Impact of Neutrino Opacities on Core-collapse Supernova Simulations

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Abstract

The accurate description of neutrino opacities is central to both the core-collapse supernova (CCSN) phenomenon and the validity of the explosion mechanism itself. In this work, we study in a systematic fashion the role of a variety of well-selected neutrino opacities in CCSN simulations where the multi-energy, three-flavor neutrino transport is solved using the isotropic diffusion source approximation (IDSA) scheme. To verify our code, we first present results from one-dimensional (1D) simulations following the core collapse, bounce, and ~250 ms postbounce of a 15 $M_\odot$ star using a standard set of neutrino opacities by Bruenn. A detailed comparison with published results supports the reliability of our three-flavor IDSA scheme using the standard opacity set. We then investigate in 1D simulations how individual opacity updates lead to differences with the baseline run with the standard opacity set. Through detailed comparisons with previous work, we check the validity of our implementation of each update in a step-by-step manner. Individual neutrino opacities with the largest impact on the overall evolution in 1D simulations are selected for systematic comparisons in our two-dimensional (2D) simulations. Special attention is given to the criterion of explodability in the 2D models. We discuss the implications of these results as well as its limitations and the requirements for future, more elaborate CCSN modeling.

Key words: hydrodynamics – neutrinos – supernovae: general

1. Introduction

A core-collapse supernova (CCSN) is triggered when the core of a massive star becomes gravitationally unstable, mainly due to electron capture on protons bound in heavy iron-group nuclei. The collapse proceeds supersonically until the central density exceeds the nuclear saturation density, when the repulsive nuclear force balances gravity such that the core bounces back with the formation of a hydrodynamics shock wave. This bounce shock propagates quickly to radii on the order of 100 km. However, the shock eventually turns into an accretion front, due to losses from neutrinos when propagating across the neutrinosphere of the last scattering as well as the continuous dissociation of infalling heavy nuclei from the still gravitationally unstable layers above the stellar core. The revival of this standing accretion shock is subject to the so-called supernova problem, i.e., the liberation of the energy available at the interior of a protoneutron star (PNS), the central hot and compact object, into a thick layer of accumulated material behind the shock front (for details, see Janka et al. 2007 for a review).

Neutrinos play crucial roles at all phases of CCSNe—from stellar core collapse, core bounce, postbounce mass accretion, onset of explosion and PNS deleptonization, until the cooling of neutron stars (NSs; e.g., Woosley et al. 2002; Langanke & Martinez-Pinedo 2003; Yakovlev & Pethick 2004; Janka 2017b for detailed reviews). In particular, during stellar core collapse, nuclear electron-capture rates determine the deleptonization (cf., Hix et al. 2003; Langanke et al. 2003), which in turn defines the location of the bounce shock. After bounce, the huge amount of gravitational energy stored in the PNS is almost completely carried away by neutrinos. A tiny fraction of the streaming neutrinos deposits energy into the postshock material via weak interactions, leading to an explosion in the neutrino-driven mechanism of CCSNe (Colgate & White 1966; Bethe & Wilson 1985; Wilson 1985).

Multidimensional (multi-D) hydrodynamic instabilities play a crucial role in the neutrino-driven mechanism. Nonlinear turbulent flows associated with convective overturn and the standing accretion shock instability (Blondin et al. 2003) increase the neutrino-heating efficiency in the gain region, essentially aiding the explosion (see Kotake et al. 2012; Burrows 2013; Foglizzo et al. 2015; Müller 2016; Hix et al. 2016; Janka 2017a for reviews). In fact, a growing number of neutrino-driven models have been reported so far in self-consistent two-dimensional (2D) simulations, which have supported the validity of the multi-D neutrino-driven mechanism (e.g., Marek & Janka 2009; Suwa et al. 2010; Müller et al. 2012a, 2017; Dolence et al. 2015; Nakamura et al. 2015; O’Connor & Couch 2015; Burrows et al. 2016; Pan et al. 2016; Summa et al. 2016; Nagakura et al. 2017).

This success, however, highlights new questions. The most challenging self-consistent three-dimensional (3D) simulations with spectral neutrino transport have failed to produce explosions for 11.2, 20.0, and 27.0 $M_\odot$ progenitors (Hanke et al. 2013; Tamborra et al. 2014; see, however, Roberts et al. 2016 for an exploding 27 $M_\odot$ model). In a few successful cases, the explosions are more delayed in 3D than in 2D (e.g., Lentz et al. 2015 and Melson et al. 2015a), leading to smaller explosion energies in 3D compared to 2D (Takiwaki et al. 2014). A few exceptions from this trend have been reported for 9.6 and 11.2 $M_\odot$ stars (Melson et al. 2015b; Müller 2015). However, these two progenitors close to the low-mass end of the SN progenitors may be rather peculiar in the sense that the...
9.6 $M_\odot$ star has a tenuous envelope and explodes even in the one-dimensional (1D) simulation, whereas the 11.2 $M_\odot$ star, with its progenitor’s compactness parameter being the smallest among the 101 solar-metallicity progenitors in Woosley et al. (2002), either very marginally produces a 3D explosion (Takiwaki et al. 2012, 2014; Müller 2015) or not at all (Hanke et al. 2013).

One of the prime candidates to enhance the “explodability” is to update the neutrino physics in the multi-D models. Horowitz (2002) pioneeredly pointed out that the contribution of strange quarks to neutrino–nucleon scattering can affect neutrino opacity by $\sim 10\% - 20\%$ (see also Kolbe et al. 1998). In fact, Melson et al. (2015a) obtained 3D explosions of a 20 $M_\odot$ star only when the strangeness effect was taken into account.6 Burrows et al. (2016) reported in their self-consistent 2D simulations of a 20 $M_\odot$ star that many-body corrections to neutrino–nucleon scattering (e.g., Horowitz et al. 2017; Burrows & Sawyer 1999, 1998, albeit in a different context) make explosions easier, increase the explosion energy, and shorten the time to explosion.

Other intriguing possibilities impacting the CCSN explodability include general relativity (GR; e.g., Kuroda et al. 2012, 2016; Müller et al. 2012a; Roberts et al. 2016), stellar rotation (e.g., Yamasaki & Foglizzo 2008; Marek & Janka 2009; Suwa et al. 2010; Nakamura et al. 2014a; Takiwaki et al. 2016; Kazeroni et al. 2017; Summa et al. 2018), magnetic fields (e.g., Kotake et al. 2006; Endeve et al. 2012; Guilet & Müller 2015; Masada et al. 2015; Obergaulinger & Aloy 2017), and inhomogeneities in the progenitor’s burning shells (e.g., Couch & Ott 2013; Couch et al. 2015; Müller 2015; AbdiKamalov et al. 2016; Müller et al. 2017).

Joining the effort to update the neutrino physics in CCSN codes, we investigate in this study the impact of neutrino opacities in 1D and 2D core-collapse simulations where the three-species neutrino transport is solved using the isotropic diffusion source approximation scheme (IDSA; Liebendörfer et al. 2009). We first start with 1D simulations, where we use the same input physics and the same equation of state (EOS) as those in Liebendörfer et al. (2005). In that seminal work, a detailed comparison between the two reference codes Agile-BOLTZTRAN (Liebendörfer et al. 2004; Mezzacappa &Bruenn 1993a) and VERTEX (-PROMETHEUS) (Rampp & Janka 2002) was made. Their results are available online,7 and so we are able to compare with their data set. In the original IDSA scheme (Liebendörfer et al. 2009), a baseline set of neutrino opacities (Bruenn 1985; often referred to as the Bruenn rate) is used. Following the implementation schemes of the microphysics updates in the literature (e.g., Buras et al. 2006b; Fischer et al. 2009; Fischer 2016), we study how individual updates to the neutrino opacity leads to differences with the baseline run with the Bruenn rate. From these systematic 1D runs, we could guess which update could potentially help (or harm) the explodability in multi-D models. Then, we move on to perform 2D simulations where several sets of neutrino opacities are included. This is because a full investigation of the individual rates is currently too computationally expensive to do in multi-D simulations.

This paper is organized as follows. In Section 2, we summarize the numerical methods, including model setup, neutrino opacities, and the three-flavor IDSA scheme. In Section 3, we first compare our results with those of Liebendörfer et al. (2005; Section 3.1). Then we proceed to study one by one the impact of the individual (updated) rates in 1D runs from Sections 3.2 to 3.6. Then, it is interesting to see which expectations regarding the explodability of the 1D models survive when we perform 2D simulations (Section 4). We summarize our results and discuss its implications in Section 5.

2. Numerical Methods

2.1. Model Setup

In our 1D simulations, we employ a standard 15 $M_\odot$ progenitor (“s15s7b2” in Woosley & Weaver 1995), following the work of Liebendörfer et al. (2005). To see the effects of individual neutrino rates, we follow the dynamics starting from core collapse, through bounce, and up to ~500 ms postbounce (pb) in each 1D run. In our 2D runs, we choose the 20 $M_\odot$ progenitor model of Woosley & Heger (2007) that has been widely used in recent multi-D simulations (e.g., Melson et al. 2015a; Burrows et al. 2016; Bollig et al. 2017). This progenitor is characterized by high explodability, where the neutrino-driven shock revival was obtained around ~250 ms (pb), the earliest in the literature (e.g., Melson et al. 2015a; O’Connor & Couch 2015; Summa et al. 2016).

Our non-relativistic hydrodynamics code employs the high-resolution shock-capturing scheme with an approximate Riemann solver of Einfeldt (1988; see Nakamura et al. 2015 for more details). Self-gravity is computed using a monopole approximation with an approximate treatment of GR gravity by the effective potential of Case A of Marek et al. (2006). Our 1D and 2D runs are computed on a spherical polar grid with a resolution of $n_r = 512$ and $n_{\phi} \times n_z = 512 \times 128$, respectively. Unequal spatial radial zones cover the center to an outer boundary of $5 \times 10^8$ cm. The radial grid is chosen such that the resolution $\Delta r$ is better than 250 m in the PNS interior and typically better than 1 km outside the PNS. Seed perturbations for aspherical instabilities are imposed by hand 10 ms after bounce by introducing random 1% perturbations to the velocity behind the stalled shock. For the spectral transport, we use 20 logarithmically spaced energy bins ranging from 3 to 300 MeV.

In our multi-D runs, we take into account the energy feedback from nuclear-burning processes in the hydrodynamic evolution by solving a thirteen-species $\alpha$-nuclei network (see Nakamura et al. 2014b for details). Throughout this paper, we use the EOS of Lattimer & Swesty (1991, hereafter LS). We set the incompressibility parameter $K = 180$ MeV (LS180) only in Section 3.1 for the sake of comparison with Liebendörfer et al. (2005); in all other models, we set $K = 220$ MeV (LS220), which can account for the 2 $M_\odot$ NS mass measurements (Demorest et al. 2010; Antoniadis et al. 2013).8

In order to clearly see the effect, Melson et al. (2015a) chose a relatively high strangeness contribution ($g_s^i = -0.2$) to the axial-vector coupling constant ($g^0_s \approx 1.26$) compared to the constraint ($g^0_s \lesssim -0.1$) proposed by Hobbs et al. (2016).

A detailed comparison of the role of the nuclear EOS, including that of LS, in CCSN simulations was reported in Hempel et al. (2012) and Fischer et al. (2014).
In set5a, inelastic contributions and weak magnetism corrections are included following Horowitz (2002) for the charged-current absorption and neutral-current scattering processes. Set5b includes a correction for the effective nucleon mass (Reddy et al. 1999). Following Buras et al. (2006b, their Equation (A.1)), we replace the nucleon mass \( m_N \) with the density-dependent nucleon mass \( m_N(\rho) \), which changes the neutrino opacities accordingly.

