Stellar Collapse Diversity and the Diffuse Supernova Neutrino Background

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Abstract

The diffuse cosmic supernova neutrino background (DSNB) is an observational target of the gadolinium-loaded Super-Kamiokande (SK) detector and the forthcoming JUNO and Hyper-Kamiokande detectors. Current predictions are hampered by our still incomplete understanding of the supernova (SN) explosion mechanism and of the neutron star (NS) equation of state and maximum mass. In our comprehensive study we revisit this problem on grounds of the landscapes of successful and failed SN explosions obtained by Sukhbold et al. and Ertl et al. with parameterized one-dimensional neutrino engines for large sets of single-star and helium-star progenitors, with the latter serving as a proxy for binary evolution effects. Besides considering engines of different strengths, leading to different fractions of failed SNe with black hole (BH) formation, we also vary the NS mass limit and the spectral shape of the neutrino emission and include contributions from poorly understood alternative NS formation channels, such as accretion-induced and merger-induced collapse events. Since the neutrino signals of our large model sets are approximate, we calibrate the associated degrees of freedom by using state-of-the-art simulations of proto-NS cooling. Our predictions are higher than other recent ones because of a large fraction of failed SNe with long delay to BH formation. Our best-guess model predicts a DSNB $\dot{\nu}_e$-flux of $28.8^{+6.6}_{-10.6} \text{ cm}^{-2} \text{s}^{-1}$ with $6.0^{+3.2}_{-2.1} \text{ cm}^{-2} \text{s}^{-1}$ in the favorable measurement interval of $[10, 30] \text{ MeV}$ and $1.3^{+1.1}_{-0.4} \text{ cm}^{-2} \text{s}^{-1}$ with $\dot{\nu}_e$ energies $> 17.3 \text{ MeV}$, which is roughly a factor of two below the current SK limit. The uncertainty range is dominated by the still insufficiently constrained cosmic rate of stellar core-collapse events.

Unified Astronomy Thesaurus concepts: Diffuse radiation (383); Supernova neutrinos (1666); Core-collapse supernovae (304); Massive stars (732); Black holes (162); Neutron stars (108)

1. Introduction

When the life of a massive star (above ~9 $M_\odot$) ends with the collapse of the inner core to a neutron star (NS) or a black hole (BH), a tremendous amount of gravitational binding energy (BE; several $10^{53}$ erg) is released, predominantly in the form of neutrinos and antineutrinos (see, e.g., Janka 2012, 2017; Burrows 2013). In 1987, when the blue supergiant Sanduleak $-69^\circ$ 202 (Walborn et al. 1987) in the Large Magellanic Cloud exploded as SN 1987A, such an associated neutrino burst was detected for the first (and so far only) time as a $\sim 10 \text{ s}$ long signal, albeit with a sparse yield of only two dozen counts (Bionta et al. 1987; Hirata et al. 1987; Alexeyev et al. 1988). Today, the size of neutrino observatories all over the world has grown significantly such that a galactic supernova (SN) would lead to a high-statistics signal (e.g., Ikeda et al. 2007; Abbasi et al. 2011), which the scientific community is eagerly waiting for.

While such a nearby SN is a rare event (Diehl et al. 2006; Ikeda et al. 2007; Agafonova et al. 2015), a vast number of massive stars have already ended their lives in cosmic history, generously radiating neutrinos. The integral flux from all those past core collapses at cosmological distances, which is steadily flooding Earth, constitutes the so-called diffuse supernova neutrino background (DSNB). It makes for a “guaranteed” (isotropic and stationary) signal of megaelectronvolt neutrinos, comprising rich information on the entire population of stellar core collapses (for dedicated reviews, see Ando & Sato 2004; Beacom 2010; Lunardini 2016; Vitagliano et al. 2020). Intriguingly, the Super-Kamiokande (SK) experiment set upper flux limits on the DSNB (Malek et al. 2003; Bays et al. 2012; Zhang et al. 2015), which are already close to theoretical predictions. This indicates the excellent discovery prospect within the next decade in the gadolinium-loaded SK detector and the forthcoming JUNO experiment (see, e.g., Beacom & Vagins 2004; Yüksel et al. 2006; Horiiuchi et al. 2009; An et al. 2016; Priya & Lunardini 2017; Möller et al. 2018) as well as in the longer term with the Hyper-Kamiokande detector (Abe et al. 2011), with DUNE (DUNE Collaboration et al. 2015), or with the proposed THEIA detector (Askins et al. 2020; Sawatzki et al. 2021).

To exploit the full potential of future observations, comprehensive theoretical models will be needed for comparison. The first predictions of the DSNB date back to the 1980s and 1990s (e.g., Bisnovatyi-Kogan & Seidov 1982; Krauss et al. 1984; Hartmann & Woosley 1997) and have been refined since. Its link to the cosmic history of star formation has been studied in detail (e.g., Ando 2004; Strigari et al. 2005; Hopkins & Beacom 2006; Mathews et al. 2014; Anandagoda et al. 2020; Riya & Rentala 2020), and the dependence on the SN source spectra, which is the focus of this paper, has been the subject of intense research. For instance, Lunardini (2007) took an analytical approach based on the work by Keil et al. (2003), while Lunardini (2006) and Yüksel & Beacom (2007) employed constraints from the measured neutrinos from SN 1987A for their DSNB predictions. The impact of the SN shock revival time has been investigated (Nakazato 2013; Nakazato et al. 2015), as has the effect of neutrino flavor conversions (Ando & Sato 2003; Chakraborty et al. 2011; Lunardini & Tamborra 2012).

In particular, the contribution from BH-forming, failed explosions to the DSNB has caught much attention in recent years. It might significantly enhance the high-energy tail of the flux spectrum, which is most relevant for detection (e.g.,
Lunardini 2009). Several studies have varied the (still unknown) fraction of failed SNe (Lunardini 2009; Lien et al. 2010; Keehn & Lunardini 2012; Priya & Lunardini 2017; Horiiuchi et al. 2018; Moller et al. 2018); in this regard, Nakazato et al. (2015) and Yüksel & Kistler (2015) further considered the cosmic evolution of stellar metallicities, and the dependence on the high-density equation of state (EoS), which is closely related to the mass limit up to which an NS can be stabilized against its own gravity, has been explored tentatively (Lunardini 2009; Keehn & Lunardini 2012; Mathews et al. 2014; Nakazato et al. 2015; Hidaka et al. 2016, 2018; Horiiuchi et al. 2018).

Detailed neutrino signals from successful and failed SNe are the premise for reliable DSNB predictions. While most previous works have employed rather approximate neutrino source spectra or spectra representative of some typical cases, numerical modeling of stellar core collapse has reached a high level of sophistication nowadays. An increasing number of 3D simulations with detailed microphysics have become available (e.g., Tak wiki et al. 2014; Tamborra et al. 2014; Lentz et al. 2015; Melson et al. 2015, 2020; Müller et al. 2017; O’Connor & Couch 2018; Ott et al. 2018; Summa et al. 2018; Burrows et al. 2019; Glas et al. 2019; Vartanyan et al. 2019). However, the high computational costs are still causing limitations. Up to now, only about 20 selected progenitors have been considered in 3D SN models; none of them have evolved longer than roughly 1 s.

At the same time, it has been shown that the outcome of a core-collapse event (successful explosion or BH formation) as well as the neutrino emission strongly depends on the progenitor structure, with large variations between different stars (O’Connor & Ott 2011; Ugliano et al. 2012; Horiiuchi et al. 2014; Nakamura et al. 2015; Pejcha & Thompson 2015; Erlt et al. 2016; Müller et al. 2016; Sukhbold et al. 2016; Ebinger et al. 2019). This has been neglected (or over-simplified) in most previous DSNB studies, which have typically employed only a few exemplary models. In particular, the signals from BH-forming, failed SNe are strongly dependent on the progenitor-specific mass-accretion rate (Fischer et al. 2009; O’Connor & Ott 2011). Comprehensive sets of neutrino signals over the entire range of pre-SN stars are therefore required to adequately account for the diversity of stellar core collapse. In light of this, Horiiuchi et al. (2018) employed a set of 101 axisymmetric (2D) SN simulations and seven models of BH formation from spherically symmetric (1D) simulations, although they needed to extrapolate the neutrino signals at times later than ~1 s. Due to the limited number of their failed explosions, they (linearly) interpolated the spectral parameters of the time-integrated neutrino emission (total energetics, mean energy, and shape parameter) of their few BH simulations as a function of the “progenitor compactness” (O’Connor & Ott 2011) to account for a larger scope of failed SNe.

In this paper, we take a different angle of approach. Referring to studies by Ugliano et al. (2012), Sukhbold et al. (2016), and Erlt et al. (2016, 2020), we use spherically symmetric simulations over a wide range of pre-SN stars that are exploded by means of a “calibrated central neutrino engine.” In this way, our analysis of the DSNB is based on detailed information about the “landscape” of successful and failed explosions with individual neutrino signals for every progenitor, including cases of long-lasting mass accretion with relatively late BH formation. Using our large sets of (approximately calculated) long-time neutrino signals, which we cross-check by comparing and normalizing them to the outcome of more sophisticated simulations (see appendices), we aim at providing refined predictions of the DSNB. In a systematic parameter study, we further investigate the impact of three critical source properties on the DSNB flux spectrum: (1) We vary the fractions of successful and failed SNe through different calibrations of the neutrino engine used for the explosion modeling of our large progenitor set. (2) We consider different values for the critical mass at which the neutrino signals stop due to BH formation and follow the continued mass accretion of failed explosions. (3) We consider different spectral shapes of the neutrino emission based on a study by Keil et al. (2003).

As in previous DSNB studies (e.g., Mathews et al. 2014; Horiiuchi et al. 2018), we also include the contribution from electron-capture SNe (ECSNe) of degenerate oxygen–neon–magnesium (ONeMg) cores (Miyaji et al. 1980; Nomoto 1984, 1987), for which we employ the neutrino signals from Hudepohl et al. (2010). Moreover, we explore other possible channels for the formation of low-mass (LM) NSs, such as accretion-induced collapse (AIC; Bailyn & Grindlay 1990; Nomoto & Kondo 1991; Ivanova & Taam 2004; Hurley et al. 2010; Jones et al. 2016; Wu & Wang 2018; Ruiter et al. 2019) and merger-induced collapse (MIC; Saio & Nomoto 1985; Ivanova et al. 2008; Schwab et al. 2016; Kashyap et al. 2018; Ruiter et al. 2019) of white dwarfs (WDs) or ultrastripped SNe from close binaries (Nomoto et al. 1994; Dewi et al. 2002; Tauris et al. 2013, 2015; Suwa et al. 2015; Müller et al. 2018). Using simplified assumptions, we estimate the flux from such a combined “LM component” and comment on its relevance.

While stellar explosion models have typically employed single-star progenitors thus far, recent observations suggest that most massive stars are in binary systems (see, e.g., Mason et al. 2009; Sana et al. 2012). In view of this we also investigate, for the first time, how the inclusion of binary models affects predictions of the DSNB using the helium-star progenitors from Woosley (2019) and the explosion models of Erlt et al. (2020).

The paper is organized as follows. In Section 2, we describe the setup of our simulations and discuss the overall properties of the neutrino signals used in our study. Section 3 is dedicated to our approach of formulating the DSNB. We present our fiducial predictions in Section 4. In Section 5, we discuss the results of our detailed parameter study: We investigate the sensitivity of the DSNB flux spectrum of electron antineutrinos to the fraction of failed explosions, the BH mass threshold, and the spectral shape of the neutrino emission. We further explore an additional contribution from LM NS–forming events (such as AIC, MIC, and ultrastripped SNe) and study the impact of including binary progenitors. In Section 6, we briefly comment on the DSNB flux spectrum of electron neutrinos. In Section 7, the effects of neutrino flavor conversions are discussed along with remaining uncertainties, followed by a comparison of our models with the SK flux limits and with previous works (Section 8). In Section 9, we categorize and rank the DSNB parameter variations and uncertainties considered in our work. We conclude in Section 10. Supplementary material can be found in the appendices.
2. Simulation Setup and Neutrino Signals

In spherical symmetry, self-consistent SN explosions have turned out to be possible only for a few low-mass stars (Kitaura et al. 2006; Janka et al. 2008, 2012; Fischer et al. 2010; Melson et al. 2015; Radice et al. 2017). To still explore the outcome of stellar core collapse in 1D over a wide range of progenitor masses, we adopt the parametric approach of Ertl et al. (2016), where a “calibrated neutrino engine” is placed in the center of all pre-SN models. By these means, we obtain neutrino signals for a large set of individual stars, in satisfactory agreement with more sophisticated simulations and including cases of long-term accretion with late BH formation, as we will elaborate in this section. For more details on our computational setup, the reader is referred to Ugliano et al. (2012), Sukhbold et al. (2016), and Ertl et al. (2016, 2020).

2.1. Pre-SN Models

In this work, we use a combined set of 200 solar-metallicity progenitor models from Woosley & Heger (2007, 2015, “WH07” and “WH15”) and Sukhbold & Woosley (2014, “SW14”), which was already applied in Sukhbold et al. (2016) and can be downloaded from the Garching Core-collapse Supernova Archive. All models are nonrotating single stars, evolved with the KEPLER code (Weaver et al. 1978) up to the onset of iron-core collapse. The resulting grid of progenitors, unevenly distributed over the zero-age main-sequence (ZAMS) mass interval of 9–120 $M_\odot$, spans the commonly assumed range of “conventional” iron-core collapse SNe (or BH-forming, failed SNe).

Below that, in the narrow band between 8.7 $M_\odot$ and 9 $M_\odot$, we additionally consider ECSNe of degenerate ONeMg cores as another channel for NS formation (Miyaji et al. 1980; Nomoto 1984, 1987); yet it should be stressed that the exact mass range of ECSNe in the local universe is not definitive according to current knowledge (see, e.g., Poelarends et al. 2008; Jones et al. 2013, 2016; Doherty et al. 2015; Kirsebom et al. 2019; Zha et al. 2019; Leung et al. 2020). We employ a simulation by Hüdepohl et al. (2010, “model SF”) for the neutrino signal of such core-collapse events. The upper-mass end of the ZAMS mass grid is similarly uncertain and depends strongly on the physics of mass loss. However, as will be detailed in Section 3.2, high-mass contributions are suppressed by the steeply declining initial mass function (IMF) and are therefore of subordinate importance for the DSNB. In Sections 5.2 and 5.3, we will further consider progenitors from binary systems. The helium-star models used in this context were published by Woosley (2019), and their explosions were investigated by Ertl et al. (2020).

2.2. SN Simulations

Our stellar collapse and explosion simulations are performed with the PROMETHEUS-HOTB code (Janka & Müller 1996; Kifonidis et al. 2003; Scheck et al. 2006; Ertl et al. 2016, 2020). The innermost 1.1 $M_\odot$ of the nascent proto-NS (PNS) is excised and replaced by a contracting inner grid boundary and an analytic one-zone core-cooling model with tunable parameters (for details, see Ugliano et al. 2012). This “central neutrino engine” is calibrated to yield explosions in agreement with the well-studied cases of SN 1987A and the Crab SN (SN 1054). More specifically, for pre-SN stars with ZAMS masses above 12 $M_\odot$, which Sukhbold et al. (2016) termed “87A-like,” a PNS core model is applied and adjusted such that a given progenitor in the range of 15–20 $M_\odot$, namely S19.8, N20, W18, W15, or W20 (as described in Sukhbold et al. 2016), reproduces the observed explosion energy (1.2–1.5) $\times 10^{51}$ erg; Arnett et al. 1989; Utrobin et al. 2015), 56Ni yield (∼0.07 $M_\odot$; Bouchet et al. 1991; Sunzettef et al. 1992), and basic neutrino emission features (Bionta et al. 1987; Hirata et al. 1987) of SN 1987A. The low-mass end (9–12 $M_\odot$) is connected to the 87A-like cases by an interpolation of the core model parameters. As a second anchor point, we use the progenitor z9.6 by A. Heger (2012, private communication), which explodes with low energy (∼1050 erg; Janka et al. 2012; Melson et al. 2015) and a small 56Ni yield (∼0.0025 $M_\odot$; Wannojo et al. 2018) in self-consistent simulations, in good agreement with the observational constraints for the Crab SN (Smith 2013; Tomina et al. 2013; Yang & Chevalier 2015). For more details on our calibration procedure, the reader is referred to Sukhbold et al. (2016).

Depending on the engine model, we obtain more or less energetic or failed explosions over the range of considered pre-SN stars, as can be seen in Figure 1. While the S19.8 and N20 calibrations lead to the largest fraction of successful SNe (red), W20 is a rather weak engine, resulting in the largest fraction of BH-forming cases (black). W18 and W15 reside between these two extremes, as can also be seen in Table 1, which shows the IMF-weighted fractions of successful and failed explosions for the different neutrino engines. The outcome in the low-mass range (9–12 $M_\odot$) is the same for all five cases, since our interpolation toward z9.6 is independent of the high-mass calibration. The nonmonotonic pattern of successful SNe and BH-forming collapses in Figure 1 has been described in previous works (Ugliano et al. 2012; Pejcha & Thompson 2015; Ertl et al. 2016; Müller et al. 2016; Sukhbold et al. 2016;

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*https://wwwmpa.mpa-garching.mpg.de/csetarchive/data/SEWBJ_2015/index.html (http://doi.org/10.17617/1.b).
Ebinger et al. 2019). It is based on the progenitor structure, which strongly varies with the ZAMS mass (O’Connor & Ott 2011; Horiiuchi et al. 2014; Nakamura et al. 2015; Sukhbold et al. 2018).

Compared to the simulations of Ertl et al. (2016) and Sukhbold et al. (2016), the neutrino transport outside of the PNS core, which is treated by a gray approximation (Scheck et al. 2006; Arcones et al. 2007), is slightly improved such that we are able to follow cases of long-lasting mass accretion until late collapse to a BH. For numerical reasons, the neutrino–nucleon scattering rate (Equation (D.68) of Scheck et al. 2006) is now split into two separate source terms, one for absorption ($\propto (\epsilon_n^2)$, with $\epsilon_n$ denoting the neutrino energy) and one for emission ($\propto T (\epsilon_n^2)$), to avoid sign fluctuations for high temperatures $T$ (for details, see Appendix B of Stockinger et al. 2020). Furthermore, an adaptive grid is implemented to better resolve the steep density gradient at the PNS surface. Our new code is applied without recalibrating the core models, which leads to slightly increased explosion energies because of decreased neutrino luminosities as compared to the models reported by Sukhbold et al. (2016). Accordingly, a few scattered progenitors that fail to explode with the old code (see Sukhbold et al. 2016, Figure 13) yield successful SNe with our new treatment. (A detailed report of the code changes and the consequences for the model results is provided in the appendices of Ertl et al. 2020.) In the work at hand, we moreover neglect the neutrino emission from the late-time fallback in so-called fallback SNe, in which the fallback matter pushes the NS beyond the BH limit after a successful explosion is initiated. This is justified because such cases turn out to be rare in the considered set of solar-metallicity progenitors (Ertl et al. 2016; Sukhbold et al. 2016) and additionally reside in the IMF-suppressed high-mass regime. In the context of our paper we therefore consider fallback SNe as successful SN events with the corresponding neutrino emission from NS formation. BH-forming events are only those cases where the BH does not form by fallback but by continuous accretion, and we use the terms “BH formation” and “failed SN” equivalently.

### 2.3. Neutrino Signals

For each progenitor, we obtain the total energy release in neutrinos through time integration of the time-dependent neutrino luminosities, $L_{\nu}(t)$, and the mean values of the energies of the radiated neutrinos by computing the (luminosity-weighted) time average of the time-dependent mean neutrino energies, $\langle E_{\nu}(t) \rangle$, for all three considered neutrino species $\nu_e = e_\nu, \bar{\nu}_e, \nu_x$, where $\nu_x$ denotes a representative heavy-lepton neutrino ($\nu_\mu, \nu_\tau, \nu_\nu, \bar{\nu}_\nu$). Successful SNe are simulated up to a post-bounce time of $t = 15$ s, when the neutrino luminosities from PNS cooling have already declined to an insignificant level (see Appendix A). In the cases of failed explosions, however, the continued accretion of infalling mass onto the PNS releases gravitational BE, leading to an ongoing accretion component of the neutrino luminosities. The signals of such cases are truncated only when the PNS is pushed beyond the (still unknown) limit of BH formation, for which we consider four different values of the baryonic mass, $M_{\text{B1}}$, namely 2.3, 2.7, 3.1, and 3.5 $M_\odot$, which are motivated as follows.

Assuming an NS radius of $(11 \pm 1)$ km for maximum-mass NSs and utilizing Equation (36) of Lattimer & Prakash (2001), which we provide as Equation (B1) in Appendix B, a baryonic NS mass of $2.3 M_\odot$ converts to a gravitating mass of $1.95^{+0.02}_{-0.03} M_\odot$. This is marginally below the largest currently measured pulsar masses of $\sim 2 M_\odot$ (Demorest et al. 2010; Antoniadis et al. 2013; Özel & Freire 2016; Cromartie et al. 2020), setting a lower limit for the maximum NS mass.

