernovae, whereas the latter is an important element for observations of supernova remnants on the timescales of tens to hundreds of years.

As mentioned in the previous sections, the nucleosynthesis of ^{44}Ti during incomplete Si-burning is quite sensitive to the peak temperature reached by each tracer. Therefore, secondary shock in multi-dimensional simulations that reheat the matter after the initial expansion can boost its production by a factor of a few (Sieverding et al. 2023; Wang & Burrows 2024b). Nonetheless, it is still insightful to compare the yields obtained in this work with those from previous 1D studies. This comparison is shown in Figure 9. Focusing on the left panel first, which shows the yields obtained in this work, we can conclude that there is no significant difference in the trend among the three different sets, and there is also no obvious dependence on the metallicity or initial rotation in the LC18 progenitors, or between single and binary stars in the F23 progenitors. This is not surprising, since only the final structure (i.e., compactness) of the star is responsible for determining the amount of mass that undergoes complete and incomplete silicon burning, as also shown in Figure 8, which in 1D is the only site of production of ^{56}Ni and ^{44}Ti .

One can immediately see that, however, the right panel exhibits a much more pronounced scatter among the points. The most strikingly different result is the one from LC18, which is however the natural consequence of imposing a fixed amount (i.e. $0.07 M_{\odot}$) of ^{56}Ni to be ejected. For WH07, both the C19 and WH07 yields of ^{44}Ti are higher compared to our findings. In the case of C19, the reason behind this was discussed in Section 4.1, and it most likely has to do with the higher NSE threshold adopted. In the case of WH07, the reason is quite similar, since they do not include neutrinos at all, and therefore all tracers have Ye of ~ 0.5 which, as illustrated in the previous section, leads to higher abundances of ^{44}Ti . For F23, the situation is instead more complicated, since there does seem to be a very large spread in their ^{44}Ti yields, although the reason behind this is unclear, since only the final yields for those models are publicly available.

It is also worth pointing out the 6 models on the top left of the left panel, which have enhanced ^{44}Ti compared to other models

with comparable amounts of ^{56}Ni . Models exhibiting this exact same behavior can also be found in the results from Sukhbold et al. (2016) and Curtis et al. (2019), as can be seen in the top left of Figure 3 in Sieverding et al. (2023). These are progenitors that experienced the so-called carbon-oxygen (C–O) shell mergers (Andrassy et al. 2020; Rizzuti et al. 2023; Roberti et al. 2025), i.e., the complete or, in some cases, incomplete (Loddo 2025) mixing of the carbon and oxygen shells a few hours to a few days before collapse, which causes an oxygen-burning-like nucleosynthesis that can produce significant amounts of ^{44}Ti as well as odd- Z (P, Cl, K, Sc) and p -nuclei (Roberti et al. 2023; Boccioli 2025). Notice that, if a C–O merger occurs, the material outside the Si-shell quickly expands, decreasing its density. Therefore, when the shock reaches these layers, it will not be energetic enough to heat up the material to temperatures high enough to reprocess the material. The net effect is therefore that in these progenitors, the pre-SN abundances in the Carbon and Oxygen shells are largely ejected without being reprocessed (Loddo 2025).

5. Results: yields and comparison with previous work

In this section, we carry out an extensive comparison with previous studies that have adopted the same progenitors analyzed in this work to perform explosive nucleosynthesis calculations. The goal is to highlight how different explodability and different explosion engines (i.e. a different explosion dynamics) can significantly affect the results. Moreover, extensive comparisons of nucleosynthesis calculations between sophisticated explosion simulations, such as the one performed in this work, and more simplistic, yet more widely used explosion models are lacking in the literature.

A sample of elemental and radioactive isotope yields is shown in Appendix D, and the full tables will be made available online.

5.1. Comparison with PUSH and WH07

The WH07 models analyzed in this work are, as mentioned above, the ones from Woosley & Heger (2007), who performed explosive nucleosynthesis calculations using the piston model to simulate the propagation of the shock (see, e.g., Woosley & Weaver 1995; Woosley et al. 2002). In particular, they provide four different explosion models for each progenitor: (i,ii) they set the mass cut to be either the edge of the iron core or the point where the specific entropy per baryon is 4; (iii,iv) they set the explosion energy (i.e. the kinetic energy at the ejecta at infinity) to be 1.4, or 2.4 B.

We compare our results to the set that has an explosion energy of 1.4 B and the mass cut set to the point where the specific entropy per baryon is 4, which roughly coincides with the Si/O interface. The reason is that, as can be seen in Figure 1, the explosion energy we obtain for all of the WH07 progenitors is smaller than 2.0 B, and therefore 1.4 B is a more reasonable choice. We stress again that for these widely used nucleosynthesis yields, the explosion energy is usually treated as a free parameter, and is usually set to some fixed value for all of the progenitors, which is, of course, nonphysical as discussed in the previous section. The mass cut is instead chosen to be the Si/O interface since, as it has extensively been shown (Boccioli et al. 2025a; Vartanyan et al. 2021; Boccioli et al. 2023; Lentz et al. 2015; Summa et al. 2016; Vartanyan et al. 2018; Ertl et al. 2016;

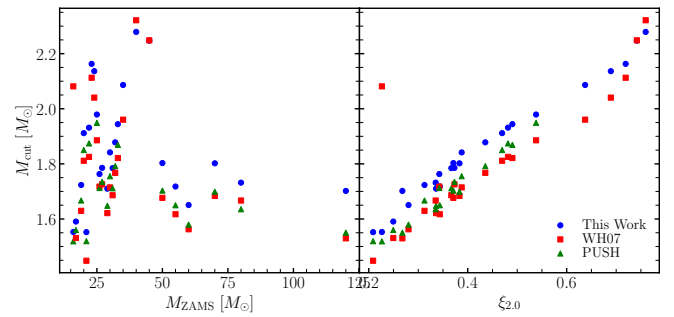


Fig. 10. Mass cuts obtained in this work (blue dots) compared to the ones obtained by WH07 (red squares) and C19 (green triangles) for the explosions of the progenitors from WH07. As expected, the mass cut correlates with compactness (except for one outlier of WH07). Notice that we find mass cuts that are larger compared to C19 and WH07, who instead find comparable values. The reason is that our explosions are triggered some small time after the Si/O interface is accreted, whereas in the case of WH07, they force the explosion exactly at the location of the interface. C19 forces the explosion by increasing the heavy neutrino heating and, as a consequence, as soon as the Si/O interface crosses the shock, the explosion occurs much more quickly compared to our simulations, and therefore mass accretion stops abruptly. The one outlier from WH07 happens to have the mass cut close to the outer Si/O interface. However, an oxygen pocket is already present much closer to the inner core, hence the large difference with respect to our models.

Wang et al. 2022; Tsang et al. 2022), the accretion of this interface through the shock often corresponds to the onset of the explosion, and therefore its location roughly corresponds to the mass of the remnant (Boccioli & Fragione 2024; Liu et al. 2021; Raithel et al. 2018; Patton & Sukhbold 2020).

The location of the mass cut is plotted for all three sets in Figure 10 as a function of both the initial ZAMS mass, as well as the pre-SN compactness of the progenitor, defined in Eq. 7. The well-known correlation between compactness and mass cut is present in all three sets and is simply the consequence of the dependence of the mass of the remnant on the Chandrasekhar mass (Timmes et al. 1996; Boccioli & Fragione 2024), which depends on the mass of the iron core and, therefore, increases with compactness since the larger the Iron core the more compact the star will be.

The pre-SN WH07 models have also been exploded in spherical symmetry with the neutrino-driven PUSH method (Perego et al. 2015; Ebinger et al. 2019, 2020). In particular, an extensive study of these same progenitors was carried out by Curtis et al. (2019) (hereafter C19). The explosion was simulated using a slightly more approximate neutrino transport compared to the one adopted in the present work. Nonetheless, their spectral neutrino transport (Liebendörfer et al. 2009) is able to accurately follow the evolution of electron neutrinos and antineutrinos. However, the much more crude advanced spectral leakage (ASL) approximation was used for the heavy-lepton neutrinos (Perego et al. 2016), causing their energies to be nonphysically large.

In PUSH, the explosion is triggered thanks to the addition of an “extra-heating” term on the RHS of the energy equation that depends on the energy and luminosity of heavy-lepton neutrinos and the compactness of the pre-SN progenitor. More details can be found in Perego et al. (2015), and a discussion on the difference with our GR1D+ model can be found in Section 5.2 of Boccioli et al. (2025b).

The other main difference with the explosion simulations carried out in this work is that the PUSH method also includes a small alpha network in the non-NSE regions. We do not expect this to impact the results too much, as also shown by Navó et al. (2023), who showed that in 1D the net effect of including a small burning network is to slightly decrease the pre-shock ram pressure, leading to a slightly stronger explosion. Given that the strength of the explosion is in any case determined by tweaking the “extra-heating” in PUSH, we do not expect this to have any impact on the comparison.

To carry out a meaningful comparison, we decided, instead of comparing each progenitor separately, to weigh the entire set by a Salpeter Initial Mass Function (IMF, Salpeter 1955) independent of metallicity:

$$P(M_{\text{ZAMS}}) \propto M_{\text{ZAMS}}^{-\alpha_{\text{IMF}}}, \quad (8)$$

with $\alpha_{\text{IMF}} = 2.35$, where $P(M_{\text{ZAMS}})$ is the probability that a star with zero age main sequence (ZAMS) mass of M_{ZAMS} forms. The reason behind this is to show how much the explodability influences the overall contribution of the population of stars to the enrichment of the interstellar medium (ISM). Therefore, the yields have been weighted by the IMF and then averaged. Finally, all of the yields have been normalized to 1 (i.e., we calculated the mass fractions of each isotope), and then converted into number fractions⁴ in order to compare our work to WH07 and C19.

As shown in Figure 4, the explodability obtained using our GRID+ simulations and the one from C19, obtained using PUSH, are quite different. Notice that WH07 assume that everything explodes, and therefore provide explosive yields for all progenitors. The main difference in the explodability is that with GRID+ stars with masses between $12 M_{\odot}$ and $15 M_{\odot}$ do not explode, whereas with PUSH they do. Moreover, stars around $25 M_{\odot}$ and $40 M_{\odot}$ do not explode with PUSH but do with GRID+. The reason behind this is that PUSH is calibrated in such a way that stars with large compactness do not explode (see Perego et al. (2015) for the exact calibration), and these stars are precisely those around $25 M_{\odot}$ and $40 M_{\odot}$.

To explicitly show the impact of explodability, in addition to comparing our yields to the original ones from WH07, we also compare them to the same exact yields with the only difference that we impose the explodability obtained with GRID+. Therefore, if a progenitor is predicted not to explode with GRID+, its contribution will only be through stellar winds. This comparison for elemental abundances after everything has decayed to stability is shown in Figure 11.

5.1.1. WH07

Overall, the yields obtained in this work agree with what was obtained in WH07 within a factor of $\sim 0.5 - 1.5$. The first effect we will analyze is the effect of explodability. If one assumes that everything explodes (i.e. the red line in Figure 11), then our simulations produce overall more light elements (except for phosphorus and fluorine, which however deserve a separate discussion) and trans-Fe elements compared to WH07. They instead produce less Fe-peak elements. When we impose our explodability on the yields of WH07 (i.e. the cyan line in Figure 11), we find that the agreement of light elements and trans-Fe elements improves significantly, whereas for Fe-peak elements it only improves very slightly.

