NEUTRINO-INDUCED NEUTRON SPALLATION AND THE SITE OF THE $r$-PROCESS

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All of the actinides and roughly half the natural abundance of elements with mass number $A > 70$ come from the rapid neutron capture process, or the $r$-process. If the $r$-process, as suggested by many, occurs deep in a supernova, then it is under the influence of an intense neutrino flux. Here we discuss the effects of both charged-current and neutral-current neutrino interactions on the $r$-process. We show that the multiple-neutron emission induced by both kinds of neutrino interactions can affect the observed $r$-process abundance pattern significantly. In particular, we find that five nuclei below the $r$-process abundance peak at $A \sim 195$ may be entirely attributed to the neutrino-induced neutron spallation processes following the $r$-process freeze-out. Furthermore, the deduced neutrino fluence agrees with the conditions in the recent supernova $r$-process model. These results strongly argue that the $r$-process occurs in supernovae. They also provide a sensitive probe for the conditions at the supernova $r$-process site.

1 Introduction

It is well known that almost all of the $^4$He nuclei in our universe were created in the Big Bang. We also know that stars shine by burning $^1$H into successively heavier elements. However, nuclear physics tells us that $^{56}$Fe is the most tightly bound nucleus. No more nuclear binding energy can be released to power the star by burning Fe. Therefore, elements heavier than $^{56}$Fe have to be made in some other processes. One such process is the rapid neutron capture process, or the $r$-process for short. This process is responsible for all of the actinides and about half the natural abundance of elements with mass number $A > 70$. Here we are mainly interested in the $r$-process nuclei near $A \sim 195$. Typical $r$-process elements in this mass region are Os, Pt, and Au.

The $r$-process mechanism is rather simple because there are only three types of reactions involved. A nucleus can increase its mass number by capturing neutrons. If the temperature is high enough, photodisintegration competes with neutron capture. In any case, if a nucleus becomes too neutron rich, it can also $\beta$-decay. By definition, neutron capture is much faster than $\beta$-decay during the $r$-process. Furthermore, the $r$-process mechanism can explain the observational data quite well. There are two prominent peaks at $A \sim 130$ and 195, respectively, in the observed solar $r$-process abundance pattern. The existence of these two peaks can be explained as follows.
At the beginning of the \( r \)-process, there are some seed nuclei and lots of neutrons. The seed nucleus then captures neutrons and moves towards the neutron-drip line. However, at some point, the binding energy of the next neutron to be captured becomes so small that it will be quickly disintegrated by the photons available in the \( r \)-process environment. At this so-called "waiting-point," the nucleus must \( \beta \)-decay to a new nuclear species before further neutron capture can proceed. Through such a series of neutron capture and \( \beta \)-decay, a distribution of progenitor nuclei far away from \( \beta \)-stability is produced on the \( r \)-process path. Clearly, the progenitor abundance at a given charge number \( Z \) is piled up at the corresponding waiting-point nucleus before it \( \beta \)-decays, and the more slowly it \( \beta \)-decays, the more abundant it will be. Due to the extreme stability of the closed neutron shells, the \( \beta \)-decay rates for the progenitor nuclei at the magic neutron numbers \( N = 82 \) and 126 are extremely small. Consequently, abundance peaks are formed at the progenitor nuclei with these magic neutron numbers. After neutron capture stops, these progenitor nuclei successively \( \beta \)-decay to stability and give rise to the peaks at \( A \sim 130 \) and 195 in the observed solar \( r \)-process abundance pattern.

Although the above \( r \)-process theory succeeds in explaining the main features of the observed solar \( r \)-process abundance pattern, it does not elaborate on where the \( r \)-process occurs. Perhaps the most popular site for the \( r \)-process is the interior of a supernova. However, a crucial step to identify supernovae as the \( r \)-process site is to find something hidden in the observed \( r \)-process abundance pattern that has the signature of a supernova. Probably one of the best known signatures of a supernova is its powerful neutrino emission, and as we show in the following discussion, we can indeed extract some fingerprints of neutrinos from the observed \( r \)-process abundance pattern. In particular, we find that five nuclei below the abundance peak at \( A \sim 195 \) may be entirely produced by the neutrino-induced neutron spallation processes after the rapid neutron capture stops (i.e., following the \( r \)-process freeze-out). Our discussion is organized as follows. In Sec. 2, we briefly describe the emission of supernova neutrinos and the characteristics of their interactions with the neutron-rich progenitor nuclei produced during the \( r \)-process. In Sec. 3, we discuss the effects of neutrino-induced neutron spallation on the \( r \)-process and identify the five nuclei that are most sensitive to these effects. We discuss the implications of our results for the \( r \)-process site in Sec. 4.

