Heavy element production during stellar helium burning

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ABSTRACT. The elements heavier than iron are efficiently produced in the so-called s-process that occurs during stellar helium burning. The currently discussed scenarios refer to massive stars in the mass range $15 < M/M_\odot < 30$ as well as to low mass stars with $1 < M/M_\odot < 3$. The relevant features of these sites are presented with emphasis on the nuclear physics information required for quantitative nucleosynthesis models. The experimental techniques for determining the corresponding neutron capture rates are outlined at the example of the stable samarium isotopes. Eventually, the analysis of the related s-process branching yields an estimate for the mean neutron density for helium shell burning in low mass stars.

RESUMEN. Los elementos más pesados que el Hierro son eficientemente producidos en los así llamados procesos-s que ocurren durante el quemado del Helio estelar. Los argumentos comúnmente discutidos se refieren a las estrellas masivas en el rango de $15 < M/M_\odot < 30$ como también a las estrellas de masa pequeña con $1 < M/M_\odot < 23$. Las características importantes de estos tamaños se presentan con énfasis en la información de la física nuclear requerida para modelos cuantitativos de nucleosíntesis. Las técnicas experimentales para determinar la taza de la captura de neutrones correspondiente son esbozadas en el ejemplo de los isótopos estables de samario. Eventualmente los análisis de bifurcación de los procesos-s dan una estimación de la densidad de neutrones medida para la capa de Helio quemado en las estrellas de masa pequeña.

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1. NEUTRON CAPTURE NUCLEOSYNTHESIS

In astrophysical scenarios, the heavy elements can only be produced by neutron capture reactions and subsequent beta decays. Fusion of charged particles no longer contributes in this mass range, since beyond $A \sim 60$, the Coulomb barriers are becoming too high and the binding energies per nucleon are decreasing. For explaining the observed abundance distribution between iron and the actinides [1], essentially three processes must be invoked as illustrated in Fig. 1. The dominating mechanisms are the slow and the rapid neutron capture processes (s- and r-process), which both account for approximately 50% of the abundances in the mass range $A > 60$. A minor part of the resulting abundance distribution is modified by the p-process at very high temperatures, which accounts for the rare proton rich nuclei.

The s-process associated with stellar helium burning is characterized by relatively low neutron densities. This implies neutron capture times of the order of several months,
much longer than typical beta decay half-lives. Therefore, in the N-Z-plane of Fig. 1 the s-process path follows the stability valley as indicated by the solid line. The developing s-abundances are determined by the respective \((n, \gamma)\) cross sections averaged over the stellar neutron spectrum, such that isotopes with small cross sections are built up to large abundances. This holds in particular for nuclei with closed neutron shells \(N = 50, 82,\) and \(126\), giving rise to the sharp s-process maxima in the abundance distribution around \(A = 88, 140,\) and \(208\).

The r-process counterparts of these maxima are caused by the effect of neutron shell closure on the beta decay half-lives. Since the r-process occurs during stellar explosions (presumably in supernovae) on time scales of a few seconds, neutron captures are much faster than beta decays. This drives the r-process path off the stability valley to where the neutron separation energies are only \(\sim 2\) MeV. At these points, \((n, \gamma)\) and \((\gamma, n)\) reactions are in equilibrium, and the reaction flow has to wait for beta decay to the next higher element. Accordingly, the r-process abundances are proportional to the half-lives of these waiting point nuclei. This means that abundance peaks accumulate again at magic neutron numbers, but at lower mass numbers compared to the related s-process maxima. The separation of the abundance peaks indicates the intersection of the r-process path with the magic neutron numbers.
The following discussion is restricted to the s-process, which is more easily accessible to laboratory experiments as well as to stellar models and astronomical observations [2]. Attempts to describe the r-process are hampered by the large uncertainties in the nuclear physics data far from stability, but also —and perhaps more severely— by the problems related to a detailed modelling of the stellar explosion [3]. These difficulties are best illustrated by the fact that the site of the r-process has not been reliably identified by now [4].