⁴ $Y^i = X^i/A^i$, where X is the mass fraction and A is the atomic mass of the nuclear species i .

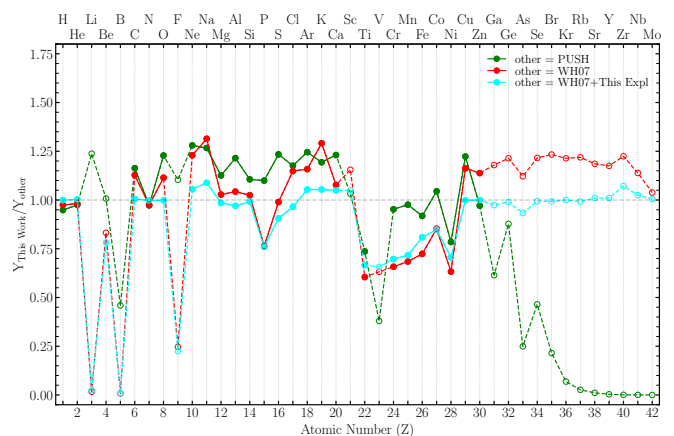


Fig. 11. Ratio of the yields obtained in this work to the yields obtained by previous works, each weighted by a Salpeter IMF and subsequently normalized to one. The red line shows the ratio of our yields to the original ones from WH07, for which every star explodes. The cyan line is the same as the red, except that we impose our explodability on the yields of WH07. This means that for the 12, 14, 15, 18, and $100 M_{\odot}$ stars, that do not explode in our simulations, we only include the winds, whereas for the rest we take the full yields from WH07. The green line shows instead the ratio of our yields to those by C19. The empty circles and the dashed line connecting them indicate elements for which the total yield is less than $10^{-7} M_{\odot}$.

When we compare our simulations with the WH07, where everything explodes, the main difference will be that less hydrogen will overall be injected into the ISM. This is because in our simulations, the 12, 13, 14, 15, 18, 28, and $100 M_{\odot}$ do not explode, and therefore will contribute to the ISM only through their winds. In the WH07 set, however, those stars will eject their H envelopes, and therefore the abundances of every other element will be comparatively reduced. This explains the larger yields of C, O, Ne, Na, Cl, Ar, and K in our simulations. Similarly, the trans-Fe elements will also be comparatively increased, since their abundance roughly scales with hydrogen because they mostly come from the initial composition of the star, and are not reprocessed during the stellar evolution or explosion phases. Therefore, when imposing our explodability on the WH07 set, the agreement for all these elements significantly improves. However, for the Fe-peak elements, the discrepancy remains. It can then be concluded that the discrepancy in the Fe-peak elements is not due to the explodability but rather to the different explosion dynamics.

As a further confirmation of this, as shown in Figure 12, if one normalizes the yields to oxygen, the discrepancy caused by a different explodability is significantly reduced for all elements except for the Fe-peak ones.

Since the amount of material synthesized in the pre-SN and in the explosive nucleosynthesis scales with O, it is not surprising that the discrepancy due to explodability is smaller. It should, however, be noted that the ratio plotted in Figure 11 is linear, whereas $[X/O]$ is a logarithmic ratio, and it therefore naturally hides small discrepancies. Nonetheless, it is quite clear that the discrepancy in light elements disappears, and the one in Fe-peak elements remains, even when imposing our explodability onto the WH07 yields, further confirming that the different explosion dynamics is responsible for it.

As outlined in Section 3.2, WH07 fixes the explosion energy at 1.4 B for all of the explosions, and as shown in Figure 1, most of our simulations have smaller explosion energies, hence

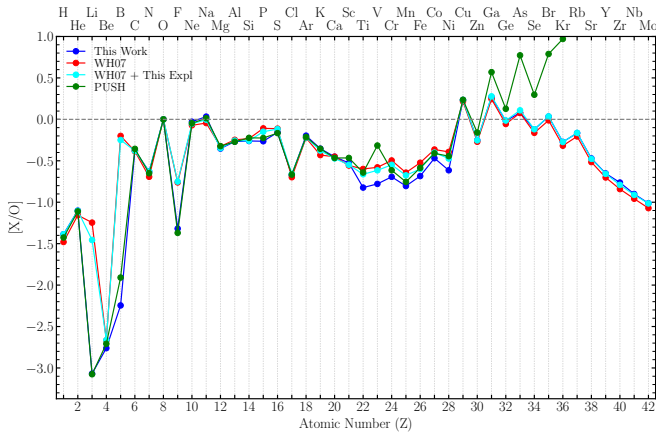


Fig. 12. Logarithmic ratio $[X/O]$ of element X and Oxygen with respect to their solar ratio. We use the solar abundances by Asplund et al. (2009). A value of zero indicates the same ratio found in the solar system. The yields of each star have been weighted by a Salpeter IMF and then normalized to one. The blue line refers to the yields obtained in this work; in red, we show the yields from WH07 assuming that everything explodes; in cyan we show the yields from WH07 adopting the explodability derived in this work; in green we show the yields from C19.

explaining the smaller amounts of Fe-peak elements, which directly correlate with the amount of ^{56}Ni ejected.

Lastly, WH07 shows much higher abundances in Fluorine and Phosphorus since they include the effects of the so-called ν -process (Woosley et al. 1990), which is responsible for synthesizing several rare isotopes through spallation reactions, including ^{19}F and ^{31}P .

5.1.2. PUSH

We also compare our yields with the work from C19, who also employed a neutrino-driven explosion from the code PUSH (Perego et al. 2015) in their calculations. From Figure 11, one can see that, overall, our yields tend to be larger for everything except for Ti, V, Ni, and all of the trans-Iron elements. For the first three elements, the isotopes causing the discrepancies are ^{47}Ti , ^{48}Ti , and ^{51}V and ^{58}Ni . Part of the discrepancy concerning these isotopes, in particular for ^{47}Ti and ^{51}V , might be caused by the weak r -process that occurs in the models of C19, which also accounts for the overproduction of trans-iron elements. However, given the more approximate neutrino transport adopted in PUSH, it is not clear how much of the neutron-rich material in the one or two zones closest to the mass cut (Ghosh et al. 2022) is a numerical artifact (Curtis, Perego, private communication). Part of the reason for the larger yields of ^{58}Ni and ^{48}Ti can however be attributed to the higher NSE threshold used in C19, which can also affect other key isotopes such as ^{44}Ti and ^{57}Ni , as discussed later in Section 4.2.

The larger yields for all other elements are then purely a result of the different explodability shown in Figure 4. If one compares each progenitor separately, the agreement is indeed remarkably good (see Appendix A for the detailed model-by-model comparison). The main difference is that in our models, the lower mass stars ($\sim 12 - 15 M_{\odot}$) do not explode, whereas the larger ones ($\sim 22 - 25 M_{\odot}$) do. Therefore, in our models, all of the elements from carbon to iron will be enhanced, whereas the hydrogen and helium will be lower, since none of our $12 - 15 M_{\odot}$

stars explode, and therefore do not eject any hydrogen in the ISM.

On the contrary, the $22 - 25 M_{\odot}$ stars explode contrary to what was assumed in C19, and since they have very large CO cores, they will comparatively contribute much more to the overall ISM abundance, explaining the larger yields of carbon, oxygen, and, in general, all other lighter elements. This is further confirmed by the fact that, if one normalizes the yields to oxygen as shown in Figure 11, the agreement is much improved except, of course, for the trans-Fe elements, vanadium, titanium, and nickel.

5.2. LC18

The LC18 models analyzed in this work are, as mentioned above, the ones from Limongi & Chieffi (2018), who performed explosive nucleosynthesis calculations using a bomb model with HYPERION (Limongi & Chieffi 2020) to simulate the propagation of the shock. In Limongi & Chieffi (2018), they launch a shock wave from the edge of the Iron core, tuning the explosion energy in such a way that the entire envelope can be unbound. This causes their explosion energies to be unphysically high, ranging from a minimum of $\sim 1.7 \text{ B}$ to $\sim 28 \text{ B}$. After having launched the shock wave, they set the mass cut manually to cause the ejection of $0.07 M_{\odot}$ of ^{56}Ni for all stars. They provide four different sets of yields for each progenitor: (i) they consider that all stars explode, or that all stars $> 25 M_{\odot}$ fail, and therefore eject only their winds; (ii) they consider that the entire envelope is ejected, or use a mixing and fallback model (although still fixing the ejection of $0.07 M_{\odot}$ of ^{56}Ni).

We compare our results to the set where everything explodes and has no mixing & fallback (their Set F⁵). We also compared with the set where everything explodes and has mixing and fallback, but qualitatively did not find significant differences since, as discussed in the remainder of this section, the mass cut and explosion energy have a much larger impact on the overall yields. We decided not to compare to the set where all stars $> 25 M_{\odot}$ fail since we are already accounting for the much more realistic explodability derived from GRID.

The location of the mass cut is plotted for all rotations and metallicities in Figure 13 as a function of both the initial ZAMS mass, as well as the pre-SN compactness of the progenitor, defined in Eq. 7. As expected, the differences in the mass cut (and, therefore, also in the mass of the resulting neutron star) are quite large (see Figure 13). This is a consequence of both the choice of launching the shock wave from the edge of the Fe core as well as the choice of ejecting a fixed amount of Nickel for all stars.

Comparing our yields with LC18 is a bit more complicated compared to the case of WH07 because some of the progenitors are rotating and, therefore, will in principle contribute differently to the IMF. To perform a global analysis, we decided, for each metallicity, to weigh the yields of each star by the IMF and by an initial distribution of rotational velocities (IDORV), as introduced in Prantzos et al. (2018). Following Rizzuti et al. (2021), we assume that the rotational velocities are sampled from a Gaussian distribution with mean:

$$\mu = \begin{cases} 300 \cdot 0.405 \cdot \exp\{-2.324 \cdot ([\text{Fe}/\text{H}] + 3)\} \text{ km/s} & \text{for } [\text{Fe}/\text{H}] \geq -3 \\ 300 \cdot 0.405 \text{ km/s} & \text{for } [\text{Fe}/\text{H}] < -3 \end{cases}, \quad (9)$$

⁵ from the ORFEO repository: <http://orfeo.iaps.inaf.it>

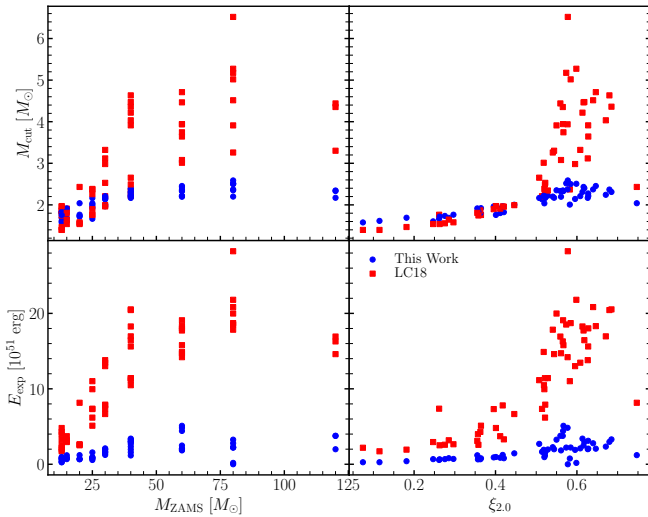


Fig. 13. Mass cut (upper panels) and explosion energies (bottom panels) for the LC18 progenitor models. The mass cut is also equal to the baryonic mass of the cold neutron star. Red squares refer to the original explosion models of LC18, and blue circles refer to the explosions carried out in this work. Left panels show each quantity as a function of initial ZAMS mass, right panels are instead as a function of pre-SN compactness. Note that all metallicities and initial rotational velocities are shown on the same plot. Both explosion energy and mass cut are larger for the LC18 explosions, especially at high compactness, although both in this work and in the original LC18 one can see a clear correlation between compactness, explosion energy, and mass cut (see the discussion in the text for more details).

and standard deviation:

$$\sigma = \begin{cases} 114.2 - 58.5 \cdot ([\text{Fe}/\text{H}] + 3) & \text{for } -3 \leq [\text{Fe}/\text{H}] \leq -1 \\ 114.2 & \text{for } [\text{Fe}/\text{H}] < -3 \\ 0 & \text{for } [\text{Fe}/\text{H}] \geq -1 \end{cases} \quad (10)$$

The ratio between our yields and the original yields from LC18 is shown in Figure 14. As done for the comparison with WH07, we impose our explodability on the yields from LC18, represented by the cyan line. As metallicity decreases, one can see that the discrepancies, particularly in the Fe-peak region, significantly increase.