In set6a, the quenching of the axial-vector coupling constant at high densities (Carter & Prakash 2002; e.g., Equation (A.9) in Buras et al. 2006b) is included but using more recent fitting formula (e.g., Equation (8) in Fischer 2016). In set6b, we employ the formulas suggested by Horowitz et al. (2017) that account for virial effects at low density and the many-body correlations at high densities (their Equations (36)–(39)) for the neutral-current axial response. Finally, in set6c, a strangeness-dependent contribution to the axial-vector coupling constant (Horowitz 2002) with \( g_5^A = -0.1 \) (Hobbs et al. 2016) is considered for neutrino–nucleon scattering.

Note that even the full set in Table 1 is in no way complete. The inclusion of muons significantly affects explosibility (Bollig et al. 2017), and a proper treatment of nucleon kinematics (Reddy et al. 1998) is not taken into account in our full set (the roles of nuclear de-excitation (Fischer et al. 2013) and light nuclear clusters (Sumiyoshi & Röpke 2008) are also not included). These updates are another major undertaking, which we leave for future work.

### 2.3. Three-flavor IDSA Scheme

The IDSA scheme splits the neutrino distribution function \( f \) into two components, \( f = f^1 + f^3 \), with \( f^1 \) and \( f^3 \) representing streaming and trapped neutrinos, respectively; both are solved using separate numerical techniques (see Liebendörfer et al. 2009 for details). In the original (two-neutrino-flavor) IDSA scheme, a steady-state approximation, \( \partial f^s(\epsilon)/\partial t = 0 \), where \( \epsilon \) represents the neutrino energy in the comoving frame, is assumed for the streaming neutrinos. Thus, we will deal with a Poisson-type equation to find the solution of \( f^s \) (e.g., Equation (10) in Liebendörfer et al. 2009). This is relatively computationally expensive, especially in multi-D simulations.

To get around this problem, we directly solve for the evolution of streaming neutrinos (e.g., Equation (1) of Takiwaki et al. 2014). In this work, we further incorporate GR effects, approximately following Rampp & Janka (2002) and O’Connor & Couch (2015), as follows:

\[
\frac{\partial \epsilon^s}{\partial \tau} + \frac{1}{r^2} \frac{\partial}{\partial r} \left( r^2 \alpha r^2 F^s \right) = S[j^s, \chi^s, \Sigma] - \alpha F^s \frac{\partial \phi}{\partial r},
\]

\[
\epsilon^s \equiv \frac{\epsilon^3}{(2\pi^2\hbar^3)^{\frac{3}{2}}} \int d\mu f^s, \quad \mu \equiv \frac{\epsilon}{\mu_N},
\]

\[
F^s \equiv \frac{\epsilon^3}{(2\pi\hbar e)^3} \int d\mu \mu f^s, \quad \Sigma \equiv -\alpha (j^s + \chi^s) \epsilon^s + \Sigma,
\]

where \( \epsilon^s \) and \( F^s \) correspond to the radiation energy and flux of the streaming particle, and \( S \) represents the source term that is a functional of the effective neutrino emissivity \( j^s \), absorptivity \( \chi^s \), and neutrino energy density \( \rho \).

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### Table 1

| Model | Weak Process or Modification | References |
|-------|-----------------------------|------------|
| set1 | \( \nu_e n = e^- p \) | Bruenn (1985) |
| set2 | \( \nu_e A = e^- A' \) | Juodagalvis et al. (2010) |
| set3a | \( \nu_e + \bar{\nu}_e = \nu_e + \bar{\nu}_e \) | Buras et al. (2003), Fischer et al. (2009) |
| set3b | \( \nu_e + \bar{\nu}_e(\epsilon_e) = \nu'_e + \bar{\nu}'_e(\epsilon'_e) \) | Buras et al. (2003), Fischer et al. (2009) |
| set4a | \( \nu_e n = e^- p, \quad \bar{\nu}_p = e^+ n \) | Martínez-Pinedo et al. (2012), Fischer (2016) |
| set4b | \( NN = \nu \bar{\nu}NN^* \) | Horowitz (2002) |
| set5a | \( \nu_e n = e^- p, \quad \bar{\nu}_p = e^+ n, \nu N = \nu N \) | Reddy et al. (1999) |
| set5b | \( m_N \rightarrow m_N^* \) | Horowitz (1999) |
| set6a | \( g_4 \rightarrow g_5^* \) | Fischer (2016) |
| set6b | \( \nu N = \nu N \) (many-body and virial corrections) | Horowitz et al. (2017) |
| set6c | \( \nu N = \nu N \) (strangeness contribution) | Horowitz (2002) |

Note: The symbols \( e^-, e^+, n, p, A \) denote electrons, positrons, free neutrons and protons, and heavy nuclei, respectively; the symbol \( N \) means \( n \) or \( p, m_N \) denotes nucleon mass, and the quantity with an \( "\cdot" \) indicates the one with an in-medium correction. \( \nu \) in the neutral-current reactions represents all species of neutrinos \( (\nu_e, \nu_x, \bar{\nu}_e) \), with \( \nu_x \) representing heavy-lepton neutrinos \( (\nu_x, \nu_\mu) \) and their antiparticles.

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**2.2. Neutrino Opacities**

With regard to the neutrino opacities, our baseline model (set1; see Table 1) employs the standard weak interaction set given in Bruenn (1985) plus nucleon–nucleon bremsstrahlung (Hannestad & Raffelt 1998; see also Rampp & Janka 2002 for detailed implementation schemes). Note that in set1, ion–ion correlations for neutrino scattering on heavy nuclei (Horowitz 1997) and the correction form factor (Mezzacappa & Bruenn 1993a; Rampp & Janka 2002) are also included. In 1D runs, all of the following updates are basically added individually to set1.

In set2 (see Table 1), the electron-capture (EC) rate on the nuclei in set1 (Fuller et al. 1982) is replaced with the currently most elaborate one by Juodagalvis et al. (2010), which is a significant extension of the EC rate by Langanke et al. (2003; see Section 2). In set3, electron–neutrino pair annihilation into \( \mu/\tau \) neutrinos (set3a in Table 1) and \( \mu/\tau \)–neutrino scattering on electron (anti)neutrinos (set3b; Buras et al. 2003) are added to set1 (see Appendices A and B for details). In set4a, medium modifications to electron/positron capture reactions on proton/(neutron) are taken into account (Martínez-Pinedo et al. 2012; Roberts et al. 2012; Hempel 2015; Roberts & Reddy 2017) at the mean-field level (Reddy et al. 1998; see Section 3.4). Set4b includes the medium-suppression of bremsstrahlung (Fischer 2016, their Equation (11)).
and the isotropic diffusion term ($\Sigma$), all defined in the laboratory (lab) frame. Note that $\phi$ is the gravitational potential and $c$ is the speed of light, is the GR correction. For closure, we use a prescribed relation between the radiation energy and flux, $F^3/E^3 = \frac{1}{2}(1 + \sqrt{1 - [R_e/\text{max}(r, R_e)]^2})$, with $R_e$ being the radius of an energy-dependent scattering sphere (see Equation (11) in Liebendörfer et al. 2009). Since the cell-centered value of the flux, $F^3$, is obtained from the prescribed relation, the cell-interface value is estimated by the first-order upwind scheme assuming that the flux is outward along the radial direction. With the numerical flux, the transport equation of $E^3$ (Equation (1)) is now expressed in hyperbolic form. The velocity-dependent terms $\langle O(\nu/c) \rangle$ are only included (up to the leading order) in the trapped part of the distribution function (Equation (15) in Liebendörfer et al. 2009).

For the three-species neutrinos considered in this work ($\nu_e$, $\nu_x$, $\bar{\nu}_x$), we take into account the collisional kernels up to the zeroth-order expansion with respect to the scattering angle (for example, $\Phi^p_{\nu_e}/\Phi^p_{\nu_x}$ in the case of neutrino pair production from pair annihilation (TP); see Equation (C62) in Bruenn 1985). To be more specific, the neutrino–electron scattering (NES) and TP with the Bruenn rate, both of which were ignored in the original IDSA scheme, are now added to the effective emissivity and absorptivity in the source terms ($S$ and $\Sigma$) as

$$\dot{j} + \dot{\chi} = j(e) + \frac{1}{\lambda(e)} - A^0_{\text{NES}}(e) - A^0_{\text{TP}}(e),$$

$$\Sigma = \min\{\alpha_{\text{diff}} + \alpha_{\bar{\nu}_x} + \alpha_{\nu_x} + \alpha_{\nu_e} - A^0_{\text{NES}}(e) - A^0_{\text{TP}}(e), 0\},$$

$$\alpha_{\nu_e} = \frac{\kappa^3}{(2\pi\hbar c)^3} \left\{ j(e) + c^0_{\text{NES}}(e) + c^0_{\text{TP}}(e) \right\},$$

where $j(e)$ and $\frac{1}{\lambda(e)}$ represent the emissivity and absorptivity of charged-current interactions (e.g., for the first three line reactions in Table 1; see also Equations (A12) of Bruenn 1985), and the exact expression of $A^0_{\text{NES}}$, $A^0_{\text{TP}}$, and $C^0_{\text{NES}}$, $C^0_{\text{TP}}$ is given in Equations (A34), (A43), (A36), and (A45) in Bruenn (1985), respectively, and that of $\alpha_{\text{diff}}$ is given in Equation (7) in Liebendörfer et al. (2009). In this work, heavy-lepton neutrino emission from TP, bremsstrahlung, electron pair neutrino annihilation (10th line reaction in Table 1), and neutrino–neutrino scattering (NNS; 11th line reaction in Table 1) are all treated as the effective emissivity and absorptivity up to the zeroth moment of the neutrino production kernels (e.g., by adding $A^0_{\nu_e}$ and $C^0_{\nu_e}$ to terms in Equations (5) and (6)). As we will show later, this approximation works well at least in the pb accretion phase. However, the consideration of higher moments (Pons et al. 1998), a full set of velocity-dependent terms (equivalently, treatment of full energy-group couplings in the transport equations), and neutrino–flavor coupling is surely needed for more sophisticated simulations (e.g., Rampp & Janka 2002; Thompson et al. 2003; Liebendörfer et al. 2004; Sumiyoshi et al. 2005; Buras et al. 2006b; Hubeny & Burrows 2007; Lentz et al. 2012; Müller et al. 2012b; Kuroda et al. 2016; Nagakura et al. 2017).

In our multi-D simulations, we apply a ray-by-ray approach where the neutrino transport is solved along a given radial direction assuming that the hydrodynamic medium for the direction is spherically symmetric. This also remains to be updated with more advanced schemes (e.g., Skinner et al. 2016; Nagakura et al. 2017).

## 3. 1D Results
Following the seminal code comparison work by Liebendörfer et al. (2005), we first make a quick comparison between our results from the three-flavor IDSA scheme and the results from the two reference codes, *Agile-BOLTZTRAN* (Liebendörfer et al. 2004) and VERTEX (Rampp & Janka 2002). *Agile-BOLTZTRAN* solves the full GR neutrino Boltzmann equation with the $\Sigma_n$ method in a spherically symmetric Lagrangian mesh, whereas VERTEX is an Eulerian code that solves the moment equations of a model Boltzmann equation using the VEF method in Newtonian hydrodynamics plus a modified GR potential with Case R of Marek et al. (2006).

