Astrophysically important reactions and nucleosynthesis processes

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Abstract.
Neutrino reactions play an important role at various stages of core-collapse supernova. During infall, neutrinos are produced by electron capture mainly on nuclei and contribute significantly to the cooling of the collapsing core. After core bounce the nascent neutron star cools by neutrino emission. It is a major goal to observe such neutrinos from a future supernova by earthbound detectors and to establish their spectra. Recently it has been shown that the spectrum of electron neutrinos from the early neutrino burst is significantly altered if inelastic neutrino-nucleus scattering is considered in supernova simulations. The neutrino wind from the nascent neutron star is the site of explosive nucleosynthesis, including the recently discovered νp process and likely the r-process. Finally spallation reactions induced by neutrinos when passing through the outer burning shells can produce certain nuclides in what is called neutrino nucleosynthesis.

1. Some introductory remarks
Massive stars end their lives as type II supernovae, triggered by a collapse of their central iron core with a mass of more than $1 M_\odot$ [1]. Despite of improved description of the microphysics entering the simulations and sophisticated neutrino transport treatment spherically symmetric simulations of core-collapse supernovae fail to explode [2, 3, 4]; i.e. the energy transported to the stalled shock, which triggers the explosion, by absorption of neutrinos on free neutrons and protons is not sufficient to revive the shock wave which has run out of energy by dissociating the matter it traverses into free nucleons. However, the revival is successful in two-dimensional simulations as convection supported by hydrodynamical instabilities increases the efficiency by which neutrinos carry energy to the shock region. Successful two-dimensional simulations are reported in [5, 6, 7].

2. Electron capture in core-collapse supernovae
During most of the collapse the equation of state is given by that of a degenerate relativistic electron gas [1]; i.e. the pressure against the gravitational contraction arises from the degeneracy of the electrons in the stellar core. However, the electron chemical potential at core densities in excess of $10^8 \text{ g/cm}^3$ is of order MeV and higher, thus making electron captures on nuclei energetically favorable. As these electron captures occur at rather small momentum transfer,
the process is dominated by Gamow-Teller transitions; i.e. by the GT\_ transitions, in which a proton is changed into a neutron.

When the electron chemical potential $\mu_e$ (which grows with density like $\rho^{1/3}$) is of the same order as the nuclear $Q$-value, the electron capture rates are very sensitive to phase space and require a description of the detailed GT\_ distribution of the nuclei involved which is as accurate as possible. Furthermore, the finite temperature in the star requires the implicit consideration of capture on excited nuclear states, for which the GT\_ distribution can be different than for the ground state. It has been demonstrated [8, 9] that modern shell model calculations are capable to describe GT\_ distributions rather well [10] (an example is shown in Figure 1) and are therefore the appropriate tool to calculate the weak-interaction rates for those nuclei ($A \sim 50 - 65$) which are relevant at such densities [11]. At higher densities, when $\mu_e$ is sufficiently larger than the respective nuclear $Q$ values, the capture rate becomes less sensitive to the detailed GT\_ distribution and depends practically only on the total GT strength. Thus, less sophisticated nuclear models might be sufficient. However, one is facing a nuclear structure problem which has been overcome only very recently. Once the matter has become sufficiently neutron-rich, nuclei with proton numbers $Z < 40$ and neutron numbers $N > 40$ will be quite abundant in the core. For such nuclei, Gamow-Teller transitions would be Pauli forbidden [12] (GT\_ transitions change a proton into a neutron in the same harmonic oscillator shell) were it not for nuclear correlation and finite temperature effects which move nucleons from the $pf$ shell into the $gds$ shell. To describe such effects in an appropriately large model space (e.g. the complete $fp$-$gds$ shell) is currently only possible by means of the Shell Model Monte Carlo approach (SMMC) [13, 14]. In [15] SMMC-based electron capture rates have been calculated for more than 100 nuclei. Together with the shell model results for $fd$ shell nuclei [9] a compilation of electron rates is available which covers the composition during the early collapse phase well (until densities to a few $10^{10}$ g/cm\(^3\)). However, at even higher densities the continuous electron capture drives the matter composition to more neutron-rich and heavier nuclei than considered in [15]. The neglect of these nuclei could lead to a systematic overestimate of the capture rates as the neglected nuclei have larger $Q$ values and enhanced Pauli blocking due to increased neutron excess than the nuclei considered. Unfortunately an SMMC evaluation of the several thousands of nuclei present in the matter composition during the late phase of collapse before neutrino trapping is numerically not feasible. Hence a simpler approach has to be adopted. This is based on the observation that the single particle occupation numbers obtained in the SMMC calculations can be well approximated by a parametrized Fermi-Dirac distribution. By adjusting the parameters of this distribution to the SMMC results for about 250 nuclei, occupation numbers were derived for nearly 3000 nuclei and the respective individual rates have been calculated within an RPA calculation based on these partial occupations. It is found that this simplified approach reproduces the SMMC results quite well at electron chemical potential $\mu_e > 15$ MeV [16]. This corresponds to the density regime where the previously neglected nuclei become abundant in the matter composition. At smaller electron chemical potentials (i.e. at lower densities) details of the GT\_ distribution are important which are not recovered by the simple parametrized approach. However, at these conditions electron capture is dominated by nuclei for which individual shell model rates exist.