As can be seen in Figure 1, the amount of ^{56}Ni ejected can vary by an order of magnitude depending on the explosion energy and compactness of the progenitor star. However, as mentioned above, in the explosions of LC18 it is fixed to $0.07M_{\odot}$, which could introduce large errors in the final abundances of Fe-peak elements. This is because all of the ^{56}Ni will decay to ^{56}Fe , which is the most abundant Fe isotope. Therefore, fixing the amount of ^{56}Ni in practice fixes the amount of Fe ejected by the supernova and, as a consequence, of all the Fe-peak elements.

For solar metallicity progenitors, $0.07M_{\odot}$ is overall a reasonable amount of ^{56}Ni ejected, and quite similar to the value obtained in our simulations when averaging over the IMF and IDORV. At lower metallicity, however, stars tend to be more compact and therefore progenitors with $M_{\text{ZAMS}} > 30M_{\odot}$ are expected to eject very large amounts of ^{56}Ni , as seen by the most energetic explosions shown in Figure 1. The important caveat to this is that some of these stars are expected to die as BHSNe, and therefore will likely experience a significant fallback that can severely decrease their explosion energy and also the mass

of Nickel ejected (Eggenberger Andersen et al. 2025). Therefore, more accurate models are needed to investigate the full impact of these high-mass stars on the overall abundance of Fe-peak elements, and we leave this to future work. Nonetheless, we do not expect our conclusions to qualitatively change even when considering the uncertainties on the explodability and ejecta mass of these BHSNe. As shown in detail in Appendix B and as it can already be seen in Figure 1, only 5 progenitors eject less than $0.07M_{\odot}$ of ^{56}Ni in our simulations. Therefore, we can robustly conclude that at low metallicity, we expect much higher abundances of Fe.

Within the Fe-peak, we find significant differences between the elements to the left (Ti, V, Cr, and Mn) and the right (Co, Ni, Cu, and Zn) of Iron. In this work, we obtain significantly more Co, Ni, Cu, and Zn compared to LC18, but less Ti, V, Cr, and Mn. As with the case of Fe, the different explodability is, however, not the main cause of this discrepancy, as can be seen by looking at the cyan and red lines in Figure 14. What determines these discrepancies boils down to the different peak temperatures of the ejecta. By starting the explosion at the edge of the iron core, the explosion energies of LC18 can be extremely large (see Figure 13), causing the mass cut to also be large. In general, the higher the explosion energy, the more ^{56}Ni can be produced during incomplete Si-burning. This is because for an explosion with a larger energy, the peak temperature decreases much more slowly as a function of mass. Therefore, especially for high-compactness progenitors, this allows larger regions of the star to reach temperatures typical of incomplete Si burning (around 4 – 5 GK) in LC18, compared to our simulations where the peak temperature decreases much faster. These regions can produce just enough ^{56}Ni to reach the threshold of $0.07M_{\odot}$, hence causing the mass cut in LC18 to be very large, at peak temperatures around ~ 5 GK. The direct consequence of this is that, contrary to our simulations, no material that reached NSE is ejected in LC18, explaining the larger yield of Co, Ni, Cu, and Zn in our models. For the same reason, since the peak temperature decreases more slowly, the region of the star experiencing incomplete silicon burning is larger in LC18, which explains the lower yields of Ti, V, Cr, and Mn in our models. As further evidence of this, one can see in Figure 14 that the elements from Si to S are also smaller in our simulations.

To summarize, LC18 tend to have very narrow (or completely absent) zones of complete Si-burning, and large zones of incomplete Si-burning due to their less steep radial dependence of T_{peak} compared to our simulations. This translates into more Co, Ni, Cu, and Zn and less Si, P, S, Ti, V, Cr, and Mn in our models. The same caveat mentioned above should however be reiterated: our models (and, to be precise, also those from LC18) ignore the fallback for high-compactness progenitors, which could significantly affect the zones of complete and incomplete Si-burning. That said, even for the lower-compactness progenitors, we find the above discussion on the Fe-peak elements to hold (see Appendix B for a star-by-star comparison), and therefore we do not expect these conclusions to qualitatively change when fallback is included.

Given the large impact that the explosion energies and mass cuts have on the final yields, the explodability in this case is of secondary importance. The main effects are on the trans-Fe elements, similarly to what was also found for the trans-Fe elements in the comparison to WH07. Since most of these elements are in the outer layers, the difference in mass cut does not significantly affect them. Interestingly, in a few cases, particularly for rotating models at low metallicity (where there is a significant s-process nucleosynthesis) and high compactness (for which the

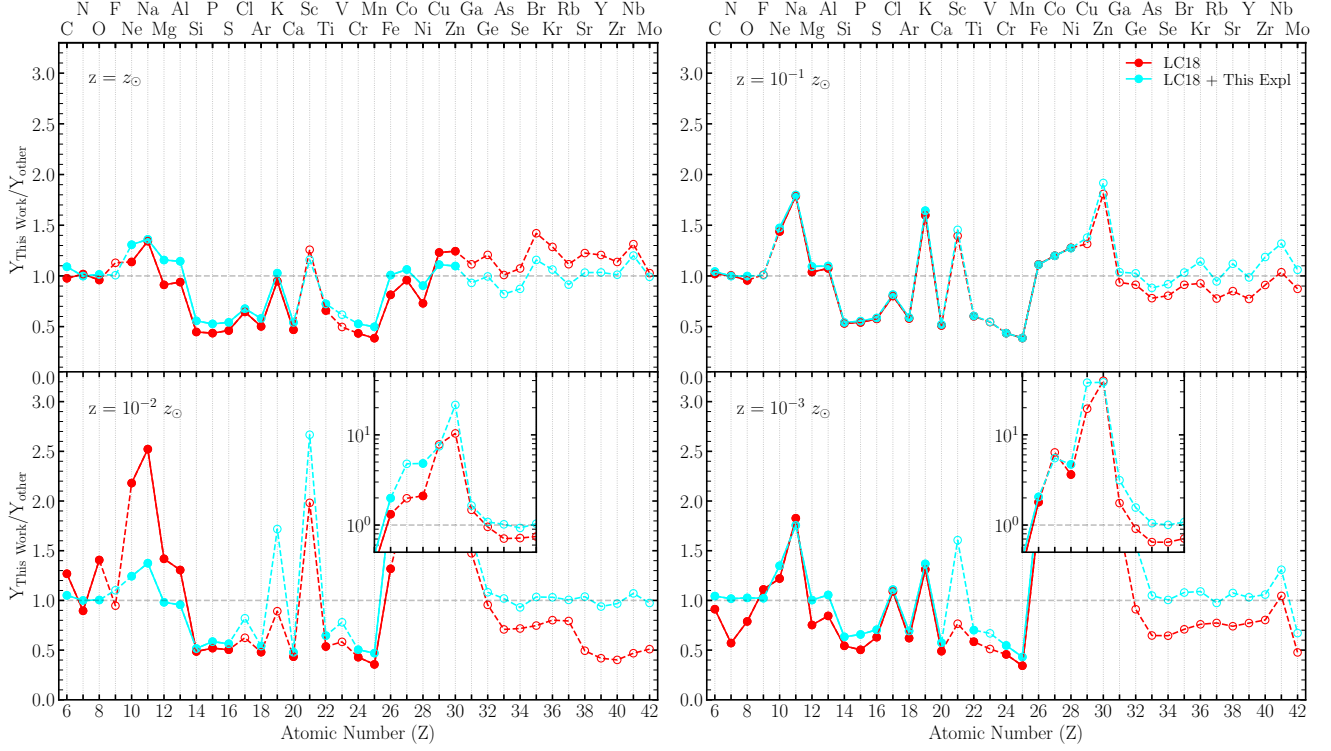


Fig. 14. Ratio of the yields obtained in this work to the yields obtained by LC18, each weighted by a Salpeter IMF and averaged over the IDORV, and subsequently normalized to one. The red line shows the ratio of our yields to the original ones from LC18, for which every star explodes. The cyan line is the same as the red, except that we impose our explodability on the yields of LC18, analogously to what was done in Figure 11. The empty circles and the dashed line connecting them indicate elements for which the total yield is less than $10^{-7} M_{\odot}$. The very large discrepancies at low metallicities are due to the fact that those are typically higher-compactness stars, and therefore are expected to eject high amounts of ^{56}Ni , which is however kept fixed at $0.07 M_{\odot}$ in LC18 (see text for more details).

explosion is strongest), some of these elements could be significantly destroyed during the explosion in LC18 (and, therefore, other trans-Fe elements could be produced). However, this is at most a 30–40% difference (see, for example, Nb and Sr) which for these *s*-process elements is not very significant.

5.3. F23

The F23 models analyzed in this work are, as mentioned above, those from Farmer et al. (2021), who performed explosive nucleosynthesis calculations using a bomb model to simulate the propagation of the shock. In particular, they launch a shock wave from the mass zone where the entropy per baryon goes above 4, tuning the explosion energy at shock breakout to be 1 B.

The location of the mass cut is plotted for single and binary-stripped stars in Figure 16 as a function of both the initial ZAMS mass, as well as the pre-SN compactness of the progenitor, defined in Eq. 7. Our models with $\xi_{2,0} < 0.5$ tend to have larger M_{cut} , whereas at higher compactnesses the F23 models have a large scatter around the value found in this work, both for single and binary stars. This is due to the fact that for high-compactness progenitors, the explosion is not usually triggered by the accretion of the Si/O interface (Boccioli et al. 2025a), but rather by the increasingly more efficient neutrino heating (Boccioli et al. 2025b).