### 3.1. Comparison with 1D Baseline Simulations
We first present Figure 1, which is plotted in a similar way to Figure 10 in Liebendörfer et al. (2005) and which shows the comparison of key neutrino quantities (left panel) and the shock radius (right panel) in IDSA (thick lines), *Agile-BOLTZTRAN* (labeled “AB,” thin lines), and VERTEX (labeled “VX,” dotted lines). Note in the left panel that all luminosities ($L_n$) and root-mean-square (rms) energies ($\langle \epsilon^2 \rangle^{1/2}$) are sampled at a radius of 500 km in the lab frame. The data from the two reference codes are originally given in the fluid frame and are converted to the lab frame using the following relations: $L_n = L_n^{\text{lab}}(1 + v_r/c)/(1 - v_r/c)$ with $v_r$ the radial velocity and $\langle \epsilon^2 \rangle^{1/2} = (\langle \epsilon^2 \rangle^{\text{lab}})^{1/2}W(1 + v_r/c)$ with $W = 1/\sqrt{1 - (v_r/c)^2}$ (i.e., Equations (56)–(58) in Müller et al. 2010). Here, we take $v_r = -0.06c$, which is the (average) infall velocity at 500 km over the entire 250 ms pb in Liebendörfer et al. (2005). Improvements to VERTEX after Liebendörfer et al. (2005) and the follow-up detailed comparison work (Müller et al. 2010) suggest that the data from *Agile-BOLTZTRAN* are currently one of the most reliable ones.

From the left panel of Figure 1, one can see that the neutrino properties from IDSA are much closer to those in AB (Agile-BOLTZTRAN) than to those in VX (VERTEX). Here, it should be emphasized that the higher neutrino luminosities and rms energies of VX are now lowered to match closely those of AB by changing the approximate GR treatment from Case R to Case A (Marek et al. 2006), the latter of which is employed in this work. In that sense, the qualitative agreement of the three codes is convincing, and we will explain this in more detail below.

More quantitatively, the peak luminosity during the electron–neutrino burst (red lines, out of the scale of the left plot) is higher by $\approx8\%$ for IDSA compared to that for AB ($L_{\text{peak}} = 3.3 \times 10^{52} \text{erg s}^{-1}$), whereas the half-width of the peak ($\approx6\ms$) agrees well with each other. After the deleptonization burst, the luminosity of $\nu_e$ (red thin line) and $\nu_x$ (green thick line) for IDSA is higher (maximally by $10\%$) than that for AB until the first $\approx160\ms$ after bounce.

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9 Consideration of up to the first-order angular expansion is technically possible, which, however, makes long-term IDSA simulations unstable at this stage.
already pointed out by Müller et al. (2010) and O’Connor (2015), this most likely comes from the higher resolution at the shock front in the Eulerian codes compared to the Lagrangian code of AB. After the neutronization burst, the maximum luminosity of $\nu_e$ and $\bar{\nu}_e$ deviates maximally by $\sim8\%$ between IDSA and AB, and the luminosity of each neutrino species between IDSA and AB points to a converged value toward the final simulation time ($250\ ms$ after bounce in this comparison).

The $\nu_e$ luminosity agrees quite well between IDSA (blue thick line) and AB (blue thin line) over the entire $250\ ms$ pb. On the other hand, IDSA fails to reproduce the spike in the $\nu_e$ rms energy, peaking at $\sim24\ MeV$, near bounce ($t \sim 0$), which is present in both AB and VX (blue thin and blue dotted lines). This is one of the limitations of the IDSA scheme (M. Liebendörfer 2018, private communication), which attempts to bridge the streaming and trapped neutrinos with the prescribed isotropic diffusion source term. For $\nu_e$ and $\bar{\nu}_e$, the neutrinosphere(s) is well-defined as the decoupling surface from thermal equilibrium mediated by charged-current interactions. IDSA works well in capturing this energy sphere. For $\nu_x$, on the other hand, it is not the energy sphere(s) but rather the scattering sphere(s) that is the decoupling surface (Raffelt 2012). In the transition region from the energy to scattering spheres, Doppler-shift terms (as well as the gravitational redshift) play an essential role in accurately determining the neutrino spectrum. By design, IDSA (in its current form) cannot treat the highly complex transport phenomena appropriately. Moreover, the energy redistribution in the neutrino phase space, such as that due to neutrino–electron scattering, cannot be treated accurately in the current effective emissivity/absorptivity approaches (see Section 2.3). All of these simplifications should potentially lead to the missing spike in the $\nu_e$ energy near bounce.

The $\bar{\nu}_e$ (rms) energy (thick green line) of IDSA is in good agreement with that of AB over the entire $250\ ms$, although IDSA underpredicts the $\nu_e$ energy (red thick line) by $\sim6\%$ compared to AB. After bounce, IDSA underestimates the $\nu_e$ energy (thick blue line) by $\lesssim10\%$ compared to AB until $\sim160\ ms$ after bounce, then closely matches AB during the simulation time. The transition timescale ($\sim160\ ms$) corresponds to the time when the silicon (Si)-rich shell accretes through the shock (e.g., middle-left panel of Figure 5). This can be seen as a hump (solid black line) in the right panel of Figure 1, which leads to a drop in the $\nu_e$ and $\bar{\nu}_e$ luminosities (see red and green solid lines in the left panel). Note that the hump, which was due to the artificial diffusion introduced in the adaptive gridding of AB as already discussed in Liebendörfer et al. (2005), is missing in AB (right panel).

More detailed code comparison not only between AB, VX, and IDSA but also including the Fornax (Skinner et al. 2016), FLASH (O’Connor & Couch 2015), GR1D (O’Connor 2015) codes is currently in progress (O’Connor et al. 2018, in preparation). Given our approximate treatment of GR, neglect of energy-bin couplings, and the partial implementation of the Doppler-shift terms, it may not be surprising that IDSA has $10\%$ levels of mismatch with the full GR and full Boltzmann result of Agile-BOLTZTRAN. In fact, such discrepancies (given the use of similar levels of approximation) have also been observed in the literature (e.g., O’Connor 2015; O’Connor & Couch 2015). Having given an overview of the validity and limitation of the current IDSA scheme, we will move on to focus on the microphysics update in the following sections.

3.2. Improved Electron-capture Rate on Heavy Nuclei (set2)

We start by describing our first update to neutrino opacity, which is electron capture on heavy nuclei. Based on detailed shell-model calculations by Langanke & Martínez-Pinedo (2000), Langanke et al. (2003) showed that it is not only a generic $0\gamma_{j2} \rightarrow 0\gamma_{j2}$ Gamow–Teller (GT) transition, but also additional GT transitions, forbidden transitions, and thermal unblocking that play a crucial role in making electron capture
Figure 2. Left panel shows the comparison of the emissivity ($j$) and absorptivity ($\chi$) of the electron capture on heavy nuclei at the thermodynamic condition of $\rho = 10^{11}$ g cm$^{-3}$, $Y_e = 0.45$, $T = 10^{10}$ K (corresponding to $\mu_e = 18.2$ MeV) between the Bruenn rate (blue dashed line) and the rate from Juodagalvis et al. (2010; red solid line). Here, $\rho$, $T$, $Y_e$, and $\mu_e$ denote the density, temperature, electron fraction, and electron chemical potential, respectively. The absorptivity (equivalently, $1/\chi$) in Bruenn (1985), i.e., Equation (C29) is calculated through detailed balance. The right panel shows the comparison of $Y_f$ and lepton fraction ($Y_l$) as a function of the central density in set1 (dashed lines) and set2 (solid lines) using either the Bruenn or Juodagalvis rate, respectively.

We employ the electron-capture rates on heavy nuclei tabulated by Juodagalvis et al. (2010), who has significantly extended the covered mass range of nuclides ($\sim$2700) compared to Langanke & Martínez-Pinedo (2000; $\sim$100). To calculate the needed abundances of heavy nuclei, a Saha-like nuclear-statistical equilibrium is assumed. The table by Juodagalvis et al. (2010) then provides an average neutrino emissivity per heavy nucleus. Following Hix et al. (2003), one can calculate the full neutrino emissivity as the product of this average neutrino emissivity and the number density of heavy nuclei calculated using the employed EOS (here for the LS EOS).

The left panel of Figure 2 compares the full neutrino emissivity/absorptivity ($j$, $\chi$; red solid line) with that of the Bruenn prescription (blue dashed line; e.g., Equation (C27) of Bruenn 1985) for a given thermodynamic condition in the core-collapse phase. From the panel, one can see a significant enhancement of the neutrino emissivity in the Juodagalvis rate (red solid line) compared to the Bruenn rate (blue dashed line) for neutrino energies above $\sim$18 MeV. This high electron-capture rate is also seen as a steeper slope in $\chi$ than that of the Bruenn rate ($\chi \propto \epsilon^2$, with $\epsilon$ representing the neutrino energy).

The right panel of Figure 2 compares $Y_f$ and lepton fraction ($Y_l$) as a function of the central density ($\rho_c$) between set1 (dashed line) and set2 (solid line) with the Bruenn or Juodagalvis rate, respectively. Deleptonization ends approximately at $\rho_c = 2 \times 10^{12}$ g cm$^{-3}$, which marks the onset of neutrino trapping (Sato 1975). The central $Y_f \sim 0.34$ at bounce (see the right edge of the blue dashed curve) using the Bruenn rate is lowered by about $\sim$10% ($Y_f \sim 0.3$; blue solid curve) with the Juodagalvis rate. For set1, the evolution of $Y_f$ and $Y_l$ matches quite nicely with that of AB (Liebendorfer et al. 2005; see their Figure 7(b)). For set2, our results are quantitatively very close to those of Hix et al. (2003) where the Langanke and Martínez-Pinedo (LMP) rate (Langanke et al. 2003) was implemented in the AB run using the same 15 $M_\odot$ progenitor model (Woosley & Weaver 1995).

To make the comparison easy, Figures 3 and 4 are plotted in a similar way to Figures 1 and 2 in Hix et al. (2003; see also Martínez-Pinedo et al. 2006). Note that LS180 was used in Hix et al. (2003); however, the differences when using a different $K$ of the LS EOS are only a few percent around core bounce and less than $\sim$10% in the first 200 ms after bounce (Thompson et al. 2003; Müller et al. 2010). So, we consider that a different choice of $K$ does not significantly affect the comparison here.

In fact, Figure 3 shows a nice agreement with Figure 1 of Hix et al. (2003). From the top panel, the use of the improved electron-capture rate leads to a $\sim$10% reduction in the central $Y_f$ and $Y_e$ compared to those with the Bruenn rate. The above match suggests that the hydrodynamics impact on the LMP rate and the Juodagalvis rate should be fairly small. From the velocity profile (bottom panel), one can see that the mass of the

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10 See Sullivan et al. (2016) for a more detailed, nucleus-by-nucleus investigation of electron capture on heavy nuclei in the CCSN context.
unshocked, homologous core at bounce is reduced by about 
\( \sim 20\% \) from \( \sim 0.62 \, M_\odot \) for set1 (see the discontinuity in the velocity plot) to \( \sim 0.5 \, M_\odot \) for set2. As already pointed out by Hix et al. (2003), the 10% reduction in \( Y_L \) leads to a \( \sim 20\% \) reduction in the homologous core because the (Chandrasekhar) mass of the unshocked core scales as \((Y_L)^2\). The smaller entropy and central density (second and third panels of Figure 2) for set2 compared to set1 are also quantitatively consistent with Hix et al. (2003).

**Figure 4.** Comparison of the neutrino luminosity (left panel) and the rms energy (right panel) measured at 500 km in the lab frame.

**Figure 5.** Comparison of several key quantities in 1D simulations with either neutrino physics set1 or set2. In the top panels, we show the neutrino luminosities (left) and the neutrino rms energy (right) measured at a radius of 500 km in the lab frame. The bottom-left panel shows the mass-accretion rate evaluated at 500 km. The bottom-right panel compares the net heating rate integrated over the gain region.