Fig. 2 compares the capture rate derived for the pool of more than 3000 nuclei (i.e. combining the rates from shell model diagonalization, SMMC, and from the parametrized approach) with those obtained purely on the basis of the shell model results. While the agreement is excellent at small electron chemical potentials (here the shell model rates dominate), the rates for the large pool are slightly smaller at higher $\mu_e$ values due the presence of neutron-rich heavy nuclei with smaller individual rates. Furthermore the new rates also include plasma screening effects which lead to an increase of the effective $Q$ values and a reduction of the electron chemical potential, which both reduce the electron capture rates. However, the effect is rather mild and does not alter the conclusion that electron capture on nuclei dominates over capture on protons during
the collapse.

**Figure 1.** Comparison of the measured $^{51}$V(d,$^2$He)$^{51}$Ti cross section at forward angles (which is proportional to the GT$_+$ strength) with the shell model GT distribution in $^{51}$V (from [21]).

**Figure 2.** Comparison of NSE-averaged electron capture rates calculated for about 3000 individual nuclei (solid, see text) with those obtained for the restricted set of nuclei (dashed) considered in [15] (from [16]).

The shell model capture rates have significant impact on collapse simulations. In the presupernova phase ($\rho < 10^{10}$ g/cm$^3$) the captures proceed slower than assumed before and for a short period during silicon burning $\beta$-decays can compete [17, 18]. As a consequence,
the core is cooler, more massive and less neutron-rich before the final collapse. However, until recently simulations of this final collapse assumed that electron captures on nuclei are prohibited by the Pauli blocking mechanism, mentioned above. However, based on the SMMC calculations it has been shown in [15] that capture on nuclei dominates over capture on free protons. The changes compared to the previous simulations are significant [15, 19, 20]. Importantly the shock is now created at a smaller radius with more infalling material to traverse, but also the density, temperature and entropy profiles are strongly modified [19].

The possibility to capture electrons on nuclei and free protons leads to a strong self-regulation during the collapse and it is observed that the collapse approaches the same core trajectory for stars between 11 and 25 solar masses before neutrino trapping and thermalization is reached at densities around $10^{12}$ g/cm$^3$ (Fig. 3, [20]). Thus it is expected that supernovae, independent of the stellar mass, have a quite similar neutrino burst spectrum. This spectrum, however, has also significantly changed after inelastic neutrino-nucleus reactions have been included in the simulations.

![Figure 3](image-url)  
**Figure 3.** The evolution of the electron-to-baryon ratio $Y_e$ and the lepton-to-baryon ratio $Y_{lep}$ as function of central density for stars with masses between 11 and 25 $M_{\odot}$. (from [20]).

![Figure 4](image-url)  
**Figure 4.** Comparison of the normalized neutrino spectra, arising from the $\nu_e$ burst shortly after bounce, without (solid) and with (dashed) consideration of inelastic neutrino-nucleus scattering in the supernova simulations. (from [25]).