From the above, it is obvious that most isotopes received abundance contributions from both, the s- and the r-process. But as indicated in Fig. 1, there are stable isotopes (marked r) that are not reached by the s-process because of their short-lived neighbors. Consequently, this species is of pure r-process origin. In turn, these nuclei shield their stable isobars against the r-process, so that there is an ensemble of s-only nuclei as well. The existence of these two groups is of vital importance for nucleosynthesis, since the credibility of any model depends on how well the abundances of these particular nuclei can be reproduced. The currently favored stellar s-process scenarios are introduced in §2.

A second observation from Fig. 1 has significant impact on the role of the s-process with respect to stellar models: At some points, the neutron capture chain encounters isotopes with long beta decay half-lives that are comparable to the neutron capture times. The resulting competition between beta decay and neutron capture causes the s-process path to split, e.g. at $^{63}$Ni, $^{79}$Se, and $^{85}$Kr, leading to abundance patterns that reflect the physical conditions of the stellar plasma. The nuclear physics data required for detailed s-process calculations are addressed in §3, and in §4, the analyses of s-process branchings are discussed.

2. s-PROCESS SCENARIOS

2.1. The classical s-process

This first attempt for a quantitative description of the s-process dates back to the fundamental paper by Burbidge et al. [5], and was later improved mainly by Seeger et al. [6]. In this approach it is empirically assumed that a certain fraction $G$ of the observed $^{56}$Fe abundance was irradiated by an exponential distribution of neutron exposures. An analytical solution can be obtained if a possible time dependence of the rates for neutron capture, $\lambda_n = n_n(\sigma)v_T$, and for beta decay, $\lambda_\beta = \ln 2/t_{1/2}$, is neglected. In other words, it is assumed that temperature and neutron density are constant. Then, the characteristic s-process quantity, stellar cross section times $s$-abundance, can be given by the simple expression

$$\langle \sigma \rangle (A) \cdot N_s(A) = \frac{G \cdot N_{56}^0}{\tau_0} \prod_{i=56}^A \left(1 + \frac{1}{\tau_i(\sigma)}\right)^{-1}.$$

Apart from the two parameters $G$ and $\tau_0$ (which are adjusted by fitting the abundances of the $s$-only nuclei), the only remaining input for the expression of $\langle \sigma \rangle \cdot N_s$ are the stellar $(n, \gamma)$ cross sections $\langle \sigma \rangle$. (For further details see Ref. [2]).
Figure 2. The characteristic product of cross section times s-process abundance plotted as a function of mass number [7]. The solid line was obtained via the classical model, and the symbols denote the empirical products for the s-only nuclei. Some important branchings of the neutron capture chain are indicated as well.

Given the very schematic nature of this classical approach, it was surprising to see that it provides for an excellent description of the s-process abundances. In fact, the agreement between observed and calculated abundances became the better the more the input data were improved. The present status is illustrated in Fig. 2, which shows the calculated $\langle \sigma \rangle \cdot N_s$ values (solid line) compared to the corresponding empirical products of the s-only nuclei (symbols) in the mass region between $A = 56$ and 209 [7]. The error bars of the empirical points reflect the uncertainties of the abundances and of the respective cross sections. One finds that equilibrium in the neutron capture flow was obtained between magic neutron numbers, where the $\langle \sigma \rangle \cdot N_s$-curve is almost constant. The small cross sections of the neutron magic nuclei act as bottlenecks for the capture flow, resulting in the distinct steps, e.g. at $A = 140$ and 208.

The global parameters, $G$ and $r_0$, that determine the overall shape of the $\langle \sigma \rangle N_s$-curve, represent a first constraint for the stellar s-process site with respect to the required seed abundance and total neutron exposure. It is found that 0.04% of the observed $^{56}$Fe abundance are a sufficient seed, and that on average about 15 neutrons are captured by each seed nucleus. These numbers refer to the main s-process component given by the solid line in Fig. 2. For $A < 90$, this line falls below the empirical points, and these discrepancies require an additional, the so-called weak, component. In terms of stellar sites, the main component can be attributed to helium shell burning in low mass stars [8-11], whereas the weak component is due to core helium burning in massive stars [12-14, 15 and references therein].
The excellent agreement between the empirical points and the data obtained with the classical model provides some confidence that at least the s-abundances can be reliably described by this schematic approach. In fact, those s-only nuclei that are not affected by branchings, are reproduced by the model with a mean square deviation of only 3% [7].