In F23, the authors comment on the role of explodability by imposing on their yields the explodability obtained by E16. Therefore, we decided to include that explodability in our analysis. As for the discussion above, we compare the overall yields

weighted by the IMF. For clarity, we compare single and binary stars separately. Since the IMF for binary-stripped stars is not known, we assume it is also a Salpeter IMF, as in the case of single stars. In Figure 17, we show the ratio between the yields obtained in this work and the ones obtained by the explosions of F23. In the top panel we assume that for the F23 set, all models explode; in the second panel, we impose the explodability from E16; in the third panel, we impose the explodability obtained in this work; for completeness, in the bottom panel we instead show the ratio between the yields from F23 with the imposed explodability from E16 and those where every star explodes.

Overall, the yields obtained in this work for elements above carbon tend to be smaller than or equal to what was obtained by F23, as seen in the upper three panels of Figure 17. Let us first analyze the third panel, i.e. the comparison between our yields and the ones from F23 on which we imposed our explodability. Therefore, only the different mass cut and explosion dynamics affect the yields. We notice that, as already pointed out in the previous sections, the different explosion dynamics mostly affects the Fe-peak elements, which are less abundant for our explosions. This is a consequence of the fixed explosion energy chosen by F23 for all their explosions, whereas in our simulations, we have a more realistic and physically consistent range of explosion energies. The value of 1B chosen by F23 is overall too large, especially for the lower mass stars, which will contribute more to the yields due to their larger IMF weight. Therefore, our models predict roughly $\sim 70\%$ less Fe-peak elements in the ISM. Notably, there is no significant difference between single and binary stars. Now we can analyze the top panel of Figure 17, i.e. the comparison between our yields and the ones from F23 where

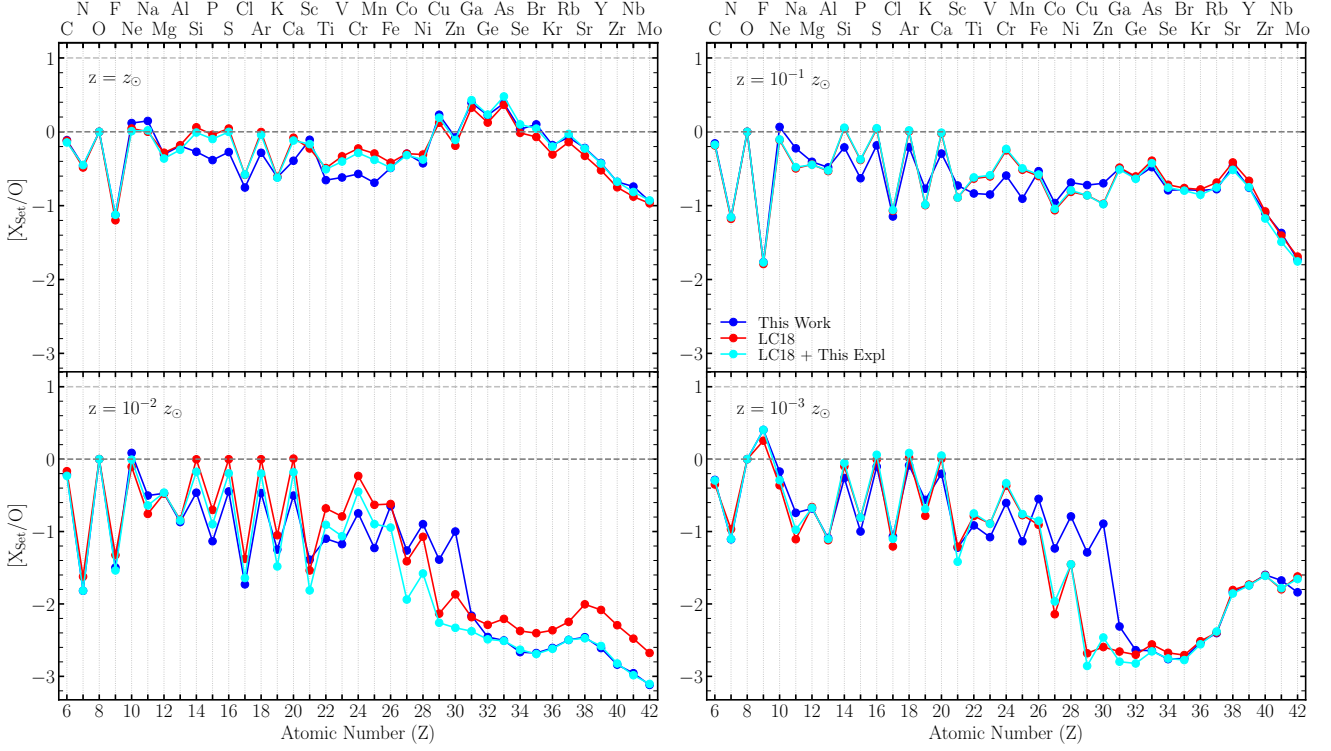


Fig. 15. Logarithmic ratio $[X/O]$ of element X and Oxygen with respect to their solar ratio. We use the solar abundances by [Asplund et al. \(2009\)](#). A value of zero indicates the same ratio found in the solar system. The yields of each star have been weighted by a Salpeter IMF, averaged over the IDORV and then normalized to one. The blue line refers to the yields obtained in this work; in red, we show the yields from LC18 assuming that everything explodes; in cyan we show the yields from LC18 adopting the explodability derived in this work.

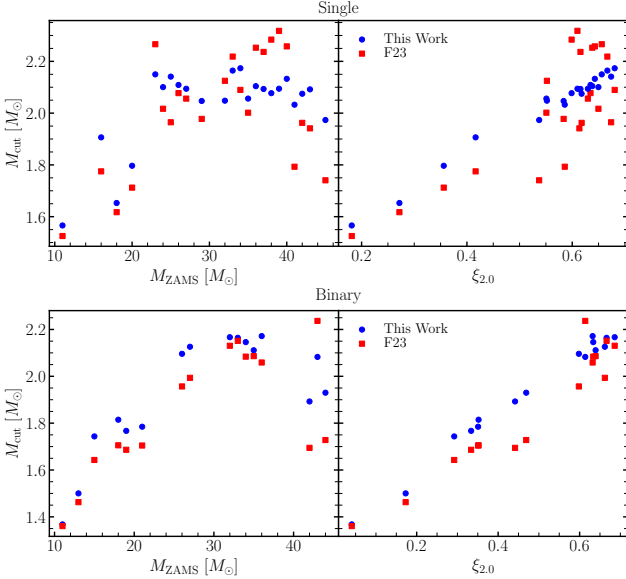


Fig. 16. Mass cut for single (upper panels) and binary (bottom panels) stars for the F23 progenitor models. The mass cut is also equal to the baryonic mass of the cold neutron star. Red squares refer to the original explosion models of F23, and blue circles refer to the explosions carried out in this work. Left panels show each quantity as a function of initial ZAMS mass, right panels are instead as a function of pre-SN compactness. The mass cut tends to be larger in the models simulated in this work, and as expected there is a clear correlation with compactness. However, the scatter in the mass cut for high compactness progenitors is much larger for the F23 single-star models (see the discussion in the text for more details).

every star explodes. The discrepancy, as expected, is larger compared to the one shown in the third panel, because of the different explodability. In particular, the discrepancy in Fe-peak elements is now even larger, and there is also a discrepancy in lighter elements. Moreover, the discrepancy in lighter elements is different for single and binary stars, and the reason is the explodability of high-mass stars. Generally speaking, the effect of explodability is to increase or decrease the weight of masses in a certain range on the overall ISM contribution. In general, fewer low-mass stars from $\sim 12M_{\odot}$ to $\sim 20M_{\odot}$ explode, less Fe-peak will be injected in the ISM. Similarly, fewer high-mass stars $\gtrsim 20M_{\odot}$ explode, less oxygen and nearby light elements will be injected into the ISM. From the explodability obtained with GRID+, shown in Figure 6, one can see that the vast majority of the high-mass single stars explode, whereas there is an island of failed explosions for binary-stripped stars with masses between $37M_{\odot}$ and $41M_{\odot}$. This causes much less Oxygen and light elements to be injected into the ISM, explaining the lower yields for these elements in binary stars compared to F23.

A similar argument can be applied to the ratio between the yields obtained in this work and those from F23 with the imposed explodability from E16, shown in the second panel of Figure 17. For binary stars, the agreement is quite good, except for the Fe-peak elements, which are lower in this work. The reason behind this has to do with both the fixed explosion energy in F23 for all stars as well as with the fact that more low-mass stars explode according to E16. For single stars, however, the yields obtained in our work are overall higher. The reason is that, according to E16, most stars (i.e. 22 out of 35) do not explode, and therefore most of the feedback into the ISM will be through winds, hence increasing the relative contribution of hydrogen

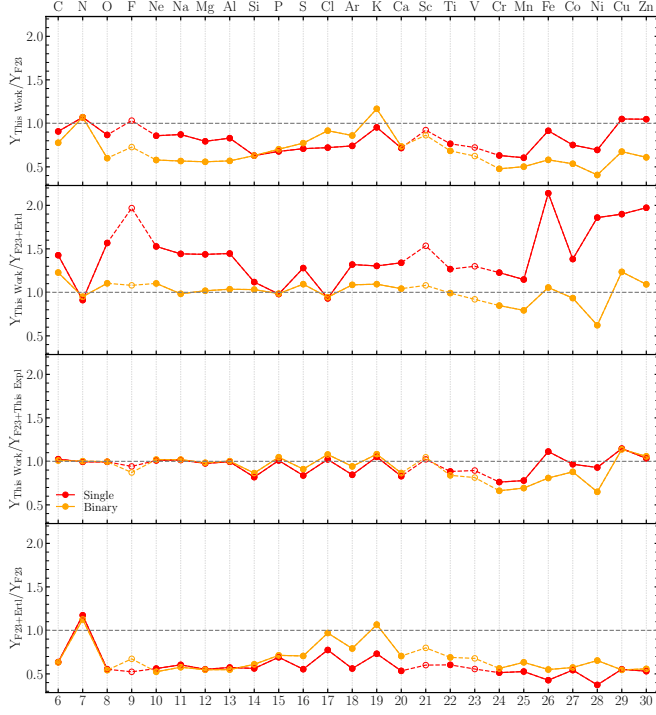


Fig. 17. Ratio of the yields obtained in this work to the yields obtained by previous works, each weighted by an IMF and subsequently normalized to one. The red lines show the results obtained for single stars, and the orange lines show the results obtained for binary stars. The empty circles and the dashed line connecting them indicate elements for which the total yield is less than $10^{-7} M_{\odot}$. In the top panel, we show the ratio of our yields with respect to the original yields of F23, assuming that everything explodes; in the second panel, we show the ratio of our yields with respect to the original yields of F23, adopting the explodability from E16; in the third panel, we show the ratio of our yields with respect to the original yields of F23, adopting the explodability found in this work; in the last panel, we show the ratio of the F23 yields adopting the explodability from E16 with the F23 yields assuming that everything explodes.

and helium, to the expense of essentially every element above carbon.

For completeness, in Figure 18 we show the abundance ratio of each element compared to Oxygen. The same conclusions reached by analyzing the yield ratios can also be reached by analyzing Figure 18. To summarize, the effect of high-mass stars not exploding is to decrease the relative contribution of Oxygen and neighboring elements to the IMF-weighted yields, since the large CO-cores of these stars are not ejected. This is confirmed by the fact that in both panels of Figure 18, $[X/O]$ for all elements from O to Al is roughly insensitive to the explosion conditions, showing that all these elements, mainly produced during oxygen burning, only depend on the Oxygen abundance. All the other elements depend instead in principle on the explodability, as proved by the large differences between the $[X/O]$ obtained in this work and the one from F23 with and without the imposed explodability from E16. When the effect of the explodability is removed (i.e. cyan versus blue lines), then the Fe-peak elements are the ones for which a significant discrepancy remains, showing that these are sensitive to how the explosion develops (i.e. how the shock propagates).