**Figure 4** also supports the correct implementation of the Juodagalvis rate in our code. The left panel shows that because of the enhanced electron-capture rate and the resulting smaller radius at the shock formation (the bottom panel of Figure 3), the shock breakout is slightly delayed and the duration becomes slightly longer (a few ms) for set2 (red solid line) compared to set1 (dashed red line). In accordance with Hix et al. (2003), this feature is also seen in other neutrino flavors near core bounce (\( \lesssim 50 \, \text{ms} \)) pb; see the green and blue curves in
the left panel). An overall trend in the rms neutrino energy (right panel of Figure 4) both in the pre- and postbounce phases is also in line with that in Hix et al. (2003); the neutrino rms energy is as much as \( \lesssim 1 \) MeV smaller for set2 than set1 over the first 50 ms after bounce (compare the red solid with the red dashed line), but lower thereafter. The difference of \( \nu_e \) is minute compared to that of \( \bar{\nu}_e \).

Figure 5 summarizes several key quantities for comparison between set1 and set2 over the first 500 ms after bounce. The top-left panel shows that the largest difference is a \( \lesssim 15\% \) reduction of the \( \nu_e \) luminosity of set2 (solid green line) compared to set1 (dashed green line) in the first \( \sim 140 \) ms after bounce, but vanishes thereafter. The bottom-left panel shows that the drop in the mass accretion rate (\( \sim 140 \) ms) corresponds to the epoch when the Si-rich layer is passing through the shock. This leads to a bump in the \( \nu_e \) and \( \bar{\nu}_e \) luminosities (see red and blue curves in the top-left panel), the peak of which is at around the \( \sim 140 \) ms after bounce. In the first 160 ms after bounce, the Si-rich layer has entirely passed though the shock as indicated by the low mass accretion in the bottom-left panel. At this phase, one can see the \( \lesssim 5\% \) reduction of the \( \nu_e \) and \( \bar{\nu}_e \) luminosities in set2 compared to set1. The rms energy (top-right panel) for all neutrino flavors becomes smaller (maximally by \( \sim 0.4 \) MeV) in set2 (solid lines) compared to set1 (dashed line). The smaller homologous mass at bounce for set2 could potentially lead to a smaller gravitational energy release in the postbounce phase compared to set1, which is reconciled with the above trends. The bottom-right panel shows that the net heating rate for set2 (red line) is smaller by \( \lesssim 7\% \) compared to set1 (blue line).

Regarding the update in this section, one would imagine from the 1D comparison that the improved electron-capture rate on nuclei might weaken the “explodability” in multi-D simulations. Readers will see whether or not this expectation is correct in Section 4 (2D results).

3.3. Electron–Neutrino Pair Annihilation (set3a) and \( \nu_e + \bar{\nu}_e (t_e) \) Scattering (set3b)

In this section, we first focus on set3a (see Table 1), where the electron–neutrino pair annihilation process \( (\nu_e + \bar{\nu}_e \rightarrow \nu_e + \bar{\nu}_e) \) for short, we call this the “nupair” process in this section) is added to set1. Buras et al. (2003) were the first to point out that as a source for \( \nu_e \), the nupair reaction is always more important than the traditional electron–positron pair annihilation process \( (e^- e^+ \rightarrow \nu_e \bar{\nu}_e) \), for short, we call this the “eepair” process in the following). The implementation scheme of the nupair process in IDSA is given in Appendix A.

In fact, the left panel of Figure 6 clearly shows the dominance of the nupair process (red solid line; labeled “\( \nu_e \bar{\nu}_e \)”) over the eepair process (blue solid line; labeled “\( e^- e^+ \)”)

In this panel, the chemical potential for electrons and neutrinos is set to zero (\( \mu_e = 0 \) and \( \mu_{\nu_e} = 0 \)), following Figure 3 of Buras et al. (2003). The peak in the spectra of the nupair process is about two times higher than that of the eepair process at the neutrino energy of \( \sim 40 \) MeV. As explained in Buras et al. (2003), this simply comes from the different weak coupling constants between the two processes. The production kernels have similar forms, \( \Phi^p = (C_V - C_A)^2 j^H + (C_V + C_A)^2 j^H \) (i.e., Equation (C63) in Bruenn 1985). For the eepair process, \( C_V = -1/2 + 2 \sin^2 \theta_w \) and \( C_A = -1/2 \) (with the weak mixing angle; \( \sin^2 \theta_w = 0.23 \)), whereas \( C_V = 1/2 \) and \( C_A = 1/2 \) for the nupair process. Given that \( j^H \sim j^p \), the two-times difference can be readily seen by putting these numbers in the above equation for \( \Phi^p \).

The right panel of Figure 6 shows the comparison between the energy production rates with (solid lines) and without chemical potentials (dashed lines; labeled with \( \mu_e = 0 \), \( \mu_{\nu_e} = 0 \)). As a reference, the production rate of nucleon bremsstrahlung is also shown (labeled “\( NN \)”)). In accordance with Fischer et al. (2009), the chemical potentials make the spectra harder and the rate smaller. Quantitatively, the \( \eta \) parameter of \( \sim 10 \) in the right panel (e.g., for electron, \( \eta_{\nu_e} = \mu_{\nu_e}/T = \sim 96 \text{ MeV}/12 \text{ MeV} \), and electron–neutrino, \( \eta_{\nu_e} = \mu_{\nu_e}/T = \sim 123 \text{ MeV}/12 \text{ MeV} \)). This leads to a \( \lesssim 1/10 \) reduction in the rates with the chemical potential.
compared to those without. This is consistent with Buras et al. (2003; their Figure 2). Comparing to Figure 16(a) of Fischer et al. (2009), the peak energy of bremsstrahlung is around \( \sim 30 \text{ MeV} \), and the spectral shape matches well together.\(^{12}\)

The top-left panel of Figure 7 compares the neutrino luminosities between set3a (solid lines) and set1 (dashed lines). Before the first 160 ms after bounce, when the mass accretion rate is still high (e.g., the bottom-left panel of Figure 5), the influence of the nupair process is biggest for the \( \nu_x \) luminosity (\( \lesssim 20\% \) increase from set1, compare the green solid with the green dashed line), followed by the \( \nu_e \) luminosity (\( \sim 4\% \) increase, compare the red solid line with the red dashed line) and by the \( \bar{\nu}_e \) luminosity (\( \lesssim 2\% \) increase, compare the blue solid line with the blue dashed line). Using the same 15 \( M_\odot \) progenitor, this is quantitatively very close to the results in Buras et al. (2003; see their Figure 7). Consistent with Buras et al. (2003) and Fischer et al. (2009), the additional source of \( \nu_L \) increases the \( \nu_x \) luminosity most significantly. In the accretion phase (\( \lesssim 160 \text{ ms pb} \)), the nupair process also leads to higher rms neutrino energies (top-right panel of Figure 7). This is due to the enhanced cooling, which makes the positions of the neutrino spheres and the PNS formed deeper inside (see Figure 17 of Fischer et al. 2009 for the corresponding \textit{Agile-BOLTZTRAN} run but for a 40 \( M_\odot \) star). After the accretion phase (\( \gtrsim 160 \text{ ms pb} \)), the luminosities of \( \nu_e \) and \( \bar{\nu}_e \) become smaller than those in set1 (top-left panel of Figure 7). As already pointed out by Buras et al. (2003), this is most likely because of the more compact neutrino spheres (e.g., smaller emission region) in response to the more accelerated PNS contraction. The (maximum) shock position in set3a is more compact compared to that in set1 by 5\% \( \sim 10\% \), which is within the change seen in Buras et al. (2003) and Fischer et al. (2009). From the bottom-right panel of Figure 7, it is interesting to note that the net heating rate of set3a (red line) dominates over that of set1 (blue line) in the accretion phase (\( \lesssim 160 \text{ ms pb} \)), which reverses thereafter (until \( \sim 400 \text{ ms pb} \)). This is in line with the higher (and lower) luminosities of \( \nu_e \) and \( \bar{\nu}_e \) of set3a compared to those of set1 in the pre- and (post-) accretion phases, respectively, as already mentioned above.

As originally pointed out by Buras et al. (2003), the cross channel of the nupair process, that is, \( \nu_e + \nu_e/\bar{\nu}_e = \nu_x + \nu_x/\bar{\nu}_x \), could be of comparable importance to \( \nu_e + e^\pm \) scattering. The top- and bottom-left panels of Figure 8 compare the (inverse) mean free path of the \( \nu_e + e^- \) scattering (red line), \( \nu_e + e^+ \) scattering (blue line), and \( \nu_x + \nu_x \) scattering (green line) for typical thermodynamics conditions in the supernova core, respectively. Note that the corresponding reactions with \( \bar{\nu}_e \) and \( e^- \) are not shown in the panels because they are much smaller compared to those with \( \nu_e \) and \( e^- \). In the prebounce phase, the top panels of Figure 8 show that \( \nu_e \nu_e \) scattering (green line) is almost comparable to \( \nu_e e^- \) scattering (blue line). But they play minor roles in the opacities (in the leptonic channels) because of the dominant contribution from \( \nu_e e^- \) scattering (red line). In the

\(^{12}\) All of the related opacities in Table 1 are compared with those in Fischer et al. (2009, 2012) and Fischer (2016). For the Bruenn rate, the opacity plots are already shown in Kuroda et al. (2016).
Figure 8. Inverse mean free path (e.g., Equation (25)) as a function of neutrino energy for three typical conditions: near neutrino trapping (top left; $\rho = 10^{12}$ g cm$^{-3}$, $T = 1.76$ MeV, and $Y_e = 0.35$), near core bounce (top-right panel; $\rho = 3 \times 10^{13}$ g cm$^{-3}$, $T = 12$ MeV, and $Y_e = 0.27$), and in the postshock region behind the shock (bottom-left panel; $\rho = 1 \times 10^{13}$ g cm$^{-3}$, $T = 7$ MeV, and $Y_e = 0.10$). Here, Fermi–Dirac final state neutrino distributions are assumed. The bottom-right panel compares the neutrino luminosities between set3a and set3ab. Note that set3ab is the run where $\nu_e + \nu_e$ scattering is added to set3a.

pb phase (bottom-left panel of Figure 8), the dominance of $\nu_e e^-$ scattering is also unchanged, but the opacity of the $\nu_x e^-$ scattering becomes higher than that of the $\nu_x \nu_e$ scattering, as previously shown in Buras et al. (2003).

The bottom-right panel of Figure 8 compares the neutrino luminosities between set3a (dashed line) and set3ab (solid line). Note that set3ab is the run where $\nu_x + \nu_x (\pi_e)$ scattering is added to set3a. The solid and dashed lines completely overlap, which confirms the expectation that $\nu_x \nu_e$ scattering plays a very minor role at least over the first 500 ms pb.

3.4. Mean-field Modifications (set4a and set4b)

Martínez-Pinedo et al. (2012) and Roberts et al. (2012) clearly pointed out that medium effects (Reddy et al. 1998) affect protons and neutrons differently (e.g., the reactions of set4a in Table 1), leading to a significant impact on the neutrino luminosities and spectra especially in the PNS cooling phase (after the onset of an explosion). By definition, our (non-explooding) 1D simulation can cover only the pre-explosion phase. Keeping in mind future applications for long-term evolution in multi-D (exploding) models, we explore in this section the impact of the mean-field corrections on the charged-current opacities treated at the elastic level (Martínez-Pinedo et al. 2012; Roberts et al. 2012).