From the first gravitational-wave observation of a binary NS merger (GW170817; Abbott et al. 2017a) and its electromagnetic counterparts (Abbott et al. 2017b), Margalit & Metzger (2017) placed a tentative upper bound on the maximum gravitational NS mass of $2.17 M_\odot$ (at a 90% confidence level), which follows from their reasoning that the merger remnant was a relatively short-lived, differentially rotating hypermassive NS, disfavoring both prompt collapse to a BH and the formation of a long-lived, supermassive NS. Their mass limit is compatible with other recent publications (e.g., Shibata et al. 2017;Alsing et al. 2018; Rezzolla et al. 2018; Ruiz et al. 2018; Lim & Holt 2019;Essick et al. 2020). Consistently, we take our case of a baryonic mass of $2.7 M_\odot$ (corresponding to a $2.23^{+0.03}_{-0.04} M_\odot$ gravitational mass), which is close to this bound, as our reference threshold for BH formation.

Nonetheless, Margalit & Metzger (2017) pointed out several uncertainties related to their analysis. For instance, they neglected the effects of thermal pressure support on the stability of the compact merger remnant, which may change their conclusions.\(^5\) Thermal effects are also important for the stability of hot PNSs on their way toward BH formation in the cases of failed SNe, possibly increasing the limiting mass as compared to the value for cold NSs (O’Connor & Ott 2011; Steiner et al. 2013; da Silva Schneider et al. 2020). For these reasons, we additionally explore two more extreme cases for the baryonic (gravitational) mass limit, namely $3.1 M_\odot$ ($2.50^{+0.04}_{-0.05} M_\odot$) and $3.5 M_\odot$ ($2.75^{+0.05}_{-0.05} M_\odot$). Eventually, further pulsar timing measurements (see Demorest et al. 2010; Antoniadis et al. 2013; Özel & Freire 2016; Cromartie et al. 2020) as well as an increased number of observed binary NS mergers (see, e.g., Abadie et al. 2010) should be able to shed more light on the maximum mass of NSs.

The most important results of the SN and BH formation simulations to be used in our DSNB calculations are the values of the time-integrated total energy release in neutrinos of all

| Engine Model | Successful SNe | Failed SNe |
|--------------|----------------|------------|
| Z9.6 and S19.8 | 82.2% | 17.8% |
| Z9.6 and N20 | 77.2% | 22.8% |
| Z9.6 and W18 | - | 26.9% |
| Z9.6 and W15 | 70.9% | 29.1% |
| Z9.6 and W20 | 58.3% | 41.7% |

**Note.** NS and BH predictions from our five engine models are IMF-weighted according to Equation (5).

\[^{4}\text{This range is motivated by recent publications constraining the NS radius from observations of the binary NS merger GW170817 (Bauswein et al. 2017; Nicholl et al. 2017; Abbott et al. 2018; Rauh et al. 2018; Capano et al. 2020), as well as by the studies of Steiner et al. (2010), Özel et al. (2016), Özel & Freire (2016), and Lattimer & Prakash (2016). For NSs at the upper-mass end, we consider (circumferential) radii of 10 km $\lesssim R_{NS} \lesssim 12$ km, while we assume $R_{NS} \geqslant 11$ km in Appendix B for “average-mass” NSs, as suggested by Bauswein et al. (2017) and in accordance with recent results by NICER (Miller et al. 2019).}

\[^{5}\text{In fact, two competing effects play a role and can dominate under different circumstances: destabilization because of an enhanced gravitational potential due to additional thermal energy or stabilization due to increased support by thermal pressure (see, e.g., Keil & Janka 1995; O’Connor & Ott 2011; Steiner et al. 2013; da Silva Schneider et al. 2020).}
species and the time-averaged mean energies of the emitted electron antineutrinos. Figure 2 provides an overview of the corresponding values over the entire range of iron-core progenitors as a function of ZAMS mass for the exemplary case of the Z9.6 and W18 engine model, whose results will serve as a reference point in our later discussion (see Section 4). The three sets of pre-SN stars, WH15, SW14, and WH07, are separated by black vertical lines. Red bars indicate successful explosions and fallback SNe, whereas the outcomes of failed SNe are marked gray, dark blue, light blue, or cyan, depending on the different choices of critical baryonic mass for BH formation.

The upper panel shows the explosion time, \( t_{\text{exp}} \), for successful SNe, defined as the time when the shock passes 500 km (and not to be confused with the termination of our successful SN simulations and neutrino-signal calculations at 15 s, which is mentioned above). In cases of failed explosions, the time of BH formation, \( t_{\text{BH}} \), is shown, which coincides with the sudden termination of the neutrino signal. Depending on the assumed NS mass limit and the progenitor-dependent mass-accretion rate, these times range from below 1 s up to 100 s in the most extreme cases (note the logarithmic scale). This illustrates the need for a large set of long-time simulations to properly sample the neutrino contribution from the BH formation events.

\[ L_{\nu_{\tau}}(t) = L_{\nu_{\tau}}(t) + L_{\bar{\nu}_{\tau}}(t) + 4L_{\nu_{e}}(t) \]

The middle panel of Figure 2 displays the total radiated neutrino energies, \( E_{\nu_{\tau}}^{\text{tot}} \), computed as the time integrals of the summed-up neutrino luminosities of all species, \( L_{\nu_{\tau}}(t) = L_{\nu_{\tau}}(t) + L_{\bar{\nu}_{\tau}}(t) + 4L_{\nu_{e}}(t) \), from core bounce (\( t = 0 \)) to the end of the simulations (\( t = 15 s \)) or the termination of the signals (\( t = t_{\text{BH}} \)) for the SN or BH formation cases, respectively. Due to the aforementioned numerical improvements in the neutrino transport, these energies are slightly lower than those in Ertl et al. (2016) and Sukhbold et al. (2016). In Appendix B, we cross-check the values of \( E_{\nu_{\tau}}^{\text{tot}} \) by comparing them to the available budget of gravitational BE released during the cooling of the PNS, as estimated by using an analytic, radius-dependent approximate fit formula from Lattimer & Prakash (2001). We find good overall agreement, although our values might overestimate the neutrino energy loss by up to about 10%–20%, depending on the NS radius. In Section 7, we will discuss this and other uncertainties related to our DSNB predictions in more detail. In our work, we neglect contributions to the neutrino loss from fallback of matter after the successful launch of an explosion, since the amount of fallback is shown to be small (typically below \( 10^{-2} M_\odot \)) for most progenitors (Ertl et al. 2016; Sukhbold et al. 2016) and since our values for the release of NS BE in neutrinos are on the high side anyway. In addition, fallback SNe with substantial late-time fallback (possibly turning NSs to BHs) are rare, as noted above.

The mean neutrino energies of the time-integrated energy emission are displayed in the bottom panel of Figure 2 for electron antineutrinos, which are most relevant for our study. Values around 15 MeV are the rather uniform outcome of successful SNe, in agreement with other publications (e.g., Mirizzi et al. 2016; Horiuchi et al. 2018). The mean energies from failed explosions, on the other hand, vary considerably between the progenitors and depend strongly on the NS mass.
neutrino species ($\nu_1 = \nu_\alpha$, $\bar{\nu}_e$, $\bar{\nu}_\alpha$):
\[
\frac{dN(t)}{dE} = \frac{L(t)}{\langle E(t) \rangle} \frac{f_\alpha(E)}{\int_0^\infty dE f_\alpha(E)}.
\]
(2)

where we assume a spectral shape $f_\alpha(E)$ according to Keil et al. (2003),
\[
f_\alpha(E) = \left( \frac{E}{\langle E(t) \rangle} \right)^\alpha e^{-\alpha+1} E/\langle E(t) \rangle.
\]
(3)

In our models, the shape parameter $\alpha$ of the spectrum is assumed to be constant over time. Although this is a simplification, sophisticated simulations show that $\alpha$ does not change dramatically with time (e.g., Tamborra et al. 2012; Mirizzi et al. 2016), justifying this approximation. Instead, we vary $\alpha$ as a free parameter over a range of values ($1 \leq \alpha \leq 4$), which we motivate in Appendix D.

For each progenitor and neutrino species, we then perform a time integration over the period of emission, from core bounce at $t = 0$ to a final time of $t = t_f$ (with $t_f = 15$ s for successful explosions and $t_f = t_{BH}$ for failed SNe):
\[
\frac{dN}{dE} = \frac{\bar{\xi}}{\xi} \int_0^{t_f} dt \frac{dN(t)}{dE}.
\]
(4)

Because the luminosities of heavy-lepton neutrinos $\nu_\alpha$ are very approximate in our sets of simulations due to the incomplete microphysics and the relatively moderate core contraction mentioned in Section 2, we rescale each time-integrated spectrum with a factor $\bar{\xi}/\xi$ (see Appendix C for details). By this procedure we adopt the total radiated neutrino energy ($E_{\nu,\text{tot}}$) from the simulated core-collapse models but redistribute them between the different neutrino species with weight factors obtained from SN and BH formation models with sophisticated neutrino treatment (see Table C1). $ar{\xi} = \xi_{95} = (E_{\nu,\text{tot}}/E_{\nu,\text{old}})^{\text{new}}$ thus constitutes the fraction of the total energy emitted in neutrino species $\nu_\alpha$. Correspondingly, $\xi = \xi_{95} = (E_{\nu,\text{tot}}/E_{\nu,\text{old}})^{\text{new}}$ stands for the relative energy as originally computed in the core-collapse models considered in our study.

In Appendix D, we compare the shapes of our time-integrated spectra with results from sophisticated simulations (with detailed microphysics) by a few exemplary cases to examine the viability of our approximate treatment. We find good agreement with these simulations for values of the instantaneous shape parameter $\alpha$ of $\sim 3$–3.5 for successful explosions and of $\sim 2$ for failed SNe. In Appendix E, we provide, for a set of representative successful and failed SN models, the total radiated neutrino energies, the mean neutrino energies, and the spectral-shape parameters of the time-integrated neutrino ($\bar{\nu}_e$, $\bar{\nu}_\alpha$ and $\nu_\alpha$) spectra.

As mentioned in Section 2, our DSNB flux calculations also include the neutrino signal of the $8.8 M_\odot$ ECSN simulated by Hüdepohl et al. (2010). The corresponding time-integrated spectra are computed according to Equations (2)–(4), but with the time-dependent shape parameters $\alpha = \alpha(t)$ given by the simulation. We use the neutrino data of "model SF,” which takes into account the full set of neutrino interactions listed in

\[\alpha \approx 2.3 \text{ corresponds to a Fermi–Dirac distribution with a vanishing degeneracy parameter; } \alpha > 2.3 \text{ to a pinched spectrum and } \alpha < 2.3 \text{ to an anti-pinched spectrum; } \alpha = 2.0 \text{ gives the Maxwell–Boltzmann distribution.} \]
Appendix A of Buras et al. (2006), including nucleon–nucleon bremsstrahlung, inelastic neutrino–nucleon scattering, and neutrino-pair conversions between different flavors, making rescaling of the spectra unnecessary, i.e., $\xi/\xi = 1$ for all flavors.

### 3.2. IMF-weighted Average

The relative number of pre-SN stars depends on their birth masses. For our DSNB flux predictions, the time-integrated neutrino spectra $dN/dE$ for each core-collapse case therefore need to be weighted by an IMF (providing the number of stars formed per unit of mass as a function of the stellar ZAMS mass $M$). As in Hopkins & Beacom (2006), Horiuchi et al. (2011), and Mathews et al. (2014), we apply the modified Salpeter-A IMF of Baldry & Glazebrook (2003),

$$\phi(M) \propto M^{-\zeta},$$

with $\zeta = 2.35$ for birth masses $M \geq 0.5 \, M_\odot$ and $\zeta = 1.5$ for $0.1 \, M_\odot \leq M < 0.5 \, M_\odot$. In our study, we consider masses up to $125 \, M_\odot$. However, due to the steep decline of Equation (5), the high-mass end is suppressed and is thus of minor relevance for the DSNB.

The IMF-weighted neutrino spectrum $dN_{\text{CC}}/dE$ of all core-collapse events can then be calculated as a sum over mass intervals $\Delta M$, associated with our discrete set of progenitor stars according to

$$\frac{dN_{\text{CC}}}{dE} = \sum_i \int_{\Delta M} dM \frac{dN_i}{dE} \phi(M),$$

where $\Delta M$ denotes the mass interval around the ZAMS mass $M_i$ with the time-integrated spectrum $dN_i/dE$ of the corresponding SN, failed SN, or ECSN simulation. Equation (6) is applied separately to the different neutrino species. As in Section 3.1, the indices $\nu_e$, $\bar{\nu}_e$, and $\nu_x$ are omitted here for the sake of clarity. In the following, we primarily focus on $\nu_e$, since the prospects for a first detection of the DSNB in upcoming detectors are the best for this species (see, e.g., Beacom & Vagins 2004; Yüksel et al. 2006; Horiuchi et al. 2009; An et al. 2016).

### 3.3. Cosmic Core-collapse Rate

Nuclear burning proceeds fast in massive stars. As a consequence, the progenitors of core-collapse SNe (and failed SNe) have relatively “short” (<10^8 yr) lives compared to cosmic timescales (see Kennicutt 1998). Therefore, the assumption is well justified that the cosmic core-collapse rate density $R_{\text{CC}}(z)$ as a function of redshift equals the birth rate density of stars in the relevant ZAMS mass range ($8.7 \, M_\odot \leq M \leq 125 \, M_\odot$), i.e.,

$$R_{\text{CC}}(z) = \psi_e(z) \int_{\Delta M} dM \phi(M) \frac{dN_i}{dE} \approx \psi_e(z) \int_{8.7 \, M_\odot}^{125 \, M_\odot} dM \phi(M),$$

where $\psi_e(z)$ describes the cosmic star formation history (SFH) in terms of the star formation rate in units of $M_\odot$ Myr^{-1}, which can be deduced from observations (e.g., Hopkins & Beacom 2006; Reddy et al. 2008; Rujopakarn et al. 2010) and is thus independent of cosmological assumptions. In our study, we adopt the parameterized description by Yüksel et al. (2008),

$$\psi_e(z) = \rho_0 \left[ \left( 1 + z \right)^{\eta_0} + \left( 1 + z \right)^{\eta_1} \right],$$

with the best-fit parameters from Mathews et al. (2014); see Table 1 therein). Note that the derivation of an SFH $\psi_e(z)$ from observational data requires the use of an IMF, which should be consistent with the one employed in Equation (7). For this reason we use the Salpeter-A IMF (Baldry & Glazebrook 2003) to be consistent with the SFH data sample compiled by Mathews et al. (2014), which is based on the data sets by Hopkins & Beacom (2006) and Horiuchi et al. (2011).

Even though the cosmic core-collapse rate is not yet known to good accuracy (its impact on the DSNB flux is discussed, e.g., by Lien et al. 2010), our work is focused on variations of the neutrino source properties. To still account for the large uncertainty of $R_{\text{CC}}$, we additionally employ the $+\pm 1\sigma$ upper and lower limits to the SFH of Mathews et al. (2014), such that we obtain $R_{\text{CC}}(0) = 8.93 \pm 2.67 \times 10^{-5} \text{ Mpc}^{-3} \text{ yr}^{-1}$ for the local universe. In Section 4, we further test parameterizations of the SFH by Madau & Dickinson (2014) and Fermi-LAT Collaboration et al. (2018). The cosmic metallicity evolution and its impact on the DSNB will be discussed briefly in Section 7.2. For our DSNB calculations, we consider contributions up to a maximum redshift of $z_{\text{max}} = 5$. This limit is justified because, as pointed out in numerous previous works (Ando 2004; Keen & Lunardini 2012; Mathews et al. 2014; Nakazato et al. 2015; Lunardini 2016), only sources at lower redshifts ($z \lesssim 1$–2) noticeably add to the high-energy part of the DSNB, which is most relevant for detection (see Figure 3). Neutrinos from higher $z$ are almost entirely shifted to energies below 10 MeV, where background sources dominate the flux and thus prevent a clear identification of the DSNB signal (see, e.g., Lunardini 2016).

### 3.4. Cosmological Model

Throughout this work we assume standard $\Lambda$CDM cosmology with the present-day mass-energy density parameters $\Omega_{\text{m}} = 0.3$ and $\Omega_{\Lambda} = 0.7$ of matter and a cosmological constant, respectively, and the Hubble constant $H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1}$. The expansion history of the universe is then given by $dz/dt = -H_0 (1 + z) \sqrt{\Omega_m (1 + z)^3 + \Omega_\Lambda}$. Using this together with Equation (1), we can write the DSNB flux

$\psi_e(z)$ describes the cosmic star formation history (SFH) in terms of the star formation rate in units of $M_\odot$ Myr^{-1}, which can be deduced from observations (e.g., Hopkins & Beacom 2006; Reddy et al. 2008; Rujopakarn et al. 2010). Even though the cosmic core-collapse rate is not yet known to good accuracy (its impact on the DSNB flux is discussed, e.g., by Lien et al. 2010), our work is focused on variations of the neutrino source properties. To still account for the large uncertainty of $R_{\text{CC}}$, we additionally employ the $+\pm 1\sigma$ upper and lower limits to the SFH of Mathews et al. (2014), such that we obtain $R_{\text{CC}}(0) = 8.93 \pm 2.67 \times 10^{-5} \text{ Mpc}^{-3} \text{ yr}^{-1}$ for the local universe. In Section 4, we further test parameterizations of the SFH by Madau & Dickinson (2014) and Fermi-LAT Collaboration et al. (2018). The cosmic metallicity evolution and its impact on the DSNB will be discussed briefly in Section 7.2. For our DSNB calculations, we consider contributions up to a maximum redshift of $z_{\text{max}} = 5$. This limit is justified because, as pointed out in numerous previous works (Ando 2004; Keen & Lunardini 2012; Mathews et al. 2014; Nakazato et al. 2015; Lunardini 2016), only sources at lower redshifts ($z \lesssim 1$–2) noticeably add to the high-energy part of the DSNB, which is most relevant for detection (see Figure 3). Neutrinos from higher $z$ are almost entirely shifted to energies below 10 MeV, where background sources dominate the flux and thus prevent a clear identification of the DSNB signal (see, e.g., Lunardini 2016).

We point out that Mathews et al. (2014) used an equality relation (instead of a proportionality) for Equation (5), which leads to a discontinuous behavior at $M = 0.5 \, M_\odot$. This seems to be in conflict with the (continuous) IMF employed in the compilations of star formation rate data by Hopkins & Beacom (2006) and Horiuchi et al. (2011), which served as a basis for the study of Mathews et al. (2014). For this reason, we construct a continuous IMF by properly choosing the normalization coefficients in the two mass intervals described by Equation (5).
The eye, the approximate detection window of (s-process elements) studies of the impact of different cosmological models on the spectrum (in units of MeV\(^{-1}\) cm\(^{-2}\) s\(^{-1}\)) as

\[
\frac{d\Phi}{dE} = \frac{c}{H_0} \int_0^{z_{\text{max}}} \frac{dN_{\text{CC}}}{dE'} \frac{R_{\text{CC}}(z)}{\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}} dz. \tag{9}
\]

We do not vary the cosmological assumptions within our work, following most publications on the DSNB topic. For recent studies of the impact of different cosmological models on the DSNB, the reader is referred to Barranco et al. (2018) or Yang et al. (2019). Having described our computational model with all of its required inputs, we now proceed to a discussion of our results.

### 4. Fiducial DSNB Model

In this section, we present our fiducial DSNB predictions and show how the single components (ECSNe, SNe, and failed SNe at various redshifts) contribute to the total flux. The following set of inputs makes up our fiducial model:

1. As in Ertl et al. (2016) and Sukhbold et al. (2016), we take the intermediate engine model Z9.6 and W18 (with 26.9% failed SNe) as our reference case.\(^{11}\)

2. Guided by Margalit & Metzger (2017), we assume a fiducial mass of \(M_{\text{NS,b}}^{\text{lim}} = 2.7 M_\odot\) for the NS (baryonic) mass limit, where a PNS is assumed to collapse to a BH and the neutrino signal is truncated (see Section 2.3).

3. According to our detailed analysis of the spectral shapes in Appendix D, we take a “best-fit” value of \(\alpha = 3.5\) for the instantaneous spectral-shape parameter of \(\tau_e\) for successful SNe with baryonic NS masses of \(M_{\text{NS,b}} \leq 1.6 M_\odot\); of \(\alpha = 3.0\) for SNe with \(M_{\text{NS,b}} > 1.6 M_\odot\), and of \(\alpha_{\text{BH}} = 2.0\) for BH-forming, failed explosions.

4. As our reference for the cosmic core-collapse rate, we take Equations (7) and (8) with the best-fit parameters for the SFH according to Mathews et al. (2014, Table 1), which yield \(R_{\text{CC}}(0) = 8.93 \times 10^{-7}\) Mpc\(^{-3}\) yr\(^{-1}\) for the local universe.

