REACTION RATES OF NEUTRON CAPTURE BY LI– AND BE–ISOTOPES

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Neutron capture by neutron-rich Li– and Be–isotopes plays a role in big–bang nucleosynthesis, especially in its inhomogeneous version and in the α–process occurring in supernovae. New reaction rates for $^7\text{Li}(n,\gamma)^8\text{Li}$ and $^9,10,11\text{Be}(n,\gamma)^{10,11,12}\text{Be}$ have been consistently calculated using direct capture for the nonresonant part and the Breit–Wigner formula for the resonant part. The spectroscopic factors, spin/parity assignments and excitation energies of the final bound and initial resonant states have been taken from existing experimental data whenever possible. For unstable nuclei where this information is not experimentally available the shell model was used to determine these quantities.

1 Introduction

Evolution of the physical conditions of the universe, galaxies and stars can be described in terms of the increase or decrease of hundreds of elemental abundances of atomic nuclides. They originate from the primordial nucleosynthesis about fifteen billion years ago and the subsequent production/destruction cycle in stars and ejection into the intergalactic space. It is therefore inevitable and even fundamental to study the nuclear processes in several astrophysical sites for a deep understanding of the evolution of the universe.

Cosmologically, the primordial nucleosynthesis provides a unique method to determine the average universal mass–density parameter $\Omega_B$. Although the homogeneous big–bang model for primordial nucleosynthesis predicts $\Omega_B h^2_0 \sim 0.04$, X–ray observations of dense clusters have indicated that $\Omega_B$ could be as large as $\lesssim 0.15$. Recent MACHO detections also suggest that there exist more baryons in our Galaxy than ever expected. There is clearly a serious potential conflict between these observations and the theoretical prediction in the homogeneous big–bang model. The situation is even crucial if high deuterium abundances, which were detected in Lyman–α absorption systems along the
line of sight to high red–shifted quasars, are presumed to be primordial. On the other hand, an inhomogeneous big–bang model, which allows inhomogeneous baryon density distribution, can predict $\Omega_B h^2_{100} \sim 0.1 - 0.2$. Among the possible observable signatures of baryon inhomogeneous cosmologies are the high abundances of heavier elements than lithium such as beryllium and boron.

In an environment of baryon inhomogeneous distribution, neutrons can easily diffuse out of the fluctuations to form high density proton–rich and low density neutron–rich regions, where a lot of proton/neutron–rich radioactive isotopes can help produce the intermediate–to–heavy mass elements.

Another astrophysical site where the neutron–rich isotopes may play a significant role in nucleosynthesis is the $\alpha$–process occurring in supernovae. The nucleosynthesis in the high–entropy bubble is thought to proceed as follows. Due to the high temperature, the previously produced nuclei up to iron will be destroyed again by photodisintegration. At temperatures of about $10^{10}$ K the nuclei would be dismantled into their constituents, protons and neutrons. At slightly lower temperatures one is still left with $\alpha$–particles. During the subsequent cooling of the plasma the nucleons will recombine again, first to $\alpha$–particles, then to heavier nuclei. Depending on the exact temperatures, densities and the neutron excess, quite different abundance distributions can be produced in this $\alpha$–rich freeze–out (sometimes also called $\alpha$–process). Temperature and density are dropping quickly in the adiabatically expanding high–entropy bubble. This will hinder the recombination of $\alpha$–particles into heavy nuclei, leading in some scenarios to a high neutron density for an $r$–process, at the end of the $\alpha$–process after freeze–out of charged particle reactions.

These astrophysical motivations have led us to critically study the role of radiative neutron capture reactions by neutron–rich Li– and Be–isotopes theoretically in explosive nucleosynthesis. Since it is the focus in recent years to study the spatial and time evolution of the nuclear reactions dynamics by the use of radioactive nuclear beams, our theoretical studies are also being tested experimentally.

2 Calculation of Radiative–Capture Cross Sections

Nuclear burning in explosive astrophysical environments produces unstable nuclei which can again be targets for subsequent reactions. In addition, it involves a very large number of stable nuclei which are not yet fully explored by experiments. Thus, it is necessary to be able to predict reaction cross sections and thermonuclear rates with the aid of theoretical models.

