Two-neutron capture reactions in supernovae neutrino bubbles

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Recent calculations suggest that the neutrino heated bubbles in the postcollapse phase of type II supernovae may be the site of the r process. The nucleosynthesis process depends on the expansion and cooling rates as well as on the rates of the α-particle and neutron recombination processes that bridge the mass gaps at $A = 5$ and 8. We have reexamined the rate for the important reaction ^{4}He$(\alpha n,\gamma)^9$Be, and estimated the rates for the two neutron capture reactions on ^{6}He and ^{8}He which could provide an alternative reaction path. We find that the most important reaction for initiating the α process is the ^{4}He$(\alpha n,\gamma)^9$Be reaction, as suggested by earlier work.

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I. INTRODUCTION

It has recently been suggested [1–3] that the postcollapse phase in type II supernovae models offers an "ideal site" for the r process whereby heavy elements of $A \geq 100$ are formed by rapid neutron capture processes in a neutron bath of 10^{20-30} neutrons per cm^3 [1–3]. The reaction path is mainly determined by nuclear statistical equilibrium (NSE), which culminates in (n,γ)-(γ,n) equilibrium [4] on the neutron-rich side of stability.

After the formation of the neutron star in the core collapse and bounce, followed by shock heating of the outer shell material, neutrino heating occurs in the wake of the shock-front, causing a high-entropy bubble to form with a high photon-to-baryon ratio. The baryonic matter is in NSE, and at these high-entropy conditions, the abundance distribution in NSE is at first dominated by α particles, protons, and neutrons. As baryonic material moves through the bubble in a neutrino-heated wind, the abundance distribution remains in NSE but shifts towards higher masses via the recombination of free α particles, neutrons, and protons. This generates an α process leading to the formation of massive isotopes (up to $A \approx 100$) largely via α capture in NSE [1,3].

This process depends on the expansion and cooling rates as well as on the rates of the various recombination paths from ^{4}He to ^{12}C. Recombination may occur both via the triple-α reaction, ^{4}He$(2\alpha,\gamma)^{12}$C, and by the reaction sequence ^{4}He$(\alpha n,\gamma)^9$Be$(\alpha,n)^{12}$C, bridging the mass 5 and mass 8 gaps. Once the α particles have recombined to form ^{12}C, higher mass nuclei are quickly built up in the α process [1,3]. For high photon-to-baryon ratios, however, the initiating recombination of α particles via three-body reactions rapidly falls out of equilibrium, delaying the formation of heavy elements. Due to the rapid cooling as material flows through the hot bubble, the charged-particle reactions first fall out of nuclear statistical equilibrium and then freeze out. This leaves free neutrons and seed nuclei for a rapid neutron capture process, the r process. The amount of the available seed material is limited by the small α-recombination rate. While the r-process path is first determined by (n,γ)-(γ,n) equilibrium, after further cooling and depletion of neutrons the neutron capture process ceases and β decay drives the reaction path closer to the line of stability.

To model this complex nucleosynthesis scenario requires an accurate knowledge of the initial α-recombination processes, the determining factor for the seed-nuclei production. Besides ^{4}He$(2\alpha,\gamma)^{12}$C and ^{4}He$(\alpha n,\gamma)^9$Be, recombination of α particles is also possible via alternative three-body reactions, depending on the initial neutron proton ratio X_n/X_p. In the case of a high neutron abundance, which is necessary for a subsequent r process, a sequence of two-neutron capture reactions, ^{4}He$(2n,\gamma)^6$He$(2n,\gamma)^8$He, could, in principle, bridge the mass 5 and 8 gaps. Further processing could then occur close to the neutron drip line via ^{8}He$(\beta)^8$Li$(n,\gamma)^9$Li and ^{8}He$(\beta n)^7$Li. The two-neutron capture cross sections on ^{4}He and ^{6}He depend strongly on the pronounced halo structure of the ^{6}He and ^{8}He compound nuclei. These nuclei may be formed preferentially by two-step resonant processes through their broad 2^+ first excited states. Detailed studies of these $2n$ configurations and the $B[E2(2_1^+ \rightarrow 0_{gs}^+)]$ γ-transition probabilities are necessary to allow reliable calculations of the rate for the two-neutron capture processes on ^{4}He and ^{6}He. In the following we present estimates of these two-neutron capture reaction rates, based on the currently available experimental data and on theoretical predictions for the strengths of the first excited levels in ^{6}He and ^{8}He. We also recalculate the rate for the competing reaction ^{4}He$(\alpha n,\gamma)^9$Be.