In Figure 3, we first illustrate how the various sources contribute to the total DSNB flux spectrum, \(d\Phi/dE\), of electron antineutrinos, using our fiducial model. The left panel shows the individual fluxes arising from ECSNe, “conventional” iron-core SNe, and BH-forming, failed SNe (light to dark solid lines). Integrated over all energies, ECSNe contribute only 2.3% (0.7 cm\(^{-2}\) s\(^{-1}\)) to the total flux (28.8 cm\(^{-2}\) s\(^{-1}\)), whose spectrum is shown by a black dashed line. This value is much lower than the \(\sim 10\%\) suggested by Mathews et al. (2014) as they assumed a considerably wider ZAMS mass range, (8–10) \(M_\odot\), compared to the (8.7–9) \(M_\odot\) applied in our work (see Jones et al. 2013; Doherty et al. 2015). Above 15 MeV, the contribution of ECSNe accounts for even less than 1% due to the more rapidly declining spectrum (recall the low mean energy of 11.6 MeV mentioned in Section 2). However, since the exact mass window of ECSNe is still unclear (see, e.g., Poelarends et al. 2008; Jones et al. 2013, 2016; Doherty et al. 2015; Kirsebom et al. 2019; Zha et al. 2019; Leung et al. 2020) and other sources such as ultrastripped SNe, AIC events, and MIC events might contribute to the DSNB flux with source spectra similar to those of ECSNe, we will consider an enhanced LM component in Section 5.2.

“Conventional” iron-core SNe and failed SNe possess comparable integrated fluxes (16.0 cm\(^{-2}\) s\(^{-1}\) and 12.1 cm\(^{-2}\) s\(^{-1}\)) in the case of our fiducial model as shown in Figure 3, yet with distinctly different spectral shapes. Below \(\sim 15\) MeV, the contribution from successful explosions is greater, whereas failed explosions dominate the flux at high energies due to their generally harder spectra (see bottom panel of Figure 2). This has been pointed out by previous works (e.g., Lunardini 2009; Keenly & Lunardini 2012; Nakazato 2013; Priya & Lunardini 2017) and can also be seen in Table 2, where we list the relative flux contributions from the various sources for different ranges of neutrino energies. Between 20 MeV and 30 MeV, failed SNe account for 62% of the total flux (at still higher energies, even 76%). Naturally, these numbers (here given for our reference

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\(^{11}\)The resulting nucleosynthesis yields show reasonable agreement with the solar element abundances (when type Ia SNe are included), and the NS mass distribution roughly fits observational data (Özel & Freire 2016), as does the distribution of BH masses (Witkowicz et al. 2014) if one assumes that only the star’s helium core collapses while its hydrogen envelope gets unbound (see Nadezhin 1980; Lovegrove & Woosley 2013; Kochanek 2014). For more details, the reader is referred to Sukhbold et al. (2016). The rather high fraction of failed explosions (26.9%: see Table 1) is not unrealistic given the large discrepancy between the observed SN rate and the SFH (Horiuchi et al. 2011). Also the recent discovery of a disappearing star (Adams et al. 2017) supports a nonzero fraction of failed explosions. Our weakest engine model, Z9.6 and W20, which yields by far the largest fraction of failed SNe (41.7%: see Table 1), is disfavored since it would lead to a significant underproduction of s-process elements (Brown & Woosley 2013; Sukhbold et al. 2016).
model set) depend strongly on the fraction of failed explosions and their neutrino emission (see Section 5.1). Compared to previous studies, we obtain a generally increased DSNB flux, advantageous for its imminent detection. We will comment on this issue more thoroughly below.

In the right panel of Figure 3, we compare the DSNB contributions from different redshift intervals (light to dark for increasing redshift). At high energies (>20 MeV), the flux mainly originates from sources below $z \sim 1$, as illustrated in several previous works (Ando 2004; Khee & Lunardini 2012; Mathews et al. 2014; Nakazato et al. 2015; Lunardini 2016). Only at lower energies does the contribution from large redshifts get increasingly important (see Table 2). In both panels of Figure 3, shaded bands bracket the approximate energy window of $\sim$(10–30) MeV, which is most relevant for DSNB detection in upcoming neutrino observatories. Beyond that, background sources (such as reactor and solar neutrinos at low energies and atmospheric neutrinos at high energies) dominate the flux and make DSNB measurement unfeasible (see, e.g., the review by Lunardini 2016).

As already pointed out in Section 3.3, the cosmic SFH constitutes one of the major uncertainties in predicting the DSNB. Thus, before we proceed to the main part of our parameter study, we test how the DSNB flux spectrum depends on the assumed parameterization of the SFH. In Figure 4, we show our fiducial DSNB model (black dashed line) together with the uncertainty corresponding to the $\pm 1$σ confidence interval of the SFH according to Mathews et al. (2014, “M+2014”; gray-shaded band). For comparison, we also employ the more conservative SFH from Madau & Dickinson (2014, “MD2014”; Equation (15); orange line) as well as the recent results of Fermi-LAT Collaboration et al. (2018) on the evolution of the extragalactic background light (EBL): an empirical EBL reconstruction (EBLr) and a physical EBL (pEBL) model (blue and green shaded bands, respectively; see their Figure 3). We already noted in Section 3.3 that, to remain consistent with the IMF employed for determining the SFH $\psi_\star(z)$, the same IMF should be taken also for the conversion of $\psi_\star(z)$ to the cosmic core-collapse rate density $R_{CC}(z)$. For this reason, we adopt a conventional Salpeter IMF (Salpeter 1955) and the one by Chabrier (2003) instead of the Salpeter-A IMF from Baldry & Glazebrook 2003) for the conversion of the SFH to the cosmic core-collapse rate (Equation (7)) when the SFHs from Madau & Dickinson (2014) and Fermi-LAT Collaboration et al. (2018) are used, respectively (see main text). As in Figure 3, vertical bands frame the approximate detection window.

![Figure 4](image-url)

**Figure 4.** Dependence of the DSNB $\bar{\nu}_e$-flux spectrum on the assumed parameterization of the cosmic SFH. In our fiducial model (black dashed line; see Figure 3), the SFH of Mathews et al. (2014) is employed; the gray-shaded band corresponds to their $\pm 1$σ upper and lower limits. The orange line indicates the DSNB spectrum for the SFH of Madau & Dickinson (2014), whereas the DSNB spectrum for the SFH from Fermi-LAT Collaboration et al. (2018) is indicated by blue (empirical EBL reconstruction) and green (physical EBL model) shaded bands (1σ confidence regions). Note that a conventional Salpeter IMF (Salpeter 1955) and the one by Chabrier (2003) are used (instead of the Salpeter-A IMF from Baldry & Glazebrook 2003) for the conversion of the SFH to the cosmic core-collapse rate (Equation (7)) when the SFHs from Madau & Dickinson (2014) and Fermi-LAT Collaboration et al. (2018) are used, respectively (see main text). As in Figure 3, vertical bands frame the approximate detection window.

5. **DSNB Parameter Study**

In this section, we present the results of our detailed DSNB parameter study. Using large grids of long-time neutrino signals (see Section 2), we probe the sensitivity of the DSNB to three critical source properties (in Section 5.1): the fraction of failed explosions (by means of our different engine models),

### Table 2

| Total DSNB Flux ($\bar{\nu}_e$) | (0–10) MeV | (10–20) MeV | (20–30) MeV | (30–40) MeV | (0–40) MeV |
|---------------------------------|------------|-------------|-------------|-------------|------------|
|                                 | 22.7 cm$^{-2}$ s$^{-1}$ | 5.4 cm$^{-2}$ s$^{-1}$ | 0.6 cm$^{-2}$ s$^{-1}$ | 0.1 cm$^{-2}$ s$^{-1}$ | 28.8 cm$^{-2}$ s$^{-1}$ |
| ECSNe                           | 2.6%       | 1.2%        | 0.5%        | 0.2%        | 2.3%       |
| Iron-core SNe                   | 57.1%      | 51.8%       | 37.5%       | 23.9%       | 55.6%      |
| Failed SNe                      | 40.3%      | 47.0%       | 62.0%       | 75.8%       | 42.1%      |

**Note.** Top row: Total DSNB flux of $\bar{\nu}_e$ for our fiducial model (Z9.6 and W18; $M_\text{lim}^{\text{CC}} = 2.7\ M_\odot$; best-fit $\alpha$), integrated over different energy intervals. Second to fourth rows: Relative contributions from the various source types (ECSNe/iron-core SNe/failed SNe with BH formation). Rows 5–9: Relative contributions from different redshift intervals (see also Figure 3).

energies (see Riya & Rentala 2020). Throughout our work, we assume an uncertainty of the cosmic core-collapse rate corresponding to the $\pm 1$σ band of Mathews et al. (2014).

Since our overall findings apply similarly to all neutrino species, we constrain our discussion to electron antineutrinos for now. In Section 6, we will briefly discuss the DSNB flux spectrum of electron neutrinos, and in Section 7.1, we will comment on the influence of heavy-lepton neutrinos in the context of neutrino flavor oscillation effects.
the threshold mass for BH formation, and the spectral shape of the neutrino emission from failed explosions. Moreover, the possible enhancement of the DSNB by an additional generic LM component is explored (Section 5.2) as well as the effect of including binary progenitor models (Section 5.3).

5.1. DSNB Parameter Dependence

First, we study the impact of our engine model (as described in Section 2.2) on the DSNB flux spectrum. In the upper left panel of Figure 5, we show $d\Phi/dE$ for the case of electron antineutrinos. In the different panels the engine models (upper left panel), the NS mass limit for BH formation (upper right panel), and the instantaneous spectral-shape parameter, $\alpha_{\rm BH}$, of the time-dependent neutrino emission from BH formation events (lower left panel) are varied, while all other parameters are kept at their reference values (Z9.6 and W18: $M_{\rm lim} = 2.7 M_{\odot}$; best-fit $\alpha$, i.e., $\alpha = 3.5$ for SNe with $M_{\rm NS,b} \leq 1.6 M_{\odot}$, $\alpha = 3.0$ for those with $M_{\rm NS,b} > 1.6 M_{\odot}$, and $\alpha_{\rm BH} = 2.0$ for failed SNe; see Section 4). In the lower right panel, the additional contribution from LM NS–forming events is shown for different constant rate densities $R_{\rm LM}$. For comparison, the pale red band marks the LM flux for an evolving rate instead (see main text for details). Our fiducial model with $R_{\rm LM} = 0$ is plotted as a dashed line. In each panel, a gray-shaded band indicates the uncertainty arising from the cosmic core-collapse rate (corresponding to the ±1σ upper and lower limits to the SFH of Mathews et al. 2014). As in Figure 3, vertical bands frame the approximate detection window.

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Figure 5. Parameter dependence of the DSNB flux spectrum, $d\Phi/dE$, for the case of electron antineutrinos. In the different panels the engine models (upper left panel), the NS mass limit for BH formation (upper right panel), and the instantaneous spectral-shape parameter, $\alpha_{\rm BH}$, of the time-dependent neutrino emission from BH formation events (lower left panel) are varied, while all other parameters are kept at their reference values (Z9.6 and W18: $M_{\rm lim} = 2.7 M_{\odot}$; best-fit $\alpha$, i.e., $\alpha = 3.5$ for SNe with $M_{\rm NS,b} \leq 1.6 M_{\odot}$, $\alpha = 3.0$ for those with $M_{\rm NS,b} > 1.6 M_{\odot}$, and $\alpha_{\rm BH} = 2.0$ for failed SNe; see Section 4). In the lower right panel, the additional contribution from LM NS–forming events is shown for different constant rate densities $R_{\rm LM}$. For comparison, the pale red band marks the LM flux for an evolving rate instead (see main text for details). Our fiducial model with $R_{\rm LM} = 0$ is plotted as a dashed line. In each panel, a gray-shaded band indicates the uncertainty arising from the cosmic core-collapse rate (corresponding to the ±1σ upper and lower limits to the SFH of Mathews et al. 2014). As in Figure 3, vertical bands frame the approximate detection window.

the threshold mass for BH formation, and the spectral shape of the neutrino emission from failed explosions. Moreover, the possible enhancement of the DSNB by an additional generic LM component is explored (Section 5.2) as well as the effect of including binary progenitor models (Section 5.3).

5.1. DSNB Parameter Dependence

First, we study the impact of our engine model (as described in Section 2.2) on the DSNB flux spectrum. In the upper left panel of Figure 5, we show $d\Phi/dE$ for the various choices of central neutrino engines for our simulations. Sets with a higher percentage of failed explosions (see Figure 1 and Table 1) yield an enhanced DSNB flux, especially in the high-energy regime. This overall picture is in line with studies by Lunardini (2009), Lien et al. (2010), and Keehn & Lunardini (2012), who varied the fraction of BH-forming collapses while applying generic neutrino spectra and thus neglecting progenitor dependences. More recently, Priya & Lunardini (2017) and Møller et al. (2018) examined the fraction of failed SNe by assuming different ZAMS mass distributions, while Horiuchi et al. (2018), for the first time, employed a larger sample of simulations including seven BH formation cases, thus taking into account progenitor-dependent variations in the neutrino emission from failed explosions (by linearly interpolating the total energetics, mean energy, and shape parameter of their time-integrated neutrino spectra as a function of the compactness parameter of O’Connor & Ott 2011; see footnote 6). They explored relative fractions of BH formation cases between 0% and 45% by taking different threshold values for the compactness above which they assumed their progenitors formed BHs.

Using our large sets of long-time simulations without predefined outcomes (also resulting in BH formation of less compact progenitors with low mass-accretion rates), we can confirm the common result of the previous studies: the larger the fraction of failed explosions, the stronger the enhancement of the DSNB at high energies. To better quantify this behavior, we follow Lunardini (2007) and fit the high-energy tail (20 MeV $\leq E \leq 30$ MeV) of our DSNB flux spectra with an exponential function:

$$
\frac{d\Phi}{dE} \simeq \phi_0 e^{-E/E_0},
$$

(10)

Our model set Z9.6 and S19.8 with the lowest fraction of failed explosions (17.8%) features the steepest decline (i.e., $E_0 = 4.5$ MeV), while Z9.6 and W20 with 41.7% BH formation
cases yields a flatter spectrum with \( E_0 = 5.1 \text{ MeV} \). The “normalization” \( \phi_0 \), on the other hand, is hardly affected by our choice of engine model. Instead, it is determined by the uncertainty arising from the cosmic core-collapse rate, which shifts the entire flux spectrum vertically without changing the slope by more than \( \sim 1\% \).\(^{12}\) The gray-shaded bands in Figure 5 indicate this severe normalization uncertainty (the \( \pm 1\sigma \) upper limit to the SFH of Mathews et al. (2014) is taken for our highest-flux model, and the \( -1\sigma \) lower limit for our lowest-flux model). The aspect that the failed-SN fraction is likely to exhibit a dependence on metallicity (and thus on redshift) was pointed out by Nakazato et al. (2015) and Yüksel & Kistler (2015). We will come back to this point in Section 7.2.

The impact of the NS mass limit on the DSNB has been discussed in the literature to some extent (Lunardini 2009; Keehn & Lunardini 2012; Nakazato et al. 2015; Hidaka et al. 2016, 2018). Commonly, the spectra from exemplary simulations of BH formation with two different EoSs have been compared: the stiff Shen EoS (Shen et al. 1998, with incompressibility \( K = 281 \text{ MeV} \)) and a softer EoS by Latimer & Swesty (1991, “LS180” or “LS220,” with \( K = 180 \text{ MeV} \) or \( K = 220 \text{ MeV} \)). Generally, a stiff EoS supports the transiently existing PNS of a failed SN against gravity up to a higher limiting mass than a soft EoS does. The final collapse to a BH therefore sets in after a longer period of mass accretion and neutrino emission with the consequence of higher spectral temperatures and an enhanced contribution to the DSNB flux.

Having a large compilation of long-time simulations at hand, we take a different (more rigorous) approach in our work: As described in Section 2, we directly vary the maximum baryonic NS mass, \( M_{\text{NS,b}} \), without applying a certain EoS. Our neutrino signals from failed explosions are then truncated when the mass accretion from the collapsing progenitor star pushes the PNS mass beyond this critical threshold for BH formation. In the upper right panel of Figure 5, we show the DSNB flux spectra for our different choices of \( M_{\text{NS,b}}^{\text{lim}} \). Raising the NS mass limit from \( 2.3 M_\odot \) to \( 3.5 M_\odot \) drastically enhances the flux at higher energies, thus lifting the value of the slope parameter, \( E_0 / \text{MeV} \) (see Equation (10)), from 4.4 to 5.6. This strong effect becomes immediately clear from Figure 2: a higher NS mass limit leads to enhanced time-integrated neutrino luminosities and generally hotter spectra, in line with studies by Lunardini (2009), Keehn & Lunardini (2012), Nakazato et al. (2015), and Hidaka et al. (2016, 2018).

We should mention that our study does not consider the possibility of a progenitor-dependent threshold mass for BH formation. O’Connor & Ott (2011) pointed out that thermal pressure support may be stronger for stars with high core compactness, lifting the maximum PNS mass to somewhat larger values. This might slightly reduce differences in the neutrino emission between individual progenitors. The results of Figure 5 should, however, remain essentially unchanged, because thermal stabilization of the PNS should be most relevant when the mass-accretion rate is high and the PNS becomes very hot. In such cases, however, the critical limit for BH formation is also reached quickly and the neutrino emission is not significantly extended. The wide range of values for \( M_{\text{NS,b}}^{\text{lim}} \) considered in our study should include the true NS mass limit, which depends on the still incompletely known high-density EoS of NS matter. Once the latter is better constrained by astrophysical observations and nuclear experiments and theory and thus once the maximum mass of cold NSs is better constrained, the question of a progenitor-dependent thermal effect on transient PNS stabilization can be addressed more thoroughly.

In our study, the spectral shape of the time-dependent neutrino emission is assumed to obey Equation (3) with a constant shape parameter \( \alpha \). Following our detailed analysis of the spectral shapes in Appendix D, we show in the lower left panel of Figure 5 how the DSNB flux spectrum changes when different values (between 1.0 and 3.0) of this instantaneous spectral-shape parameter \( \alpha = \alpha_{\text{BH}} \) are taken for the emission from failed explosions. For successful SNe, \( \alpha \) is not varied but kept constant at the best-fit values of 3.0 and 3.5 (for \( M_{\text{NS,b}} > 1.6 M_\odot \) and \( M_{\text{NS,b}} \leq 1.6 M_\odot \), respectively). Similarly, its influence on individual failed-SN source spectra, a small value of \( \alpha_{\text{BH}} \) broadens the shape of the DSNB such that its high-energy tail gets lifted relative to the peak (see Keil et al. 2003; Lunardini 2007, 2016). For \( \alpha_{\text{BH}} = 1.0 \) (i.e., anti-pinched failed-SN source spectra), the exponential fit of Equation (10) yields \( E_0 = 5.4 \text{ MeV} \) and \( \phi_0 = 6.3 \text{ MeV}^{-1} \text{ cm}^{-2} \text{s}^{-1} \) in the range of neutrino energies 20 MeV \( \leq E \leq 30 \text{ MeV} \). Choosing \( \alpha_{\text{BH}} = 3.0 \), on the other hand, results in a more prominent peak at the cost of a suppressed flux at high energies (\( E_0 = 4.5 \text{ MeV} \); \( \phi_0 = 12.3 \text{ MeV}^{-1} \text{ cm}^{-2} \text{s}^{-1} \)). Because the instantaneous spectral-shape parameter \( \alpha \) is only varied for failed SNe while the contribution from successful SNe is unchanged, a slight “kink” becomes visible in the overall DSNB flux spectrum for cases of small \( \alpha_{\text{BH}} \), unveiling its “two-component” nature. Notice the crossings of the different curves at \( \sim 3 \text{ MeV} \) and \( \sim 15 \text{ MeV} \). Accordingly, we construct the shaded band for the uncertainty of \( R_{\text{CC}} \) such that the lowest-flux and highest-flux models are considered in each segment.

In Table 3, we provide an overview of the two fit parameters \( \phi_0 \) and \( E_0 \) for all models discussed in this section. We use the following naming convention for our DSNB models: “W18-BH2.7-α2.0” corresponds to our fiducial model with the Z9.6 and W18 neutrino engine (“W18”), with a baryonic NS mass limit of \( M_{\text{NS,b}}^{\text{lim}} = 2.7 M_\odot \) (“BH2.7,”) and with the best-fit choice for the instantaneous spectral-shape parameter (“α2.0;” i.e., \( \alpha_{\text{BH}} = 2.0 \)). The two models “W20-BH3.5-α1.0” and “S198-BH2.3-α3.0,” which employ the most extreme parameter combinations, yield the largest and smallest slope parameters \( E_0 \) among all our models and thus the highest and lowest fluxes at high energies, respectively.