In astrophysically relevant nuclear reactions two important reaction mechanisms take place. These two mechanisms are compound–nucleus reactions (CN) and direct reactions (DI). The reaction mechanism and therefore also
the reaction model depends on the number of levels in the CN. If one is consid-
ering only a few CN resonances the R–matrix theory is appropriate. In the case of single resonance the R–matrix theory reduces to the simple phenomenological Breit–Wigner formula. If the level density of the CN is so high that there are many overlapping resonances, the CN mechanism will dominate and the statistical HF–model can be applied. Finally, if there are no CN resonances in a certain energy interval the DI mechanism dominates and one can use DI models, like Direct Capture (DC).

In the case of a single isolated resonance the resonant part of the cross section is given by the well–known Breit–Wigner formula:

\[
\sigma_{r}(E) = \frac{\pi \hbar^{2}}{2 \mu E} \frac{(2J + 1)}{(2j_{p} + 1)(2j_{t} + 1)} \frac{\Gamma_{in} \Gamma_{out}}{(E_{r} - E)^{2} + \frac{\Gamma_{tot}^{2}}{4}},
\]

where \(J\) is the angular momentum quantum number and \(E_{r}\) the resonance energy. The partial widths of the entrance and exit channels are \(\Gamma_{in}\) and \(\Gamma_{out}\), respectively. The total width \(\Gamma_{tot}\) is the sum over the partial widths of all channels. One important aspect is that the particle width \(\Gamma_{p}\) can be related to spectroscopic factors \(S\) and the single–particle width \(\Gamma_{s,p}\).

\[
\Gamma_{p} = C^{2} S \Gamma_{s,p},
\]

where \(C\) is the isospin Clebsch–Gordan coefficient. The single–particle width \(\Gamma_{s,p}\) can be calculated from the scattering phase shifts of a scattering potential with the potential depth determined by matching the resonance energy.

The nonresonant part of the cross section can be obtained using the DC model:

\[
\sigma_{nr} = \sum_{c} C^{2} S_{c} \sigma_{c}^{DC},
\]

The sum extends over all bound states in the final nuclei. The DC cross sections \(\sigma_{c}^{DC}\) are essentially determined by the overlap of the scattering wave function in the entrance channel, the bound–state wave function in the exit channel and the multipole transition–operator.

The total cross section can be calculated by summing over the resonant (Eq. 1) and nonresonant parts (Eq. 3) of the cross section (if the widths of the resonances are broad, also an interference term has to be added). For both parts the spectroscopic factors have to be known. They can be obtained from other reactions, e.g., the spectroscopic factors necessary for calculating \(A(n,\gamma)B\) can be extracted from the reaction \(A(d,p)B\). The \(\gamma\)–widths can be extracted from reduced electromagnetic transition strengths. For unstable nuclei
where only limited or even no experimental information is available, the spectroscopic factors and electromagnetic transition strengths can also be extracted from nuclear structure models like the shell model (SM).

The most important ingredients in the potential models are the wave functions for the scattering and bound states in the entrance and exit channels. This is the case for the DC cross sections $\sigma_{\text{DC}}^c$ in Eq. 3 as well as for the calculation of the single–particle width $\Gamma_i$ in Eq. 2. For the calculation of these wave functions we use real folding potentials which are given by:

$$V(R) = \lambda V_F(R) = \lambda \int \rho_a(r_1) \rho_A(r_2) v_{\text{eff}}(E, \rho_a, \rho_A, s) \, dr_1 \, dr_2 ,$$  \tag{4}

with $\lambda$ being a potential strength parameter close to unity, and $s = |\mathbf{R}+\mathbf{r}_2-\mathbf{r}_1|$, where $R$ is the separation of the centers of mass of the projectile and the target nucleus. The density can be derived from measured charge distributions or from nuclear structure models (e.g., Hartree–Fock calculations) and the effective nucleon–nucleon interaction $v_{\text{eff}}$ has been taken in the DDM3Y parametrization. The imaginary part of the potential is very small because of the small flux into other reaction channels and can be neglected in most cases involving neutron capture by neutron–rich target nuclei.