2.2. Helium shell burning in AGB-stars

Most of their life, stars are quietly burning their central hydrogen inventory. The stellar structure during that phase is characterized by the equilibrium between the gravitational pressure exerted by the outer layers and the radiation pressure created by the energy production in the interior, a situation that can be described, for example, by the standard solar model [16]. After hydrogen exhaustion in the central region, stellar evolution speeds up and becomes much more violent. While hydrogen burning continues in a shell around the center, the inert helium core is no longer producing the energy to maintain the radiation pressure. Therefore, it shrinks and heats up by the released gravitational energy until He burning is ignited at temperatures in excess of ~ 150 million K. The high central temperatures cause the stars to expand and to cool on the outside — they are becoming Red Giants.

This stage of evolution is schematically sketched in Fig. 3 for a star during its Red Giant phase when it has already exhausted its central helium. Energy is produced in a narrow double shell on top of an inert core consisting of $^{12}\text{C}$ and $^{16}\text{O}$. The inset shows that hydrogen burning at the bottom of the deeply convective envelope produces helium that accumulates in a thin layer around the core. As soon as this layer exceeds a critical mass, helium burning is ignited, leading to a highly unstable situation because of the large associated energy production. As a result, strong convective motions cause practically instantaneous mixing in the helium burning layer, and eventually freshly synthesized matter is dredged up to the surface. It is important to note that helium burning lasts for only about 500 yr, while it takes about $2 \times 10^5$ yr for hydrogen burning to replenish the consumed helium. These helium burning episodes can repeat up to about 20 times. As a consequence of the comparably short time scale and the strong impact of convection, stellar models for the Red Giant phase are very difficult to quantify and often have to resort to uncertain approximations [8,9].

An alternative way for gaining insight in that state of stellar evolution is to investigate the abundance pattern in the various s-process branchings, which still carry the information about the physical conditions at which they were formed (§4). Neutrons are liberated during helium burning by side reactions, i.e. by $^{13}\text{C}(\alpha, n)^{16}\text{O}$ and to a smaller fraction by $^{22}\text{Ne}(\alpha, n)^{25}\text{Mg}$.

2.3. Core helium burning in massive stars

Towards the end of core helium burning in massive stars the temperatures become sufficiently high for neutron production via the $^{22}\text{Ne}(\alpha, n)^{25}\text{Mg}$ reaction. The resulting neutron density is comparably weak and lasts for about $3 \times 10^4$ yr, resulting in the modest neutron exposure of the weak s-process component. This scenario is discussed in the following contribution [17].
FIGURE 3. Schematic sketch of the stellar structure during the helium shell burning phase in low mass stars. Energy is produced in a narrow zone on top of the inert core consisting essentially of \(^{12}\text{C}\) and \(^{16}\text{O}\). The envelope is deeply convective, so that freshly synthesized matter is transported to the surface. Helium burning occurs in repeated flashes, which renders the associated s-process difficult to model.

3. LABORATORY STUDIES

This section contains a brief discussion of those quantities which determine the s-process abundance patterns: the stellar \((n, \gamma)\) cross sections and \(\beta\)-decay rates. The charged particle reactions during helium burning, and in particular the neutron producing \((\alpha, n)\) reactions are the subject of the following contribution [17].

3.1. Neutron capture cross sections

3.1.1. The Karlsruhe 4\(\pi\) BaF\(_2\) detector

The best signature for the identification of neutron capture events is the total energy of the gamma cascade by which the product nucleus deexcites to its ground state. Hence, the accurate measurement of \((n, \gamma)\) cross sections should use a detector that operates as a calorimeter with good energy resolution. In the gamma spectrum of such a detector, all capture events would fall in a line at the neutron separation energy (typically between 5 and 10 MeV), well separated from the gamma-ray backgrounds that are inevitably in neutron experiments.