II. REACTION RATES

Two-particle capture sequences on a nucleus 1 that proceeds through an intermediate state in the unstable nucleus 2 to form the final nucleus 3 are historically included under the heading of three-body reactions. For successive capture of two particles, the three-body rate $\langle 1pp\rangle$ is defined by

$$\dot{Y}_3 = (1/2)\rho^2 N_A^2 \langle 1pp\rangle Y_p^2 Y_1 , \quad (1)$$

where $\dot{Y}_3$ denotes the change in time of the abundance by

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number Y_3 of nucleus 3, Y_1 and Y_p are the abundance of the target nucleus 1 and the particle abundance, respectively, ρ is the stellar density, and N_A is Avogadro's number. The factor 1/2 accounts for the double counting of collisions involving identical particles. The three-body rate $\langle 1pp \rangle$ is then given by [5,6]

$$N_A^2\langle 1pp\rangle = N_A^2 \int_{E_1} \frac{d\langle (p,p)\rangle(E_1)}{dE_1} \frac{2\hbar}{\Gamma_2(E_1)} \left[\int_{E_2} \frac{d\langle p,\gamma\rangle(E_1,E_2)}{dE_2} dE_2 \right] dE_1 . \quad (2)$$

Here, $\Gamma_2(E_1)$ is the energy-dependent width of the intermediate state in nucleus 2, and E_1 (E_2) denotes the collision energy of the first (second) neutron (plus nucleus) in the center-of-mass system. The differentials represent the integrands of the $\langle \sigma v \rangle$ integration for the single-step reaction rates [7] and are described by

$$\frac{d\langle \sigma v\rangle}{dE} = \sqrt{\frac{8}{\pi\mu}} \frac{1}{(kT)^{3/2}} \sigma(E) E \exp\left(-\frac{E}{kT}\right). \quad (3)$$

For resonant capture the energy dependence of the cross section $\sigma(E)$ is described by the Breit-Wigner formula [7]

$$\sigma(E)_{ab} = \pi \lambda\!\!\!^{-2} \frac{(2J+1)}{(2J_p+1)(2J_t+1)} \frac{\Gamma_a(E)\Gamma_b(E)}{(E-E_R)^2+[\Gamma(E)/2]^2} . \quad (4)$$

$\Gamma_a(E)$, $\Gamma_b(E)$, and $\Gamma(E)$ are the energy-dependent widths of the entrance and the exit channels and the total width. J_t, J_p, and J are the spins of the target, the incident particle, and the resonance, respectively, and E_R denotes the resonance energy. For the first step, (p,p), the widths are determined by the particle width Γ_p, $\Gamma=\Gamma_a=\Gamma_b=\Gamma_p$, and in the second step, (p,γ), by $\Gamma_a=\Gamma_p$, $\Gamma_b=\Gamma_\gamma$, and $\Gamma=\Gamma_p+\Gamma_\gamma$, where Γ_γ is the γ width of the excited state of the final nucleus.

The second step of the reaction sequence can also proceed through the nonresonant direct radiative capture (DC) process [8]. The cross section is given by the product of the spectroscopic factor of the final state and the theoretical cross section, which was calculated in a Woods-Saxon potential. The cross section, especially for s-wave neutron capture, is sensitive to the phase shift of the unbound wave function. Therefore, the phase shifts were calculated in a Woods-Saxon potential with a diffuseness of $a=0.65$ fm, and the potential depths were chosen to reproduce the energies of the first excited states of the final nuclei.