5.2. Additional LM Component

As we mentioned in Section 2.1, the low-mass range of core-collapse SN progenitors is rather uncertain. It is widely believed that in degenerate ONeMg cores electron-capture reactions on \(^{20}\text{Ne}\) and \(^{24}\text{Mg}\) can counter the effects of oxygen deflagration, initiating collapse to an NS rather than thermonuclear runaway (Miyaji et al. 1980; Nomoto 1984, 1987). However, the conditions for the occurrence of such an ECSN in nature are still being debated (see, e.g., Jones et al. 2016; Kirsebom et al. 2019; Zha et al. 2019; Leung et al. 2020). Moreover, observations suggest that most massive stars are in binary systems (see, e.g., Mason et al. 2009; Sana et al. 2012), and evolution in binaries might lead to a larger population of

\(^{12}\) The fact that \( E_0 \) is not entirely unaffected by changes of \( R_{\text{CC}} \) is due to the different functional dependences of the \( \pm 1\sigma \) upper/lower limits to the cosmic SFH on the redshift \( z \) (see Table 1 of Mathews et al. 2014).
Table 3
Exponential-fit Parameters of Equation (10) for a Subset of Our DSNB Models

| Model           | $\alpha_0$ (MeV$^{-1}$ cm$^2$ s$^{-1}$) | $E_0$ (MeV) |
|-----------------|----------------------------------------|-------------|
| W18-BH2.7-α2.0 (fiducial) | 9.4±2.3, 4.8, 6.7, 4.4 | 4.8±0.4, 4.4, 4.4, 5.2 |
| W20-BH3.5-α1.0 (max.) | 6.6±1.5, 4.9, 3.3 | 6.7±0.5, 6.4, 7.1, 7.1 |
| S19.8-BH2.3-α3.0 (min.) | 12.6±10.5, 9.6, 8.7, 5.6 | 4.9±0.3, 4.3, 4.4, 4.4 |
| S19.8-BH2.7-α2.0 | 10.6±2.2, 8.5, 7.5, 4.9 | 4.5±0.4, 4.6, 4.8, 4.9 |
| N20-BH2.7-α2.0 | 9.3±3.2, 7.4, 6.7, 4.3 | 4.7±0.4, 4.7, 4.5, 5.1 |
| W15-BH2.7-α2.0 | 9.1±3.3, 7.0, 6.5, 4.2 | 4.9±0.4, 4.9, 5.2, 5.3 |
| W20-BH2.7-α2.0 | 9.6±3.4, 6.7, 6.8, 4.4 | 5.1±0.5, 5.4, 5.5, 5.5 |
| W18-BH2.3-α2.0 | 9.9±3.4, 7.7, 7.0, 4.5 | 4.4±0.4, 4.5, 4.7, 4.8 |
| W18-BH2.7-α2.0 | 9.4±3.3, 7.3, 6.7, 4.4 | 4.8±0.4, 4.8, 5.0, 5.3 |
| W18-BH3.1-α2.0 | 9.0±3.1, 6.9, 6.4, 4.3 | 5.2±0.5, 5.1, 5.0, 5.6 |
| W18-BH3.5-α2.0 | 8.7±3.1, 6.6, 6.3, 4.2 | 5.5±0.5, 5.4, 5.0, 6.0 |
| W18-BH2.7-α1.0 | 6.3±3.1, 5.5, 4.7, 3.2 | 5.4±0.4, 5.2, 5.0, 5.7 |
| W18-BH2.7-α1.5 | 7.8±3.2, 6.5, 5.7, 3.8 | 5.0±0.4, 5.0, 5.2, 5.4 |
| W18-BH2.7-α2.0 | 9.4±3.3, 7.3, 6.7, 4.4 | 4.8±0.4, 4.8, 5.1, 5.2 |
| W18-BH2.7-α2.5 | 10.9±3.3, 8.1, 7.6, 5.0 | 4.6±0.4, 4.7, 4.9, 5.0 |
| W18-BH2.7-α3.0 | 12.3±3.5, 8.9, 8.5, 5.5 | 4.4±0.4, 4.5, 4.8, 4.9 |
| W18-BH2.7-α3.0-He33 | 8.0±3.4, 6.3, 5.7, 3.7 | 4.7±0.4, 4.7, 5.0, 5.1 |
| W18-BH2.7-α3.0-He100 | 5.5±4.4 | 4.5±0.4, 4.5, 4.8 |
| S19.8-BH2.3-α2.0 | 11.1±3.0, 8.8, 7.8, 5.0 | 4.2±0.3, 4.4, 4.5, 4.6 |
| W18-BH3.5-α1.0 | 5.8±3.0, 4.9, 4.4, 3.0 | 6.3±0.4, 5.9, 6.1, 6.7 |
| W15-BH3.5-α1.0 | 5.7±4.6, 4.7, 4.3, 2.9 | 6.4±0.4, 6.0, 6.8, 6.8 |
| W20-BH2.7-α1.0 | 6.2±4.9, 4.8, 4.5, 3.1 | 5.9±0.5, 5.7, 6.2, 6.3 |
| W20-BH3.1-α1.0 | 6.3±5.1, 4.7, 4.7, 3.2 | 6.3±0.5, 6.1, 6.7, 6.7 |
| W20-BH3.5-α2.0 | 10.2±5.6, 6.7, 7.2, 4.8 | 5.8±0.6, 5.7, 6.2, 6.3 |

Note. The fits are applied in the energy region 20 MeV ≤ $E$ ≤ 30 MeV. The listed values correspond to the unoscillated $\beta_0$ DSNB flux spectra using the SFH from Mathews et al. (2014) with its associated ±1σ uncertainty. In parentheses, the values for the case of a complete flavor swap ($\beta_0 \leftrightarrow \bar{\beta}_0$) are provided as well as the results for the SFH according to the EBL reconstruction model by Fermi-LAT Collaboration et al. (2018) and for the SFH of Malou & Dickinson (2014). The one-sided error intervals of $E_0$ in the cases with the SFH from Mathews et al. (2014) are caused by the fact that the functional fits to the SFH scale slightly differently with redshift (see footnote 12), with the best-fit case by Mathews et al. (2014) yielding the largest relative contribution from high-redshift regions and thus the smallest value of $E_0$ compared to both the +1σ and the −1σ limits.

degenerate ONeMg cores that produce ECSNe (Podsiadlowski et al. 2004).

In addition to these uncertain SN progenitors, three channels are being discussed that may lead to preferentially rather LM NSs, whose formation might contribute to the DSNB: Electron capture–initiated collapse may also occur when an ONeMg WD is pushed beyond the Chandrasekhar mass limit due to Roche-lobe overflow from a companion. Such an NS-forming event is referred to as a NIC (see, e.g., Bailyn & Grindlay 1990; Nomoto & Kondo 1991; Ivanova & Taam 2004; Hurley et al. 2010; Jus et al. 2016; Wu & Wang 2018; Ritter et al. 2019). Similarly, Saio & Nomoto (1985) suggested the NIC of two WDs as another possible scenario for the formation of a single NS (also see Ivanova et al. 2008; Schwab et al. 2016; Ritter et al. 2019). Moreover, close-binary interaction might in some cases lead to the stripping of a star’s hydrogen envelope and (most of) its helium envelope onto a companion NS, leaving behind a bare carbon–oxygen core (Nomoto et al. 1994; Dewi et al. 2002) and causing subsequent iron-core collapse. The explosion of such ultrastripped SNe (Tauris et al. 2013, 2015; Suwa et al. 2015; Müller et al. 2018) is discussed as the most likely evolutionary pathway leading to the formation of double-NS systems (Tauris et al. 2017; Mandel et al. 2021).

Previous works (e.g., Mathews et al. 2014; Horiuchi et al. 2018) have considered the contribution from ECSNe to the DSNB flux, and in a footnote, Lien et al. (2010) have already mentioned that, to a minor degree, neutrinos from the AIC of WDs might also add to the DSNB.

In our study we explore the consequences of additional formation channels of (rather) LM NSs on our DSNB predictions in a quantitative and systematic way, subsuming the possible contributions from ultrastripped SNe, NIC events, and MIC events in addition to the contribution from ECSNe that is included in our standard models. To this end, we employ a generic neutrino spectrum ($dN_{\nu}/dE$') adopted from the ECSN calculations of Hüdepohl et al. (2010, "model S") since neutrino signals from sophisticated long-time simulations of AIC, MIC, and ultrastripped SNe are still lacking. We expect the neutrino emission properties of all three additional formation channels of LM NSs to be fairly similar to the case of ECSNe. Our approach is therefore meant to serve as an order-of-magnitude estimate, but it cannot capture any details connected to differences in the individual event rates and in the neutrino signals of the three channels of ultrastripped SNe, NIC events, and MIC events, which we combine to a single, additional LM NS formation component.
It should be mentioned here that the cosmic rates of such events are highly uncertain, because a large parameter space in the treatment of binary interaction (especially common-envelope physics) makes precise predictions difficult. Using population synthesis methods, Zapartas et al. (2017) found that core-collapse events in binary systems are generally delayed as compared to those of single stars. More particularly, Ruiter et al. (2019) showed that AIC and MIC can proceed in various evolutionary pathways, featuring a variety of delay times (from below $10^2$ Myr up to over $10^7$ yr) between starburst and eventual stellar collapse. For simplicity, we thus explore on the one hand different values of the comoving rate density, $R_{LM}(z) = R_{LM}$, which does not change with cosmic time (“LM$_{com}$”). On the other hand, we examine how our DSNB results differ in the case of an evolving rate for additional LM NS formation events (“LM$_{evolv}$”). The DSNB flux spectrum (Equation (9)) can be rewritten in the generalized form

$$\frac{d\Phi}{dE} = \frac{c}{H_0} \int_0^5 dz \frac{R_{CC}(z) \frac{dn_{CC}}{dz} + R_{LM}(z) \frac{dn_{LM}}{dz}}{\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}}. \quad (11)$$

In the lower right panel of Figure 5, we separately plot our fiducial DSNB prediction (dashed line; see Section 4) and the additional contribution from LM events for four different constant rate densities $R_{LM}$ (solid lines), which we take as multiples of the local stellar core-collapse rate, $R_{CC}(0) = 8.93 \times 10^{-3}$ Mpc$^{-3}$ yr$^{-1}$. However, since $R_{CC}(z)$ varies strongly with redshift (it increases by over an order of magnitude from $z = 0$ to $z = 1$), we also consider the ratio of the comoving rate densities of LM NS formation events to those of “conventional” core-collapse SNe, both integrated over the cosmic history:

$$\chi = \frac{\int_0^5 dz R_{LM}(z) \frac{dt_c}{dz}}{\int_0^5 dz R_{CC}(z) \frac{dt_c}{dz}}. \quad (12)$$

This serves as a measure of the relative importance of how both types of neutrino sources contribute to the DSNB from the time of the highest considered redshifts ($z_{max} = 5$) to the present day. In Table 4, we show the ratios $\chi$ for different energy intervals $[E_1, E_2]$ for our different choices of $R_{LM}$. To see an effect of at least 10% within the detection window (10–30 MeV), an additional (constant) LM rate $R_{LM} = 1.55 \times 10^{-4}$ Mpc$^{-3}$ yr$^{-1}$ is required, which is nearly twice the local stellar core-collapse rate, $R_{CC}(0)$, and corresponds to $\chi = 0.20$. Such a fraction is well above present estimates for both AIC/MIC events (Metzger et al. 2009; Ruiter et al. 2019) and ultrastripped SNe (Tauris et al. 2013) of at most a few percent of the “conventional” core-collapse SN population. However, due to large uncertainties in the physics of binary interaction, the possibility of such a large proportion of LM NS formation events may not be ruled out completely.

As a sensitivity check, we additionally consider a comoving rate density, $R_{LM}(z)$, that linearly increases by a factor of 4 between $z = 0$ and $z = 1$ and stays constant at even larger redshifts, roughly following the observationally inferred rate of type Ia SNe (e.g., Graur et al. 2011). In the lower right panel of Figure 5, the LM flux contribution resulting from such an evolving rate is indicated by the pale red band (defined by $0.11 \leq \chi \leq 1.00$, as for the four cases of constant rates $R_{LM}$). The spectra are shifted toward lower energies, as expected from the relatively increased contribution from events at high redshifts. This can also be seen in Table 4 (values in parentheses). An enhancement of the DSNB flux by 10% at energies above 10 MeV would even require $\chi = 0.26$ for an evolving LM rate, which would mean that, for example, roughly more than half of WD mergers lead to NS formation instead of a type Ia SN, if merging WDs explain the majority of SN Ia events. Although this seems to be disfavored on grounds of current observations and population synthesis models (e.g., Metzger et al. 2009; Ruiter et al. 2019), it might not be entirely impossible. Nevertheless, within the relevant detection window, the contribution from AIC/MIC events and ultra-stripped SNe to the DSNB is likely to be hidden by the current uncertainty of the cosmic core-collapse rate (gray-shaded band in Figure 5). Only when this dominant uncertainty is reduced significantly may there be a chance to uncover a contribution to the neutrino background from such LM NS–forming events.

### Table 4: DSNB Contribution from Additional LM NS Formation Events

| Energy Range | Flux (Flux Units) |
|--------------|-------------------|
| (0–10) MeV   | 22.7 cm$^{-2}$ s$^{-1}$ |
| (10–20) MeV  | 5.4 cm$^{-2}$ s$^{-1}$ |
| (20–30) MeV  | 0.6 cm$^{-2}$ s$^{-1}$ |
| (30–40) MeV  | 0.1 cm$^{-2}$ s$^{-1}$ |
| (0–40) MeV   | 28.8 cm$^{-2}$ s$^{-1}$ |

Note. First row: DSNB $\Phi$-flux for our fiducial model (with $R_{LM} = 0$; see Table 2), integrated over different energy intervals. Rows 2–5: Flux contributions for four different choices of the constant (LM$_{com}$) rate density $R_{LM}$. In parentheses, the values of $x$ for an evolving LM NS formation rate (LM$_{evolv}$) with the same value of $\chi$ (Equation (12)) are given (see main text for details).

5.3. Inclusion of Binary Models

A large fraction of massive stars are expected to undergo binary interaction with a companion, possibly shedding their hydrogen envelopes (e.g., via Roche-lobe overflow or common-envelope ejection) and leaving behind bare helium stars (Sana et al. 2012). Taking this as a motivation, we explore how the inclusion of binary models affects our DSNB
predictions. To this end, we employ a set of 132 helium stars with initial masses in the range of $2.5-40\, M_\odot$, originating from hydrogen burning in nonrotating, solar-metallicity stars, as reported in Woosley (2019). According to Equations (4) and (5) therein, this range of initial helium-core masses converts to ZAMS masses of $13.5-91.7\, M_\odot$. Stars with masses lower than those are assumed to form WDs, and thus do not contribute to the DSNB.\textsuperscript{13} For the details of the pre-SN evolution (which includes wind mass loss), the reader is referred to Woosley (2019).

We use these progenitor models and perform SN simulations with the PROMETHEUS-HOTB code as done for single-star progenitors (see Section 2.2). A detailed and dedicated analysis of the explosions of these helium stars can be found in a recent paper by Ertl et al. (2020). In Figure 6, we show, for engine model Z9.6 and W18, the landscape of NS and BH formation events with the basic properties of the neutrino emission of relevance for our DSNB calculations. Compared to Figure 2, the range of stars experiencing core collapse is shifted toward higher ZAMS masses, starting only at $13.5\, M_\odot$. Moreover, there are no cases of BH formation below a ZAMS mass of $33\, M_\odot$. This can be understood as a consequence of mass loss by stellar winds during the pre-SN evolution of the helium stars, yielding less compact cores compared to stars that still possess their hydrogen envelopes (see Figures 1 and 10 in Woosley 2019).

We note in passing that the values for the neutrino energy loss used in our present study differ in details from the numbers shown in Figure 5 of Ertl et al. (2020). First, we do not consider the additional neutrino energy loss from fallback accretion and consistently treat fallback SNe as NS formation events, whereas Ertl et al. (2020) took fallback into account in their estimates of the compact remnant masses and the associated release of gravitational BE through neutrinos.\textsuperscript{14} Second, in our present study we extrapolate the neutrino emission of non-explooding cases until the accreting PNS in our PROMETHEUS-HOTB runs reaches the assumed and parametrically varied baryonic mass limit of stable, cold NSs, $M_{\text{lim}}^{\text{NS,b}}$, and therefore collapses to a BH. In contrast, Ertl et al. (2020) employed for BH cases (with $t_{\text{BH}} > 10$ s) the radius-dependent fit formula of Lattimer & Prakash (2001) for the gravitational BE of an NS with the maximum mass assumed in their work. The energy release ($\epsilon_{\text{nu}}^{\text{lim}}$) estimated that way is somewhat larger than our accretion-determined estimates (see Figure B2 in Appendix B).

Figure 7 illustrates how the inclusion of binary models impacts our DSNB predictions. In the left panel, we separately show the contributions from successful and failed explosions to the DSNB flux spectrum of electron antineutrinos for our fiducial model parameters ($Z9.6$ and W18; $M_{\text{lim}}^{\text{NS,b}} = 2.7\, M_\odot$; best-fit $\alpha$; SFH from Mathews et al. 2014), assuming that all (100\%) of progenitors evolve as helium stars (“W18-BH2.7-\alpha2.0-He100”). Compared to single stars (Figure 3 and black dashed line in the right panel of Figure 7), the overall DSNB flux is reduced by a factor of ~2 owing to the smaller fraction of stars experiencing core collapse. At the same time, the less frequent failed explosions produce a lower high-energy tail of the spectrum compared to our fiducial DSNB spectrum based

\textsuperscript{13} Consistently, the lower integration bounds in Equations (6) and (7) are raised from $8.7\, M_\odot$ to $13.5\, M_\odot$.

\textsuperscript{14} Note that five such progenitors, which explode at relatively late times ($\sim2$ s) and consequently reach high PNS masses, are treated in this work either as “normal” successful SNe (without fallback) or, if the PNS mass exceeds $M_{\text{lim}}^{\text{NS,b}}$ at any time during the post-bounce evolution, as failed explosions. In the latter case, the neutrino signals are truncated at this time, $t_{\text{diff}}$. 

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{landscape.png}
\caption{Figure 6. Landscape of NS or BH formation for the set of helium-star progenitors from Woosley (2019) as obtained in simulations with the engine model Z9.6 and W18 (see Figure 2 for single-star progenitors). From top to bottom: Time of explosion or BH formation, total energy radiated in all species of neutrinos, and mean neutrino energy of electron antineutrinos vs. ZAMS mass of the progenitors. Note the different mass range for stellar core-collapse progenitors compared to Figure 2. Red bars indicate successful SN explosions (and fallback SNe), while the outcomes of BH-forming, failed SNe are shown for the different baryonic NS mass limits in gray ($2.3\, M_\odot$), dark blue ($2.7\, M_\odot$), light blue ($3.1\, M_\odot$), and cyan ($3.5\, M_\odot$). Five special progenitors yield successful or failed explosions depending on the NS mass limit (see footnote 14).}
\end{figure}
on single stars (see also Table 3). If we assume that only 33% of all massive stars strip their hydrogen envelopes (as suggested by Sana et al. 2012; “W18-BH2,7-α2.0-He33”), the effects of helium stars on the DSNB spectrum are less dramatic, and the shifted spectrum lies within the uncertainty band associated with the SFH (gray-shaded band; see Figure 4).

Applying the other neutrino engines considered in our work to the helium-star models, we obtain relative changes of the DSNB spectra similar to those in the case of our fiducial engine model Z9.6 and W18. We should stress at this point that if progenitors do not lose their entire hydrogen envelopes, end stages of stellar evolution more similar to those of single-star evolution can be expected (Woosley 2019).

### 6. DSNB Spectrum of Electron Neutrinos

Although the main focus of our study lies on the DSNB’s $\bar{\nu}_e$ component, we briefly comment on the flux spectrum of $\nu_e$, which is an observational target of DUNE (DUNE Collaboration et al. 2015). Combining future DSNB $\nu_e$-flux measurements by DUNE with the $\nu_e$-flux data gathered by the gadolinium-loaded SK (and Hyper-Kamiokande) and by JUNO will yield complementary constraints on the DSNB precursor space (see, e.g., Möller et al. 2018) and will help in testing different neutrino oscillation scenarios or nonstandard-model physics, such as neutrino decays (see, e.g., Fogli et al. 2004; de Gouvêa et al. 2020; Tabrizi & Horiuchi 2020).

In Figure 8, we show our predictions for the DSNB flux spectrum for $\nu_e$ in comparison to the $\bar{\nu}_e$ component for our fiducial model parameters (Z9.6 and W18 neutrino engine; $M_{\text{lim}}^{\text{He}} = 2.7 M_\odot$; $\alpha_{\text{BH}} = 2.0$; best-fit SFH from Mathews et al. 2014). The main differences are a more prominent spectral peak (at energies $E \lesssim 8$ MeV) and a faster decline of the spectrum toward high neutrino energies for the case of $\nu_e$ compared to $\bar{\nu}_e$. The exponential fit of Equation (10) yields a value of the “slope parameter” $E_0$ of 4.49 MeV for the $\nu_e$ spectrum (compared to $E_0 = 4.82$ MeV for $\bar{\nu}_e$). This is a consequence of generally lower mean neutrino energies of $\nu_e$ compared to $\bar{\nu}_e$ (see Table E1 in Appendix E). Note that for $\nu_e$ a DSNB detection will not be possible below $\sim 17$ MeV due to the overwhelming solar hep (and $^8$B) neutrino flux (see, e.g., Figure 8 of Zhu et al. 2019).

To give an impression of the spectral DSNB variability for $\nu_e$, we also show the $\nu_e$-flux spectra for models with different neutrino engines applied and thus varied fractions of failed SNe with BH formation (left panel), as well as for a model that includes 33% hydrogen-stripped helium-star progenitors (as suggested by Sana et al. 2012; right panel). The overall trends (i.e., enhanced high-energy tail of the DSNB spectrum for a larger fraction of failed SNe and reduced DSNB flux for inclusion of helium stars) are similar to the case of $\bar{\nu}_e$.

### 7. Neutrino Flavor Conversions and Remaining Uncertainties

#### 7.1. Neutrino Flavor Conversions

So far we have not taken neutrino flavor oscillations into account but have identified the emission of electron antineutrinos (or neutrinos) by the considered astrophysical sources with the measurable DSNB flux of $\bar{\nu}_e$ (or $\nu_e$). However, on their way out of a collapsing star, neutrinos (and antineutrinos) undergo collective and matter-induced (MSW) flavor conversions (Wolfenstein 1978; Mikheyev & Smirnov 1985; Duan et al. 2010; Mirizzi et al. 2016). Hereafter, we discuss how such oscillations can affect our DSNB flux predictions.