### 3. Inelastic neutrino-nucleus reactions

When, with progressing collapse, matter gets denser, neutrino interactions with matter become increasingly important and require careful treatment which is usually done within Boltzmann transport theory [3, 4]. Here elastic neutrino-nucleus and inelastic neutrino-electron scattering have been identified as the sources for neutrino trapping (which occurs at densities around a few $10^{11}$ g/cm$^3$) and neutrino thermalization. Recently inelastic neutrino-nucleus scattering, as another mode of energy exchange between neutrinos and matter, has been included for the first time in supernova simulations. The relevant cross sections have been calculated based on large-scale shell model calculations for the allowed GT transitions and within the random phase approximation for forbidden transitions [22], taking special care of finite temperature effects [23]. At low and modest neutrino energies $E_\nu$, the cross sections are dominated by GT0 contributions for which the shell model has been validated by detailed comparison to precision M1 data derived from electron scattering on spherical nuclei which are mainly due to the same isovector response [24].

Although inelastic neutrino-nucleus scattering contributes to the thermalization of neutrinos with the core matter, the inclusion of this process has no significant effect on the collapse
trajectories. However, it increases noticeably the opacity for high-energy neutrinos in the $\nu_e$ neutrino burst just after the bounce. As these neutrinos excite the nuclei, they are down-scattered in energy, in this way significantly reducing the high-energy tail of the spectrum of emitted supernova neutrinos (see Fig. 4, [25]). This makes the detection of supernova neutrinos by earthbound detectors more difficult, as the neutrino detection cross section scales with $E_{\nu}^2$.

4. Supernova nucleosynthesis

When in a successful explosion the shock passes through the outer shells, its high temperature induces an explosive nuclear burning on short time-scales. This explosive nucleosynthesis can alter the elemental abundance distributions in the inner (silicon, oxygen) shells. More interestingly the shock has sufficiently high temperatures that nuclear binding cannot withstand and matter is ejected as free protons and neutrons from the neutron star surface. As the outflow is adiabatic, the ejected matter reaches cooler regions at larger distances and nuclei can be assembled from free nucleons. The abundance outcome obviously depends strongly on the ratio of neutrons and protons which is set by the competition of neutrino and anti-neutrino absorption on free nucleons.

Recently explosive nucleosynthesis has been investigated consistently within supernova simulations, where a successful explosion has been enforced by slightly enlargening the neutrino absorption cross section on nucleons or the neutrino mean-free path, which both increase the efficiency of the energy transport to the stalled shock. The results presented in [26, 27] showed that in an early phase after the bounce the ejected matter is actually proton-rich as the $\nu_e$ neutrinos have sufficiently large energies to drive the matter proton-rich. In later stages, i.e. a few seconds after bounce, the neutrino opacities in the neutron-rich matter ensure that $\bar{\nu}_e$ have larger average energies than $\nu_e$, and it is believed that $\bar{\nu}_e$ absorption on protons dominates driving the matter neutron-rich (allowing for the r-process to occur).

4.1. The $\nu p$ process

The studies presented in [26, 27, 28] show that matter with $Y_e$ larger than 0.5 will always be present in core-collapse supernovae explosions with ejected matter irradiated by a strong neutrino flux, independently of the details of the explosion. As this proton-rich matter expands and cools, nuclei can form resulting in a composition dominated by $N = Z$ nuclei, mainly $^{56}\text{Ni}$ and $^4\text{He}$, and protons. Without the further inclusion of neutrino and antineutrino reactions the composition of this matter will finally consist of protons, alpha-particles, and heavy (Fe-group) nuclei (in nucleosynthesis terms a proton- and alpha-rich freeze-out), with enhanced abundances of $^{45}\text{Sc}$, $^{49}\text{Ti}$, and $^{64}\text{Zn}$ [26, 27]. In these calculations the matter flow stops at $^{64}\text{Ge}$ with a small proton capture probability and a beta-decay half-life (64 s) that is much longer than the expansion time scale ($\sim 10$ s) [26].