The three-body rate can be calculated by numerical integration of Eq. (2). Input parameters are the excitation energies, the spins, and the partial widths of the states involved. If in one step or in both steps several resonances contribute, the corresponding integration in Eq. (2) has to extend over the sum of the Breit-Wigner terms [6].

The nuclei ^{6}He and ^{8}He are currently being investigated by many groups, both experimentally and theoretically, because of their neutron halo structure and the three-body structure of some states (see, e.g., Ref. [9]). For the resonant contributions in the second step, only the $E2$ γ-transition strengths of the first excited states in ^{6}He and ^{8}He are of importance because the two-neutron capture rates are directly proportional to Γ_γ [see Eq. (1)]. To study these contributions, the $B(E2)$ strengths of the excited states in ^{6}He and ^{8}He were calculated in the framework of a three-body model.

The ^{6}He nucleus can be well described in a three-body model as two neutrons plus an α-particle core [10]. The 0^+ ground state is bound, and the 2^+ excited state a resonance. Both can be described by solving the hyperspherical harmonics equations [10] with a local Woods-Saxon n-core potential [11] and the n-n potential of Ref. [12]. The n-core potential has a $0s$ bound state at ≈ -11 MeV. This was taken as the "occupied" level of the neutrons in the alpha-particle core, and all the three-body wave functions are constrained to be orthogonal to this level. The $dB(E2)/dE$ distribution is then integrated over a range of continuum energies including the resonance. The integrated $B(E2)$ may be compared with the limit from the non-energy-weighted sum rule (NEWSR), of $B(E2,2^+\rightarrow 0^+) \leq Z^2e^2/4\pi\langle r_\alpha^4\rangle$, where r_α is the distance of the (charged) α core from the center of mass of the whole nucleus, and $\langle\ \rangle$ indicates the ground-state expectation value.

The structure of ^{8}He in a three-body model is more complicated than that of ^{6}He because the ^{6}He core is no longer a closed shell and the Pauli orthogonalization is no longer so straightforward. We therefore estimated $B(E2,2^+\rightarrow 0^+)$ by finding an effective n-^{6}He potential, which reproduces the 0^+ and 2^+ energies. We compare this result with that from the Slater-determinant wave function [13] for ^{8}He. In the following subsections the reaction rates for the individual reactions will be determined and discussed in detail.

A. ^{4}He$(2n,\gamma)^6$He

The first step of this reaction sequence, ^{4}He$(n,n)^4$He, has a Q value of $Q=-0.89$ MeV [14], and the cross section is dominated by a p-wave resonance corresponding to the ground state of ^{5}He. We have adopted the width of 1.2 MeV observed in elastic neutron scattering [15]. This value is in good agreement with the most recent width of 1.36 ± 0.19 MeV determined in the ^{3}H$(d,\gamma)^5$He reaction [16]. A second resonance has been observed at $E_n=22.13$ MeV but is too high in energy to contribute to the reaction rate.

The second step proceeds through the first excited state of ^{6}He at 1.797 MeV. This $J^\pi=2^+$ state is below the one-neutron threshold by 0.060 MeV [14] but is unbound against two-neutron decay by 0.824 MeV [14]. Therefore, the ^{5}He $(n,\gamma)^6$He reaction has to proceed through the high-energy tail of this subthreshold resonance. Since the one-neutron

TABLE I. Reaction rate for the ^{4}He$(2n,\gamma)^6$He reaction.