Following Chakraborty et al. (2011) and Lunardini & Tamborra (2012), we write the DSNB flux spectrum of electron antineutrinos after including the effect of flavor conversions as

$$d\Phi_{\bar{\nu}_e}/dE = \tilde{\rho} \frac{d\Phi_{\bar{\nu}_e}^0}{dE} + (1 - \tilde{\rho}) \frac{d\Phi_{\nu_e}^0}{dE}, \quad (14)$$

where $d\Phi_{\bar{\nu}_e}^0/dE$ and $d\Phi_{\nu_e}^0/dE$ are the unoscillated spectra for electron antineutrinos ($\bar{\nu}_e$) and a representative heavy-lepton neutrino ($\nu_e$). $\tilde{\rho} \simeq 0.7$ or $\tilde{\rho} \simeq 0$ denotes the survival probability of $\bar{\nu}_e$ in cases of normal (NH) or inverted (IH)
mass hierarchy, respectively. Recently, Møller et al. (2018) confirmed by numerically solving the neutrino kinetic equations of motion that (matter-induced) neutrino flavor conversions can be well approximated by the simplified analytic expression of Equation (14) for the small set of PROMETHEUS-VERTEX simulations that they used in their study and that we also employ in our work as reference cases to calibrate some degrees of freedom in our modeling approach (see Table C1 in Appendix C). We already mentioned earlier that the large sets of core-collapse simulations underlying our DSNB calculations do not provide reliable information on the heavy-lepton neutrino source emission, which is why we use PROMETHEUS-VERTEX SN and BH formation models to rescale the neutrino energy release in the different neutrino species (see Section 3.1). For the same reason, we also adjust the spectral parameters, \( E_{\nu_0} \) and \( \sigma_{\nu_e} \), of the time-integrated \( \nu_e \) emission (the bar in the symbol \( \sigma_{\nu_e} \) indicates that the shape parameter refers to the time-integrated spectrum rather than the instantaneous spectrum), guided by the sophisticated PROMETHEUS-VERTEX models listed in Table C1, to get a useful representation of the unoscillated DSNB spectrum of heavy-lepton neutrinos, \( d\Phi_{E_0}^{\nu_e}/dE \) (see Appendix C for details).

In the left panel of Figure 9, we show our unoscillated, fiducial DSNB spectrum for \( \bar{\nu}_e \), \( d\Phi_{E_0}^{\bar{\nu}_e}/dE \) (black dashed line), and the corresponding unoscillated DSNB spectrum for \( \nu_e \), \( d\Phi_{E_0}^{\nu_e}/dE \) (red solid line), for our fiducial model parameters (see Section 4). According to Equation (14), the latter represents the case of IH, where a complete flavor swap \( \bar{\nu}_e \leftrightarrow \nu_e \) takes place. If, instead, the case of NH is realized in nature, an outcome between the two plotted extremes can be expected. The uncertainty arising from the cosmic core-collapse rate (corresponding to the \( \pm 1\sigma \) interval of the SFH from Mathews et al. 2014) is indicated by shaded bands. In Table 5, we additionally provide the integrated \( \bar{\nu}_e \)-flux for different energy intervals and a complete flavor swap \( \bar{\nu}_e \leftrightarrow \nu_e \) in analogy to what is given in Table 2 for the case of no flavor oscillations (\( \bar{\nu}_e \)). The most important difference is a reduced contribution from failed SNe. This can be understood by the small relative fraction of heavy-lepton neutrino emission, \( \bar{\xi}_e \), in our two PROMETHEUS-VERTEX reference models for BH formation, which we employ for our rescaling (Appendix C). At the same time, the contribution from successful explosions (including ECSNe) is largely unchanged, which reflects the approximate flavor equipartition in their neutrino emission.

Despite the less relevant contribution from failed SNe, the slope parameter \( E_0/\text{MeV} \) of the exponential fit of Equation (10) is increased marginally from 4.82 to 4.84 in the case of a complete flavor swap (see Table 3) because smaller values of the spectral-shape parameter \( \sigma_{\nu_e} \) for heavy-lepton neutrinos (see Table C1; \( \lambda_{\nu_e}^{\nu_0} < 1 \)) partly compensate for the reduced flux of \( \nu_e \) in the high-energy region associated with BH cases. The mean energies of the time-integrated neutrino signals are fairly similar for \( \nu_e \) and \( \bar{\nu}_e \) (see Table C1; \( \lambda_{\nu_e}^{\nu_0} \sim 1 \)), as suggested by state-of-the-art simulations (e.g., Marek et al. 2009; Møller & Janka 2014) and as a consequence of the inclusion of energy transfers (non-isooenergetic effects) in the neutrino–nucleon scattering reactions (see Keil et al. 2003; Hudevågh 2014). In conflict with this result of modern SN models with state-of-the-art treatment of the neutrino transport, several previous DSNB studies have employed spectra with \( E_{\nu_0} \) being considerably higher than \( \langle E_{\nu_0} \rangle \) (particularly for the emission from failed explosions).

In line with recent studies by Priya & Lunardini (2017) and Møller et al. (2018), we find that neutrino flavor conversions have a fairly moderate influence on the DSNB (for \( \bar{\nu}_e \)), which is well dominated by other uncertainties. Nonetheless, for our highest-flux models (with a weak central engine and a high maximum NS mass), which possess a large DSNB contribution
from BH-forming events, the oscillation effects become more pronounced. We will further comment on this in Section 8.1.

For the DSNB $\nu_e$ flux the effects of neutrino flavor oscillations can be described in an analogous manner (see, e.g., Chakraborty et al. 2011; Lunardini & Tamborra 2012). In the most extreme case of NH (and purely MSW-induced flavor conversions), a complete flavor swap ($\nu_e \leftrightarrow \nu_x$) can take place, whereas for IH a measurable DSNB $\nu_x$-flux spectrum between the unoscillated spectra of $\nu_e$ and $\nu_x$ can be expected.

7.2. Tests of Remaining Uncertainties

As we point out in Appendix B, the total radiated neutrino energies ($E_{\nu}^{\text{tot}}$) of our successful SNe might, on average, be overestimated by a few percent, whereas the neutrino emission from failed explosions could be slightly underestimated in our modeling approach for the neutrino signals. In the right panel of Figure 9, we therefore compare our fiducial DSNB prediction (black dashed line) with a spectrum where the $E_{\nu}^{\text{tot}}$ of all exploding progenitors is reduced by 15% (lower edge of the red band). This choice of reduction is guided by a comparison of $E_{\nu}^{\text{tot}}$ with the gravitational BEs (BE12) of the corresponding NS remnants (Equation (B1) with $R_{\text{NS}} = 12$ km; see Figure B1 and Table B1), consistent with the cold-NS radius suggested by recent astrophysical observations and constraints from nuclear theory and experiments (see footnote 4). Analogously, the upper edge of the red band in Figure 9 indicates a model where the $E_{\nu}^{\text{tot}}$ of all failed explosions is increased by 15%. This case is motivated by the circumstance wherein the maximum neutrino emission in our failed-SN models with late BH formation lies $\sim$10%–20% below the maximally available gravitational BE according to Equation (B1) of an NS at its mass limit (see Figure B2 and Table B2). Any mix of changes of the NS and BH energy release will lead to intermediate results. Note that the corresponding red uncertainty band is hardly visible on the logarithmic scale.

A somewhat stronger effect can be seen when we vary the mean energies, $\langle E \rangle$, of the time-integrated spectra by $-10\%$ for successful SNe or by $+10\%$ for failed explosions (lower or upper edge, respectively, of the blue shaded band). Particularly at high energies, the spectra fan out noticeably. Such an uncertainty range cannot be ruled out according to present knowledge. Again, changing $\langle E \rangle$ for both successful and failed SNe at the same time yields a result in between the given limits. In Appendix D, we show that the outcome of our simplified approach is in reasonable overall agreement with results from the sophisticated PROMETHEUS-VERTEX simulations; however, the mean energies of the time-integrated spectra do not match perfectly (they lie $\sim 1$ MeV higher/lower than in the VERTEX models for successful/failed SNe; see Figures D1 and D2). Besides this fact, we should emphasize that the neutrino emission characteristics depend considerably on the still incompletely known high-density EoS (e.g., Steiner et al. 2013; Schneider et al. 2019) and also on the effects of muons, which have been neglected in most previous core-collapse models but can raise the mean energies of the radiated neutrinos (Bollig et al. 2017).

Despite these uncertainties associated with the neutrino source, the cosmic core-collapse rate $R_{\text{CC}}$ still constitutes the largest uncertainty affecting the DSNB, especially at lower energies (see Figure 4). Accordingly, the gray-shaded band in the right panel of Figure 9 indicates the $\pm 1\sigma$ variation of $R_{\text{CC}}$ for the SFH from Mathews et al. (2014). Upcoming wide-field surveys such as LSST (Tyson 2002) should be able to pin down the visible SN rate (below redshifts of $z \sim 1$) to good accuracy, opening up a chance for DSNB measurements to specifically probe the contribution from faint and failed explosions (Lien et al. 2010).

Finally, one should keep in mind that we only employ solar-metallicity progenitor models in our simulations. Obviously, this is a simplification, because the distribution of metals in the universe is spatially nonuniform (see, e.g., the low metallicities in the Magellanic Clouds) and evolves with cosmic time. Since the fraction of failed explosions depends on metallicity (e.g., Woosley et al. 2002; Heger et al. 2003; Langer 2012), Nakazato et al. (2015) and Yüksel & Kistler (2015) considered a failed-SN fraction that increases with redshift. On the other hand, Panter et al. (2008) suggested that the average metallicity

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**Figure 9.** Effects of neutrino flavor conversions on the DSNB flux spectrum and remaining modeling uncertainties for our fiducial model parameters (see Section 4). The left panel shows the unoscillated DSNB spectrum of electron antineutrinos ($d\Phi^{\bar{\nu}_e}_{\nu}/dE$; black dashed line) and the predicted DSNB spectrum for one species of heavy-lepton neutrinos ($d\Phi^{\bar{\nu}_x}_{\nu}/dE$; red solid line), which would become the measurable $\bar{\nu}_e$ spectrum in the case of a complete flavor swap $\bar{\nu}_e \leftrightarrow \bar{\nu}_x$ (see Equation (14)). The uncertainty arising from the cosmic core-collapse rate $R_{\text{CC}}$ (represented by the $\pm 1\sigma$ limits to the SFH from Mathews et al. 2014) is indicated by shaded bands. In the right panel, our fiducial model (black dashed line; unoscillated $\bar{\nu}_e$) is compared to DSNB flux spectra where the total radiated neutrino energy, $E_{\nu}^{\text{tot}}$, is reduced by 15% for successful SNe or increased by 15% for BH formation cases. The corresponding red band is partly covered by the blue band, which marks the DSNB variation when the time-integrated mean $\bar{\nu}_e$ energies, $\langle E \rangle$, are shifted by $-10\%$ for successful SNe or by $+10\%$ for failed explosions (see main text for details). Changing $E_{\nu}^{\text{tot}}$ or $\langle E \rangle$ for successful and failed SNe at the same time yields spectra within the uncertainty bands shown. The uncertainty of the fiducial spectrum due to $R_{\text{CC}}$ is indicated by the gray band.
Figure 10. Comparison of our most extreme DSNB predictions with the upper flux limit from SK: $\Phi_{>17.3} \equiv \Phi(E > 17.3 \text{ MeV}) \lesssim (2.8$–$3.1) \text{ cm}^{-2} \text{s}^{-1}$ (Bays et al. 2012). The shaded bands in the left panel show the spread between the flux spectra $d\Phi/dE$ of electron antineutrinos, resulting from various combinations of the source parameters considered in Section 5.1 (see Figure 5). Our fiducial model (W18-BH2.7-$\alpha$2.0; Section 4) is displayed as a dashed line. To guide the eye, we discriminate the approximate ranges for models that yield an integrated flux $\Phi_{>17.3}$ below 3.1 cm$^{-2}$ s$^{-1}$ (gray) or exceed this limit (red); see the main text for details. As in the previous figures, vertical bands frame the approximate detection window. In the right panel, $\Phi_{>17.3}$ is shown for a selection of models (including our fiducial case; black cross) that reach close to or beyond the SK limit (light and dark shaded regions for 2.8 and 3.1 cm$^{-2}$ s$^{-1}$, respectively) as a function of the fit parameter $E_0$ (Equation (10)). Both vertical and horizontal error bars indicate the uncertainty connected to the cosmic SFH (+1σ limits of Mathews et al. 2014). The one-sided horizontal error intervals are caused by the fact that the functional fits to the SFH scale slightly differently with redshift (see footnote 12), with the best-fit case by Mathews et al. (2014) yielding the largest relative contribution from high-redshift regions and thus the smallest value of $E_0$ compared to both the +1σ and the −1σ limits.

### Table 5

| Energy Interval | Total DSNB Flux (x) |  (0–10) MeV | (10–20) MeV | (20–30) MeV | (30–40) MeV | (0–40) MeV |
|----------------|---------------------|-------------|-------------|-------------|-------------|------------|
|                | cm$^{-2}$ s$^{-1}$  | 19.3        | 4.3         | 0.5         | 0.1         | 24.2       |
| ECSNe          | 3.0%                | 1.5%        | 0.8%        | 0.4%        | 2.7%        |
| Iron-core SNe  | 68.4%               | 65.4%       | 52.3%       | 37.9%       | 67.5%       |
| Failed SNe     | 28.6%               | 33.2%       | 47.0%       | 61.8%       | 29.9%       |

Note. First row: DSNB $\bar{\nu}_e$-flux for the case of a complete flavor swap ($\bar{\nu}_e \leftrightarrow \bar{\nu}_\mu$), integrated over different energy intervals. Rows 2–4: Relative contributions from the various source types (ECSNe/iron-core SNe/failed SNe). Our fiducial model parameters (Z9.6 and W18; $M_{\nu_{\text{NS,b}}} = 2.7 \ M_\odot$; best-fit $\alpha$) are used. Compare with Table 2, where values for the unoscillated $\bar{\nu}_e$-flux are provided.

does not decline dramatically up to $z \sim 2$. Assuming solar metallicity should therefore be a sufficiently good approximation, in view of the fact that the DSNB flux in the energy window favorable to DSNB detection is produced almost entirely by sources at moderate redshifts (see Figure 3).

At this point we should also remind the reader that a core-collapse SN is an inherently multidimensional phenomenon (see, e.g., Müller 2016). While our simplified 1D approach should be able to capture the overall picture of the progenitor-dependent neutrino emission, the increasing number of fully self-consistent 3D simulations will have to validate our results eventually.

### 8. Comparison with the SK Flux Limits and Previous Works

#### 8.1. Comparison with the SK Flux Limits

After discussing the dependence of the predicted DSNB spectrum on different inputs in Sections 5 and 7, we compare our results now with the most stringent $\bar{\nu}_e$-flux limit set by the SK experiment (Bays et al. 2012): $\Phi_{>17.3} \equiv \Phi(E > 17.3 \text{ MeV}) \lesssim (2.8$–$3.1) \text{ cm}^{-2} \text{s}^{-1}$.

The various parameter combinations considered in our study lead to a wide spread between the DSNB flux spectra, as can be seen in the left panel of Figure 10. At high energies, the spectral tails of our different models fan out over more than an order of magnitude, with our most extreme cases yielding an integrated flux $\Phi_{>17.3}$ that clearly exceeds the SK limit. To guide the eye, we roughly mark the region of such disfavored models (with $\Phi_{>17.3} \gtrsim 3.1 \text{ cm}^{-2} \text{s}^{-1}$) by a red shaded band, while flux spectra with $\Phi_{>17.3} \lesssim 3.1 \text{ cm}^{-2} \text{s}^{-1}$, including our fiducial prediction (dashed line; see Section 4), lie in the gray band. We take the specific model “W20-BH3.5-$\alpha$2.0” (i.e., Z9.6 and W20 neutrino engine, $M_{\nu_{\text{NS,b}}} = 3.5 \ M_\odot$; $\alpha_{\text{BH}} = 2.0$) with the best-fit parameters taken for the SFH from Mathews et al. (2014) as a bounding case; it yields an integrated flux $\Phi_{>17.3} = 3.09 \text{ cm}^{-2} \text{s}^{-1}$, just within the uncertainty range of the SK limit (2.8–3.1 cm$^{-2}$ s$^{-1}$). We should emphasize, however, that this does not define a rigorous borderline, since spectra with quite different values of the slope parameter $E_0$ (Equation (10)) can yield similar integrated fluxes in the energy range above 17.3 MeV.

In the right panel of Figure 10, we therefore plot $\Phi_{>17.3}$ as a function of the fit parameter $E_0$ for a selection of models reaching close to (or beyond) the SK bound, which is marked by the red shaded region (with its uncertainty indicated by two slightly shifted lines). This plot is also intended for facilitating comparison with other works (see, e.g., Table 1 in Lunardini & Peres 2008 and Figure 19 in Bays et al. 2012). The tendency of greater integrated fluxes $\Phi_{>17.3}$ for higher values of $E_0$ is obvious, yet there is significant scatter. In particular, the large uncertainty connected to the cosmic core-collapse rate ($\pm 1\sigma$
interval from Mathews et al. 2014, indicated by error bars) impedes definite conclusions. Nonetheless, models of our study with the most extreme combinations of parameters, such as the different cases of W20-BH3.5, which possess a strong contribution from failed SNe and thus large values of $E_{\nu}$ (see Section 5.1 and Table 3), are already disfavored, because their fluxes $\Phi_{\nu,17.3}$ reach beyond the SK limit (unless a minimal $R_{\text{CC}}$ is taken). Also a less extreme value of the NS mass limit or a neutrino energy with a lower fraction of BH formation events can lead to an integrated flux close to the SK bound: models W20-BH3.1-0.1 and W15-BH3.5-0.1 (not shown in Figure 10) yield $\Phi_{\nu,17.3} = 2.7 \pm 0.3$ cm$^{-2}$ s$^{-1}$ and $\Phi_{\nu,17.3} = 2.6 \pm 0.5$ cm$^{-2}$ s$^{-1}$, respectively, with a dominant fraction (85% and 81%, respectively) of $E_{\nu}$ above 17.3 MeV originating from BH formation events. In Table 6, we provide the total integrated fluxes ($\Phi_{\text{tot}}$), the fluxes within the observational window of 10–30 MeV ($\Phi_{10-30}$), and the flux integrals above 17.3 MeV ($\Phi_{\nu,17.3}$) for a subset of our DSNB models.

Unlike the experimental DSNB flux limits of Malek et al. (2003), those provided by Bays et al. (2012) depend on the DSNB model employed. However, for an energy threshold close to $\sim 20$ MeV, the flux limits are rather insensitive to the shape of the DSNB spectrum as pointed out by Lunardini & Peres (2008). In any case, the (Fermi–Dirac) spectral temperatures (3 MeV $\leq T_{\nu} \leq 8$ MeV) that Bays et al. (2012) considered for their modeling of a “typical” SN source spectrum lead to DSNB spectra with slope parameters $E_0$ that cover the range of values obtained in our work.\footnote{In an analytic study, Lunardini (2007) showed that the spectral temperature of the employed SN source spectrum (before integration over redshifts) translates into the slope parameter $E_0$ of the DSNB spectrum up to some tens of percent.} Repeating their analysis of computing upper DSNB flux limits with our DSNB models should therefore lead to comparable bounds. Instead, we simply compare the experimental flux limit of (2.8–3.1) cm$^{-2}$ s$^{-1}$ to a subset of our model predictions in Figure 10. Naturally, this cannot replace a sophisticated statistical analysis, which is beyond the scope of this work.

Our fiducial model (W18-BH2.7-0.2) yields an integrated flux of $\Phi_{\nu,17.3} = 1.3^{+1.1}_{-0.9}$ cm$^{-2}$ s$^{-1}$, which is just below the SK bound, possibly not even by a factor of 2. Intriguingly, Bays et al. (2012) pointed out that there might already be a hint of a signal in the SK-II and SK-III data, giving hope that the first detection of the DSNB is within close reach now (see Beaumont & Vagins 2004; Yüksel et al. 2006; Horiiuchi et al. 2009; Keen & Lunardini 2012; An et al. 2016; Priya & Lunardini 2017).