3 Reaction Rates for Li– and Be–Isotopes

The parameters for the resonant and nonresonant contributions to the reaction rates are listed in Tables III and IV respectively. In the tables we give experimental values if available. Otherwise the excitation energies, spectroscopic factors, neutron– and $\gamma$–widths were calculated with the shell model. We used the code OXBASH for the calculations. For normal parity states we employed the interaction (8–16)POT of Cohen and Kurath. For nonnormal parity states we used the WBN interaction of Warburton and Brown.

With Eq. 3 the resonant reaction rate can be derived as

$$N_A \langle \sigma v \rangle_r = 1.54 \times 10^5 \mu^{-3/2} T_9^{-3/2} \sum_i \omega \gamma_i \exp(-11.605E_r/T_9) \, \text{cm}^3 \, \text{mole}^{-1} \, \text{s}^{-1} ,$$  \tag{5}

where $T_9$ is the temperature in $10^9 \text{K}$, $E_r$ the resonance energy in the c.m. system (in MeV), and the resonance strength $\omega \gamma$ (in eV) is given by

$$\omega \gamma = \frac{2J+1}{(2J_p+1)(2J_1+1)} \frac{\Gamma_{\text{in}} \Gamma_{\text{out}}}{\Gamma_{\text{tot}}} .$$  \tag{6}
Table 1: Resonance parameters

| Reaction          | $E_x$  | $E_n$  | $J^\pi$ | $\Gamma_n$ | $\Gamma_\gamma$ | $\omega_\gamma$ |
|-------------------|--------|--------|---------|-------------|------------------|-----------------|
| $^7\text{Li}(n,\gamma)^8\text{Li}$ | 2.26   | 0.227  | 3$^+$   | $3.1 \times 10^4$ | 0.07             | 0.061           |
| $^8\text{Li}(n,\gamma)^9\text{Li}$ | 4.31   | 0.247  | $5/2^-$ | $1 \times 10^5$  | 0.11             | 0.066           |
| $^9\text{Be}(n,\gamma)^{10}\text{Be}$ | 7.371  | 0.559  | 3$^-$   | $1.57 \times 10^4$ | 0.661            | 0.578           |
|                   | 7.542  | 0.73   | 2$^+$   | $6.3 \times 10^4$ | 0.814            | 0.509           |

The partial widths of the entrance and exit channel, $\Gamma_in$ and $\Gamma_out$, are in the case of $(n,\gamma)$–reactions the neutron– and $\gamma$–widths. Since the neutron width is usually much larger than the $\gamma$–width, the total width $\Gamma_{tot}$ is practically identical with the neutron–width.

In Table 1 we list the excitation energies, resonance energies, neutron– and $\gamma$–widths and the resonance strengths of the resonances.

The nonresonant capture cross section is parametrized as

$$\sigma_{nr}(E) = A/\sqrt{E} + B\sqrt{E} - CE^D,$$

with $[A] = \mu b \text{MeV}^{1/2}$, $[B] = \mu b \text{MeV}^{-1/2}$, and $[C] = \mu b \text{MeV}^{-D}$. The parameters $A, B, C$ and $D$ are listed in Table 2. Using this equation, we obtain for the reaction rate

$$N_A\langle \sigma v \rangle_{nr} = \left( 836.565A\mu^{-1/2} + 108.130B\mu^{-1/2}T_9 \right. \\
\left. -277.097C\mu^{-1/2}\frac{\Gamma(2+D)}{11.605^{D/2}}T_9^{D+1/2} \right) \text{cm}^3 \text{s}^{-1} \text{mole}^{-1},$$

where $\mu$ is the reduced mass in units of the atomic mass unit and $\Gamma(z)$ is the Euler gamma function.

The total reaction rate is given as the sum of the resonant (Eq. 5 and nonresonant (Eq. 8) part.