T (GK)	$N_A^2\langle^4\text{He}2n\rangle$		
	Present		
	DC	Resonant	[18]
0.5	4.00×10^{-11}	3.04×10^{-16}	8.17×10^{-19}
0.8	8.00×10^{-11}	2.75×10^{-15}	4.39×10^{-16}
1.0	1.15×10^{-10}	8.18×10^{-15}	3.16×10^{-15}
1.5	2.20×10^{-10}	3.71×10^{-14}	3.64×10^{-14}
2.0	3.29×10^{-10}	7.06×10^{-14}	1.07×10^{-13}
2.5	4.25×10^{-10}	9.37×10^{-14}	1.88×10^{-13}
3.0	5.01×10^{-10}	1.05×10^{-13}	2.60×10^{-13}
5.0	6.49×10^{-10}	9.33×10^{-14}	4.02×10^{-13}

decay is energetically not allowed at the resonance energy, the observed width of the state of 0.113 ± 0.020 MeV [15] is entirely due to the two-neutron structure of the state. However, the shape of the high-energy tail above the one-neutron threshold will also be influenced by the contribution of the one-neutron structure, $\Gamma_{\text{total}}=\Gamma_{1n}+\Gamma_{2n}$. Γ_{1n} has been calculated from the well-known relation

$$\Gamma_{1n}=\frac{3\hbar}{\mu R^2}P_l C^2 S_{1n}. \tag{5}$$

P_l is the neutron penetrability and $R=1.26(A_n^{1/3}+A_t^{1/3})$ the interaction radius. A value of $C^2S_{1n}=0.5$ has been assumed for the one-neutron spectroscopic factor. This value is suggested by shell model calculations using the (6–16) general two-body matrix elements of Cohen and Kurath [17]. No experimental information is available for the γ width of this state. Using the three-body model for the ^{6}He 0^+ and 2^+ states, the integrated $B(E2,2^+\rightarrow 0^+)$ is predicted to be 0.57 e^2fm^4, which leads to Γ_γ = 8.8 μeV for the 1.80 MeV γ-ray decay. The NEWSR limit is $B(E2,\downarrow)\leq 1.74$ e^2fm^4, using $\langle r_\alpha^4\rangle=5.47$ fm^4 from the three-body ground state (g.s.), corresponding to an upper limit on Γ_γ of 26.8 μeV.

The second step can also proceed by an s-wave direct radiative capture into the ground state of ^{6}He, and its cross section was calculated in a Woods-Saxon potential with a radius $R=2.6$ fm and a diffuseness $a=0.65$ fm. For the calculations of the phase shifts, the potential depth was chosen to reproduce the energy of the first excited state at 1.797 MeV. The theoretical cross section was weighted with a spectroscopic factor of 0.5 to account for the single-particle structure of the ground state.

Equation (2) was evaluated numerically with these input parameters, and the resulting three-particle reaction rate $\langle^4\text{He}2n\rangle$ is listed in Table I as a function of temperature, together with the previous rate of [18]. The nonresonant and resonant contributions of the rate are listed in the second and third columns.

The nonresonant rate is much larger than the resonant contribution and completely determines the reaction rate. The present resonant rate is comparable to the previous rate of Fowler, Caughlan, and Zimmerman [18] who made only a rough calculation of the rate by analogy with ^{6}Li$(\gamma,np)^4$He to set a reasonable limit on the rate.

TABLE II. Reaction rate for the ^{6}He$(2n,\gamma)^8$He reaction.

T (GK)	$N_A^2\langle^6\text{He}2n\rangle$	
	DC	Resonant
0.5	8.40×10^{-11}	9.75×10^{-17}
0.8	2.75×10^{-10}	7.93×10^{-16}
1.0	4.75×10^{-10}	2.24×10^{-15}
1.5	1.01×10^{-9}	1.74×10^{-14}
2.0	1.40×10^{-9}	7.61×10^{-12}
2.5	1.61×10^{-9}	1.90×10^{-11}
3.0	1.71×10^{-9}	3.35×10^{-11}
5.0	1.59×10^{-9}	7.41×10^{-11}