These arguments point to a 4\(\pi\) detector of high efficiency, made of a scintillator with reasonably good time and energy resolution. In addition, this detector should be insensitive to scattered neutrons, since —on average— the scattering cross sections are about 10 times larger than the capture cross sections. The combination of these aspects is best covered by a 4\(\pi\) detector of BaF\(_2\). Most important is the weak neutron sensitivity of that scintillator, which keeps the background from scattered neutrons at a manageable level.
The concept of the Karlsruhe $4\pi$ BaF$_2$ detector and the way it is operated is sketched in Fig. 4, showing the setup at the accelerator. The detector is indicated schematically by a computer simulation, which presents only the spherical BaF$_2$ shell (15 cm thickness and 20 cm inner diameter) consisting of 42 individual BaF$_2$ crystals, and the supporting structure. All other details are omitted for the sake of clarity.

Neutrons are produced by means of a pulsed Van de Graaff accelerator (repetition rate 250 kHz, average beam current 2 $\mu$A, pulse width 0.7 ns) via the $^7$Li$(p, n)^{7}$Be reaction. The collimated neutron beam hits the sample in the centre of the detector at a flight path of 77 cm. Up to 8 samples can be used in the measurements, all mounted on a vertical ladder. One position is occupied by a gold sample for determination of the neutron flux, while an empty sample position and a scattering sample are routinely used for background determination. The samples are cycled into the measuring position in intervals of about 10 min that are defined by integrating the beam current on target. Additional neutron monitors are used to check for equal neutron exposure per sample and to collect data for the correction of backgrounds due to scattered neutrons.

The essential features of the detector are: a resolution in gamma-ray energy ranging from 14\% at 662 keV to 6\% at 6.13 MeV, a time resolution of 500 ps, and an overall efficiency for capture events of about 96\%. For a detailed description of the detector and a typical measurement see Wisshak et al. [18,19].

With this new detector, $(n, \gamma)$ cross sections can be determined with an accuracy of 1 to 2\%. Recently, it was used for a precise measurement of the cross sections of the s-only samarium isotopes [19], which define the branchings at $A = 147, 148, 149$ ($\S$4). This experiment has been performed on highly enriched samarium oxide samples in the neutron energy range from 3 to 225 keV. The efficiency for capture events was evaluated by means of the experimental pulse height spectra shown in Fig. 5. Most of the events are, indeed, recorded in a line at the neutron separation energy, but there is also a tail towards lower energies. This tail corresponds to events where not all gamma-rays of the cascade
could be detected, mainly due to losses through the openings for the neutron beam and the sample changer. The related systematic uncertainty of 0.7% is the major contribution to the overall uncertainty.

The stellar cross sections were obtained by folding the differential data, $\sigma(E_n)$, with the stellar neutron spectra for various temperatures. The Maxwellian averages for the "standard" thermal energy of $kT = 30$ keV exhibit uncertainties of $\sim 1\%$, 4 times smaller than reported for the best previous data [20]. The precision of the present results provides for a new quality that is indispensable in the astrophysical discussion of the branchings at $A = 147-149$ (see §4).

3.1.2. The activation technique

The activation technique represents an alternative method for the determination of stellar $(n, \gamma)$ rates. Compared to the techniques using the detection of the prompt capture $\gamma$-rays, this method offers the advantages of higher sensitivity (which means that much smaller cross sections can be measured reliably — an important aspect for the investigation of radioactive isotopes on the $s$-process path), and of selectivity (which means that isotope mixtures can be studied via the characteristic $\gamma$-ray energies of the respective decay products). However, it is restricted to those cases, where neutron capture produces an unstable nucleus, and it yields the stellar rate only for two thermal energies at $kT = 25$ and 52 keV.