Since the SN neutrino emission is different for heavy-lepton neutrinos compared to electron antineutrinos (see Section 7.1), a complete (or partial) flavor swap ($\nu_{\ell} \leftrightarrow \bar{\nu}_{\ell}$) would affect our

### Table 6

| Model                      | $\Phi_{\text{tot}}$ (cm$^{-2}$ s$^{-1}$) | $\Phi_{10-30}$ (cm$^{-2}$ s$^{-1}$) | $\Phi_{\nu,17.3}$ (cm$^{-2}$ s$^{-1}$) |
|----------------------------|----------------------------------------|------------------------------------|---------------------------------------|
| W18-BH2.7-0.2 (fiducial)   | 24.3 $^{+2.9}_{-2.4}$                  | 21.7 $^{+2.4}_{-2.3}$               | 17.4 $^{+2.1}_{-2.0}$                 |
| W20-BH3.5-0.1 (max.)       | 41.7 $^{+3.5}_{-3.0}$                  | 39.0 $^{+3.3}_{-3.2}$               | 31.9 $^{+3.0}_{-2.9}$                 |
| S19.8-BH2.3-0.3 (min.)     | 24.3 $^{+2.2}_{-1.8}$                  | 22.8 $^{+2.0}_{-1.7}$               | 19.4 $^{+1.8}_{-1.6}$                 |

Note. The given values correspond to the unoscillated $\nu_{\ell}$ DSNB flux spectra using the SFH from Mathews et al. (2014) with its associated $\pm 1\sigma$ uncertainty. In parentheses, the values for the case of a complete flavor swap ($\nu_{\ell} \leftrightarrow \bar{\nu}_{\ell}$) are provided as well as the results for an SFH according to the EBL reconstruction model by Fermi-LAT Collaboration et al. (2018) and for the SFH of Madau & Dickinson (2014).
previous conclusions: In the case of IH (\(\rho \approx 0\)), the integrated flux above 17.3 MeV of our most extreme model (W20-BH3.5-α-1.0) decreases by 39% from 3.5^{+2.0}_{-1.2} \text{ cm}^{-2} \text{s}^{-1} \text{ to } 2.1^{+1.5}_{-0.7} \text{ cm}^{-2} \text{s}^{-1}, which is just below the SK bound (however, note the large uncertainties due to the cosmic core-collapse rate). For the case of NH (\(\rho \approx 0.7\)), we obtain \(\Phi_{\nu,17.3} = 3.1^{+0.5}_{-1.1} \text{ cm}^{-2} \text{s}^{-1}\), which is still somewhat above the SK limit. Applying neutrino flavor conversions to our fiducial model, the effects are much reduced, as described in Section 7.1 (see left panel of Figure 9 and Table 5): \(\Phi_{\nu,17.3}\) decreases by only 7% (21%) from 1.3^{+0.4}_{-0.4} \text{ cm}^{-2} \text{s}^{-1} \text{ to } 1.2^{+0.1}_{-0.3} \text{ cm}^{-2} \text{s}^{-1} \text{ for NH (IH), still reaching close to the SK bound. In Table 6, the first values in parentheses correspond to the case of a complete flavor swap.}

Taking alternative SFHs such as the ones from Madau & Dickinson (2014) or Fermi-LAT Collaboration et al. (2018), which we discussed in Section 4, leads to lower predictions of the DSNB flux compared to our fiducial model, which employs the SFH from Mathews et al. (2014). This implies weaker constraints by the experimental limit, as can be seen in Table 6, where the second and third values in parentheses show the results for an SFH according to the EBL reconstruction by Fermi-LAT Collaboration et al. (2018) and for the SFH from Madau & Dickinson (2014), respectively. Independently of the chosen model parameters, the integrated fluxes are reduced as compared to the cases with the SFH from Mathews et al. (2014). The flux values of \(\Phi_{\nu,17.3}\) for the case of the SFH from Madau & Dickinson (2014) lie about one-third below the ones when taking the best-fit SFH from Mathews et al. (2014; roughly corresponding to their \(-1\sigma\ lower limit case), whereas there is still significant overlap between the flux values for the SFH of Fermi-LAT Collaboration et al. (2018) and our fiducial flux values (also note the large uncertainty ranges). At the same time, the values of the slope parameter \(E_0\) are increased (i.e., the spectral tails are lifted) when taking the SFH of Fermi-LAT Collaboration et al. (2018) or the one from Madau & Dickinson (2014) (see Table 3 and Figure 4). Apparently, the large degeneracy between the parameters entering the flux calculations impedes both precise predictions and the exclusion of models.

8.2. Comparison with Previous Works

Finally, we compare our DSNB flux predictions with the results of other recent works. For instance, Priya & Lunardini (2017) found a \(\Phi_{\nu}\)-flux above 11 MeV in the range of (1.4–3.7) \text{ cm}^{-2} \text{s}^{-1}, with their highest-flux model being a factor of \(\sim 3\) below the SK limit of Bays et al. (2012). In contrast, our fiducial model yields flux values of 4.6^{+3.9}_{-1.3} \text{ cm}^{-2} \text{s}^{-1} (3.9^{+1.3}_{-1.3} \text{ cm}^{-2} \text{s}^{-1}) above 11 MeV in the case where we consider neutrino oscillations for NH (IH) (to follow Priya & Lunardini 2017), reaching very close to the SK bound (see Section 8.1). Likewise, a recent study by Møller et al. (2018) suggested a clearly lower DSNB flux compared to our result (see their Figures 3 and 10). These differences between our DSNB estimates and previous results can be understood by the large variations of the neutrino outputs between the different core-collapse events in our sets of SN and BH formation models, as shown in Figure 2. While progenitors at the low end of the considered ZAMS mass range radiate \(E_{\nu}^{\text{tot}} \approx 2 \times 10^{53} \text{ erg}\), the emission increases to values of \((3–4) \times 10^{53} \text{ erg}\) for progenitors above about (11–12) M\(_\odot\). On the other hand, Priya & Lunardini (2017) and Møller et al. (2018) applied the low-energy neutrino signals (\(E_{\nu}^{\text{tot}} \approx 2 \times 10^{53} \text{ erg}\)) of the s11.2c and s9.6co models they considered for the entire mass interval between \(\sim 8 M_\odot\) and \(\sim 15 M_\odot\), which received a high weight by the IMF in the integration over all core-collapse events. Moreover, both studies made use of failed-SN models that formed BHs relatively quickly (within \(< 2\) s after bounce) and therefore radiated less energy (\(< 3.7 \times 10^{53} \text{ erg}\)) than most of our failed explosions. Each of these two aspects accounts for the reduction of the integral flux by several tens of percent as compared to our work.

Horiiuchi et al. (2018) for the first time employed a larger set of neutrino signals in their DSNB study, including seven models of BH-forming, failed explosions. However, the total neutrino energies \(E_{\nu}^{\text{tot}}\) radiated from their failed SNe are in general below \(\sim 3.5 \times 10^{53} \text{ erg}\) (see their Figure 5). In contrast, we find total neutrino energies in the cases of failed explosions up to \(5.2 \times 10^{53} \text{ erg}\) for an NS mass limit of \(M_{\text{lim,NS,b}} = 2.3 M_\odot\) and of up to \(6.7 \times 10^{53} \text{ erg}\) when using our fiducial value, \(M_{\text{lim,NS,b}} = 2.7 M_\odot\) (see Figures 2 and B2), enhancing the integral flux by a few tens of percent as compared to Horiiuchi et al. (2018). Accordingly, our study suggests that in particular the inclusion of slowly accreting progenitors that lead to late BH formation (not considered in previous works) is responsible for a significant contribution to the DSNB.

9. Summary of Uncertainties

Having discussed numerous dependences of the DSNB, we summarize our main results with their corresponding uncertainties in this section. Again, these uncertainties are considered in reference to our fiducial DSNB spectrum, which is based on the Z9.6 and W18 neutrino engine with 26.9% BH formation cases, a baryonic NS mass limit of 2.7 M\(_\odot\), a value of \(\alpha_{\text{BH}} = 2.0\) for the instantaneous neutrino emission spectrum of failed SNe, no additional contribution from LM NS formation events (i.e., \(\chi = 0\); Equation (12)), only single-star progenitors (i.e., no hydrogen-stripped helium stars), no neutrino flavor oscillations, and the best-fit SFH of Mathews et al. (2014). The corresponding DSNB uncertainties can be grouped into the following four categories:

(1) \textbf{Stellar-diversity uncertainties (see Sections 5.1–5.3; Figures 5, 7): These include the still undetermined fraction of BH-forming stellar core-collapse events; a possible, still poorly understood contribution from LM NS formation events (AIC, MIC, or ultrastripped SNe); and the relative fraction of helium stars, which serve as a proxy for SN progenitors that have stripped their hydrogen envelopes as a consequence of binary interaction at the end of core-hydrogen burning.}

(2) \textbf{Microphysical uncertainties (see Sections 5.1, 7.1; Figures 5, 9): These concern, on the one hand, the still incompletely known high-density EoS of NS matter with the corresponding NS mass limit and, on the other hand, possible effects of neutrino flavor conversions.}

(3) \textbf{Modeling uncertainties (see Sections 5.1, 7.2; Figures 5, 9): These are connected to our numerical description of the neutrino emission from successful and failed SNe. Here we subsume approximations of the spectral-shape parameter (\(\alpha_{\text{BH}}\)) for the instantaneous neutrino emission spectrum, of the total neutrino energy loss from NS and BH formation events, and of the mean energy of the time-integrated \(\Phi_{\nu}\) spectrum.}

(4) \textbf{Astrophysical uncertainties (see Section 4; Figure 4): These refer to the still insufficiently constrained cosmic SFH, for which we test different representations.}
The upper left panel of Figure 11 shows our fiducial DSNB $\bar{\nu}_e$-flux spectrum with its major uncertainties stacked on top of each other (shaded/hatched bands): the failed-SN (fSN) fraction (17.8%–41.7% of core-collapse progenitors depending on the strength of the neutrino engine); the NS baryonic mass limit (2.3–3.5 $M_\odot$); the instantaneous spectral-shape parameter for the emission from failed SNe ($1.0 < \alpha_{BH} < 3.0$); and the uncertainty connected to the cosmic SFH (±1σ limits of Mathews et al. 2014). The resulting “total” uncertainty band is the same as that in Figure 10. The lower left panel and the right panels show the residuals of our DSNB models where only one parameter is changed relative to the fiducial model, while all other parameters are kept at their default values, grouped into stellar-diversity uncertainties, microphysical uncertainties, modeling uncertainties, and astrophysical uncertainties (see Section 9 for more details). L_{core} and L_{evol} denote cases where the rate densities of LM NS formation events (AIC, MIC, ultrastripped SNe) are constant and evolve with redshift, respectively (both for a value of $\chi = 0.34$, which corresponds to a relative abundance of LM NS formation events of 34% as compared to “conventional” core-collapse SNe plus failed SNe; Equation (12)). In each panel, gray-shaded vertical bands frame the approximate detection window.

Figure 11. Overview of DSNB uncertainties. The upper left panel shows the $\bar{\nu}_e$-flux spectrum, $d\Phi/dE$, of our fiducial DSNB model (dashed line) together with its major uncertainties stacked on top of each other (shaded/hatched bands): the failed-SN (fSN) fraction (17.8%–41.7% of core-collapse progenitors depending on the strength of the neutrino engine); the NS baryonic mass limit (2.3–3.5 $M_\odot$); the instantaneous spectral-shape parameter for the emission from failed SNe ($1.0 < \alpha_{BH} < 3.0$); and the uncertainty connected to the cosmic SFH (±1σ limits of Mathews et al. 2014). The resulting “total” uncertainty band is the same as that in Figure 10. The lower left panel and the right panels show the residuals of our DSNB models where only one parameter is changed relative to the fiducial model, while all other parameters are kept at their default values, grouped into stellar-diversity uncertainties, microphysical uncertainties, modeling uncertainties, and astrophysical uncertainties (see Section 9 for more details). L_{core} and L_{evol} denote cases where the rate densities of LM NS formation events (AIC, MIC, ultrastripped SNe) are constant and evolve with redshift, respectively (both for a value of $\chi = 0.34$, which corresponds to a relative abundance of LM NS formation events of 34% as compared to “conventional” core-collapse SNe plus failed SNe; Equation (12)). In each panel, gray-shaded vertical bands frame the approximate detection window.

The upper left panel of Figure 11 shows our fiducial DSNB $\bar{\nu}_e$-flux spectrum with its main uncertainties (failed-SN fraction, NS baryonic mass limit, and spectral shape of the neutrino emission from failed SNe in terms of $\alpha_{BH}$) stacked on top of each other. The uncertainty of the SFH is additionally applied to the upper and lower limits of the uncertainty range. The impact of the different uncertainties according to the four categories listed above is illustrated by their corresponding residuals relative to the fiducial spectrum in the four additional panels of Figure 11.

Concerning stellar-diversity uncertainties, a large failed-SN fraction can enhance the DSNB spectrum by up to ~50%, whereas a considerable fraction of helium stars can shift the spectrum in the opposite direction by about the same margin. Among the microphysical uncertainties, the NS baryonic mass limit has the major impact, but an assumed value of 3.5 $M_\odot$ appears to be on the extreme side in view of current gravitational-wave and kilonova constraints, which seem to point to a mass limit around 2.7 $M_\odot$ (e.g., Margalit & Metzger 2017), which we apply to our fiducial spectrum.

Future gravitational-wave and kilonova measurements as well as astrophysical observations by NICER (Miller et al. 2019) are likely to constrain this mass limit with increasingly better precision. Among the modeling uncertainties, which are specific to our approach based on large sets of core-collapse simulations with approximate neutrino treatment, the spectral-shape parameter $\alpha_{BH}$ has the dominant influence (up to ~35% enhancement of the DSNB $\bar{\nu}_e$ spectrum at a neutrino energy of 30 MeV seems possible). However, this uncertainty as well as the (subdominant) ones connected to the total gravitational BE release and the mean energy of the radiated neutrinos will also be reduced once the NS EoS is better determined and neutrino-signal predictions from detailed transport calculations for large sets of NS and BH formation events become available.

Finally, the SFH can make changes to the DSNB $\bar{\nu}_e$ spectrum by up to a factor of two and is certainly a much desirable aspect for further improvements through astronomical observations. If this can be achieved, DSNB measurements will provide an interesting approach to deducing information on the stellar core-collapse diversity, whose effects are the main focus of our...
work. Conversely, if theoretical and observational advances lead to a better understanding of the population of core-collapse progenitors and their final destinies (i.e., their fates as successful or failed SNe), the forthcoming detection of the DS NB flux spectrum. Our study is based on large sets of single-star models (Sukhbold et al. 2016) and helium-star models (Ertl et al. 2020) for successful and failed SNe. The helium-star progenitors from Woosley (2019) are considered as a proxy of massive stars that evolved to the onset of stellar core collapse after stripping their hydrogen envelopes at the end of core-hydrogen burning through binary interaction, e.g., by common-envelope evolution or Roche-lobe overflow (see Sana et al. 2012). The progenitor sets contain between 100 and 200 stellar models with ZAMS masses between $<9 M_{\odot}$ and $120 M_{\odot}$. These models are exploded (or fail to explode) in spherically symmetric simulations with the PROMETHEUS-HoTB code, employing a parameterized neutrino engine that is calibrated to reproduce the basic properties of the well-studied SNe of SN 1987A and the Crab Nebula.

Our stellar core-collapse models provide the total energy output in neutrinos from NS and BH formation events as well as the time-dependent mean energies of the radiated neutrinos, specifically of $\bar{\nu}_e$. Since the treatment of the PNS cooling and its neutrino emission in these large model sets is only approximate, we compare our estimates of the total neutrino energy loss with the gravitational BEs of NSs (up to their mass limit) as given by the radius-dependent fit formula of Lattimer & Prakash (2001). We find good agreement for NS radii of 11–12 km, which is the range favored by recent astrophysical observations and nuclear theory and experiments. Moreover, we use NS and BH formation simulations with the PROMETHEUS-VERTEX code (which employs a state-of-the-art treatment of neutrino transport based on a Boltzmann moment closure scheme and a mixing-length treatment of PNS convection) to calibrate degrees of freedom in our approximate neutrino signals—for example, the shape of the time-dependent neutrino spectrum, which we characterize by the widely used $\alpha$-fit of Keil et al. (2003). We note that our treatment of the neutrino emission by successful and failed SNe is not based on a detailed microphysical PNS model, but nevertheless our procedure of combining information from PROMETHEUS-HoTB simulations with neutrino data from PROMETHEUS-VERTEX models enables our study to capture the generic properties of neutrino signals radiated from NS and BH formation cases.

In the course of our investigation we vary the neutrino engine, whose power is connected to the properties of the progenitor model considered for SN 1987A, yielding different relative fractions of successful SN events in contrast to failed explosions with BH formation. Moreover, we explore the effects of alternative paths to NS formation besides the stellar core-collapse channel, which could be associated with the AIC or MIC of WDs or with ECSN and ultrastripped core-collapse progenitors in close-binary systems. All of these cases would preferentially lead to the formation of rather LM NSs with little postshock accretion, for which reason we treat this component in analogy to the ECSNe (Hüdepohl et al. 2010) that are included in our standard set of stellar core-collapse models. We also vary the still uncertain NS mass limit (above which a transiently stable, accreting PNS collapses to a BH) between the currently measured largest masses of galactic NSs ($2.3 M_{\odot}$, baryonic and $\sim 2.0 M_{\odot}$ gravitational) and the maximum mass that can be stabilized by still viable microphysical EoSs ($3.5 M_{\odot}$, baryonic and $\sim 2.75 M_{\odot}$ gravitational). Moreover, we vary the shape parameter, $\alpha_{\text{BH}}$, of the time-dependent neutrino emission spectrum from failed explosions and consider, in a standard way, the effects of neutrino flavor oscillations.

Our fiducial case employs a neutrino engine that is fully compatible with observationally determined NS and BH masses as well as chemogalactic constraints on SN nucleosynthesis, an NS mass limit of $2.7 M_{\odot}$ baryonic and $\sim 2.25 M_{\odot}$ gravitational mass (compatible with recent limits from GW170817), and a best-fit $\alpha$-spectrum for the time-dependent neutrino emission. With the SFH adopted from Mathews et al. (2014), it yields a total DS NB $\bar{\nu}_e$-flux of $28.8^{+24.6}_{-10.5} \text{cm}^{-2} \text{s}^{-1}$ with a contribution of $6.0^{+5.1}_{-2.0} \text{cm}^{-2} \text{s}^{-1}$ in the energy interval of $[10, 30] \text{MeV}$, which is most favorable for measurements. Our best value of the predicted flux for $\bar{\nu}_e$ energies $>17.3 \text{MeV}$ is $1.3^{+0.4}_{-0.2} \text{cm}^{-2} \text{s}^{-1}$, which is slightly lower than the result of $1.6 \text{cm}^{-2} \text{s}^{-1}$ published by Ando et al. (2003, with an update at NNN05) and about a factor of two below the current SK limit (see Bays et al. 2012; preliminary updated value of $2.7 \text{cm}^{-2} \text{s}^{-1}$ at 90% confidence level by the Super-Kamiokande Collaboration; El Hedri et al. 2020; Nakajima 2020).

Because of the currently expected narrow mass range of ECSNe from single stars, these events yield a negligible contribution to the DS NB. Similarly, the tested alternative LM NS formation channel via AIC and MIC events or SNe from ultrastripped progenitors can contribute on a significant level ($>10\%$) only in the case of an implausibly high constant event rate or in the case of an evolving rate on the level of the cosmic SN Ia rate. But even then the enhancement of the DS NB spectrum would happen mainly at low neutrino energies $\lesssim 10 \text{MeV}$ and thus outside of the most favorable energy window for detection.

Our study confirms previous results (e.g., Lunardini 2009; Keehn & Lunardini 2012; Nakazato et al. 2015; Hidaka et al. 2016, 2018), which were based on the consideration of exemplary cases of BH formation, that an increased fraction of failed SNe flattens the exponential-like decline of the DS NB spectrum beyond its peak and lifts the high-energy tail of the spectrum. This effect can be observed both in our model sets with weaker neutrino engines, where a larger fraction of stars collapse to BHs, and, particularly strongly, in those model sets where we assume a high value for the maximum NS mass. The rise of the high-energy spectrum is mainly connected to core-collapse events with a long delay time until BH formation, where the mass-accreting PNS radiates harder neutrino spectra and releases a considerably higher total BE. Correspondingly, the high-energy tail of the DS NB spectrum varies by a factor of 6.6 at 30 MeV and the DS NB flux values above 17.3 MeV, $\Phi_{\bar{\nu}_e, 17.3}$, by a factor of 3.9 between the limits of 0.8 cm$^{-2}$s$^{-1}$ and 3.1 cm$^{-2}$s$^{-1}$ (for the model S19.8-BH23.0-2.0 compared to W20-BH3.5-2.0). A similar effect, though considerably weaker (about a 14% increase of $\Phi_{\bar{\nu}_e, 17.3}$ relative to our fiducial case), can be seen when the radiated neutrino spectra from
failed explosions are considered to be anti-pinched ($\alpha_{\text{BH}} = 1$) at all times instead of being Maxwell–Boltzmann like ($\alpha_{\text{BH}} = 2$).

A larger population of hydrogen-stripped binary progenitors of SNe can have a significant impact on the DSNB spectrum, because compared to single stars, the ZAMS mass range of stars that experience stellar core collapse is shifted upward by $\sim 5 M_\odot$ to the more IMF-suppressed high-mass regime (compare Figure 6 with Figure 2). At the same time, a lower fraction of BH formation events reduces the high-energy tail. Correspondingly, we find a reduction of the total DSNB $\bar{\nu}_e$-flux by $\sim 18\%$ (53%) and a reduction of $\Phi_{\gamma,17.3}$ by $\sim 20\%$ (60%) if $33\%$ (100%) of the core-collapse progenitors evolve as helium stars (see Figure 7). Neutrino flavor oscillations have an effect that is, at most, of roughly comparable magnitude. A complete swap of $\bar{\nu}_e$ and $\nu_x$ (the most extreme case) reduces our predictions of the total DSNB $\bar{\nu}_e$-flux again by $\sim 16\%$ and of $\Phi_{\gamma,17.3}$ by $\sim 21\%$ relative to our fiducial case.