B. ^{6}He$(2n,\gamma)^8$He

The ground state of ^{7}He is neutron unbound by 0.44 MeV [14] with a width of 0.16 ± 0.03 MeV [15] and is populated by p-wave elastic neutron scattering on ^{6}He. ^{8}He is subsequently produced by neutron radiative capture through the first excited state as a resonance. Various energies for the first excited state have been reported [15,19–21]. We have adopted the results of Bohlen *et al.* [20], who report an energy of 3.70 ± 0.06 MeV and a width of 0.20 ± 0.05 MeV for this state in good agreement with the results of Ref. [21]. This state is unbound against one neutron decay by 1.12 MeV [14] and against two-neutron decay by 1.56 MeV [14]. The invariant mass spectrum for the ^{6}He+n coincidences [21] shows a peak at an energy of 0.45 ± 0.01 MeV, which corresponds to the ground-state energy of ^{7}He. This indicates a dominant sequential neutron decay of the excited state. For this reason, we assumed $\Gamma_{\text{tot}}\approx\Gamma_{1n}$ and $\Gamma_{2n}\approx 0$. No experimental information is available for the γ width of this state. The Slater-determinant model for ^{8}He of Ref. [13] gives, for the g.s. expectation value, $\langle r_\alpha^4\rangle=0.877$ fm^4, which leads to upper limits $B(E2,\downarrow)\leq 0.279$ e^2fm^4 and $\Gamma_\gamma\leq 160$ μeV. A dynamical model for ^{8}He similar to that for ^{6}He gives an integrated $B(E2,\downarrow)=0.115$ e^2fm^4 and Γ_γ = 65 μeV, which is again about 1/3 of the sum-rule limit.

The cross section of the s-wave direct capture to the ground state of ^{8}He has been calculated in a Woods-Saxon potential with radius $R=2.6$ fm and diffuseness $a=0.65$ fm. For the calculations of the phase shifts the potential depth was chosen to reproduce the energy of the first excited state at 3.70 MeV. The theoretical cross section was weighted with a spectroscopic factor of 0.5 to account for the single-particle structure of the ground state.

The resulting three-body reaction rate $\langle^6\text{He}2n\rangle$ is listed in Table II as a function of temperature. Again the DC contribution is much larger than the resonant rate and determines the total rate.

C. ^{4}He$(\alpha n,\gamma)^9$Be

The first step of this reaction sequence is identical to that of the triple-α process [5]. It proceeds through a narrow resonance at 0.091 78 MeV with a width Γ = 6.8 ± 1.7 eV [15]. The final compound nucleus, ^{9}Be, is neutron bound by 1.665 MeV [14], and the excited states contribute as resonances for the ^{8}Be(n,γ) rate. We have included only the first excited state at E_x = 1.684 MeV (J^π = 3/2$^-$, Γ = 0.217

TABLE III. Reaction rate for the ^{4}He$(\alpha n,\gamma)^9$Be reaction (for details see text).

T (GK)	$N_A^2\langle\alpha\alpha n\rangle$		
	Present		
	Destructive	Constructive	[18]
0.5	5.54×10^{-7}	6.10×10^{-7}	1.06×10^{-6}
0.8	4.14×10^{-7}	5.23×10^{-7}	8.40×10^{-7}
1.0	3.23×10^{-7}	4.39×10^{-7}	6.66×10^{-7}
1.5	1.84×10^{-7}	2.86×10^{-7}	3.74×10^{-7}
2.0	1.17×10^{-7}	1.98×10^{-7}	2.25×10^{-7}
2.5	8.09×10^{-8}	1.46×10^{-7}	1.46×10^{-7}
3.0	5.97×10^{-8}	1.12×10^{-7}	9.93×10^{-8}
5.0	2.59×10^{-8}	5.19×10^{-8}	3.08×10^{-8}