As noted by Martinez-Pinedo and explored in [29] the synthesis of nuclei with $A > 64$ can be obtained, if one explores the previously neglected effect of neutrino interactions on the nucleosynthesis of heavy nuclei. $N \sim Z$ nuclei are practically inert to neutrino capture (converting a neutron into a proton), because such reactions are endoergic for neutron-deficient nuclei located away from the valley of stability. The situation is different for antineutrinos that are captured in a typical time of a few seconds, both on protons and nuclei, at the distances at which nuclei form ($\sim 1000$ km). This time scale is much shorter than the beta-decay half-life of the most abundant heavy nuclei reached without neutrino interactions (e.g. $^{56}\text{Ni}$, $^{64}\text{Ge}$). As protons are more abundant than heavy nuclei, antineutrino capture occurs predominantly on protons, causing a residual density of free neutrons of $10^{14}$–$10^{15}$ cm$^{-3}$ for several seconds, when the temperatures are in the range 1–3 GK. The neutrons produced via antineutrino absorption on protons can easily be captured by neutron-deficient $N \sim Z$ nuclei (for example $^{64}\text{Ge}$), which have large neutron capture cross sections. The amount of nuclei with $A > 64$ produced is
Figure 5. Elemental abundance yields (normalized to solar) for elements produced in the proton-rich environment shortly after the supernova shock formation. The matter flow stops at nuclei like $^{56}\text{Ni}$ and $^{64}\text{Ge}$ (open circles), but can proceed to heavier elements if neutrino reactions are included during the network (full circles) (from [29]).

then directly proportional to the number of antineutrinos captured. While proton capture on $^{64}\text{Ge}$ takes too long, the $(n, p)$ reaction dominates (with a lifetime of 0.25 s at a temperature of 2 GK), permitting the matter flow to continue to nuclei heavier than $^{64}\text{Ge}$ via subsequent proton captures.

Figure 6. Light $p$-nuclei abundances in comparison to solar abundances as a function of $Y_e$. The $Y_e$-values given are the ones obtained at a temperature of 3 GK that corresponds to the moment when nuclei are just formed and the $\nu p$-process starts to act. (from [29]).
Fröhlich et al. argue that all core-collapse supernovae will eject hot, explosively processed matter subject to neutrino irradiation and that this novel nucleosynthesis process (called $\nu p$ process) will operate in the innermost ejected layers producing neutron-deficient nuclei above $A > 64$. However, how far the mass flow within the $\nu p$ process can proceed, strongly depends on the environment conditions, most notably on the $Y_e$ value of the matter. Obviously the larger $Y_e$, the larger the abundance of free protons which can be transformed into neutrons by antineutrino absorption. Fig. 6 shows the dependence of the $\nu p$-process abundances as a function of the $Y_e$ value of the ejected matter. Nuclei heavier than $A = 64$ are only produced for $Y_e > 0.5$, showing a very strong dependence on $Y_e$ in the range 0.5–0.6. A clear increase in the production of the light $p$-nuclei, $^{92,94}\text{Mo}$ and $^{96,98}\text{Ru}$, is observed as $Y_e$ gets larger. Thus the $\nu p$ process offers the explanation for the production of these light $p$-nuclei, which was yet unknown. It might also explain the presence of strontium in the extremely metal-poor, and hence very old, star HE 1327-2326 [30]. The $\nu p$ process results of Fröhlich et al. have been recently confirmed by Pruet et al. [31], in a study of the nucleosynthesis that occurs in the early proton-rich neutrino wind.

4.2. The r-process
About half of the elements heavier than mass number $A \sim 60$ are made within the $r$-process, a sequence of rapid neutron captures and $\beta$ decays [32, 33]. The process is thought to occur in environments with extremely high neutron densities. Then neutron captures are much faster than the competing decays and the $r$-process path runs through very neutron-rich, unstable nuclei. Once the neutron source ceases, the process stops and the produced nuclides decay towards stability producing the neutron-rich heavier elements. Many parameter studies of the reaction network, aimed at reproducing the observed $r$-process abundances, have progressed our general understanding. So it is generally accepted that the $r$-process operates in an environment with temperatures of order $10^{9}$ K and with neutron densities in excess of $10^{20}/\text{cm}^3$ [34]. Simulations of the $r$-process reaction network indicate that under such conditions for temperature and neutron densities the reaction path runs through extremely neutron-rich nuclei with separation energies, $S_n < 2 - 3$ MeV [34], for which most of the required properties (i.e. mass, lifetime and neutron capture cross sections) are experimentally unknown. Thus, these quantities have to be modelled, based on experimental guidance. As we will discuss in the next subsection some progress has been achieved recently, but the real breakthrough is expected from future radioactive ion-beam facilities. Finally we mention that the observed peaks in the $r$-process abundance distribution are related to the magic neutron numbers, $N = 50, 82, 126$, where the $r$-process runs through several, rather long $\beta$ decays at $N = 50, 82$ and 126 (the respective nuclei are called $r$-process waiting points) and a relatively large amount of matter is built up.