A major uncertainty in all predictions of the DSNB, however, is the still insufficiently constrained stellar core-collapse rate. With a defined form for the stellar IMF this refers to uncertainties in the cosmic SFH, which render all estimates uncertain within a factor of roughly 3 (considering the $\pm 1\sigma$ range of Mathews et al. 2014). Rigorous constraint of individual inputs of the DSNB by measurements is further hampered by the existing large degeneracies between different effects of relevance. Nevertheless, the most extreme cases included in our study, which combine a very large fraction of BH-forming core-collapse events (up to an IMF-weighted fraction of 42%) and/or the highest considered value of the NS mass limit ($3.5 M_\odot$ baryonic and $\sim 2.75 M_\odot$ gravitational mass), seem to be ruled out by the current SK limit already.

Some of the physical quantities entering the DSNB calculations can be expected to be better constrained in the not too distant future. The increasing number of gravitational-wave detections from binary NS mergers (Abadie et al. 2010) will yield more information on the maximum NS mass and NS radii, placing tighter constraints on the high-density EoS; steadily improved statistics of binary BH mergers might lead to better constraints on BH formation events and progenitors (see, e.g., Woosley et al. 2020); long-baseline neutrino oscillation experiments should be able to determine the neutrino mass hierarchy (e.g., LBNE Collaboration et al. 2013); and upcoming wide-field surveys such as LSST (Tyson 2002) will measure the rate of visible SNe (below $z \sim 1$) to good accuracy.

Complementary to these perspectives, future observations of the DSNB will probe the entire population of stellar core-collapse events with its full diversity, particularly including faint and failed explosions (see Lien et al. 2010). This opens up the chance of better constraining the cosmic core-collapse rate as well as the fraction of BH-forming, failed SNe (Müller et al. 2018). Moreover, the DSNB may even carry the imprint of new physics (e.g., Fogli et al. 2004; Farzan & Palomares-Ruiz 2014; Jeong et al. 2018; de Gouvêa et al. 2020; Tabrizi & Horiuchi 2020). These exciting prospects for both particle physics and astrophysics motivate ongoing efforts to steadily improve theoretical predictions of the DSNB. The next upgrade in this direction should be fully self-consistent successful and failed SN simulations with a detailed modeling of the neutrino signal radiated by the forming compact remnant.

Our results are made available for download upon request on the following website: https://wwwmpa.mpa-garching.mpg.de/ccsnarchive/archive.html.

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**Software:** PROMETHEUS-HoTB (Janka & Müller 1996; Kifonidis et al. 2003; Scheck et al. 2006; Ertl et al. 2016, 2020), NumPy (Harris et al. 2020), SciPy (Virtanen et al. 2020), IPython (Pérez & Granger 2007), Matplotlib (Hunter 2007), Bibmanager (Cubillos 2020).

*Note added in proof.* When our paper was already in the production process, we got notice of a new arXiv posting by Horiuchi et al. (2021), dealing with the impact of mass transfer and mergers during binary evolution on the DSNB spectrum. We agree that this progenitor component of core-collapse SNe, which we did not take into account in our study, can potentially increase the DSNB flux, because SN progenitors arising from a population of binary interacting systems originate on average from lower ZAMS masses (Zapartas et al. 2021). This fact can partially compensate for the reducing influence of stripped progenitors discussed in our work. However, the relevant effects of mass transfer and mergers depend on a variety of uncertain processes during stellar evolution and are hard to assess in quantitative detail (see the detailed discussion by Zapartas et al. 2021).

**Appendix A**

Extrapolation of Neutrino Signals

In our analysis as described in Sections 2 and 3, we employ the neutrino signals from successful SNe (including rare cases of fallback SNe) up to 15 s post-bounce, at which time their luminosities have declined to a level that is not relevant for our purpose of estimating the DSNB; moreover at late times the NS temperature drops and therefore the mean spectral energies of the emitted neutrinos shift out of the DSNB detection window. In contrast, the signals from failed explosions have to be followed until the accreting NS reaches the mass limit for BH formation, $M_{\text{lim,NS}_{15}}$, which may take tens of seconds in cases of low mass-accretion rates and high $M_{\text{lim,NS}_{15}}$ (see upper panel of Figure 2). Not all of our successful or failed SN simulations could be carried out long enough, because of rising computational costs or due to numerical problems emerging at late times (after several seconds). We thus extrapolate these neutrino signals after the computational end at post-bounce time $t_0$. In the upper panel of Figure A1, $t_0$ is plotted against the ZAMS mass for our reference engine model Z9.6 and W18. Typically, our extrapolation starts at around $8–10$ s, whereas no extrapolation is needed for a few successful SNe near the low-
Figure A1. Systematics of our signal extrapolation over the range of progenitor models from the WH15, SW14, and WH07 sets for the Z9.6 and W18 engine (see Figure 2). In the upper panel, the starting time of our extrapolation, \( t_0 \), is given. The lower panel shows the relative fraction of the total radiated neutrino energy arising from the extrapolation (note the logarithmic scale). Both quantities are plotted vs. the ZAMS mass. Red bars indicate successful SN explosions (including rare fallback SNe), while BH-forming, failed SNe are marked dark blue, light blue, and cyan corresponding to baryonic NS mass limits of \( 2.7 \, M_\odot \), \( 3.1 \, M_\odot \), and \( 3.5 \, M_\odot \), respectively (no extrapolation is needed for the case of \( 2.3 \, M_\odot \)). The time \( t_0 \) is independent of the mass limit. Progenitors below a ZAMS mass of \( 10 \, M_\odot \) as well as fast-accreting BH cases do not require extrapolation.

The cooling phases of our successful, NS-forming SNe can be described approximately by an exponential decline of the neutrino signal at sufficiently late times after shock revival, when mass accretion onto the hot PNS has ceased and diffusion of neutrinos from the core defines the emission (Burrows & Lattimer 1986; Keil & Janka 1995; Pons et al. 1999). We thus extrapolate the signals of our successful SNe according to

\[
L_{\text{core}}(t) = L_0^{\text{core}} e^{-(t-t_0)/\tau}
\]

for all neutrino species \( \nu_i \), with \( L_0^{\text{core}} = L_0(0) \) being the corresponding luminosity at the end of our simulations at time \( t_0 \) and \( \tau = \tau_0 \) being the core-cooling timescale, which we obtain from least-squares fits over the last 2 s of the computed neutrino signals. Our values for \( \tau \) typically range between 1 and 4 s, in agreement with the work by Hüdepohl (2014) (also see Müller et al. 2016, their Table 1). The lower panel of Figure A1 shows the relative contributions to the total radiated neutrino energies from our extrapolations (in the time interval \( t_0 \leq t \leq 15 \text{ s} \) for the cases of successful explosions). They lie below \( \sim 1\% - 2\% \) for all successful SNe, which illustrates that further extrapolation of exponentially declining signals beyond 15 s is not necessary. Similar results are obtained for all our engine models. The mean neutrino energies, \( \langle E_\nu(t) \rangle \), are simply extrapolated by keeping them constant at their final values at \( t_0 \), which, because of the small contribution from the neutrino emission at late times \( (t > t_0) \) to the time-integrated signals, has no significant influence on our DSNB predictions and is therefore unproblematic.

In the cases of BH-forming, failed SNe, on the other hand, the continued infall of matter feeds the accretion luminosity in addition to the diffusive flux from the core (Burrows 1988). Therefore, we describe the total neutrino emission (of all species) as the sum of a core and an accretion component, \( L_{\text{tot}}(t) = L_{\text{core}}(t) + L_{\text{acc}}(t) \). For the accretion luminosity, we follow the description by Burrows (1988),

\[
L_{\text{acc}}(t) = \frac{\eta GM_{\text{NS,b}}(t) M_{\text{NS,b}}(t)}{R_{\text{NS}}(t)}
\]

with the gravitational constant \( G \) and an adjustable efficiency parameter \( \eta \) (see Fischer et al. 2009; Hüdepohl 2014; Müller & Janka 2014). For computational reasons, we take the late-time evolution of the progenitor-dependent (baryonic) PNS mass, \( M_{\text{NS,b}}(t) \), and accretion rate, \( \dot{M}_{\text{NS,b}}(t) \), from pure hydrodynamic simulations with the neutrino engine switched off and an open inner boundary of the computational grid placed in the supernovically infalling matter exterior to the stalled accretion shock. As we do not have model-based information on the time-dependent radius, \( R_{\text{NS}}(t) \), of the contracting PNS, we adopt Equation (9) of Müller et al. (2016),

\[
R_{\text{NS}}(t) = \left[ R^3 \left( \frac{M_{\text{NS,b}}(t)}{M_\odot s^{-1}} \right)^{3/4} + R_2^3 \right]^{1/3}
\]

We find that the late phases of those failed-SN simulations that are carried on beyond 10 s (21 cases in the N20 set and 72 in the W20 set, 17 of them beyond 20 s) are reproduced by Equations (A2) and (A3) with an accuracy of a few percent when we choose the parameter values \( R_1 = 40 \text{ km} \), \( R_2 = 11 \text{ km} \), and accretion efficiency \( \eta = 0.51 \). Similar values for \( \eta \) were found by Fischer et al. (2009), Hüdepohl (2014), and Müller & Janka (2014). We apply this description of the accretion luminosity to all of our extrapolated failed-SN signals, independently of the engine model. For the core luminosity (of all neutrino species) in failed explosions, we employ Equation (A1) with an initial value \( L_0^{\text{core}} = L_{\text{tot}}(t_0) - L_{\text{acc}}(t_0) \).

\(^{17}\) The absolute values of \( R_1 \) and \( R_2 \) can be chosen somewhat arbitrarily since the adjustable parameter \( \eta \) compensates for shifts of \( L_{\text{acc}} \) in Equation (A2). For consistency with the measured NS radii, we take \( R_2 = 11 \text{ km} \) (see footnote 4). The resulting best-fit value of \( R_1 = 40 \text{ km} \) is much smaller than the 120 km in Müller et al. (2016), which reflects the moderate core contraction in our simulations.
and a core-cooling timescale $\tau = \tau_{\nu_i} \approx 1 \text{s}$ from a least-squares fit of the heavy-lepton neutrino signal between 3 s and 6 s after bounce in each model. During this phase, $L_{\nu_i}$ is dominated by its core component and can be well approximated by an exponential decline. We hence adopt this prescription also for the core luminosities of electron-type neutrinos, which are not as readily accessible (see Hudepohl 2014; Müller & Janka 2014). In the extrapolation, the relative contributions of the different neutrino species to the total emission are kept constant at their final values obtained at the end of the simulations (i.e., $L_{\nu_i}(t) \approx f_i L_{\text{tot}}(t)$, with the factor $f_i = L_{\nu_i}(t_f)/L_{\text{tot}}(t_f)$ equally applied to the core and accretion components).

As can be seen in the lower panel of Figure A1, our extrapolation accounts for up to $\sim 40\%$ of the total radiated neutrino energy for the case of an NS mass limit of $3.5 \, M_\odot$ in the most extreme conditions, while no extrapolation is required for a limiting NS mass of $2.3 \, M_\odot$. This is true for all of our engine models. The mean neutrino energies from slowly accreting failed SNe, where the extrapolation has the biggest influence, flatten to rather constant values ($\sim 20 \, \text{MeV}$) at late times in simulations that could be carried on for more than $\sim 10 \, \text{s}$. We thus extrapolate the mean neutrino energies in failed SNe by keeping them constant at their final values at $t_f$, in analogy to what we do in the cases of successful SNe. We test other extrapolation schemes but find that the time-integrated spectra are largely insensitive to the late-time description of the mean energies.

Appendix B

Total Energies of Radiated Neutrinos

In both successful and failed core-collapse SNe, the neutrino emission is fed by the release of gravitational BE from an assembling PNS, which either cools down to become a stable NS or further collapses to a BH. To assess the viability of our DSNB flux predictions, we compare the total radiated neutrino energy, $E_{\nu}^{\text{tot}}$, obtained from our simulations with an analytic estimate of the BE. For this purpose, we adopt Equation (36) of Lattimer & Prakash (2001), which connects the PNS’s baryonic mass, $M_{\text{NS,b}}$, with its gravitating mass, $M_{\text{NS,g}}$, assuming a final (cold) NS radius $R_{\text{NS}}$:

$$\frac{\text{BE}}{c^2} = \frac{0.6 \beta}{1 - 0.5 \beta} M_{\text{NS,g}},$$

with $\text{BE}/c^2 = M_{\text{NS,b}} - M_{\text{NS,g}}$ and the dimensionless parameter $\beta = GM_{\text{NS,b}}/R_{\text{NS}}^2$.

In the left panel of Figure B1, the $E_{\nu}^{\text{tot}}$ of our successful explosions in the Z9.6 and W18 set is plotted against the baryonic mass of the relic NS (turquoise dots). We compare these values with the corresponding gravitational BEs $\text{BE}_{11}$ and $\text{BE}_{12}$ (gray and red dashed lines), computed with Equation (B1) for an assumed final NS radius of 11 km and 12 km, respectively. The shaded bands indicate deviations of $\pm 10\%$ from the analytic relations. In the right panel, we also show the ratio of the total radiated neutrino energy to the BE for the case of $R_{\text{NS}} = 11 \, \text{km}$, plotted against the ZAMS mass $M_{\text{ZAMS}}$ of the progenitors.

Our simulations feature good overall agreement with Equation (B1), compatible with the PNS of a successful SN radiating essentially its entire gravitational BE in the form of neutrinos. Assuming an NS radius of 11 km, $93\%$ of the successful explosions in our Z9.6 and W18 set deviate by less than $15\%$ from the analytic fit provided by Lattimer & Prakash (2001). Most of our simulations overestimate the total radiated neutrino energy on the order of $10\%$, but for the majority of low-mass progenitors the values of $E_{\nu}^{\text{tot}}$ are close to or below $\text{BE}_{11}$, which leads to an IMF-weighted mean deviation of $+7.1\%$. If we assume $R_{\text{NS}} = 12 \, \text{km} (13 \, \text{km})$ instead, the deviation increases to a value of $+15.6\% (+24.1\%)$ above the analytic description. In Table B1, we show the IMF-weighted mean deviations for all of our engine models.

Compared to successful explosions, the total energy reservoir that could be released in neutrinos by BH-forming, failed SNe is generally higher if the PNS at the limiting mass remains stable until it has emitted its entire gravitational BE before it collapses to a BH (see Table B2). However, the BE of a maximum-mass NS constitutes just an upper limit for the radiated neutrino energy $E_{\nu}^{\text{tot}}$, because BH formation typically occurs before the NS has cooled to a cold state, terminating the neutrino emission before the total gravitational energy is carried away by neutrinos. This can be seen in the left panel of Figure B1. Comparison of the total neutrino energies, $E_{\nu}^{\text{tot}}$, radiated by the successful explosions of our reference set (Z9.6 and W18) with the gravitational BEs of the relic NSs as estimated with an analytic expression from Lattimer & Prakash (2001). In the left panel, the relation between $E_{\nu}^{\text{tot}}$ and the baryonic NS mass, $M_{\text{NS,b}}$, is shown (turquoise dots). The gray (red) dashed line indicates the NS’s BE as a function of $M_{\text{NS,b}}$, computed with Equation (B1), assuming an NS radius of 11 km (12 km). The shaded bands correspond to deviations of $\pm 10\%$. In the right panel, the ratio of the total radiated neutrino energy to the BE is plotted vs. the ZAMS mass for an NS radius of 11 km. The dashed turquoise line additionally indicates the IMF-weighted mean value, which deviates by $+7.1\%$ from the BE. Note the scale break at $M_{\text{ZAMS}} \approx 30 \, M_\odot$. 

Figure B1.
available reservoir of gravitational BE as given by the analytic formula of Lattimer & Prakash (2001). The left panel shows \( E^{\text{tot}}_n \) vs. the time until BH formation (turquoise dots). The three dashed lines (in blue, gray, and red) indicate the total gravitational BEs, according to Equation (B1), of an NS with an assumed maximum baryonic mass of \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \), and an assumed radius of 10 km, 11 km, and 12 km, respectively. In the right panel, the ratio of the radiated neutrino energy to the maximally available gravitational BE is plotted vs. the progenitor’s ZAMS mass for an assumed NS radius of 11 km. Note the scale break at \( M_{\text{ZAMS}} \sim 30 \, M_\odot \).

![Figure B2](image)

**Figure B2.** Comparison of the total neutrino energies, \( E^{\text{tot}}_n \), radiated by the failed explosions of our reference set (Z9.6 and W18, \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \)) with the maximally available reservoir of gravitational BE as given by the analytic fit formula of Lattimer & Prakash (2001). The left panel shows \( E^{\text{tot}}_n \) vs. the time until BH formation (turquoise dots). The three dashed lines (in blue, gray, and red) indicate the total gravitational BEs, according to Equation (B1), of an NS with an assumed maximum baryonic mass of \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \), and an assumed radius of 10 km, 11 km, and 12 km, respectively. In the right panel, the ratio of the radiated neutrino energy to the maximally available gravitational BE is plotted vs. the progenitor’s ZAMS mass for an assumed NS radius of 11 km. Note the scale break at \( M_{\text{ZAMS}} \sim 30 \, M_\odot \).

**Table B1**

| Engine Model | \( R_{\text{NS}} = 11 \, \text{km} \) | \( R_{\text{NS}} = 12 \, \text{km} \) | \( R_{\text{NS}} = 13 \, \text{km} \) |
|--------------|-------------------------------|-------------------------------|-------------------------------|
| Z9.6 and S19.8 | +11.8% | +20.7% | +29.6% |
| Z9.6 and N20 | +6.1% | +14.6% | +23.0% |
| Z9.6 and W18 | +7.1% | +15.6% | +24.1% |
| Z9.6 and W15 | +5.1% | +13.5% | +21.8% |
| Z9.6 and W20 | +7.0% | +15.6% | +24.1% |

**Note.** For the computations of the BE, Equation (B1) is used with final NS radii of 11, 12, or 13 km.

**Table B2**

| Baryonic NS Mass Limit | \( E^{\text{tot}}_n \) (10^{53} \text{erg}) | \( \text{BE}_{10} \) (10^{53} \text{erg}) | \( \text{BE}_{11} \) (10^{53} \text{erg}) | \( \text{BE}_{12} \) (10^{53} \text{erg}) | \( \text{BE}_{13} \) (10^{53} \text{erg}) |
|------------------------|---------------------------------|---------------------------------|---------------------------------|---------------------------------|---------------------------------|
| \( M_{\text{lim}}^{\text{NS,b}} = 2.3 \, M_\odot \) | 6.8 (77.8%) | 6.3 (84.1%) | 5.9 (90.5%) | 5.5 (96.8%) |
| \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \) | 9.1 (77.3%) | 8.4 (83.4%) | 7.8 (89.5%) | 7.3 (95.6%) |
| \( M_{\text{lim}}^{\text{NS,b}} = 3.1 \, M_\odot \) | 11.6 (76.0%) | 10.8 (81.8%) | 10.1 (87.6%) | 9.4 (93.5%) |
| \( M_{\text{lim}}^{\text{NS,b}} = 3.5 \, M_\odot \) | 14.4 (74.2%) | 13.4 (79.7%) | 12.5 (85.3%) | 11.8 (90.8%) |

**Note.** The values of the BE are computed according to Equation (B1) for different values of the baryonic NS mass limit, \( M_{\text{lim}}^{\text{NS,b}} \), and for different NS radii (10, 11, 12, and 13 km). In parentheses we give the largest value of the ratio \( E^{\text{tot}}_n / \text{BE} \) for all of our failed explosion models and all our neutrino engines. Note the slightly larger value of 83.4% for \( E^{\text{tot}}_n / \text{BE} \) in the case of \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \) compared to the ~80% in the right panel of Figure B2, which shows the case of the Z9.6 and W18 neutrino engine.

Figure B2, where we plot \( E^{\text{tot}}_n \) for the failed SNe of our reference set (Z9.6 and W18, \( M_{\text{lim}}^{\text{NS,b}} = 2.7 \, M_\odot \)) against the time until BH formation (turquoise dots). Only the slowly accreting cases (with \( t_{\text{BH}} \gtrsim 3 \, t_s \)) come close to the maximally available BE according to Equation (B1), which is indicated by a blue, gray, or red dashed line for an NS radius of 10 km, 11 km, or 12 km, respectively. In the right panel, we show the ratio of the radiated to the maximally available energy for an assumed NS radius of 11 km versus the ZAMS mass range of the corresponding progenitors.

For all of our simulations the neutrino emission from failed SNe lies well below the analytically computed energy limit. For our reference set shown in Figure B2, at most 80% of \( \text{BE}_{11} \) is radiated before a BH forms, while the progenitors at around 22–25 \( M_\odot \) and \( \sim 40 \, M_\odot \), which exhibit very high mass-accretion rates (see footnote 6 and upper panel of Figure 2), feature considerably lower percentages (~30%–60%). The results for the other neutrino engines are very similar, because the emission from a failed SN is dominated by the progenitor-dependent accretion component rather than the PNS core emission. For larger NS radii applied in Equation (B1), the ratio \( E^{\text{tot}}_n / \text{BE} \) tends toward unity, as can be seen in Table B2 (values in parentheses).