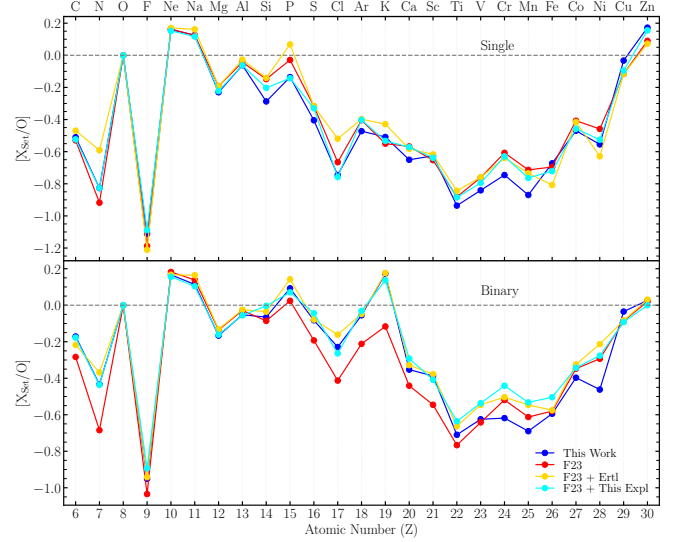


Fig. 18. Logarithmic ratio $[X/O]$ of element X and Oxygen with respect to their solar ratio. We use the solar abundances by [Asplund et al. \(2009\)](#). A value of zero indicates the same ratio found in the solar system. The yields of each star have been weighted by a Salpeter IMF and then normalized to one. The top and bottom panels show the results for single and binary stars, respectively. The blue line refers to the yields obtained in this work; in red we show the yields from F23 assuming that everything explodes; in orange we show the yields from F23 adopting the explodability from E16; in cyan we show the yields from F23 adopting the explodability derived in this work.

6. Conclusions

In this paper, we presented nucleosynthetic yields for three different sets of CCSN progenitors. The explosion was carried out using state-of-the-art 1D+ simulations that are much more sophisticated than the more commonly used piston and bomb models. Therefore, we were able to obtain ejected ^{56}Ni masses and explosion energies fully compatible with observations as well as state-of-the-art 3D CCSN simulations, and remnant masses that were obtained without the need to impose an artificial mass cut, as most current studies akin to this do.

We showed that when carrying out nucleosynthetic studies in CCSNe, it is important to track the evolution of the electron fraction caused by neutrino interactions. This can only be done if the temperature threshold for NSE is not too high, which instead causes the start of the post-process calculation to be too early, hence underestimating the proton richness of the ejecta.

Production of selected elements, such as ^{44}Ti , can have significant contributions from multi-dimensional effects ([Sieverding et al. 2023](#); [Wang & Burrows 2024a](#)). However, there can be significant pre-supernova contributions for selected progenitors (for example the one that undergo mergers of the convective carbon and oxygen shells a few hours before collapse), which is not usually captured by multi-dimensional studies, given the limited amount of simulations that can be run. Moreover, for selected progenitors, 3D simulations predict moderate to significant neutron-rich neutrino-driven winds, not found in 1D or 1D+ simulations. Likewise, the amount of proton-rich material (also present in late-time outflows from the PNS) is usually a few times smaller in 1D compared to 3D simulations ([Wang & Burrows 2024b](#)).

We showed that the explodability can have a significant impact on the nucleosynthesis. Depending on which stars explode

and which ones do not, the amount of light versus heavier elements can vary significantly, and it is therefore important to consider realistic explodability criteria when performing population studies. Moreover, we showed how piston and bomb models systematically overestimate the amount of Fe-peak elements produced. This is usually due to the fact that all stars explode with roughly the same explosion energy, whereas, as we showed in the paper, there is a clear trend of increasing explosion energy with the pre-supernova compactness of the star.

We also showed that 1D+ explosion models have a significant impact on the extent of different explosive burning regions compared to bomb and piston models. This obviously results in a significant difference also in the production of nuclear species that are produced far enough from the PNS, such as the short- (and long-) lived radioactive nuclei (e.g., ^{26}Al , ^{36}Cl , ^{40}K , ^{41}Cl , ^{60}Fe , see [Curtis et al. 2019](#); [Lawson et al. 2022](#); [Battino et al. 2024](#); [Falla et al. 2025](#), and Table D.2) as well as the p-nuclei (n-deficient isotopes of elements heavier than Fe, like, e.g., ^{92}Mo or ^{130}Ba , see [Roberti et al. 2023, 2024](#), and references therein). We aim to extend the discussion on the production of radioactive species in a forthcoming paper.

For future studies, it is therefore crucial to consider the uncertainties related to the explodability and to the explosion dynamics (i.e. mass cut and explosion energy). Until large sets of 3D explosion models become available, 1D+ simulations can be a great tool to explore these uncertainties while at the same time relying on much more robust explosion mechanisms compared to the old piston and bomb models.

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Appendix A: Comparison for each progenitor (WH07)

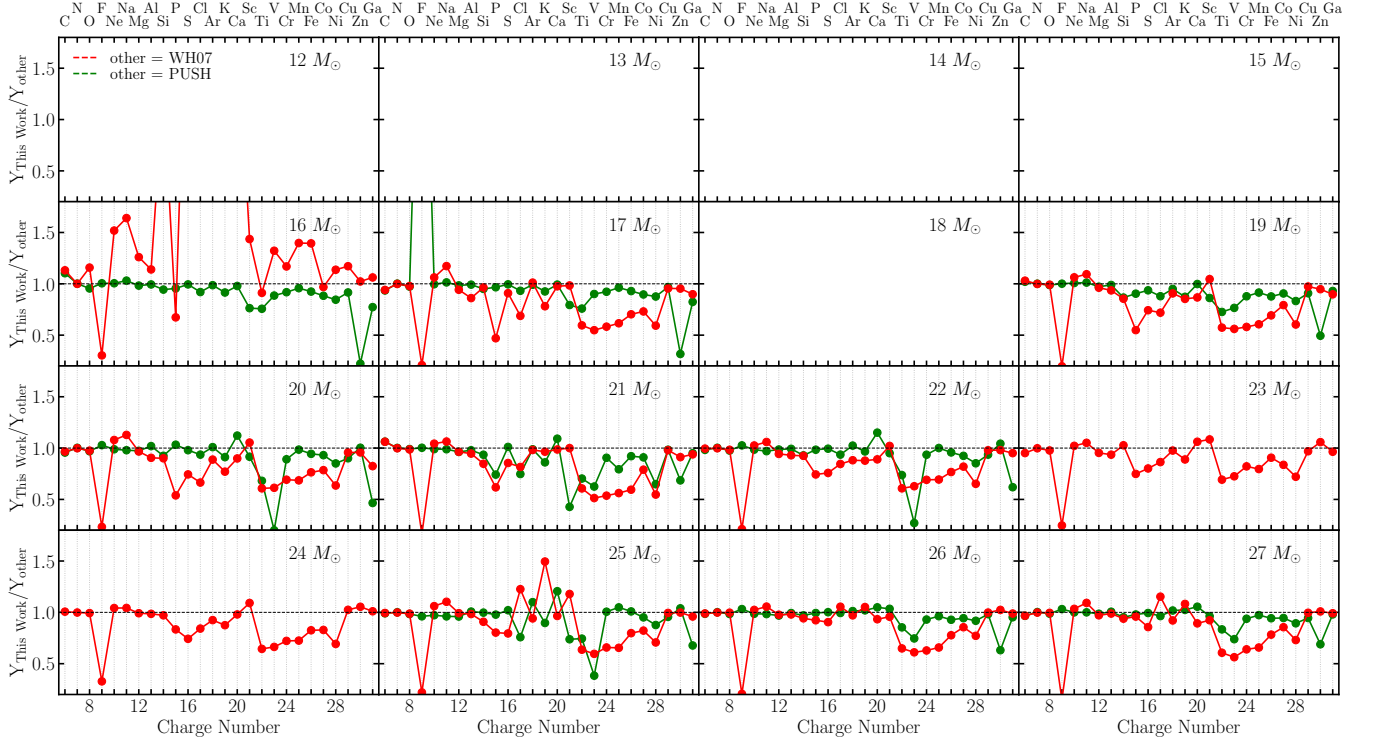


Fig. A.1. Ratio of our yields to the yields of WH07 (red) and C19 (green) for all progenitors from WH07. Blank panels indicate stars that did not explode in our simulations (see explodability in Figure 4). The 16 $M_{\odot}$ progenitor has significantly different yields due to the much larger mass cut in one is an outlier since the remnant masses (and therefore the ejecta masses) are significantly different. If C19 predicts that a specific progenitor does not explode, we do not include it in the comparison (e.g. the case of the 23 and 24 $M_{\odot}$ progenitors).

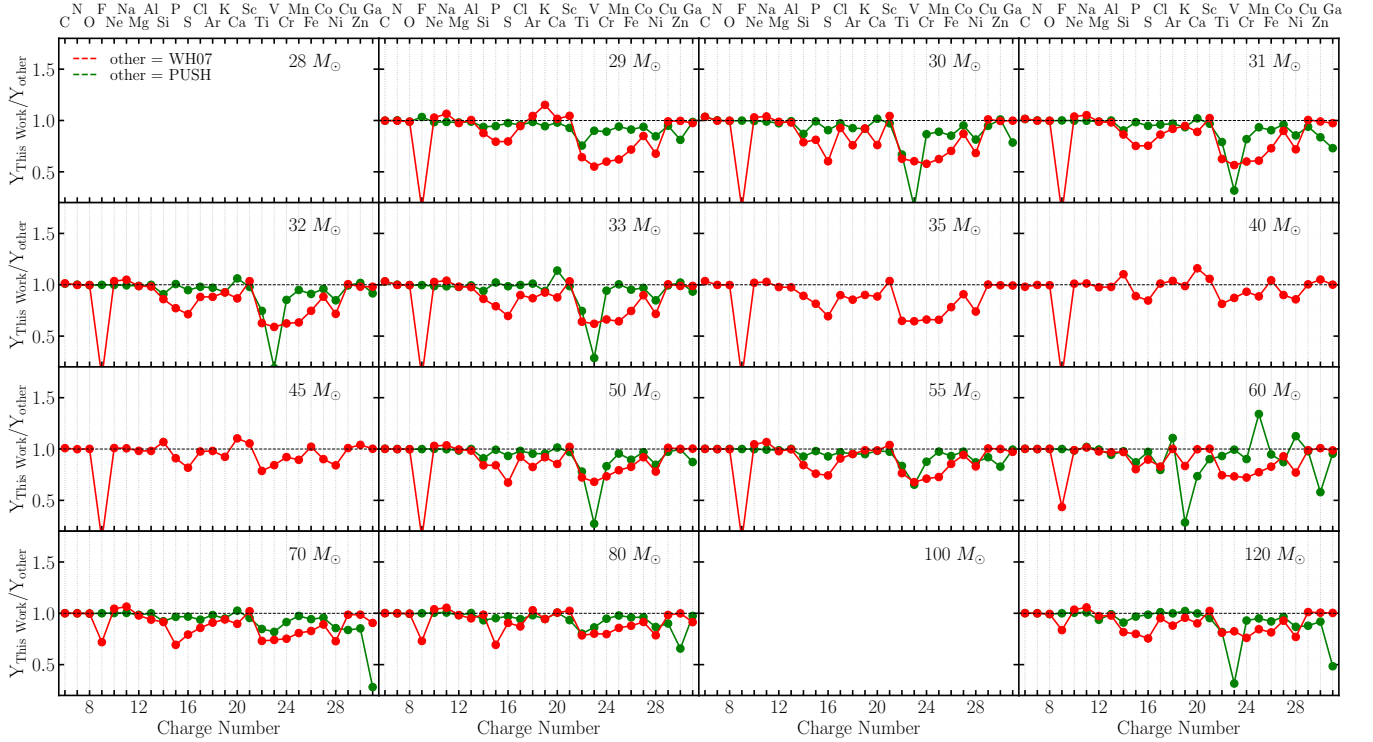


Fig. A.2. Same as Figure A.1.