MeV, $\Gamma_\gamma = 0.30$ eV) [15]. Calculations showed that the contributions of the states at higher excitation energies can be neglected. Since the first step proceeds through a very narrow resonance, the expression for the three-particle rate, Eq. (2), reduces to

$$N_A^2\langle 1pp\rangle = N_A^2\frac{2\hbar}{\Gamma_2}\langle\alpha\alpha\rangle\langle{}^8\mathrm{Be}+n\rangle, \tag{6}$$

where $\langle\alpha\alpha\rangle$ is given by [7]

$$\langle\alpha\alpha\rangle = \left[\frac{2\pi}{\mu kT}\right]^{3/2}\hbar^2\Gamma_2\exp\left(-\frac{E_R}{kT}\right). \tag{7}$$

In addition, the second step can also proceed through the s-wave direct capture to the ground state of ^{9}Be. For the calculations of the phase shifts, the potential depth was chosen to reproduce the ground state. The theoretical cross section was weighted with a spectroscopic factor of 0.5 to take account of the single-particle structure of the ground state.

Since both reaction contributions involve an incoming s wave, interference between the resonance and the DC process is possible. Therefore, the rates were calculated assuming constructive (destructive) as well as destructive (constructive) interference at energies above (below) the resonance energy. The resulting reaction rate is listed in Table III as a function of the temperature together with the previous rate of [18]. The present rate agrees with the previous rate of [18] within a factor of 2.

III. CONCLUSION

The present rate calculations show the importance of the nonresonant direct radiative capture process, especially in the case of the two-neutron capture reaction on ^{4}He, where the DC rate is approximately four orders of magnitude larger than the resonance rate. The calculation of the s-wave DC cross section is particularly sensitive to the choice of the phase shift because of the absence of any barriers. This uncertainty in the cross section was investigated by varying the phase shifts. In the case of the DC to the ground states of ^{6}He and ^{8}He the phase shifts were also calculated with various potential depths chosen to reproduce the ground state. This led to variations of the DC cross sections up to a factor of 10. For the neutron capture to the ground state of ^{9}Be, the phase shift was varied by $\pm 100\%$ leading to a change of the cross section of $\pm 35\%$. Therefore, we believe that the present rates are reliable order-of-magnitude estimates. This is supported by a recent analysis of the nonresonant contribution of the ^{4}He$(2n,\gamma)^6$He reaction in terms of the three-body problem in the framework of the hyperspherical method [22]. These calculations yield values similar to the present work (within a factor of 2) for the nonresonant reaction rate.

The reaction flow in the α process in nuclear statistical equilibrium during the postcollapse phase in type II supernovae models is determined by the reactions rates of the competing reactions ^{4}He$(2\alpha,\gamma)^{12}$C, ^{4}He$(\alpha n,\gamma)^9$Be, and ^{4}He $(2n,\gamma)^6$He as well as the relative abundance of neutrons and α particles. The last parameter is listed for a $20M_\odot$ delayed supernova explosion in Table III of Ref. [3] for 40 trajectories in $\rho(t)$, $T(t)$, and $Y_e(t)$ for a logarithmic grid of mass elements for the last ≈ 0.03 $M_\odot$ to be ejected by the proto-neutron star. For the first 24 trajectories the flow in the α process is determined by the triple-α reaction, and the neutron abundance is too small for the r process to take place. However, for the last 16 trajectories the neutron-to-seed ratio is sufficient for the r process [3]. Under these conditions the ^{4}He$(\alpha n,\gamma)$ reactions is approximately three orders of magnitude larger than the triple-α reaction. With the present increased rate, the ^{4}He$(2n,\gamma)$ reaction is also stronger than the triple-α reaction by about a factor of 5, but it is approximately two orders of magnitude weaker than the ^{4}He$(\alpha n,\gamma)$ reaction.

The key reaction for initiating the α process remains the ^{4}He$(\alpha n,\gamma)$ reaction. The present reevaluation establishes a firm base for this rate. The remaining uncertainties are caused by the choice of the phase shift for the DC calculations and the question of whether the DC interferes with the s-wave resonant capture constructively or destructively. An experimental study of the Coulomb dissociation of ^{9}Be could possibly reduce the remaining uncertainty substantially.

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