This method is based on the fact that quasi-stellar neutron spectra can be produced in the laboratory via the $^7\text{Li}(p, n)^7\text{Be}$ and the $^3\text{He}(p, n)^3\text{He}$ reactions [21–23]. With the first reaction, the proper neutron spectrum is obtained by bombarding thick metallic lithium targets with protons of 1912 keV, only 31 keV above the reaction threshold. The resulting neutrons exhibit a continuous energy distribution with a high-energy cutoff at
$E_n = 106$ keV and a maximum emission angle of $60^\circ$. The angle-integrated spectrum corresponds closely to a Maxwell-Boltzmann distribution for $kT = 25$ keV. Hence, the reaction rate measured in such a spectrum yields immediately the proper stellar cross section. The second reaction can be applied analogously. The $^{197}\text{Au}(n,\gamma)$ cross section can be used as a standard in both cases.

Recent applications of this technique can be found in the literature [24,25].

3.2. Weak interaction rates

The competition between neutron capture and beta decay at s–process branchings requires the consideration of thermal effects on the beta decay rates at stellar conditions. That isotopes, which are exposed to the high temperatures and densities in a He burning plasma, may experience a dramatic enhancement of their decay rate, will be illustrated at the example of the branch point isotope $^{79}\text{Se}$. Such temperature-dependent branchings are important, since their abundance pattern can be considered as a thermometer for the s–process site.

The inset in Fig. 6 illustrates the mechanism responsible for the enhancement of the $\beta^-$–rate of $^{79}\text{Se}$ under stellar conditions. The decay of the ground state is unique first forbidden with a terrestrial half-life of about $6.5 \times 10^4$ yr. The first excited state is a 3.9 min isomer at 96 keV, which decays by internal transitions but can also undergo allowed beta decay. From log $ft$ systematics one expects a half-life of about 20 d for this decay branch. The intense photon bath in the stellar plasma gives rise to thermally induced transitions, resulting in a finite population probability of low lying excited nuclear states. At typical s–process temperatures, the isomer in $^{79}\text{Se}$ is populated to $\sim 1\%$, corresponding to an effective stellar half-life of about 2000 d—an enhancement by a factor $10^4$ compared to the terrestrial value.

The most important nuclear physics parameter in this problem is the log $ft$ value for the isomer in $^{79}\text{Se}$. Since theoretical estimates were too uncertain, this quantity was determined experimentally by producing $^{79m}\text{Se}$ via activation with thermal neutrons [26]. With the experimental log $ft$ value of $4.70\pm0.10$ for the beta decay branch of the isomer, the temperature dependence of the effective stellar half-life of $^{79}\text{Se}$ was recalculated with the formalism of Takahashi and Yokoi [27].

The resulting half-life is plotted in Fig. 6 with the error band according to the experimental range of log $ft$ values. This figure indicates also a remaining ambiguity in the half-life of the ground state: While only an upper limit of $6.5 \times 10^4$ yr is known experimentally, a lower limit of 1700 yr would also be consistent with the log $ft$ systematics for unique first forbidden transitions. At temperatures above $200 \times 10^6$ K, however, the effective half-life is completely insensitive to this ambiguity. This leads to the rather strange consequence that, at present, the stellar half-life of $^{79}\text{Se}$ is much better known than the terrestrial one!

4. Branchings in the s–process path

In addition to the parameters $G$ and $\tau_0$, which were deduced from the overall shape of the $(\sigma)N_s$–curve, analyses of the abundance patterns in the various branchings can be
expected to yield more detailed information on the stellar s-process site. The strength of such a branching is preserved in the respective abundance pattern and can be used to determine information on the physical conditions during the s-process. As an example, Fig. 7 shows the s-process flow in the mass region between neodymium and samarium, with the possible branchings at $^{147}$Nd and $^{147-149}$Pm. Note that $^{148}$Sm and $^{150}$Sm are shielded against the r-process by their isobars in neodymium. As the result of a significant branching at $A = 147-149$, the $\langle \sigma \rangle N_s$-value of $^{150}$Sm will, therefore, be always larger than that of $^{148}$Sm.