From Equation (3) of Martínez-Pinedo et al. (2012), the opacity of $\nu_e$ absorption on neutrons ($\nu_e n \rightarrow e^- p$) is expressed as

$$\frac{1}{\lambda_{\nu_e}} \propto E_{\nu_e}^2 \left[1 - f_e(E_{\nu_e})\right] \frac{n_n - n_p}{\Gamma - \exp \beta (\mu_p - \mu_n + \Delta U)},$$

(7)

where $E_{\nu_e}$ is the electron energy, $f_e$ is the electron distribution function, $n_i$ and $\mu_i$ are the number density and chemical potential (without rest mass) for $i = n, p$ (neutrons and protons), respectively, and $\beta$ is the inverse temperature. At the level of elastic approximation (Reddy et al. 1998; Martínez-Pinedo et al. 2012), the following relation holds:

$$E_{\nu_e} = E_{\nu_{\bar{e}}} + Q + \Delta U,$$

(8)

where $E_{\nu_{\bar{e}}}$ is the $\nu_e$ energy, $Q = m_n - m_p$ is the so-called $Q$ value with $m_i$ the rest mass for $i = n, p$, and $\Delta U = U_n - U_p$ is the difference of the mean-field potentials of neutrons and protons.\textsuperscript{13}

From Equation (8), the $E_{\nu_e}^2$ in Equation (7) becomes larger due to $\Delta U$, which leads to an increase in the $\nu_e$ opacity compared to the free gas case ($\Delta U = 0$) at lower neutrino

\textsuperscript{13} Note that for a neutron-rich environment (like in the pre-explosion phase), $\Delta U > 0$ (e.g., Roberts et al. 2012).
energies. At larger neutrino energies, Pauli blocking disappears, making the opacity with and without the mean-field effects approach each other closely (e.g., Martínez-Pinedo et al. 2014 for more details). For $\bar{\nu}_e$, the positron energy becomes $E_e^\nu = E_{\nu_e} - Q - \Delta U$. This leads to the reduction of the opacity at lower neutrino energies. Note also that the $Q$ value of this reaction increases from $E_{\nu_e} > Q$ to $E_{\nu_e} > Q + \Delta U$.

Figure 9 is consistent with the above explanations, which compare the inverse mean free path for $\nu_e$ (red lines) and $\nu_e$ (blue lines) with (dashed lines, labeled “in-medium”) and without (solid lines) mean-field corrections. The thermodynamics conditions of $T = 8$ MeV, $n_B = 0.02$ fm$^{-3}$, and $Y_e = 0.027$ are chosen with $n_B$ the baryon number density, which corresponds to Figure 3 of Roberts et al. (2012). Note that the nucleon potential difference is $\Delta U = 9$ MeV for the EOS used in Roberts et al. (2012), whereas $\Delta U = 7.67$ MeV for LS220 in this work, leading to a slight difference quantitatively.

The bigger mean-field effects observed in this study, such as those on the $\nu_e$ luminosity compared to those on the $\bar{\nu}_e$ luminosity and the same for the $\bar{\nu}_e$ rms energy compared to the $\nu_e$ rms energy, are consistent with those in Horowitz et al. (2012). Note that Horowitz et al. (2012) observed a stronger impact of the mean-field effects, especially on the increase of the $\bar{\nu}_e$ rms energy and the reduction of the $\nu_e$ luminosity (see their Figure 4). The employed progenitor (15 $M_\odot$) and EOS (LS220) are the same as those in this work. However, the quantitative differences with Horowitz et al. (2012) could originate from their use of a $\Delta U$ obtained by a virial expansion calculation (not from the LS220 EOS as in this work), the inclusion of the weak magnetism correction (not included in our set4a and set1), and GR hydrodynamics (essentially Newtonian hydrodynamics in this work).14

The left panel of Figure 11 compares the shock radius between set4a and set1. The shock radius becomes slightly bigger for set4a (red line) compared to that for set1 (blue line) for a short period ($\sim$120–160 ms after bounce) when the Si-rich layer is advecting through the shock, but the difference disappears thereafter. Note that in this period the $\bar{\nu}_e$ luminosity is bigger than the $\nu_e$ luminosity (left panel of Figure 10). As mentioned above, the mean-field effects of set4a lead to a higher $\bar{\nu}_e$ luminosity than in set1, which is consistent with the bigger shock radius, albeit transiendly.

The right panel of Figure 11 compares the net heating rate in the gain region, suggesting that the mean-field effects, as previously reported (e.g., Horowitz et al. 2012; Martínez-Pinedo et al. 2012), would not have a significant impact on the onset of an explosion. A similar comparison between set4b and set1 shows that the difference with set1 due to the in-medium suppression of bremsstrahlung (see Table 1) is much smaller compared to the mean-field effects mentioned above. The comparison plots between set4b (not shown) and set1 almost completely overlap (like in the middle-left panel of Figure 12 or in the bottom-right panel of Figure 8). It was shown (e.g., Fischer 2016 and Bartl et al. 2016) that the suppression of bremsstrahlung by the medium only clearly affects the neutrino properties after the onset of an explosion and the subsequent PNS cooling phase. In this respect, our results showing a negligible impact on the pre-explosion phase are in line with the literature.

3.5. Weak Magnetism and Recoil (set5a) and Nucleon Effective Mass (set5b)

In order to take into account the effects of weak magnetism and recoil both on the charged-current (CC) and neutral-current (NC) reactions, we follow Horowitz (2002; their Equations (22) and (32)). The top-left panel of Figure 12 corresponds to Figure 1 (Horowitz 2002), showing that the main effect from weak magnetism is to reduce the $\bar{\nu}_e$ opacity (solid line) by a large amount ($\sim$15% reduction at a neutrino energy of 20 MeV). In comparison, the $\nu_e$ opacity is enhanced only by a small amount (dashed line). Regarding NC reactions, Figure 2 of Horowitz (2002) shows that the reduction of the opacity is slightly higher for $\nu p$ scattering than for $\nu n$ scattering, and that the reduction of the $\bar{\nu}_e$ reactions is higher than the corresponding $\nu$ reactions ($\sim$10% reduction for $\bar{\nu}_e$ at a neutrino energy of 20 MeV).

14 Note also that a comparison with Martínez-Pinedo et al. (2012) is more difficult because they focused on the later postbounce evolution (after $\sim$500 ms) of a different progenitor (18 $M_\odot$ star) using a different EOS (Shen et al. 1998) in the Agile-BOLTZTRAN run.
The top-right panel of Figure 12 compares the neutrino luminosities between set5a (solid lines) and set1 (dashes line) with and without weak magnetism and recoil, respectively. One can clearly see the enhancement of the $\bar{\nu}_e$ luminosity (blue solid line) for set5a, which is $\sim$8% bigger than that in set1 (blue dashed line). This comes from the reduction of the $\bar{\nu}_e$ opacity as mentioned above. The difference, however, becomes very small after $\sim$340 ms pb. The $\nu_e$ luminosities (red solid line and red dashed line) are hardly affected, which is in line with the very small change in the $\nu_e$ opacity. The $\nu_x$ luminosity (green solid line and green dashed line) is enhanced by up to $\sim$10% for set5a compared to set1. Regarding the rms neutrino energies (middle-right panel), the reduced opacities of $\bar{\nu}_e$ and $\nu_x$ result in the higher $\bar{\nu}_e$ (blue solid line) and $\nu_x$ energies (green solid lines) up to $\sim$1 MeV, compared to those of set1 (blue dashed line and green dashed line).

From the bottom-left panel of Figure 12, one can see that the shock radius of set5a is transiently bigger than that of set1 for the epoch when the Si-rich shell is passing. This most likely comes from the higher $\bar{\nu}_e$ luminosity and energy (top-right and middle-right panels). In fact, the net heating rate is slightly higher for set5a (red line) compared to set1 (blue line) up to the first $\sim$340 ms after bounce. Thereafter, the $\bar{\nu}_e$ luminosities with and without the weak magnetism correction approach each other (top-right panel). Set5b (effective mass correction; see Table 1) does not exhibit visible changes from set1 (as was the case for set3b and set4b); only the comparison plot of the neutrino luminosities (middle-left panel) is shown as a reference.

### 3.6. Quenching of $g_a$ (set6a), Many-body Effect (set6b), Strangeness Contribution (set6c), and the Whole Set

Finally, the model series with “set6” includes modifications to the axial-vector currents in the weak interactions either from in-medium effects (set6a), many-body effects (set6b), or strangeness-dependent contributions (set6c), respectively (see Section 2.2 and Table 1 for details).

From the top panels of Figure 13, it is very hard to see significant differences between set6a and set1. This suggests that the quenching of $g_a$ plays a negligible role in the first 500 ms after bounce covered in our 1D run. For set6b, the middle-left panel shows that the $\nu_e$ luminosity is higher by $\sim$10% (green solid line) compared to set1 (green dashed line). The relative difference becomes larger in the later pb phase predominantly because many-body effects reduce the opacity of the $\nu N$ scattering at high densities (Horowitz et al. 2017). This is also the case for the $\nu_x$ and $\bar{\nu}_x$ luminosities, where the luminosities become higher by $\sim$3%–4% for set6b toward the final simulation time. The clearer impact of the many-body effects on $\nu_x$ compared to that on $\nu_e$ and $\bar{\nu}_e$ is also seen in the middle-right panel, showing an increase of $\sim$1 MeV in the $\nu_x$ energy for set6b (green solid line) compared to set1 (green dashed line).

The bottom-left panel of Figure 13 shows a clear increase of the $\nu_x$ and $\bar{\nu}_x$ luminosities (by $\sim$4%) for set6c compared to set1 in the first $\sim$160 ms after bounce. At this epoch, the increase in the $\nu_x$ luminosity is bigger ($\sim$9%). The bottom-right panel shows that the strangeness effects lead to a slight increase in the rms neutrino energies where the maximum upshift is $\sim$0.2 MeV in the $\nu_x$ energy (green solid line and green dashed line). These trends with the strangeness contribution are qualitatively consistent with those in Melson et al. (2015a). In the full-scale 3D simulations by Melson et al. (2015a), they observed much bigger effects from the strangeness effects, such as $\sim$30% and 10%–15% increases in the $\nu_x$ and $\nu_e/\nu_x$ luminosities, respectively, and a $\sim$1 MeV increase in the mean neutrino energies. Note that in Melson et al. (2015a) the use of the larger value of $g_a^s = -0.2$ and the choice of the more massive progenitor with the higher mass accretion rate (a 20 $M_\odot$ star) could potentially have led to the clearer impact of the strangeness effect compared to this work (see also Bollig et al. 2017 for 2D results using $g_a^s = -0.1$).

The top-left panel of Figure 14 shows that the maximum shock extent becomes $\sim$5% bigger for set6b and set6c compared to set1 near the hump region ($\sim$160 ms after bounce). The mentioned higher $\nu_e$ and $\bar{\nu}_e$ luminosities in the accretion phase (e.g., Figure 13) are in line with this feature. In fact, the right panel of Figure 14 shows that the net heating rate...
Figure 11. Comparison of the shock radius (left panel) and the net heating rate in the gain region (right panel) for set1 and set4a.

for set6b (red line) and set6c (green line) is bigger than that for set1 (blue line). For set6b, the increase from set1 (blue line) is \( \sim 6\% \) around 100 ms after bounce and higher at later times. This is in good agreement with Horowitz et al. (2017; see their Figure 3). With respect to the strange-quark contribution, the heating rate of set6c (green line) becomes larger than that of set1 by \( \sim 12\% \). This is in accordance with Horowitz et al. (2017). Note that a significantly bigger impact (\( \sim 20\% \) increase) was observed by Horowitz et al. (2017) probably because of the larger value of \( g_s = -0.2 \) and the use of a more massive 20 \( M_\odot \) progenitor.