**Appendix C**

**Flavor Rescaling**

The approximate treatment of the microphysics and the relatively modest contraction of our inner grid boundary in the considered core-collapse simulations result in underestimated
Table C1

| Model                  | $\xi_{\nu_e}$ | $\xi_{\nu_x}$ | $\xi_{\nu_\tau}$ | $\lambda^\nu_{\nu_e}$ | $\lambda^\nu_{\nu_x}$ | $\lambda^\nu_{\nu_\tau}$ | Compact Remnant | $M_{\text{NS,b}}$ (M$_\odot$) |
|------------------------|---------------|---------------|-------------------|------------------------|------------------------|------------------------|-----------------|-------------------------------|
| VERTEX, s11.2co, LS220 | 0.166         | 0.194         | 0.160             | 0.990                  | 0.808                  | 0.960                  | NS              | 1.366                         |
| VERTEX, s27.0co, LS220 | 0.155         | 0.173         | 0.168             | 0.992                  | 0.810                  | 0.963                  | NS              | 1.361                         |
| Average ("L")          | 0.159         | 0.181         | 0.165             | 0.991                  | 0.803                  | ...                    | ...             | ...                           |
| VERTEX, s20.0, SFHo     | 0.172         | 0.176         | 0.163             | 0.965                  | 0.813                  | 0.947                  | NS              | 1.947                         |
| VERTEX, s27.0co, LS220  | 0.172         | 0.181         | 0.162             | 0.957                  | 0.807                  | 1.776                  | NS              | 1.772                         |
| VERTEX, s27.0co, SFHo   | 0.170         | 0.179         | 0.160             | 0.973                  | 0.810                  | NS                     | NS              | 1.772                         |
| Average ("H")          | 0.171         | 0.179         | 0.163             | 0.965                  | 0.810                  | ...                    | ...             | ...                           |
| VERTEX, s40.78bc, LS220 ("F") | 0.212 | 0.257       | 0.133             | 1.068                  | 0.639                  | BH (fast; 0.57 s)       | (2.320)         |
| VERTEX, s40.0c, LS220 ("S") | 0.231 | 0.251       | 0.129             | 0.940                  | 0.724                  | BH (slow; 2.11 s)       | (2.279)         |

Note. Relative fractions $\xi_{\nu_e}$, $\xi_{\nu_x}$, and $\xi_{\nu_\tau}$ of the total energy $E^\text{tot}_\nu$ radiated in the neutrino species $\nu_e$, $\nu_x$, and $\nu_\tau$ (Equation (C2)) and conversion factors $\lambda^\nu_{\nu_e}$, $\lambda^\nu_{\nu_x}$, and $\lambda^\nu_{\nu_\tau}$ (Equation (C4)) are listed for eight models that are simulated with the 1D version of the PROMETHEUS-VERTEX code. The lines are grouped according to the masses of the compact remnants. For models s40.78bc and s40.0c, the values in parentheses indicate the post-bounce times of BH formation and the corresponding baryonic PNS masses at these times. The conversion factors applied to our PROMETHEUS-HOTB models (according to the four cases “L,” “H,” “F,” and “S”) are listed in Table C1: z9.6co and s27.0co, both simulated with the LS220 EoS (Lattimer & Swesty 1991) and the SFHo EoS (Steiner et al. 2013) and discussed in detail in Mirizzi et al. (2016); the unpublished model s20.0 of a 20 M$_\odot$ progenitor of Woosley & Heger (2007), computed with the SFHo EoS in the same way as the four models mentioned before (R. Bollig 2018, private communication); and the three models s11.2co, s40.0c, and s40.78bc from Hudepohl (2014), all of them computed with the LS220 EoS. The suffix “c” of the model names indicates the use of a mixing-length treatment for PNS convection, and the suffix “o” indicates that mean-field potentials are taken into account in the charged-current neutrino–nucleon interactions (see Mirizzi et al. 2016 for details). The neutrino signals of all eight models can be found in the Garching Core-collapse Supernova Archive.18

Although we constrain our analysis in most parts on the emitted $\bar{\nu}_e$ signals, we need information on the time-integrated spectra, $dN_{\nu}_e/dE$, also for heavy-lepton neutrinos in our discussion of flavor oscillation effects in Section 7.1. Instead of taking the outcome of (too approximate) SN and BH formation models, we directly employ the spectral shape from Keil et al. (2003),

\[
\frac{dN_{\nu_e}}{dE} = \frac{(\bar{\nu}_e + 1)(\bar{\nu}_e + 1)}{\Gamma(\bar{\nu}_e + 1) \langle \bar{\nu}_e \rangle} \frac{E^{\nu_e}}{\langle E^{\nu_e} \rangle} \exp \left[ - \left( \bar{\nu}_e + 1 \right) E / \langle E^{\nu_e} \rangle \right],
\]

with the total energy radiated in a single heavy-lepton neutrino species $E^\text{tot}_{\nu_e} = \bar{\nu}_e E^{\nu_e}$, the time-averaged mean neutrino energy $\langle E^{\nu_e} \rangle = \lambda^\nu_{\nu_e} E^{\nu_e}$, and the spectral-shape parameter $\bar{\nu}_e = \lambda^\nu_{\nu_e} / \lambda^\nu_{\bar{\nu}_e}$ of the time-integrated $\nu_e$ spectrum. Here, $\langle E^{\nu_e} \rangle$ and $\bar{\nu}_e$ are computed from the time-integrated spectra of $\bar{\nu}_e$ obtained in our large set of core-collapse simulations. For the conversion factors

\[
\lambda^\nu_{\nu_e} \equiv \langle \nu_e / \nu_{\bar{\nu}_e} \rangle^{\nu_e}
\]

and

\[
\lambda^\nu_{\bar{\nu}_e} \equiv \langle \bar{\nu}_e / \nu_{\bar{\nu}_e} \rangle^{\bar{\nu}_e},
\]

we take the values from the PROMETHEUS-VERTEX models in Table C1. The shape parameters, $\bar{\nu} = \bar{\nu}_e$, and mean neutrino energies, $\langle E \rangle = \langle E^{\nu_e} \rangle$, of the time-integrated spectra, $dN_{\nu_e}/dE = dN_{\bar{\nu}_e}/dE$, are computed as

\[
\bar{\nu} = \frac{2 \langle E^2 \rangle - \langle E \rangle^2}{\langle E^2 \rangle - \langle E \rangle^2},
\]

18 https://wwwmpa.mpa-garching.mpg.de/csnarchive/archive.html (access provided upon request).

19 The bar in the symbol $\bar{\nu}_e$ indicates that the shape parameters refer to the time-integrated spectra rather than the instantaneous spectra (see Section 3.1 and Appendix E).
The neutrino spectra of our successful explosions that form NSs with baryonic masses of \( M_{\text{NS,b}} \leq 1.6 \) \( M_\odot \) are rescaled by the average conversion factors of the s11.2co model and the two z9.6co models (upper part of Table C1; case “L”). For SNe with \( M_{\text{NS,b}} > 1.6 \) \( M_\odot \), we apply the average values of the s20.0 and s27.0co models (middle part of Table C1; case “H”). In cases of failed explosions with BH formation (lower part of Table C1), we distinguish between fast-accreting (\( t_{\text{BH}} < 2 \) s; “F”) and slowly accreting (\( t_{\text{BH}} \geq 2 \) s; “S”) cases. The spectra of our fast-accreting models (progenitors with high core compactness; see footnote 6) are rescaled according to VERTEX model s40s7b2c, which forms a BH after 0.57 s. For our slowly accreting cases with long delays until BH formation, which correlate with higher maximum NS masses and with progenitors that have relatively lower core compactness, we employ the rescaling factors of model s40.0c, where BH formation occurs at \( t_{\text{BH}} = 2.11 \) s. For completeness, we also give the baryonic PNS masses just before the PNSs collapse to BHs in Table C1. Note that approximate flavor equipartition (\( \tilde{\xi}_e \cong \tilde{\xi}_\nu \cong \tilde{\xi}_{\bar{\nu}} \)) is realized for successful SNe, whereas \( \tilde{\nu}_e \) and \( \nu_e \) dominate over heavy-lepton neutrinos in cases of failed explosions. This can be understood by the continued accretion of infalling matter, which is accompanied by \( e^\pm \) captures on free nucleons in the PNS’s accretion mantle (Janka 2012), giving rise to an enhanced accretion luminosity of electron-flavor neutrinos and antineutrinos (see Equation (A2)).

Appendix D
Spectral Shapes

Our simplified approach does not provide information on the spectral shape of the neutrino emission. As described in Section 3, we therefore assume a spectral-shape parameter \( \alpha \), which is constant in time and for which we adopt different values in Section 5. Here, we examine how well our time-integrated spectra match the outcome of more sophisticated simulations with time-dependent \( \alpha \). Moreover, the range of values for the instantaneous shape parameters used in our study shall be motivated in this context.

In Figure D1, we compare the time-integrated spectra, \( dN/dE \), of electron antineutrinos, obtained from exemplary SN simulations of our Z9.6 and W18 set for different values of the instantaneous spectral-shape parameter \( \alpha \) with the spectra from four SN models that are computed with the 1D version of the PROMETHEUS-VERTEX code (also see Appendix C and Table C1). We take models in the same range of ZAMS masses as the VERTEX models to compare with, and with NS baryonic masses \( M_{\text{NS,b}} \) similar to those of the VERTEX calculations. Because neutrinos emitted with energies less than

\[
\langle E \rangle = \frac{\int dE \: E \: dN/dE}{\int dE \: dN/dE},
\]

\[
\langle E^2 \rangle = \frac{\int dE \: E^2 \: dN/dE}{\int dE \: dN/dE}.
\]

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\[ \text{Figure D2. Time-integrated spectra, } \frac{dN}{dE}, \text{ of electron antineutrinos (normalized by } N_{15} = \int_{15 \text{ MeV}}^{\infty} N_{15} \, dE(\frac{dN}{dE}) \text{) from selected BH formation simulations of our Z9.6 and W18 set (with } M_{\text{BH,lim}} = 2.7 M_\odot \text{; blue solid lines), compared to two BH formation models that are computed with the PROMETHEUS-VERTEX code (dashed lines). In the upper panels the spectra from two exemplary progenitors (s22.1 and s18.0) for different values of the instantaneous spectral-shape parameter } \alpha_{\text{BH}} \text{ are compared to the VERTEX models s40s7b2c and s40.0c, which form BHs relatively “fast” (after 0.57 s at } M_{\text{BH,lim}} = 2.32 \text{ M}_\odot \text{) and “slowly” (after 2.11 s at } M_{\text{BH,lim}} = 2.27 M_\odot \text{), respectively. The lower left panel shows the spectra from the exemplary s28.0 progenitor for different choices of the NS baryonic mass limit } M_{\text{NS,b}} \text{; the lower right panel shows the spectra for eight different progenitors with increasing accretion times until BH formation (between 1.0 s and 9.1 s for the shown case of } M_{\text{NS,b}} = 2.7 M_\odot \text{). In both lower panels } \alpha_{\text{BH}} = 2.0 \text{ is taken. The arrows at the bottom of each panel mark the mean energies of the spectra (Equation (C7)). The gray-shaded vertical bands edge the most relevant energy region of } E \gtrsim 15 \text{ MeV for most SNe after accounting for the cosmological redshift, we restrict our comparison to the (most relevant) high-energy range of } E \gtrsim 15 \text{ MeV. Accordingly, we normalize the spectra by } N_{15} = \int_{15 \text{ MeV}}^{\infty} dE(\frac{dN}{dE}) \text{ for better comparability of their shapes.}

We find good overall agreement of our models with the VERTEX simulations at energies } E \gtrsim 15 \text{ MeV. There is a noticeable mismatch to the left of the spectral peak, which is connected to slightly higher mean neutrino energies (by } \sim 1 \text{ MeV) compared to the reference models computed with VERTEX (see arrows in Figure D1 and the values of } (E_{\text{BH}}) \text{ in Table E1). This cannot be avoided with our chosen normalization, but it is of no relevance as pointed out above. The best fits are achieved when we take an instantaneous shape parameter of } \alpha = 3.5 \text{ for SNe with low-mass NSs (} M_{\text{NS,b}} \lesssim 1.6 M_\odot \text{; upper panels) and } \alpha = 3.0 \text{ for SNe with } M_{\text{NS,b}} > 1.6 M_\odot \text{ (lower panels). This parameter choice is largely insensitive to variations of the progenitors or of our engine model. We thus use these “best-fit” values of } \alpha \text{ for successful SNe in all of our DSNB calculations (see Section 4).}

For the case of BH-forming, failed SNe we provide an analogous comparison of our time-integrated spectra with two PROMETHEUS-VERTEX models in Figure D2. The two upper panels show the spectra from the exemplary models s22.1 and s18.0 (with our standard value of the NS baryonic mass limit of } M_{\text{NS,b}} = 2.7 M_\odot \text{) for different values of the instantaneous spectral-shape parameter (} \alpha = \alpha_{\text{BH}} \text{). These two models are chosen such that the mean neutrino energies of their spectra (Equation (C7)) are not too different from the ones of the two VERTEX models to compare with (see Table E1). We find a best fit of the spectra for } \alpha_{\text{BH}} = 2.0 \text{ in both cases (i.e., when the instantaneous spectra are Maxwell–Boltzmann like at all times of emission). This value is used as our fiducial case for failed SNe (see Section 4).}

However, as the neutrino emission from BH formation events is strongly dependent on the maximum stable mass of cold NSs as well as on the progenitor-specific accretion rates, we also investigate how the spectral shapes change under variation of } M_{\text{BH,lim}} \text{ (lower left panel) or of the chosen progenitor (lower right panel). When the NS mass limit is increased from } 2.3 M_\odot \text{ to } 3.5 M_\odot \text{, the mean energies of the time-integrated spectra rise by roughly } 4 \text{ MeV from } \sim 15 \text{ MeV to } \sim 19 \text{ MeV for the exemplary case of model s28.0 (see arrows in the lower left panel of Figure D2), which leads to a flattened spectral slope. The same trend, i.e., higher mean energies and thus flatter spectral slopes, can also be seen for spectra from progenitors with increasingly late BH formation (compare, e.g., case s18.0...
with the rapid BH formation case s40 in the lower right panel of Figure D2). Because the comparison to only two VERTEX reference models cannot be ultimately conclusive for the optimal choice of instantaneous $\alpha_{\text{BH}}$ values, we perform a set of DSNB calculations with varied choices of $\alpha_{\text{BH}}$ between 1 and 3 for our BH formation cases in Section 5.1. In doing this, we intend to test the uncertainties connected to the spectral variations of the neutrino emission from failed explosions in a systematic way.

### Appendix E

#### Spectral Parameters

For the sake of completeness and as service to the community for use in future studies, we provide in Table E1 the spectral parameters (i.e., the total radiated neutrino energies, $E_{\nu}^{\text{tot}}$, mean neutrino energies, $\langle E_{\nu} \rangle$, and spectral-shape parameters, $\pi_{\nu}$) for the time-integrated neutrino emission of all neutrino species (\(\bar{\nu}_e, \nu_e, \nu_x, \nu_x, \bar{\nu}_x, \nu_x\)) and the same PROMETHEUS-VERTEX reference models listed in Table C1 as well as for the 8.8\,M$_\odot$ ECSN (model “Sf”) from Hüdepohl et al. (2010). Moreover, we list the values for a selected set of exemplary PROMETHEUS-HOTB models (as shown in Figures D1 and D2). Note that the values of $\pi_{\nu}$ (Equation (C6)) are generally somewhat (~5%–15%) smaller than the $\alpha$ parameters of the instantaneous neutrino emission (Equation (3)), i.e., the time-integrated spectra are slightly wider than the instantaneous ones.

We also point out that in the PROMETHEUS-VERTEX simulations the total energy released in neutrinos by the cooling PNSs is EoS-dependent and for BH-forming cases in particular, it depends on the accretion time until the NS collapses to a BH. For a comparison of the results obtained from our PROMETHEUS-HOTB models with generic values based on the (EoS-independent but NS radius–dependent) fit formula of Lattimer & Prakash (2001; see Equation (B1)), we refer the reader to Appendix B. For a comparative discussion of the time-integrated neutrino spectra of the PROMETHEUS-VERTEX simulations and our best-fit spectra for the PROMETHEUS-HOTB models (obtained by suitable choices of values of the instantaneous shape parameter $\alpha$), we refer the reader to Appendix D.

| Model | $E_{\nu}^{\text{tot}}$ (10$^{50}$ erg) | $\langle E_{\nu} \rangle$ (MeV) | $\pi_{\nu}$ | $\pi_{\nu}$ | $\pi_{\nu}$ | $\pi_{\nu}$ | Remnant | $M_{\text{NS,B}}$ (M$_\odot$) |
|-------|---------------------------------|-----------------|---------|---------|---------|---------|---------|----------------|
| VERTEX, 8.8\,M$_\odot$, ECSN (“Sf") | 2.67 | 3.20 | 2.62 | 11.6 | 9.5 | 11.5 | 2.49 | 3.06 | 2.10 | NS | 1.366 |
| VERTEX, 29.6\,M$_\odot$, SS220 | 2.93 | 3.28 | 3.17 | 12.4 | 9.7 | 12.4 | 2.51 | 2.82 | 2.03 | NS | 1.361 |
| VERTEX, z9.6\,M$_\odot$, SFHo | 3.13 | 3.49 | 3.31 | 12.1 | 9.6 | 12.0 | 2.83 | 3.03 | 2.24 | NS | 1.363 |
| VERTEX, z11.2\,M$_\odot$, SS220 | 3.09 | 3.56 | 3.02 | 13.7 | 10.6 | 13.6 | 2.90 | 2.76 | 2.34 | NS | 1.366 |
| VERTEX, s27.0\,M$_\odot$, SS220 | 5.72 | 5.99 | 5.37 | 13.7 | 10.9 | 13.1 | 2.25 | 2.15 | 1.82 | NS | 1.776 |
| VERTEX, s27.0\,M$_\odot$, SFHo | 5.91 | 6.24 | 5.68 | 13.6 | 10.9 | 13.2 | 2.61 | 2.50 | 2.11 | NS | 1.772 |
| VERTEX + S220, SS220 | 7.36 | 7.53 | 6.96 | 14.0 | 11.3 | 13.5 | 2.48 | 2.31 | 2.02 | NS | 1.947 |
| VERTEX + s407\,b2c, SS220 | 4.49 | 5.44 | 2.81 | 17.6 | 14.4 | 18.8 | 2.52 | 2.08 | 1.61 | BH (0.57 s) | (2.320) |
| VERTEX + s407\,b2c, SS220 | 8.62 | 9.38 | 4.83 | 18.7 | 15.7 | 17.6 | 1.95 | 1.58 | 1.41 | BH (2.11 s) | (2.279) |

**Table E1** Spectral Parameters for Selected PROMETHEUS-VERTEX and PROMETHEUS-HOTB Models

**Note.** Total energies, $E_{\nu}^{\text{tot}}$, radiated in the neutrino species $\nu_i = (\bar{\nu}_e, \nu_e, \nu_x)$, mean energies, $\langle E_{\nu} \rangle$, and shape parameters, $\pi_{\nu_i}$, of the time-integrated $\nu_i$ spectra according to Equations (C7) and (C6), for the same PROMETHEUS-VERTEX models employed in Appendices C and D, for the 8.8\,M$_\odot$ ECSN (model “Sf”) from Hüdepohl et al. (2010), and for the same subset of PROMETHEUS-HOTB models shown in Figures D1 and D2. Note that $E_{\nu}^{\text{tot}} = E_{\nu}^{\text{SSNe}} + E_{\nu}^{\text{SN}} + 4\,E_{\nu}^{\text{SN}}$. For our PROMETHEUS-HOTB models, we take $E_{\nu}^{\text{SSNe}} = \xi_1 E_{\nu_{\odot}}^{\text{teff}}, (E_{\nu_{\odot}} = \lambda_{\nu_{\odot}} (E_{\nu_{\odot}}),$ and $\pi_{\nu_{\odot}} = \lambda_{\nu_{\odot}} \pi_{\nu_{\odot}}$ with the conversion factors $(\xi_1, \lambda_{\nu_{\odot}}, \lambda_{\nu_{\odot}})$; Equations (C2), (C4), (C5) according to the four cases “L,” “H,” “F,” and “S” in Table C1 (see Appendix C). The values of $\pi_{\nu_{\odot}}$ are given for our best-fit choices of the instantaneous spectral-shape parameter $C_{\text{inst}}$; i.e., $\alpha = 3.5$ for SNes with $M_{\text{SN,B}} < 1.6\,M_{\odot}$, $\alpha = 3.0$ for SNes with $M_{\text{SN,B}} > 1.6\,M_{\odot}$, and $\alpha_{\text{BH}} = 2.0$ for failed SNes; see Appendix D), as well as for the choices of $\alpha_{\text{inst}} + 0.5$ (in superscript) and $\alpha_{\text{inst}} - 0.5$ (in subscript). The NS formation cases are sorted according to the baryonic masses of the remnant NSs ($M_{\text{NS,B}}$; last column), and the failed-SN cases according to the times of BH formation ($t_{\text{BH}}$; second to last column); the baryonic NS masses at these times are listed in the last column (values in parentheses).