Appendix B: Comparison for each progenitor (LC18)

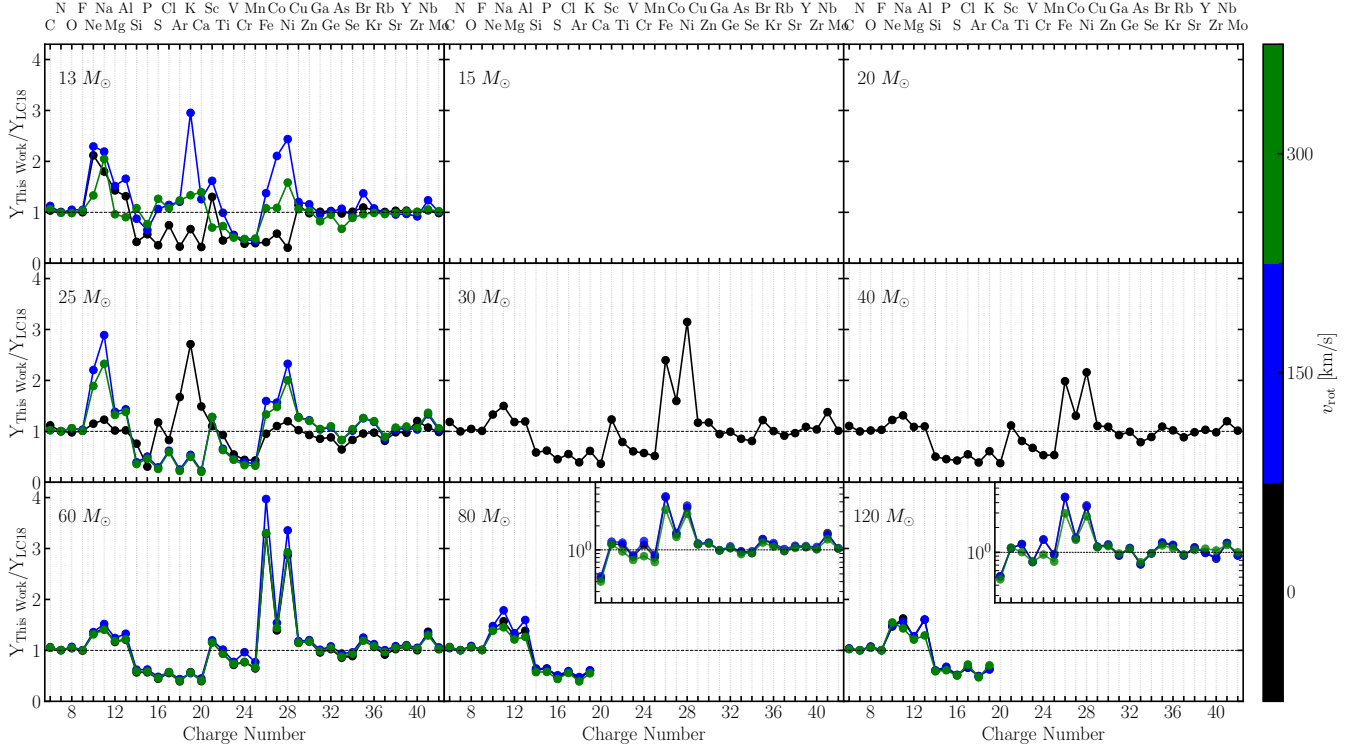


Fig. B.1. Ratio of our yields to the yields of solar metallicity progenitors from LC18, color coded by initial rotations. If a star did not explode, it is not included in the comparison (e.g., of the 30 $M_{\odot}$ stars at solar metallicity, only the non-rotating one exploded). Blank panels indicate cases where none of the stars with that particular mass exploded.