Finally, Figure 15 compares the model including all of the updates to set1 (from set2 to set6c in Table 1, labeled “set-all” in the panels\(^{[5]} \) with set1. Comparing with set1, the top-left panel shows that the largest increase of the neutrino luminosities is for \( \nu_e \) (by \( \sim 31\% \)), which is followed in order by \( \bar{\nu}_e \) (\( \sim 14\% \)) and \( \nu_x \) (\( \sim 11\% \)). Among the individual updates, the increase of the \( \nu_x \) luminosity is the biggest (\( \sim 20\% \)) due to the inclusion of electron–neutrino pair annihilation (set3a) as shown in the top-left panel of Figure 7. Note that the increase from the weak magnetism and recoil is \( \sim 10\% \) in set5b (top-right panel of Figure 12), from the many-body effects is \( \sim 8\% \) in set6b (middle-left panel of Figure 13), and \( \sim 9\% \) from the strangeness contribution (middle-left panel of Figure 13). As is expected, simply adding the contributions from these individual rates (for the \( \nu_x \) luminosity) does not explain the total increase (\( \sim 31\% \)). This is not surprising because the individual update nonlinearly affects the postbounce evolution, which is governed by nonlinear neutrino-radiation hydrodynamics.

The \( \nu_x \) luminosity is bigger (\( \sim 5.4\% \)) than the \( \nu_e \) luminosity at the peak around 100 ms after bounce; thereafter, the difference becomes smaller toward the final simulation time. As already seen from Figure 12 (top-left panel), the dominance of the \( \nu_x \) over the \( \nu_e \) luminosity is mainly due to the inclusion of weak magnetism and recoil (set5a). Using the same 15 \( M_\odot \) progenitor (Woosley & Weaver 1995), this feature is also seen in Müller & Janka (2014; see their Figure 1, the panel labeled “s15z7b2”), where neutrino signals from 2D GR simulations using the Vertex–CoCoNuT code were investigated. In Müller & Janka (2014), the main difference regarding the microphysics inputs from this work is the use of the LS180 EOS and the inclusion of non-elastic effects in the charged-current absorption reactions (Burrows & Sawyer 1998, 1999).

The energy redistribution from the latter would make the recoil effect smaller, which would explain the smaller difference between the \( \nu_e \) and \( \nu_x \) luminosities in Müller & Janka (2014). Their 2D run of the 15 \( M_\odot \) star (G15) starts to explode \( \sim 570 \) ms after bounce, and the postbounce dynamics deviates from 1D around 100 ms after bounce (Müller et al. 2013). At 100 ms after bounce, their \( \nu_e \) and \( \bar{\nu}_e \) luminosities are \( \sim 5 \times 10^{52} \) erg s\(^{-1} \), which is slightly lower than those in this study, \( \sim 5.3 \times 10^{52} \) erg s\(^{-1} \). Note that the neutrino luminosities in Müller et al. (2013) take into account the GR effects (their Equations (2) and (3)), which could potentially lead to a \( \sim 10\%–20\% \) reduction compared to those without the GR corrections. Regarding the \( \nu_e \) luminosity, the peak value is \( \sim 2.4 \times 10^{52} \) erg s\(^{-1} \) in Müller et al. (2013), which is \( \sim 25\% \) lower than that in this work. Although we do not have a clear-cut answer, we consider that the reduction of the \( \nu_e \) luminosity due to GR redshift effects could partly explain the discrepancy. This is because the neutrinospheric radii of \( \nu_e \) are formed deeper inside where the GR correction becomes more significant among the other neutrino species.

The top-right panel of Figure 15 shows that all of the rms neutrino energies become higher for set-all (solid lines) than set1 (dashed lines) over the first 500 ms after bounce. The enhancement due to the updated opacity is bigger for \( \nu_x \) and \( \nu_e \) by \( \sim 2 \) MeV compared to \( \bar{\nu}_e \) by \( \sim 1 \) MeV. Note that in Müller et al. (2013) it was not the rms but the mean neutrino energy that was plotted. At 100 ms after bounce, the mean energy is (probably incidentally) very close, and \( \nu_x, \bar{\nu}_x \), and \( \nu_e \) are \( \sim 9 \) MeV, \( \sim 13 \) MeV, and \( \sim 15 \) MeV in Müller et al. (2013), and which are \( \sim 10 \) MeV, \( \sim 13 \) MeV, and \( \sim 15 \) MeV in our work, respectively.

The middle-left panel of Figure 15 shows that the shock position where the bounce shock stalls (at 100 ms after bounce) is \( \sim 3\% \) smaller for set-all (red line) compared to set1 (blue line). At the hump that marks the passing of the Si-rich layer through the shock (\( \sim 160 \) ms after bounce), the difference in the shock between set-all and set1 becomes largest, \( \sim 9\% \). After \( \sim 300 \) ms pb, the two shock radii approach very closely, but the shock radius of set-all becomes \( \sim 5\% \) bigger compared to set1 toward the final simulation time. The enhanced \( \nu_x \) and \( \bar{\nu}_x \) luminosities due to the many-body effects (set6b; see the middle-left panel \( \sim 300 \) ms pb) and the extended shock radius

\(^{[5]} \) Note that the Pauli blocking factor of nucleons has already been included in Horowitz et al. (2017), which one should not double count.
(red line in the left panel of Figure 14) are reconciled with the above features seen in set-all. The middle-right panel of Figure 15 shows the more compact PNS radius for set-all (red line) compared to set1 (blue line). Note that the PNS radius is estimated at a fiducial density of $10^{15}$ g cm$^{-3}$. The difference is biggest at the hump seen in the shock evolution ($\sim$160 ms after bounce), when the PNS radius is smaller by $\sim$17% for set-all relative to set1. Although the difference becomes smaller toward the final simulation time, the PNS radius is always smaller for set-all over the entire 500 ms after bounce.

The bottom panel of Figure 15 shows that the maximum enhancement of the net heating rate of set-all (red line) is $\sim$30% compared to set1 (blue line) around 100 ms after bounce. After 160 ms pb, the net heating rate in the gain region becomes $\sim$10%–24% higher for set-all. This is predominantly because of the higher $\nu_e$ and $\bar{\nu}_e$ luminosities and rms energies (top-left and right panels) and the smaller PNS radius (middle-right panel). As already discussed above, the improved opacities add nonlinearly and synergistically to increase the net heating rate, where each of the
individual update amounts to the \( \lesssim 10\% \) level (e.g., see Figures 5–14).

4. 2D Results

In this section, we present the results of our 2D core-collapse simulations where the selected sets of neutrino opacities shown in Table 1 are included in set1. As already mentioned in Section 2.1, we choose a 20 \( M_\odot \) star (Woosley & Heger 2007) in our 2D runs.

Our 2D run with the Bruenn rate (set1) is now called “G1” (meaning group one). “G2” is the model that is equivalent to set2 in the previous section. “G3” is the model where set3a and set3b are added to G2. In the same way, for “G4,” set4a and set4b are added to G3, and for “G5,” set5a and set5b are added to G4. For “G6ab,” set5b, set6a, and set6b are added to G5. Note that for G6ab we collectively add the three updates to G5 because set5b and set6a have no visible impact on our 1D runs. The model difference between G6ab and G5 is the inclusion of the many-body effects of Horowitz et al. (2017). For “G6abc,” set5b, set6a, and set6c are added to G5. Finally “G6abc” corresponds to set-all in our 1D models where the strangeness contribution is added to G6ab.
Figure 14. Comparison of the shock radius and the net heating rate in the gain layer between set6b, set6c, and set1. Note that set6a is not shown because of the overlap with set1, which makes the differences in the plots difficult to see.

Figure 16 shows a compact overview of all of the 2D runs. Up to the final simulation time of \(\sim 600\) ms after bounce, we observe the shock revival only for G6ab and G6abc. This may not be very surprising because in 1D, the many-body effects (set6b) and the strangeness contribution (set6c) are expected to primarily enhance the explosability (see, e.g., the right panel of Figure 14). We furthermore explain in detail the reason why G6ac, which simply includes the strangeness correction to G6a, does not lead to explosion.

Figures 17–19 show several key quantities useful for our 2D model comparison. In each model in Figure 17, the top-left panel shows the time evolution of the average shock radius (black line) with the mass accretion rate \((M_{\text{dot}})\) at 500 km (green line), the top-right panel shows the net heating rate in the gain region, the bottom-left panel is the PNS radius, and the bottom-right panel is the diagnostic explosion energy. The model name is indicated in the upper-right corner of the top-left panel of each subfigure.

The top two panels of Figure 17 compare G1 (left panel) and G2 (right panel), where the Juodagalvis (electron-capture) rate is implemented in G2 instead of the Bruenn prescription of G1. These two panels are almost identical with respect to the shock evolution (top left), PNS contraction (bottom left), and the non-explodability \((E_{\text{dia}} \approx 0, \text{ bottom right})\). Here, \(E_{\text{dia}}\) denotes the diagnostic (explosion) energy that is calculated following the literature (e.g., Buras et al. 2006a; Suwa et al. 2010; Bruenn et al. 2013). Note also from the comparison of the neutrino luminosities and rms energies that no clear differences between G2 (black line) and G1 (blue line) can be seen (top panels of Figure 19). The only exception is the reduction of the net heating rate in the gain region for G2 (Figure 17) compared to G1. The reduction of the net heating rate for G2 is more apparent in the accretion phase, namely before \(\sim 300\) ms pb. Note that the timescale (\(\sim 300\) ms pb) coincides with the sudden drop in the mass accretion rate (green line). Our 1D comparison between set2 and set1 (e.g., the bottom panel of Figure 5) suggests that the improved electron-capture rate on heavy nuclei lowers the explosability. Using the different progenitor models (note again the use of 20 \(M_{\odot}\) in 2D and 15 \(M_{\odot}\) in 1D), our results show that this feature still remains in 2D.

Three panels in the second and third rows of Figure 17 show more detailed comparisons between G3 and G1 (labeled with G3/G1), G4 and G3 (with G4/G3), and G5 and G4 (with G5/G4), respectively. Regarding the shock evolution (second rows), the average shock radii of G3 and G4 (black lines) show no remarkable difference from G1 (green line). This is in line with the lack of significant change in the net heating rate (top right) relative to G1. As expected in 1D, the inclusion of the nupair reaction (Section 3.3) makes the PNS radius\(^16\) more compact compared to G1 (see panel labeled G3/G1). The PNS radius of G4 somehow becomes slightly bigger between \(\sim 375–600\) ms compared to G3; however, it comes closer to G3 thereafter. From Figure 19, the neutrino luminosities and rms energies show no clear discrepancies between G3 and G4 compared to those already observed in the corresponding 1D models.

More big change can be seen in the shock evolution of G5 (black line; see panel with G5/G4 in Figure 17). The shock of G5 (black line) starts to expand \(\sim 214\) ms after bounce (see the hump in the shock evolution), maximally reaching \(\sim 163\) km, but returns to closely match the shock trajectory of G4 (green line) 400 ms pb (see also Figure 16). In fact, the rms energies of \(\nu_e\) and \(\bar{\nu}_e\) are higher for G5 than for G1 (see the panel with G5/G1 in Figure 19), and the net heating rate (top panel of Figure 18) becomes clearly larger for G5 (blue line) compared to G1 (green line) during the transient shock expansion phase. This is reconciled with the (slightly) higher heating rate due to the weak magnetism and recoil effect (as we saw in our 1D model, set5a). The higher heating rate is also fingerprinted in the diagnostic explosion energy (black line; see the panel with G5/G4 in Figure 17).