2 Neutrino Emission and Interactions in Supernovae

A supernova occurs when the core of a massive star collapses into a compact neutron star with a mass of \( \sim 1 \, M_\odot \) and a radius of \( \sim 10 \, \text{km} \). The grav-
itational binding energy of the final neutron star is \( \sim 10^{53} \) erg. Due to the high temperatures and densities encountered during the collapse, the most efficient way to release this energy is to emit neutrinos and antineutrinos mainly through electron-positron pair annihilation. In fact, because of the intense elastic neutral-current scatterings on free nucleons for all neutrino species, even neutrinos have to diffuse out of the neutron star on a timescale of \( \sim 10 \) s, as confirmed by the detection of neutrinos from SN1987a. Since these neutrinos are in thermal equilibrium with the neutron star matter and with each other for most part of the diffusion process, the individual neutrino luminosities are about the same and have an average value of \( L_\nu \sim 10^{51} \) erg s\(^{-1}\).

However, the average neutrino energies are very different. This is because these neutrinos have different abilities to exchange energy with matter. The elastic neutral-current scatterings on free nucleons, which govern the diffusion of all neutrinos, essentially cost neutrinos no energy because the nucleon rest mass is much higher than the typical neutrino energies. All neutrino species can exchange energy with matter through scatterings on electrons. However, only \( \nu_e \) and \( \bar{\nu}_e \) are energetic enough to have charged-current capture reactions on free neutrons and protons, respectively. Therefore, \( \nu_\mu(\tau) \) and \( \bar{\nu}_\mu(\tau) \) have the weakest ability to exchange energy with matter and decouple from thermal equilibrium at the highest temperatures and densities as they diffuse towards the neutron star surface. Correspondingly, they have the highest average energy. Between \( \nu_e \) and \( \bar{\nu}_e \), the \( \nu_e \) have more chances to exchange energy with matter because there are more neutrons than protons inside a neutron star. Consequently, the \( \nu_e \) decouple at the lowest temperatures and densities, and have the lowest average energy. The average \( \bar{\nu}_e \) energy lies in the middle. Typically, the average neutrino energies are \( \langle E_{\nu_e} \rangle \approx 11 \) MeV, \( \langle E_{\bar{\nu}_e} \rangle \approx 16 \) MeV, and \( \langle E_{\nu_\mu(\tau)} \rangle = \langle E_{\bar{\nu}_\mu(\tau)} \rangle \approx 25 \) MeV.

With such powerful neutrino emission, neutrino interactions with the progenitor nuclei can be very important if the \( r \)-process indeed occurs in supernovae. The charged-current \( \nu_e \) capture reaction and the inelastic neutral-current scatterings for \( \nu_\mu(\tau) \) and \( \bar{\nu}_\mu(\tau) \) are of particular interest. For the average supernova neutrino energies given previously, charged-current capture reactions are energetically forbidden for \( \nu_\mu(\tau) \) and \( \bar{\nu}_\mu(\tau) \). However, these neutrinos dominate the inelastic neutral-current scatterings because they have the highest average energy and the inelastic neutral-current scattering cross section of interest is strongly energy dependent. On the other hand, the charged-current \( \nu_e \) capture reaction is unimportant because it changes a proton into a neutron, and all the allowed transitions are Pauli blocked for the extremely neutron-rich progenitor nuclei made in the \( r \)-process. Obviously, the neutrino interaction rates scale with the neutrino flux (i.e., \( \propto L_\nu/\tau^2 \)). For the progenitor nuclei
at $A \sim 195$, the charged-current $\nu_e$ capture rate is $\sim 8 \text{ s}^{-1}$ and the total neutral-current scattering rate is $\sim 12 \text{ s}^{-1}$ at a distance of $r = 100$ km from the neutron star for a luminosity of $L_\nu = 10^{51} \text{ erg s}^{-1}$ per neutrino species.

With an average energy of $\sim 10$ MeV, the $\nu_e$ mainly put the daughter nucleus to the excited states at $\sim 10$ MeV above the ground state of the parent nucleus through the charged-current capture reaction. A typical ground state energy difference between the parent and daughter nuclei is also $\sim 10$ MeV. So typically the daughter nucleus has an excitation energy of $\sim 20$ MeV above its ground state. Because the progenitor nuclei made in the $r$-process are extremely neutron rich, their daughter nuclei deexcite almost exclusively through neutron emission. The neutron binding energy ranges from $\sim 3$ MeV for the progenitor nuclei to $\sim 8$ MeV for the stable nuclei. On the average, $\sim 4$ neutrons are emitted after each $\nu_e$ capture reaction, and the branching ratios for emitting different numbers of neutrons can be estimated using statistical techniques.