The strength of a branching can be expressed in terms of the rates for beta decay and neutron capture of the branch point nuclei as well as by the $\langle \sigma \rangle N_s$-values of the involved s-only isotopes,

$$f_\beta = \frac{\lambda_\beta}{\lambda_\beta + \lambda_n} \approx \frac{\langle \sigma \rangle N_s}{\langle \sigma \rangle N_s^{148\text{Sm}}}.$$

Inserting the equations for the decay rates this expression can be solved for the neutron density. If—for simplicity—only the branching at $^{148}$Pm is considered, one obtains

$$n_n = \frac{1 - f_\beta}{f_\beta} \cdot \frac{1}{\nu T \langle \sigma \rangle N_s^{148\text{Pm}}} \cdot \frac{\ln 2}{t_{1/2}^{148\text{Pm}}}.$$

This equation demonstrates the input data that are important for reliable branching analyses.
The first term depends on the cross sections for the $s$-only nuclei, which define the branching factor $f_\beta$. Since the neutron density is required to $\sim 10\%$, the branching factor, and, hence, the cross sections for the $s$-only nuclei need to be known to about $1\%$ in many cases. While conventional techniques are limited to uncertainties $\geq 4\%$, the new $4\pi$ BaF$_2$ detector allows now to reach the required accuracy.

The second term contains the stellar cross sections of the radioactive branch point isotopes. Since there are practically no measurements for the unstable branch point nuclei, only calculated cross sections are available at present. But even the most careful statistical model calculations are limited to uncertainties of 20–30%, not sufficient for deducing the entire information contained in the abundance patterns. For some of these short-lived nuclei, experimental cross section studies aiming at a 5 to 10% uncertainty were recently suggested [28,29].

The last term denotes the stellar decay rate of the branch point isotope. While there is no difference between the stellar and the terrestrial rate for some of the branch points, a variety of examples exhibit sometimes drastic changes under the high temperatures and densities of the stellar plasma. For the branchings at $A=147–149$, these effects are of minor importance [29].

Since the relevant beta decay rates in Fig. 7 are not significantly affected by temperature, the branchings at $A=147–149$ can be used for obtaining an improved estimate of the $s$-process neutron density. Compared to the previous result of Winters et al. [20] ($f_\beta = 0.92 \pm 0.04$), the measurement with the $4\pi$ BaF$_2$ detector yields a considerably more precise value, $f_\beta = 0.870 \pm 0.009$, with a 4 times smaller uncertainty [19]. This branching factor implies a neutron density of $n_n = (3.8 \pm 0.6) \times 10^8$ cm$^{-3}$, in agreement with the previously estimated $(3.4 \pm 1.1) \times 10^8$ cm$^{-3}$ [7].

Along the $s$-process path there are about 15 to 20 significant branchings, which can be studied with respect to the physical conditions at the stellar site. In a first step, the neutron

\begin{figure}[h]
\centering
\includegraphics[width=0.7\textwidth]{s-process.png}
\caption{The $s$-process flow between neodymium and samarium.}
\end{figure}
density must be obtained from those branchings which are not affected by temperature. With this information, the effective branching factor of the remaining examples can be derived and the mean stellar decay rates be determined. Eventually, the dependence of these rates on temperature and/or electron density can then be used to estimate the s-process temperature and mass density [7].

The present status of branching analyses is characterized by the neutron density from the branchings of Fig. 7, \( n_n = (3.8 \pm 0.6) \times 10^8 \text{ cm}^{-3} \), and a mean temperature, \( T_s = (3.1 \pm 0.6) \times 10^8 \text{ K} \). A first attempt to determine the mass density in the s-process was carried out by Beer et al. [30] by analyzing a small branching to \(^{163}\text{Ho}\). Their value, \( \rho_s = (8 \pm 5) \times 10^3 \text{ g cm}^{-3} \), is in reasonable agreement with stellar models for helium shell burning [8,9].

So far, the results from the different branchings are still consistent with each other. This means that they still satisfy the assumption of the classical approach, that neutron density and temperature are constant during the s-process. In contrast, the stellar models suggest these parameters to vary strongly with time. Consequently, the freeze-out behavior at the end of each neutron exposure may well cause a distortion of the abundance pattern predicted by the classical approach. Since these effects are still obscured by the present uncertainties of the input data, future efforts have to concentrate on improving the relevant cross sections and beta decay rates, which characterize the various branchings.

References