Using the same progenitor, LS220 EOS, and a similar set of neutrino opacities, our G5 run is close to model s20-2007 of Summa et al. (2016). Their s20-2007 model starts to explode \(\sim 200–300\) ms after bounce, whereas our G5 model does not. For a quantitative discussion, we choose to compare the neutrino luminosities and the rms energies 100 ms after bounce, which closely corresponds to the peak of the luminosities (e.g., their Figure 3). For the progenitor, the luminosity of \(\nu_e, \bar{\nu}_e\), and \(\nu_x\) is \(\sim 7.0, 6.5, \text{ and } 4.2 \times 10^{52} \text{ erg s}^{-1}\) for s20-2007, and 7.5, 7.5, and 4.0 \(\times 10^{52} \text{ erg s}^{-1}\) for G5, respectively. The rms energy of \(\nu_e, \bar{\nu}_e\), and \(\nu_x\) is \(\sim 12.2, 14.4, \text{ and } 16.2 \text{ MeV for s20-2007, and 11.0, 13.8, and } 14.9 \times 10^{52} \text{ erg s}^{-1}\) for G5, respectively. The lower \(\nu_e\) and \(\bar{\nu}_e\) rms energies would explain the more difficult explosion for G5

\(^{16}\) Note that the PNS radius is estimated at a fiducial density of \(10^{11} \text{ g cm}^{-3}\).
compared to s20-2007 of Summa et al. (2016). As already mentioned, our neglect of non-elastic effects in charged-current reactions, in the simplified transport schemes could explain such ∼10% level of discrepancies. We cannot unambiguously identify which of the missing improvements in this work could explain the above difference. This apparently shows that the implementation of detailed neutrino opacities and the accurate treatment of neutrino transport as well as GR are mandatory for quantitative studies of the CCSN mechanism.

G6ab shows a shock revival ∼500 ms after bounce (Figure 16), which we can also see from the clear deviation of the diagnostic energy from zero (see the panel labeled G6ab in Figure 17). On the other hand, G6ac does not explode during the simulation time (Figure 16). The neutrino luminosities and rms energies show fast time variations in the accretion phase (Figure 19); it is not easy to clearly see the increase for G6ab relative to G5.

In order to understand the above trend, we show the ratio of advection to heating timescale in the gain region (Buras et al. 2006b; the bottom-right panel in Figure 18). From the panel, one can see that the strangeness contribution (magenta line, G6ac) does enhance the chance of explosion at around 300 ms pb, which can be seen as peaks in the ratio ($t_{adv}/t_{heat}$) exceeding unity. In fact, the shock radius (the top-left panel)

Figure 15. Similar to Figure 5, but for the comparison between set-all and set1. Note that set-all includes all of the updates to set1 (from set2 to set6c in Table 1).
and the net heating rate in the gain region (the top-right panel) become slightly bigger for G6ac (magenta lines) compared to those of G6ab (black lines) at the same time.

The enhanced chance of explosion in models with the strangeness correction (G6ac and G6abc) originates from the reduction of the neutrino opacity, which is most significant in the accretion phase. This was already seen in our 1D runs of the 15 $M_\odot$ star (see the green line in the right panel of Figure 14 in the accretion phase). Our 2D results show that the inclusion of only the strangeness correction (relative to the standard rates) is not sufficient to trigger the onset of explosion for the 20 $M_\odot$ star. As one can see from the red line (G6abc, the bottom panel in Figure 18), our 2D run demonstrates that the combination of the strangeness and the many-body corrections makes the onset of the explosion easier.

As already mentioned in our 1D comparison (Section 3.6), the many-body correction of Horowitz et al. (2017) is expected to enhance the explodability primarily after the accretion phase because the many-body correction reduces the opacity at high densities. Our 2D results are in line with the 1D expectation. The net heating rate in the gain region is higher and the PNS radius is slightly smaller for G6ab compared to those for G5 and G1 in the post-accretion phase (e.g., top panel of Figure 18). These results demonstrate that the many-body correction mainly impacts the explodability in 2D after the accretion phase. Using the same EOS (LS220) and the same progenitor, model s20.0-LS220 in Bollig et al. (2017), which includes the many-body correction in addition to their standard neutrino opacities, leads to an explosion $\sim$400 ms pb (see the sky-blue line in their Figure 18), whereas the corresponding model which further includes the strangeness correction ($g_s' = -0.1$) leads to an earlier explosion at $\sim$200 ms pb (see the blue line in the figure). These features are basically consistent with our 2D runs.

Finally, G6abc (Figure 16) shows the earliest runaway shock expansion (starting $\sim$200 ms after bounce) among our 2D models. As already seen in set6c (Section 3.6), the strangeness effects contribute to the enhancement of the $\nu_e$ and $\bar{\nu}_e$
Figure 17. Summary of our 2D models (see the text).
luminosities before the accretion phase ends (∼300 ms after bounce). Quantitatively, the peak of the accretion luminosity of $\nu_e$ and $\bar{\nu}_e$ is enhanced by ∼3.8% and ∼1.3% for G6abc compared to G6ab (Figure 19). The slightly enhanced heating rate in the accretion phase due to the strangeness correction works synergetically with the many-body correction to revive the stalled shock into explosion for the 20 $M_\odot$ star (e.g., Figure 18). The diagnostic energy when the shock reaches 1000 km is 0.24 B and 0.2 B, with B representing “Bethe” = $10^{51}$ erg. To obtain the saturated value, long-term simulations are needed, which is beyond the scope of this work. Using the same progenitor and LS220 EOS, the onset time of an explosion (∼200 ms after bounce) is close to that seen in the 2D model of Bollig et al. (2017) with the strangeness contribution ($q_s' = -0.1$). However, the match may be simply coincidental because the contributions from muons are not yet included in this work.

5. Conclusions and Discussion

In this study, we have explored impact of updated neutrino opacities in CCSN simulations where spectral neutrino transport is solved by the three-flavor IDSA scheme. To verify our code, we first presented our 1D results following the core collapse, bounce, and up to ∼250 ms pb of a 15 $M_\odot$ star using the standard set of neutrino opacities by Bruenn (1985) and made a comparison with the seminal work by Liebendörfer et al. (2005). The good agreement from the code comparison supports the reliability of our three-flavor IDSA scheme with the standard opacity set. We then investigated in the 1D runs how the individual updated rate could lead to the difference between the baseline run and the standard opacity set. Through a detailed comparison with previous literature, we checked the validity of each implementation in a step-by-step manner. As previously noted, we have confirmed that adding up the individual rates impacts the neutrino luminosities and energies nonlinearly. In our 2D runs, we implemented selected sets of the neutrino opacities because a full investigation of the individual rates is currently too computationally expensive to do even in 2D with our improved IDSA scheme. Regarding the explodability, our results showed that several expectations from the individual updates in 1D are indeed correct in 2D. Among the updates considered in this work, the inclusion of both the strangeness-dependent contribution and the many-body correction to the neutrino–nucleon scattering has the largest impact on enhancing the explodability in our 2D models.

Using the same progenitor, the same EOS, and a similar set of neutrino opacities, there are ∼10% levels of discrepancies in the neutrino luminosities and the rms energies between our results and the results from the codes with more accurate neutrino transport schemes (e.g., Agile-BOLTZTRAN and Vertex). Our neglect of energy-bin/flavor coupling in the transport equation, non-isoeenergetic effects in the charged-current reactions, and the partial implementation of the Doppler-shift terms could solve the mismatch. Moreover, our approximate GR treatment in neutrino transport should be also improved. In this respect, the microphysics update we have done in this work is nothing but the first steps toward more sophisticated CCSN modeling.

We are thankful to M. Hempel for providing a table for calculating the nucleon potential difference ($\Delta U$) for the LS220 EOS. K.K. is wholeheartedly thankful to H.T. Janka, E. Müller, R. Bollig, A. Lohs, T. Foglizzo, T. Kuroda, E. Abdikamalov, and R. Kaceroni for stimulating discussions during his six-month stay at the Max Planck Institute for Astrophysics in 2017 that was supported by JSPS KAKENHI Grant Number JP15KK0173. T.F. acknowledges support from the Polish National Science Center (NCN) under grant number UMO-2016/23/B/ST2/00720. G.M.P. acknowledges partial support by the Deutsche Forschungsgemeinschaft through grant SFB 1245 (“Nuclei: From Fundamental Interactions to Structure and Stars”). Numerical computations were carried out in part on XC30 and general common use computer system at the center for Computational Astrophysics, CICA, the National Astronomical Observatory of Japan, and also on XC40 at YITP at Kyoto University. This study was also supported by JSPS KAKENHI grant numbers JP15H00789, JP15H01039, JP17H01130, JP17H06364, by the Central Research Institute of Fukuoka University (Nos. 171042, 1717103), and JICFuS as a priority issue to be tackled by using the Post “K” Computer.
Figure 19. Comparison of neutrino luminosities (upper panels) and rms neutrino energies (lower panels) for all 2D models.
Appendix A
Implementing Electron–Neutrino Pair Annihilation

A.1. $\nu_e + \bar{\nu}_e \rightarrow \nu_x + \bar{\nu}_x$: Evolution Equation of $\nu_x$

Following Buras et al. (2003), the scattering kernels of electron–neutrino pair annihilation can be calculated essentially in the same way as that for electron–positron annihilation, $e^- + e^+ \rightarrow \nu_x + \bar{\nu}_x$. It is convenient to define the $\nu_x$-pair production kernels labelled (p):

$$ \mathcal{R}^{p}_{\nu_x, \nu_{x'}}(\cos \theta_{\nu_x, \nu_{x'}}, E_{\nu_x} + E_{\nu_{x'}}) = \int \frac{d^3p_{\nu_x}}{(2\pi \hbar)^3} \frac{d^3p_{\nu_{x'}}}{(2\pi \hbar)^3} \times 2f_{\nu_x}(p_{\nu_x})f_{\nu_{x'}}(p_{\nu_{x'}})|M|^2 \delta^4(p_{\nu_x} + p_{\nu_{x'}} - p_{\nu_x} - p_{\nu_{x'}}), $$

with the initial-particle’s distribution functions $f_{\nu_x}$ and $f_{\nu_{x'}}$, for which we assume local thermodynamic equilibrium, i.e., $\mu_{\nu_x} = \mu_{\nu_{x'}} = (\mu_{\nu_x} - \mu_{\nu_{x'}})$. The pair-production kernel (22) depends on the incident scattering angle $\theta_{\nu_x, \nu_{x'}}$ between $\nu_x$ and $\bar{\nu}_x$ (for the definition, see Mezzacappa & Bruenn 1993b) as well as on the sum of the $\nu_x$ and $\bar{\nu}_x$ energies, $E_{\nu_x}$ and $E_{\bar{\nu}_x}$, respectively. Moreover, the spin-averaged and squared matrix element, $|M|^2$, is obtained from $e^- - e^+ \rightarrow \nu_x - \nu_{x'}$ (cf. Bruenn 1985) with the following replacements for the weak coupling constants, $\Gamma_v = \Gamma_0 = + 1/2$. Since within the IDSA no explicit angle dependence of weak processes is employed, we perform a Legendre expansion of the pair-production kernel (22) in terms of $\cos \theta$,

$$ \mathcal{R}^{p}_{\nu_x, \nu_{x'}}(\cos \theta_{\nu_x, \nu_{x'}}, E_{\nu_x} + E_{\nu_{x'}}) \rightarrow \frac{1}{2} \Phi^{p}_{0, \nu_x, \nu_{x'}}(E_{\nu_x} + E_{\nu_{x'}}) + \mathcal{O}(\cos \theta), $$

such that the corresponding collision term for the the zeroth component of the distribution function, $f^{(0)}_{\nu_x}$, reads as follows:

$$ \left. \frac{\partial f^{(0)}_{\nu_x}(E_{\nu_x})}{\partial t} \right|_{\text{coll}} = \frac{2\pi}{c(2\pi \hbar)^3} (1 - f^{(0)}_{\nu_x}(E_{\nu_x})) \times \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} (1 - f^{(0)}_{\bar{\nu}_x}(E_{\bar{\nu}_x})) \Phi^{p}_{0, \nu_x, \nu_{x'}}(E_{\nu_x} + E_{\bar{\nu}_x}) $$

$$ - \frac{2\pi}{c(2\pi \hbar)^3} f^{(0)}_{\nu_x}(E_{\nu_x}) \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} f^{(0)}_{\nu_{x'}} $$

$$ \times \Phi^{q}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}), $$

where $\Phi^{q}_{0, \nu_{x'}, \nu_x}$ denotes the zeroth-order Legendre coefficient of the $\nu_{x'}$ pair absorption kernel. It is related to $\Phi^{p}_{0, \nu_x, \nu_{x'}}$ via the relation of detailed balance,

$$ \Phi^{q}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}) = \exp \left\{ \frac{E_{\nu_{x'}} + E_{\bar{\nu}_x}}{T} \right\} \times \Phi^{p}_{0, \nu_x, \nu_{x'}}(E_{\nu_x} + E_{\bar{\nu}_x}), $$

which is realized straightforwardly in a similar fashion to Equation (22). Now, Equation (11) can be rewritten as follows,

$$ \left. \frac{\partial f^{(0)}_{\nu_x}(E_{\nu_x})}{\partial t} \right|_{\text{coll}} = C^{0}_{\text{nopair}}(E_{\nu_x}) + A^{0}_{\text{nopair}}(E_{\nu_x}) f^{(0)}_{\nu_x}(E_{\nu_x}), $$

with

$$ C^{0}_{\text{nopair}}(E_{\nu_x}) = \frac{2\pi}{c(2\pi \hbar)^3} \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} $$

$$ \times (1 - f^{(0)}_{\bar{\nu}_x}(E_{\bar{\nu}_x})) \Phi^{p}_{0, \nu_x, \nu_{x'}}(E_{\nu_x} + E_{\bar{\nu}_x}), $$

$$ A^{0}_{\text{nopair}}(E_{\nu_x}) = - \frac{2\pi}{c(2\pi \hbar)^3} \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} $$

$$ \times [(1 - f^{(0)}_{\bar{\nu}_x}(E_{\bar{\nu}_x})) \Phi^{p}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}) ] $$

$$ + f^{(0)}_{\nu_x}(E_{\nu_x}) \Phi^{q}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}), $$

which denote the production (Equation (14)) and annihilation rates (Equation (15)) of this process. In IDSA, the terms $C^{0}_{\text{nopair}}$ and $A^{0}_{\text{nopair}}$ are added to Equations (5) and (6) for the streaming neutrinos and to Equation (15) of Liebendörfer et al. (2009) for the trapped neutrinos.

A.2. $\nu_x + \bar{\nu}_x \rightarrow \nu_x + \bar{\nu}_x$: Evolution Equations of $\nu_x$ and $\bar{\nu}_x$

The collision term for $f^{(0)}_{\nu_x}$ associated with this process is obtained in a similar fashion to the production of $\nu_x$ pairs (Appendix A.1),

$$ \left. \frac{\partial f^{(0)}_{\nu_x}(E_{\nu_x})}{\partial t} \right|_{\text{coll}} = \frac{2\pi}{c(2\pi \hbar)^3} (1 - f^{(0)}_{\nu_x}(E_{\nu_x})) \times \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} (1 - f^{(0)}_{\bar{\nu}_x}(E_{\bar{\nu}_x})) \Phi^{p}_{0, \nu_x, \nu_{x'}}(E_{\nu_x} + E_{\bar{\nu}_x}) $$

$$ - \frac{2\pi}{c(2\pi \hbar)^3} f^{(0)}_{\nu_x}(E_{\nu_x}) \int E_{\bar{\nu}_x}^2 dE_{\bar{\nu}_x} f^{(0)}_{\nu_{x'}} $$

$$ \times \Phi^{q}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}). $$

Note that the expression for the collision integral for $\bar{\nu}_x$ is obtained equivalently by replacing the labels $\nu_x \leftrightarrow \bar{\nu}_x$ in the above expression. Moreover, due to the following symmetry considerations,

$$ \Phi^{q}_{0, \nu_{x'}, \nu_x}(E_{\nu_{x'}} + E_{\bar{\nu}_x}) = \Phi^{p}_{0, \nu_x, \nu_{x'}} $$

expression (16) can be reduced to a simple form similarly to expression (13), with equivalent definitions for $C^0$ and $A^0$. Due to the relation of detailed balance (24), which holds here as well, it becomes clear that it is necessary to obtain only one pair-reaction kernel, e.g., $\Phi^{p}_{0, \nu_x, \nu_{x'}}$. Therefore, we follow Equations (C62)–(C74) in Bruenn (1985) for the computation presented in this work. Note that here we also assume that $\nu_x$ obeys local thermodynamic equilibrium, i.e., $\mu_{\nu_x} = 0$ (this was also assumed in Buras et al. 2003; Fischer et al. 2009). For practical reasons, we monitor the $\nu_x$ distribution function, which must not differ by more than 10% from the corresponding Fermi–Dirac distribution. This treatment was tested to work well in Agile–BOLTZTRAN simulations by Fischer et al. (2009).

A.3. Integrated Pair Production Rates

Here we provide definitions of the quantities shown in the main part of this paper. Therefore, the $(\nu_x, \bar{\nu}_x)$ pair-production

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rate is defined as follows:

\[
Q_{\nu_i,\bar{\nu}_i}(E_{\nu_i}) = \frac{1}{(2\pi\hbar)^3} \int E_{\nu_i}^2 dE_{\bar{\nu}_i} (1 - f^{(0)}_{\nu_i}(E_{\nu_i})) \\
\times 2\pi \Phi_{\nu_i,\bar{\nu}_i}^0(E_{\nu_i} + E_{\bar{\nu}_i}) \quad [\text{s}^{-1}],
\]

(18)
such that the total number production rate of \(\nu_i\) is given as follows:

\[
\frac{\partial n_{\nu_i}}{\partial t} = \frac{1}{(2\pi\hbar)^3} n_{\nu_i} \int E_{\nu_i}^2 dE_{\nu_i} \\
\times (1 - f^{(0)}_{\nu_i}(E_{\nu_i})) 4\pi Q_{\nu_i,\bar{\nu}_i}(E_{\nu_i}),
\]

(19)
which is normalized to the number density of neutrinos, \(n_{\nu_i}\). Then, we obtain the number production spectra,

\[
\frac{\partial^2 n_{\nu_i}}{\partial E_{\nu_i} \partial t} = \frac{1}{(2\pi\hbar)^3} E_{\nu_i}^2 (1 - f^{(0)}_{\nu_i}(E_{\nu_i})) \\
\times 4\pi Q_{\nu_i,\bar{\nu}_i}(E_{\nu_i}) \quad [\text{s}^{-1}\text{MeV}^{-1}\text{cm}^{-3}],
\]

(20)
where \(\nu_i \leftrightarrow \bar{\nu}_i\) in the above expressions. For the calculations of the scattering kernels, we follow Equation (C50) in Bruenn (1985) for the computation presented in this work. Then, the inverse mean free path \(1/\lambda_{\text{mfp}}\) shown in Figure 8 is obtained as follows,

\[
A^0_{\text{NNS}, \nu_i}(E_{\nu_i}) = -\frac{2\pi}{c(2\pi\hbar)^3} \int E_{\nu_i}^2 dE_{\nu_i} \\
\times (f^{(0)}_{\nu_i}(E_{\nu_i}) \Phi_{\nu_i,\bar{\nu}_i}^0(E_{\nu_i} - E_{\bar{\nu}_i})) \\
+ (1 - f^{(0)}_{\nu_i}(E_{\nu_i})) \Phi_{\nu_i,\bar{\nu}_i}^0(E_{\nu_i} - E_{\bar{\nu}_i}),
\]

(25)
where \(A^0_{\text{NNS}, \nu_i}\) is in units of \([\text{s}^{-1}]\), which also holds for the other processes (cf., Equation (139) of Kuroda et al. 2016).

Appendix B

To calculate the scattering kernel, we follow the same procedure as for neutrino–electron/positron scattering (Buras et al. 2003),

\[
R_{\nu_i}^{\text{in}}(\cos \theta_{\nu_i,\nu_i'}, E_{\nu_i} - E_{\nu_i'}) = \int \frac{d^3\nu}{(2\pi\hbar)^3} \frac{d^3\nu'}{(2\pi\hbar)^3} 2f_{\nu_i}(p_{\nu_i}) \\
\times (1 - f_{\nu_i}(p_{\nu_i})) |M|^2 \delta^4(p_{\nu_i} + p_{\nu_i'} - p_{\nu_i'}),
\]

(22)
which is evaluated under the assumption of local thermodynamic equilibrium, exactly as for the neutrino-pair processes in Appendix A. Here, the spin-averaged and squared matrix element, \(|M|^2\), is obtained from neutrino-electron scattering (cf., Bruenn 1985) with the following replacements for the weak coupling constants, \(C_6 = C_{6a} = +1/2\) for \(\nu_e\) scattering on \(\nu_e\), and \(C_7 - 1\) and \(C_{6a} - 1\) for \(\nu_e\) scattering on \(\bar{\nu}_e\). With the Legendre expansion of the scattering kernel in terms of \(\cos \theta_{\nu_i,\nu_i'}\), we obtain the zeroth-order term of the collision integral for \(f^{(0)}_{\nu_i}\) as follows,

\[
\frac{\partial f^{(0)}_{\nu_i}(E_{\nu_i})}{\partial t}_{\text{coll}} = \frac{2\pi}{c(2\pi\hbar)^3} (1 - f^{(0)}_{\nu_i}(E_{\nu_i})) \\
\times \int E_{\nu_i}^2 dE_{\nu_i'} f^{(0)}_{\nu_i}(E_{\nu_i'}) \Phi_{\nu_i,\nu_i'}^{0\text{in}}(E_{\nu_i} - E_{\nu_i'}) \\
- \frac{2\pi}{c(2\pi\hbar)^3} f^{(0)}_{\nu_i}(E_{\nu_i}) \int E_{\nu_i}^2 dE_{\nu_i'} \\
\times (1 - f^{(0)}_{\nu_i}(E_{\nu_i'}) \Phi_{\nu_i,\nu_i'}^{0\text{out}}(E_{\nu_i} - E_{\nu_i'}),
\]

(23)
where the out-scattering kernel, \(\Phi_{\nu_i,\nu_i'}^{0\text{out}}\), is obtained via the relation of detailed balance,

\[
\Phi_{\nu_i,\nu_i'}^{0\text{out}}(E_{\nu_i} - E_{\nu_i'}) = \exp \left\{ \frac{E_{\nu_i} - E_{\nu_i'}}{T} \right\} \\
\times \Phi_{\nu_i,\nu_i'}^{0\text{in}}(E_{\nu_i} - E_{\nu_i'}).
\]

(24)

Similarly to Equations (14) and (15), one can obtain \(C^0_{\text{NNS}}\) and \(A^0_{\text{NNS}}\) for the NNS processes here. These are added to the evolution equation of the streaming and trapped neutrinos accordingly. Note that the expression for \(\nu_i\) scattering on \(\bar{\nu}_i\) is obtained equivalently by replacing the labels \(\nu_i \leftrightarrow \bar{\nu}_i\) in the above expressions.