As for the $\nu_\mu(\tau)$ and $\bar{\nu}_\mu(\tau)$, they can excite the progenitor nucleus to the states near the neutron emission threshold or to those at $\sim 15$ MeV above the ground state through the inelastic neutral-current scattering. On the average, $\sim 2$ neutrons are emitted after each inelastic neutral-current scattering, and various neutron emission branching ratios can be estimated as in the charged-current case.

### 3 Effects of Neutrino-Induced Neutron Spallation on the $r$-Process

During the dynamic phase of the $r$-process, neutrino-induced neutron spallation is unimportant because the typical rates for neutron capture and photodisintegration are faster than those for the neutrino reactions by orders of magnitude. When the neutron number density drops below a critical level, rapid neutron capture stops and the progenitor abundance pattern freezes out. Ideally, this progenitor abundance pattern would lead to the observed abundance pattern. However, standard $r$-process calculations without neutrino effects have a major deficiency. For example, these calculations greatly underproduce the nuclei at $A \sim 182$–187 below the abundance peak at $A \sim 195$. This is because these calculations give a generic progenitor abundance pattern that lacks the corresponding progenitor nuclei.

Due to the competition between neutron capture and photodisintegration, it turns out that the progenitor nuclei on the $r$-process path have roughly the same neutron binding energy. In general, at a given charge number $Z$ the neutron binding energy decreases with increasing neutron number $N$, and at a given $N$ it increases with increasing $Z$. Therefore, the same neutron binding energy is reached at a larger $N$ as $Z$ increases. Due to the existence
of the closed neutron shell at $N = 126$, the neutron binding energy for a given $Z$ decreases very slowly as $N$ increases towards 126. As a result, when $Z$ increases from 63 to 64, $N$ jumps from 118 to 124 for the progenitor nuclei on the $r$-process path, and the progenitor nuclei at $A \sim 182–187$ are not produced in significant abundance during the $r$-process. This then leads to the deficiency in the corresponding mass region of the final abundance pattern.

However, if neutrino reactions are significant after the $r$-process freeze-out, the final abundance at a given $A$ will receive contributions from progenitor nuclei at higher mass numbers due to the neutrino-induced neutron spallation. We refer to this effect of neutrino interactions on the $r$-process abundances as “neutrino postprocessing.” With neutrino postprocessing, the final abundance $Y_f(A)$ at mass number $A$ is given by

$$Y_f(A) = \sum_{n=0}^{n_{\text{max}}} P_n Y_{\text{pro}}(A + n),$$

where $Y_{\text{pro}}(A + n)$ is the progenitor abundance at mass number $A + n$, and $P_n$ is the probability for emitting a total of $n$ neutrons during postprocessing. Provided that we have the right form of these postprocessing neutron emission probabilities, it is possible that we can reproduce the observed abundance pattern everywhere despite the deficiency in the progenitor abundance pattern.

The postprocessing neutron emission probability $P_n$ can be calculated as follows. First of all, we note that there are different ways to emit a total of $n$ neutrons. For example, a total of two neutrons can be emitted by one event emitting two neutrons or by two events each of which emits one neutron. The occurrence of an event emitting $i$ neutrons is governed by the Poisson distribution

$$P(N_i) = \frac{\langle N_i \rangle^N_i}{N_i!} \exp(-\langle N_i \rangle),$$

where $\langle N_i \rangle$ is the average number of events emitting $i$ neutrons. The post-processing neutron emission probability $P_n$ can then be obtained by listing the ways to emit a total of $n$ neutrons and summing up the corresponding probabilities calculated from the Poisson distributions.

In Eq. (2), $\langle N_i \rangle$ depends on the average number of neutrino interactions during postprocessing and the branching ratio for emitting $i$ neutrons after each neutrino interaction. The neutrino interaction rates and the associated neutron emission branching ratios have been discussed in Sec. 2. More details were given elsewhere. The average number of neutrino interactions also depends on the neutrino fluence

$$\mathcal{F} = \int_{t_{10}}^{\infty} \frac{L_\nu(t)}{r(t)^2} dt,$$
where $t_{fo}$ is the time at the $r$-process freeze-out.

The postprocessing neutron emission probabilities for the progenitor nuclei at $A \sim 195$ are plotted for $F/(10^{47} \text{ erg km}^{-2}) = 0.015, 0.030, \text{ and } 0.060$ in Fig. 1 (hereafter, the values of $F$ are always given in unit of $10^{47} \text{ erg km}^{-2}$). Two features of these probabilities are worth noticing. The first feature is that for all three chosen neutrino fluences, the probability for emitting zero neutron is the largest. This is because for these fluences, the average number of neutrino interactions is $\sim 1$ or less. Consequently, according to Poisson statistics, the probability for no interaction at all is large. The second feature is that these probabilities have bumps at $n = 4$ and 5. This is because the neutron emission characteristics of charged-current and neutral-current neutrino interactions are very different. On the average, $\sim 4$ neutrons are emitted after each charged-current interaction, twice as many as in the neutral-current case. The superposition of these two kinds of interactions then leads to the bumps at $n = 4$ and 5.