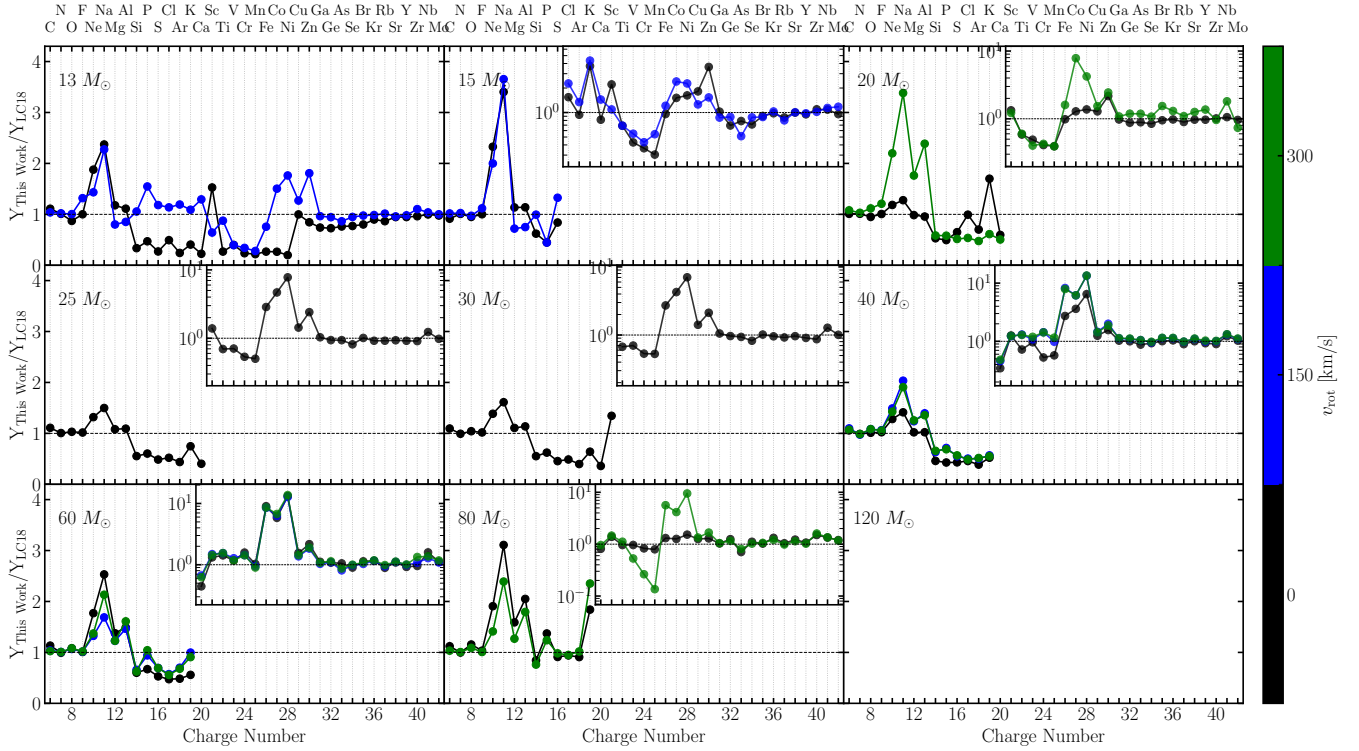


Fig. B.2. Same as Figure B.1, but for $[\text{Fe}/\text{H}] = -1$

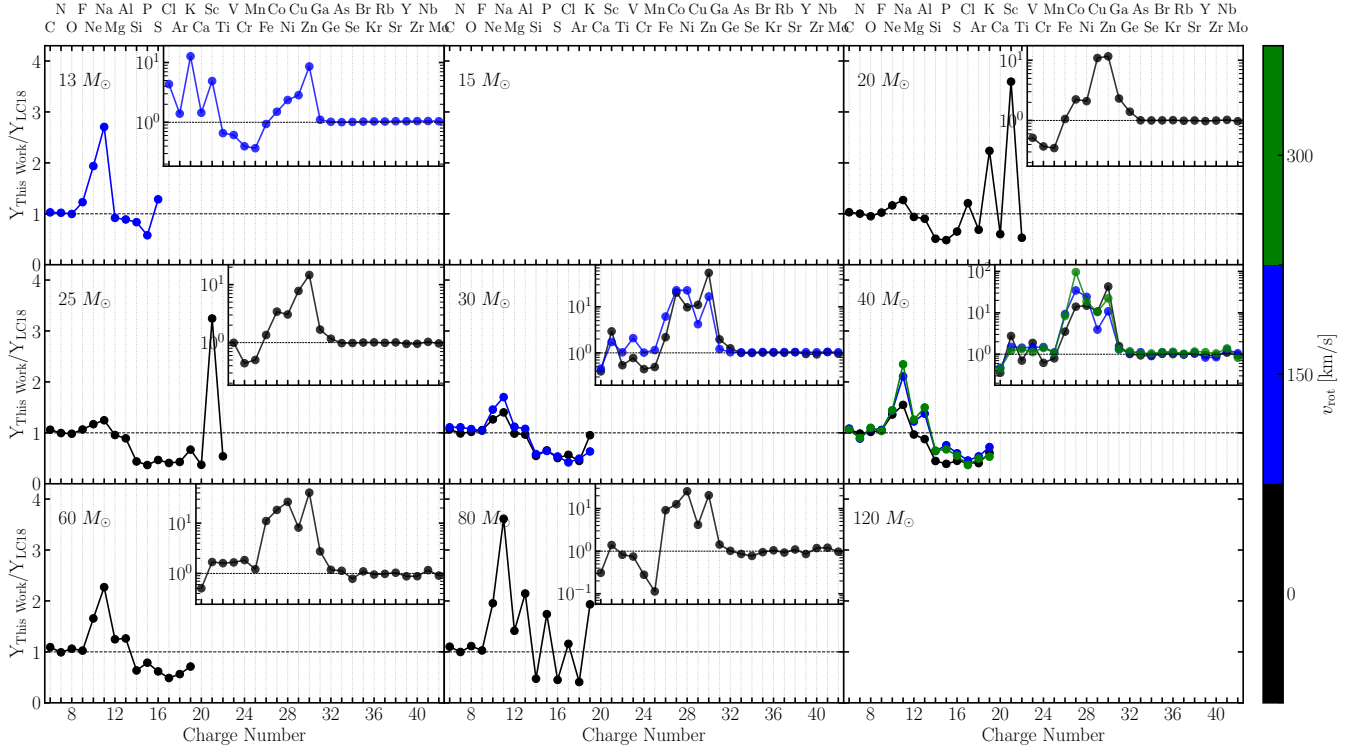


Fig. B.3. Same as Figure B.1, but for $[\text{Fe}/\text{H}] = -2$

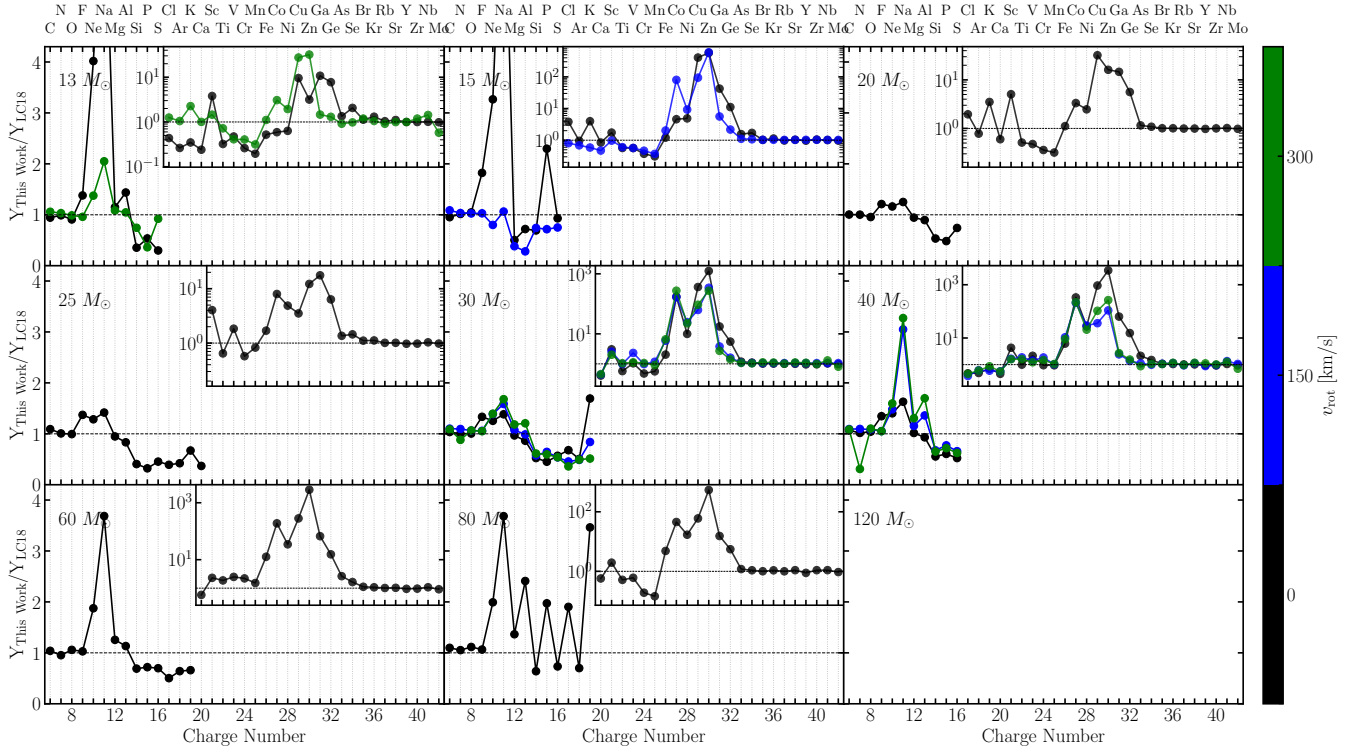


Fig. B.4. Same as Figure B.1, but for $[\text{Fe}/\text{H}] = -3$

Appendix C: Comparison for each progenitor (F23)

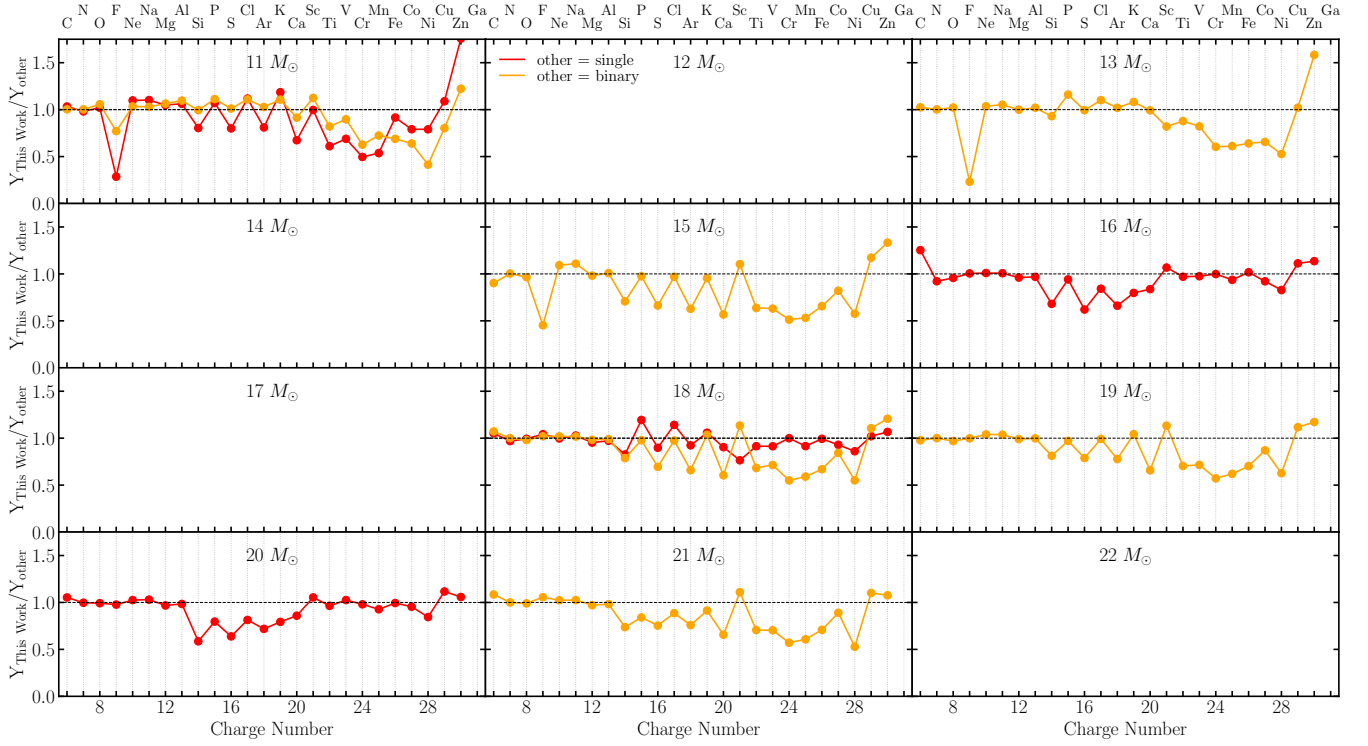


Fig. C.1. Ratio of our yields to the yields of single-star (red) and binary-stripped (orange) progenitors from F23. If a star did not explode, it is not included in the comparison (e.g., of the 15 $M_{\odot}$ stars, only the binary-stripped one exploded). Blank panels indicate cases where none of the stars with that particular mass exploded.