Despite many promising attempts the actual site of the $r$-process has not been identified yet. However, parameter studies have given clear evidence that the observed $r$-process abundances cannot be reproduced at one site with constant temperature and neutron density [35]. Thus the abundances require a superposition of several (at least three) $r$-process components. This likely implies a dynamical $r$-process in an environment in which the conditions change during the duration of the process. The currently favored $r$-process sites (type II supernovae [36] and neutron-star mergers [37]) offer such dynamical scenarios. However, recent meteoritic clues might even point to more than one distinct site for our solar $r$-process abundance [38]. The same conclusion can be derived from the observation of $r$-process abundances in low-metallicity stars (e.g. [39]), a milestone of $r$-process research.

4.2.1. Nuclear r-process input
Arguably the most important nuclear ingredient in $r$-process simulations are the nuclear masses as they determine the flow-path. Unfortunately nearly all
of them are experimentally unknown and have to be theoretically estimated. Traditionally this is done on the basis of parametrizations to the known masses. Although these empirical mass formulae achieve rather remarkable fits to the data (the standard deviation is of order 700 keV), extrapolation to unknown masses appears less certain and different mass formulae can predict quite different trends for the very neutron-rich nuclei of relevance to the r-process. The most commonly used parametrizations are based on the finite range droplet model (FRDM), developed by Möller and collaborators [40], and on the ETFSI (Extended Thomas-Fermi with Strutinski integral) model of Pearson [41]. A new era has been opened very recently, as for the first time, nuclear mass tables have been derived on the basis of nuclear many-body theory (Hartree-Fock-Bogoliubov model) [42, 43, 44] rather than by parameter fit to data.

The nuclear half-lives strongly influence the relative r-process abundances. In a simple $\beta$-equilibrium picture the elemental abundance is proportional to the half-life, with some corrections for $\beta$-delayed neutron emission [45]. As r-process half-lives are longest for the magic nuclei, these waiting point nuclei determine the minimal r-process duration time; i.e. the time needed to build up the r-process peak around $A \approx 200$. We note, however, that this time depends also crucially on the r-process path and can be as short as a few 100 milliseconds if the r-process path runs close to the neutron dripline.

There are a few milestone half-life measurements including the N=50 waiting point nuclei $^{80}\text{Zn}$, $^{79}\text{Cu}$, and $^{78}\text{Ni}$, and the N=82 waiting point nuclei $^{130}\text{Cd}$ and $^{129}\text{Ag}$ [46, 47]. Although no half-lives for $N = 126$ waiting points have yet been determined, there has been decisive progress towards this goal recently [48].

These data play crucial roles in constraining and testing nuclear models which are still necessary to predict the bulk of half-lives required in r-process simulations. It is generally assumed that the half-lives are dominated by allowed Gamow-Teller (GT) transitions, with forbidden transitions contributing noticeably for the heavier r-process nuclei [49]. The $\beta$ decays only probe the weak low-energy tail of the GT distributions and provide quite a challenge to nuclear modeling as they are not constrained by sumrules. Traditionally the estimate of the half-lives are based on the quasiparticle random phase approximation on top of the global FRDM or ETFSI models. Recently half-lives for selected (spherical) nuclei have been presented using the QRPA approach based on the microscopic Hartree-Fock-Bogoliubov method [50] or a global density functional [51, 52]; in particular the later approach achieved quite good agreement with data for spherical nuclei in different ranges of the nuclear chart. Applications of the interacting shell model [53, 54, 55] have yet been restricted to waiting point nuclei with magic neutron numbers. Here, however, this model which accounts for correlations beyond the QRPA approach, obtains quite good results (for an example see Fig. 7). We remark that, besides a good description of the allowed (and forbidden) transition matrix elements, an accurate reproduction of the $Q_n$ values is required from the models.