These postprocessing neutron emission probabilities, which only depend on the neutrino fluence, are needed to test if the observed abundances at $A \sim 182–187$ can be produced by neutrino postprocessing. The progenitor abundances at $A > 187$ are also needed for this test. However, since the observed abundances at $A > 187$ should be reproduced with the same neutrino fluence, a given number of these progenitor abundances can be obtained from the same number of constraints on the final abundances after postprocessing [see Eq. (1)]. Therefore, we are trying to find one neutrino fluence that can fit a number of abundances at $A \leq 187$ in this test.

Our results are shown in Fig. 2. In this figure, the data points, some with error bars, give the observed solar $r$-process abundances. The solid line gives the progenitor abundance pattern, and the dashed line gives the final abundance pattern after neutrino postprocessing. These results correspond to a neutrino fluence of $F = 0.015$. The results for $A = 183–187$ are highlighted in the inset. For the best-fit fluence of $F = 0.015$, the observed abundances at these five mass numbers are reproduced within $1 \sigma$, although the progenitor abundance pattern is deficient at these mass numbers. Because the final abundances at these mass numbers entirely come from neutrino postprocessing, they are the fingerprints of neutrinos on the $r$-process abundance pattern.

### 4 Implications for the Site of the $r$-Process

It was suggested that the deficiency of the standard $r$-process calculations in the mass region of $A = 183–187$ could result from the deficiency of nuclear mass models for the extremely neutron-rich nuclei far away from stability.
Figure 1: Postprocessing neutron emission probabilities for the progenitor nuclei at $A \sim 195$. The points connected by the long-dashed, dot-dashed, and short-dashed lines are for neutrino fluences of $F = 0.015, 0.030$, and $0.060$, respectively.
Figure 2: Neutrino postprocessing results in the $A \sim 195$ region for a fluence of $F = 0.015$. The data points, some with error bars, give the observed solar $r$-process abundances. The solid (dashed) line gives the abundance pattern before (after) neutrino postprocessing.
It was shown that a similar deficiency below the abundance peak at \( A \sim 130 \) no longer exists after the strength of the \( N = 82 \) closed neutron shell is quenched. However, quenching the strength of closed neutron shells may not be a general way to eliminate this kind of deficiencies. As explained in Sec. 1, the abundance peaks at \( A \sim 130 \) and 195 result from the sharp contrast in the \( \beta \)-decay rates for the nuclei with and without closed neutron shells. Clearly, quenching the strength of closed neutron shells tends to diminish this contrast. Therefore, it may not be always possible to make the deficiencies below the abundance peaks disappear by quenching the strength of closed neutron shells without changing the distinct features of the abundance peaks at the same time.

We have shown in Sec. 3 that for a neutrino fluence of \( \mathcal{F} = 0.015 \), the deficiency of the standard \( r \)-process calculations in the mass region of \( A = 183–187 \) can be remedied completely by neutrino postprocessing without quenching the strength of closed neutron shells. Furthermore, we can show that this fluence is consistent with the conditions in the recent supernova \( r \)-process model. The postprocessing neutrino fluence is approximately determined by the neutrino flux at the \( r \)-process freeze-out and the typical timescale over which the neutrino flux significantly decreases. The neutrino flux decreases either due to the decay of the neutrino luminosity or due to the increasing distance of the \( r \)-process material from the neutron star. A typical decay timescale for the neutrino luminosity is \( \sim 3 \) s, whereas the \( r \)-process material is ejected with a dynamic timescale of \( \sim 0.1–1 \) s. In addition, the \( r \)-process typically freezes out at \( \sim 300–1000 \) km in the recent supernova \( r \)-process model of Woosley et al. For a luminosity of \( L_\nu = 10^{51} \text{ erg s}^{-1} \) per neutrino species, these conditions correspond to a neutrino fluence of \( \mathcal{F} \sim 0.01 \), which is very close to the best-fit value found in Sec. 3.

To summarize, we find that the deficiency of the standard \( r \)-process calculations in the mass region of \( A = 183–187 \) can be remedied completely by neutrino-induced neutron spallation following the \( r \)-process freeze-out. In addition, the neutrino fluence needed for this postprocessing agrees with the conditions in the recent supernova \( r \)-process model. These results strongly argue that the \( r \)-process occurs in supernovae. Furthermore, they provide a sensitive probe for the conditions at the supernova \( r \)-process site because the freeze-out radius and the dynamic timescale for the \( r \)-process are now severely constrained by the postprocessing neutrino fluence.

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