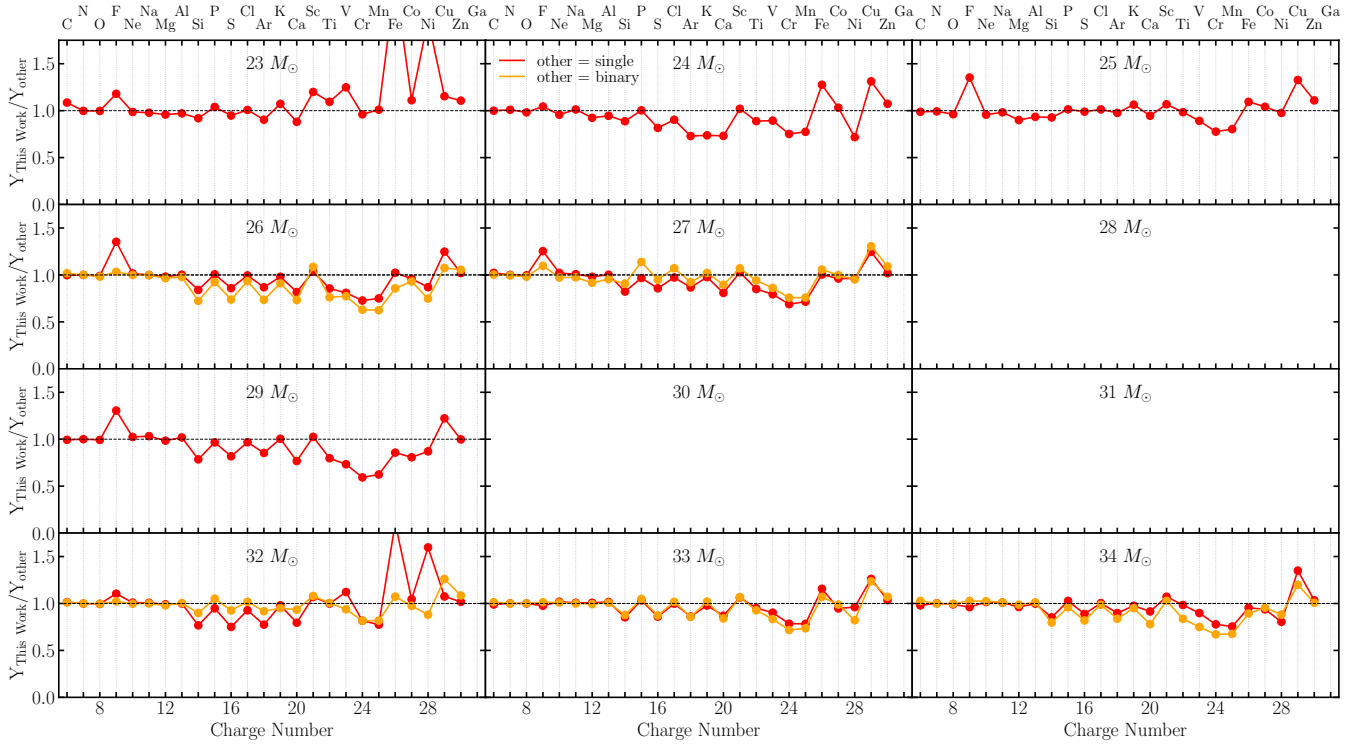


Fig. C.2. Same as Figure C.1

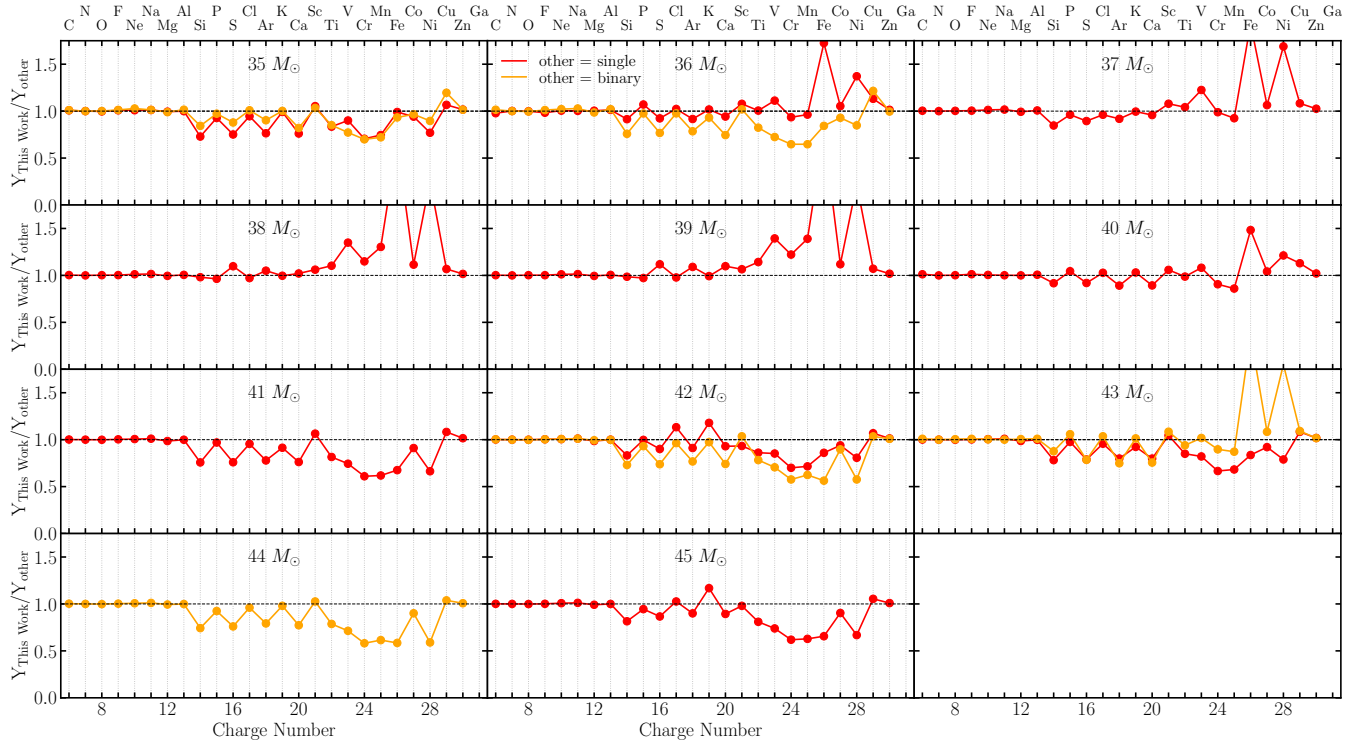


Fig. C.3. Same as Figure C.1

Appendix D: Yield Tables

Complete Tables of yields will be made available online upon acceptance.

Table D.1. Stable Elements for a $25M_{\odot}$ single, non-rotating, solar metallicity progenitor from WH07, LC18, and F23. Post refers to the post-explosion total ejected yield. Wind refers to the total yield ejected through winds during the pre-collapse evolution. f_{prod} is the production factor calculated as the ratio between the post-explosion and pre-collapse yield (i.e., not including winds). Complete tables in machine-readable form are available online.

	WH07			LC18			F23		
	Post ($M_{\odot}$)	Wind ($M_{\odot}$)	f_{prod}	Post ($M_{\odot}$)	Wind ($M_{\odot}$)	f_{prod}	Post ($M_{\odot}$)	Wind ($M_{\odot}$)	f_{prod}
C	3.70×10^{-1}	1.83×10^{-2}	1.04	6.94×10^{-1}	2.30×10^{-2}	1.07	3.17×10^{-1}	2.03×10^{-2}	1.03
N	6.36×10^{-2}	1.94×10^{-2}	1.44	5.94×10^{-2}	5.53×10^{-2}	1.45×10^1	5.22×10^{-2}	3.54×10^{-2}	3.11
O	3.34	5.30×10^{-2}	9.69×10^{-1}	2.43	6.73×10^{-2}	9.97×10^{-1}	3.76	6.47×10^{-2}	9.55×10^{-1}
F	9.09×10^{-6}	3.61×10^{-6}	1.72	6.29×10^{-6}	5.72×10^{-6}	1.09×10^1	2.44×10^{-5}	3.31×10^{-6}	1.13
Ne	5.10×10^{-1}	1.15×10^{-2}	9.46×10^{-1}	1.02	2.03×10^{-2}	1.01	5.95×10^{-1}	1.80×10^{-2}	8.70×10^{-1}
Na	7.26×10^{-3}	4.96×10^{-4}	9.50×10^{-1}	2.66×10^{-2}	1.14×10^{-3}	1.01	3.97×10^{-3}	7.43×10^{-4}	1.11
Mg	2.29×10^{-1}	6.65×10^{-3}	9.89×10^{-1}	1.63×10^{-1}	1.17×10^{-2}	1.07	3.63×10^{-1}	7.00×10^{-3}	8.24×10^{-1}
Al	2.01×10^{-2}	6.10×10^{-4}	1.02	1.80×10^{-2}	9.90×10^{-4}	1.04	3.39×10^{-2}	6.36×10^{-4}	8.43×10^{-1}
Si	2.72×10^{-1}	7.53×10^{-3}	9.09×10^{-1}	1.35×10^{-1}	1.11×10^{-2}	4.52×10^{-1}	5.12×10^{-1}	7.72×10^{-3}	7.40×10^{-1}
P	2.72×10^{-3}	6.96×10^{-5}	9.41×10^{-1}	4.74×10^{-4}	9.70×10^{-5}	7.70×10^{-1}	5.71×10^{-3}	8.51×10^{-5}	8.82×10^{-1}
S	1.41×10^{-1}	3.83×10^{-3}	7.88×10^{-1}	9.61×10^{-2}	5.15×10^{-3}	5.33×10^{-1}	1.64×10^{-1}	3.85×10^{-3}	5.61×10^{-1}
Cl	1.31×10^{-3}	4.98×10^{-5}	7.21×10^{-1}	6.29×10^{-4}	1.37×10^{-4}	7.40×10^{-1}	9.12×10^{-4}	5.11×10^{-5}	4.28×10^{-1}
Ar	2.57×10^{-2}	9.98×10^{-4}	9.49×10^{-1}	2.67×10^{-2}	1.11×10^{-3}	5.54×10^{-1}	2.68×10^{-2}	7.60×10^{-4}	4.22×10^{-1}
K	1.08×10^{-3}	3.88×10^{-5}	8.31×10^{-1}	6.23×10^{-4}	5.11×10^{-5}	5.73×10^{-1}	3.18×10^{-4}	3.99×10^{-5}	1.64×10^{-1}
Ca	1.33×10^{-2}	6.76×10^{-4}	1.69	1.86×10^{-2}	1.07×10^{-3}	5.66×10^{-1}	2.00×10^{-2}	6.79×10^{-4}	4.93×10^{-1}
Sc	9.81×10^{-6}	4.13×10^{-7}	1.53	1.26×10^{-5}	7.74×10^{-7}	4.71×10^{-1}	7.89×10^{-6}	4.29×10^{-7}	2.08×10^{-2}
Ti	2.70×10^{-4}	3.12×10^{-5}	2.40	2.51×10^{-4}	5.20×10^{-5}	4.17×10^{-3}	4.39×10^{-4}	3.16×10^{-5}	7.73×10^{-3}
V	2.83×10^{-5}	3.94×10^{-6}	3.31	2.17×10^{-5}	5.29×10^{-6}	7.44×10^{-4}	6.12×10^{-5}	4.03×10^{-6}	8.99×10^{-4}
Cr	2.01×10^{-3}	1.80×10^{-4}	7.09	1.28×10^{-3}	2.77×10^{-4}	4.81×10^{-3}	5.03×10^{-3}	1.93×10^{-4}	1.26×10^{-2}
Mn	9.00×10^{-4}	1.35×10^{-4}	5.50	6.58×10^{-4}	1.80×10^{-4}	5.76×10^{-3}	2.30×10^{-3}	1.41×10^{-4}	9.62×10^{-3}
Fe	1.56×10^{-1}	1.26×10^{-2}	1.01×10^1	1.04×10^{-1}	2.15×10^{-2}	1.51×10^{-1}	5.13×10^{-1}	1.34×10^{-2}	5.87×10^{-1}
Co	7.97×10^{-4}	3.68×10^{-5}	1.21	9.60×10^{-4}	7.02×10^{-5}	2.01×10^{-2}	1.41×10^{-3}	3.63×10^{-5}	2.39×10^{-2}
Ni	1.03×10^{-2}	7.53×10^{-4}	2.80	7.73×10^{-3}	1.19×10^{-3}	4.43×10^{-2}	3.34×10^{-2}	7.90×10^{-4}	2.03×10^{-1}
Cu	6.80×10^{-4}	8.99×10^{-6}	9.72×10^{-1}	4.54×10^{-4}	1.20×10^{-5}	2.13×10^{-2}	5.34×10^{-4}	9.38×10^{-6}	4.88×10^{-2}
Zn	5.41×10^{-4}	2.16×10^{-5}	1.18	3.53×10^{-4}	2.89×10^{-5}	2.22×10^{-1}	1.41×10^{-3}	2.32×10^{-5}	9.54×10^{-1}

Table D.2. Selected long-lived radioactive isotopic yields for a $25M_{\odot}$ single, non-rotating, solar metallicity progenitor from WH07, LC18, and F23. Post refers to the post-explosion total ejected yield. f_{prod} is the production factor calculated as the ratio between the post-explosion and pre-collapse yield (i.e., not including winds). Complete tables in machine-readable form are available online.

	WH07		LC18		F23	
	Post ($M_{\odot}$)	f_{prod}	Post ($M_{\odot}$)	f_{prod}	Post ($M_{\odot}$)	f_{prod}
^{26}Al	8.77×10^{-5}	1.42	1.11×10^{-5}	2.47	4.95×10^{-4}	9.41×10^{-1}
^{36}Cl	7.68×10^{-6}	1.48	3.12×10^{-6}	9.19×10^{-1}	2.46×10^{-5}	3.94×10^{-1}
^{40}K	1.03×10^{-5}	1.69	2.37×10^{-6}	9.97×10^{-1}	5.52×10^{-6}	2.61×10^{-1}
^{41}Ca	5.02×10^{-5}	1.65	3.96×10^{-5}	5.98×10^{-1}	1.27×10^{-5}	8.88×10^{-2}
^{44}Ti	3.97×10^{-5}	2.62×10^1	2.15×10^{-5}	2.42	9.64×10^{-5}	8.54×10^{-1}
^{48}V	2.80×10^{-7}	1.75	1.63×10^{-7}	2.31×10^{-2}	1.51×10^{-7}	4.93×10^{-4}
^{53}Mn	1.36×10^{-5}	2.96	8.82×10^{-6}	4.93×10^{-4}	3.21×10^{-5}	5.24×10^{-4}
^{56}Ni	1.19×10^{-1}	5.66×10^6	6.97×10^{-2}	4.21×10^1	4.61×10^{-1}	5.46
^{60}Fe	1.20×10^{-4}	9.69×10^{-1}	3.76×10^{-5}	1.42×10^{-3}	4.85×10^{-4}	7.93×10^{-2}
^{63}Ni	4.64×10^{-4}	9.08×10^{-1}	2.76×10^{-4}	2.40×10^{-2}	2.86×10^{-4}	3.21×10^{-2}