Further r-process simulations require rates for neutron capture, and for the various fission processes (neutron-induced, beta-delayed, spontaneous, perhaps neutrino induced) together with the corresponding yield distributions. If the r-process occurs in strong neutrino fluences, different neutrino-induced charged-current (e.g. ($\nu_e$, $e^-$)) and neutral-current (e.g. ($\nu$, $\nu'$)) reactions, which are often accompanied by the emission of one or several neutrons [56, 57, 58], have to be modelled and included as well. We mention that the occurrence of low-lying dipole strength in neutron rich nuclei, as predicted by nuclear models [59], can have noticeable effects on neutron capture cross sections [60, 61].

### 4.3. R-process simulations and the role of fission

Recently r-process simulations, which for the first time included detailed fission yield distributions derived on the basis of the ablation-ablation model code ABLA [62], have been performed [63, 64]. The studies have been performed assuming the neutrino-driven wind scenario...
Figure 7. top: Comparison of various theoretical halflife predictions with data for the $N = 82$ r-process waiting points. Bottom: Ratio of forbidden and allowed Gamow-Teller contributions to the halflives calculated with the Fayans density functional (courtesy I.N. Borzov).

(i.e. matter ejected from the neutron star in a core-collapse supernova) as the r-process site. The entropy of the ejected matter, or equivalently the neutron-to-seed ($n/s$) ratio, has been treated as a free parameter. To explore the potential effect of fission on the r process, this ratio has to be relatively large, ($n/s > 250$). If this condition is fulfilled, then the r-process runs quite fast along a path close to the dripline reaching heavy nuclei with neutron numbers close to the magic number $N = 184$ and above. Once the neutron source is used, i.e. the ejected matter has reached radii where the neutron number density is too low to support fast neutron captures, the produced nuclides decay. For the heavy nuclei the dominant decay mode
Figure 8. R-process abundances obtained in simulations which explore the potential impact of fission. The calculations are performed for different neutron-to-seed ratios and two different mass tabulations. The solid line is the solar r-process abundance distribution. (from [64])

is neutron-induced fission which via their fission decay products produce peaks around mass numbers \( A \sim 130 \) and 195, corresponding to the second and third peaks in the observed r-process abundance distributions (Fig. 8). These simulations clearly stress the importance of the nuclear masses. While the studies, which use the FRDM masses, give quite similar abundance distributions for different \( n/s \) ratios once a threshold value is overcome, the calculations based on the ETFSI masses yield abundance results which vary strongly with the assumed \( n/s \) value. This strikingly different behavior can be traced back to differences in the predicted masses for very neutron rich nuclei. The FRDM mass tabulation predicts that certain nuclei just above the magic neutron number \( N = 82 \) act as obstacles for the r-process matter flow holding back material and ensuring that more free neutrons are available when part of the matter reaches heavy nuclei with \( N = 184 \) to guarantee production of nuclides even beyond these waiting points. The ETFSI mass tabulation does not identify such obstacles on the r-process path and, as no matter is kept back at relatively small mass numbers, less neutrons are available once the matter flow reaches \( N = 184 \) suppressing matter flow beyond this waiting point [64]. We will know which of the two mass models is more realistic once masses of very neutron-rich nuclei will be measured; this is one of the important aims at radioactive ion-beam facilities like FAIR. It should be stressed that observations of r-process abundances in old halo stars in our galaxy show always the same pattern between mass numbers \( A = 130 \) and 195 (Fig. 9), while the patterns differ for \( A < 130 \). It is intriguing that these old stars in completely different locations of the Milky Way, which have witnessed only a few, and importantly not the same, supernova outbursts, show identical r-process patterns between the second and third peaks as obtained when averaged over the chemical history of the galaxy (solar r-process abundance). The fact that r-process simulations including fission and based on the FRDM mass model yield just such a behavior is certainly interesting. However, the nuclear and astrophysical input in these simulations is yet too uncertain to draw any conclusions. including fission and based on the FRDM mass model yield just such a behavior is certainly interesting. However, the nuclear and astrophysical input in these simulations is yet too uncertain to draw any conclusions.

4.4. The neutrino process

When neutrinos, produced in the hot supernova core, pass through the outer shells of the star, they can induce nuclear reactions and in this way contribute to the elementsynthesis (the \( \nu \)-process, [65]). For example, the nuclides \(^{11}\text{B}\) and \(^{19}\text{F}\) are produced by \((\nu,\nu'\bar{n})\) and \((\nu,\nu'p)\) reactions on the quite abundant nuclei \(^{12}\text{C}\) and \(^{20}\text{Ne}\). These reactions are dominantly induced by \(\nu_\mu\) and \(\nu_\tau\) neutrinos and their antiparticles (combined called \(\nu_x\) neutrinos) which have larger
average energies (about 20 MeV) than $\nu_e$ and $\bar{\nu}_e$ neutrinos. As found in detailed stellar evolution studies [66] the rare odd-odd nuclides $^{138}$La and $^{180}$Ta are mainly made by the charged-current reaction $^{138}$Ba($\nu_e$, $e^-$) $^{138}$La and $^{180}$Hf($\nu_e$, $e^-$) $^{180}$Ta. Hence, the $\nu$-process is potentially sensitive to the spectra and luminosity of $\nu_e$ and $\nu_x$ neutrinos, which are the neutrino types not observed from SN1987a. Both charged-current cross sections are dominated by the low-energy tail of the GT$^-$ strength distribution which is notoriously difficult to simulate. As large-scale shell model calculations are yet not possible for $^{138}$Ba or $^{180}$Hf, the charged-current cross sections had been evaluated within the random phase approximation [66]. As a major improvement it has been possible last year to measure the GT$^-$ strengths on $^{138}$Ba and $^{180}$Hf below the particle thresholds and to convert these data into the relevant ($\nu_e$, $e^-$) cross sections [67]. It is found that the new cross sections are somewhat larger than the RPA predictions. Upon using these results in stellar evolution models, slightly larger $^{138}$La and $^{180}$Hf abundances are found, which agree reasonably well with observation [67].

We mention that neutrino nucleosynthesis studies are quite evolved requiring state-of-the-art stellar models with an extensive nuclear network [66]. In the first step stellar evolution and nucleosynthesis is followed from the initial hydrogen burning up to the presupernova models. The post-supernova treatment then includes the passage of a neutrino fluence through the outer layers of the star, followed by the shock wave which heats the material and also induces noticeable nucleosynthesis, mainly by photodissociation (the $\gamma$ process). Modelling of the shock heating is quite essential as the associated $\gamma$ process destroys many of the daughter nuclides previously produced by neutrino nucleosynthesis.

Neutrino nucleosynthesis has recently also been studied as a possible thermometer for supernova neutrinos, which is particularly interesting under the assumption of certain oscillation scenarios [68].

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Figure 9. Observed r-process abundances in three old stars in the galactical halo, compared to the solar r-process abundance (blue line) which reflects the chemical history of the galaxy. For better viewing the different curves are off-set by arbitrary numbers. (courtesy John Cowan)
5. Conclusion

Supernova simulations are a quite interdisciplinary task. In recent years significant progress has been achieved due to improved astrophysical modelling and based on more reliable nuclear inputs. The latter is mainly concerned with weak-interaction processes like electron capture on nuclei, beta decays and neutrino-induced reactions. Here it has been possible to derive the relevant rates on the basis of modern nuclear many-body models, constrained and guided by recent measurements of nuclear Gamow-Teller and M1 strengths functions. The improved nuclear inputs have led to quite significant changes in the core-collapse dynamics and the supernova neutrino spectra. Nuclear network simulations, consistently coupled to trajectories of state-of-the-art supernova simulations, have yielded new insights in supernova nucleosynthesis, resulting for example to the discovery of the $\nu p$ process which produces proton-rich nuclides and might be the searched-for source of the light p-nuclides. Progress in the understanding of the r-process is expected from future radioactive ion-beam facilities which will help to remove the large uncertainties on the nuclear input data which still hamper current r-process simulations.

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