3.1 $^7\text{Li}(n,\gamma)^8\text{Li}$

The cross section of the reaction $^7\text{Li}(n,\gamma)^8\text{Li}$ is well known (see, e.g., 4). The cross section is dominated by s–wave capture to the $^8\text{Li}$ ground state and a resonance at 227 keV neutron energy. Using the spectroscopic factors of Cohen and Kurath 5 yields a thermal cross section of $8.2 \times 10^{-2} \text{b}$, which is a factor 1.8 higher than the experimental value of $4.54 \times 10^{-2} \text{b}$. The shell model calculation is purely p–shell and does not include excitations to other oscillator
Table 2: Parametrization of the nonresonant cross section (see text).

|       | A   | B      | C   | D   | \(\sigma_{nr}(\mu b)\) at 30 keV |
|-------|-----|--------|-----|-----|----------------------------------|
| \(^7\)Li(n, \(\gamma\))\(^8\)Li | 6.755\(^a\) | —     | —   | —   | 39.000                          |
| \(^8\)Li(n, \(\gamma\))\(^9\)Li | 2.909 | —     | —   | —   | 16.795, 30.392                  |
| \(^9\)Be(n, \(\gamma\))\(^{10}\)Be | 1.147\(^a\) | 11.000 | 6.815 | 0.962 | 8.294, 6.622                   |
| \(^{10}\)Be(n, \(\gamma\))\(^{11}\)Be | 0.132 | 24.000 | 15.725 | 0.914 | 4.281, 3.943                   |
| \(^{11}\)Be(n, \(\gamma\))\(^{12}\)Be | —     | 7.000  | 4.851 | 0.887 | 0.996, 2.373                   |

\(^a\)extracted from experimental thermal cross section\(^14\).

shells. Therefore the spectroscopic amplitude of 0.977 for a \(p_{3/2}\)–transition to the ground state of \(^8\)Li might be too high.

For the resonance, however, we find excellent agreement between calculation and experiment. The calculated width — using the folding potential and spectroscopic amplitudes from Cohen and Kurath\(^15\) — is 28.9 keV, almost identical to the known value of 31 ± 7 keV.

3.2 \(^8\)Li(n, \(\gamma\))\(^9\)Li

The resonance at 247 keV is a \(5/2^-\) state\(^13\). With a total width of 100 keV the resonance strength is determined by the \(\gamma\)–width which was previously estimated with 0.56 eV\(^18\). A shell model calculation yielded a width \(\Gamma_\gamma = 0.11\) eV. Therefore the resonance strength is a factor 5 smaller than previously assumed\(^19\).

The calculated thermal cross section, resulting from \(s\)–wave capture to the ground state and first excited state in \(^9\)Li, is \(1.94 \times 10^{-2}\) b and is smaller than the value of \(3.51 \times 10^{-2}\) b given by Rauscher \textit{et al.}\(^19\).

3.3 \(^9\)Be(n, \(\gamma\))\(^{10}\)Be

Like in the reaction \(^7\)Li(n, \(\gamma\))\(^8\)Li the spectroscopic factors of Cohen and Kurath\(^13\) are a little too high. The thermal cross section is dominated by the transition to the \(^{10}\)Be ground state with a theoretical spectroscopic factor of 2.36. With this value the calculated thermal cross section is \(1.06 \times 10^{-2}\) b, compared to experimental cross section of \(7.6 \times 10^{-3}\) b. With the spectroscopic factor given by Mughabghab\(^20\) of 1.45 the calculated cross section would be close to the experimental value. For high temperatures the \(p\)–wave capture to excited states has to be taken into account.
Two resonances are known at 559 keV and 730 keV. The total widths are known experimentally. We have calculated the $\gamma$–widths which were only estimated previously. Both resonance strengths are larger than the previous estimates, for the 559 keV resonance the enhancement is one order of magnitude.

With the higher resonance strengths and the p–wave contribution the reaction rate is clearly higher compared to Ref. 19.

3.4 $^{10}\text{Be}(n,\gamma)^{11}\text{Be}$

Cross section and reaction rate of this reaction were recently determined experimentally with the help of the inverse Coulomb dissociation. They supported their experimental values by a direct capture calculation. In order to reproduce the experimental data they enhanced the spectroscopic factors to the $^{11}\text{Be}$ ground state by 20%. Our calculation confirms the results. Using the spectroscopic factors from the $(\text{d},p)$–reaction the calculated cross section is a little smaller than the experimental. The results are grossly different from the rate given in Ref. 19.

3.5 $^{11}\text{Be}(n,\gamma)^{12}\text{Be}$

There is no resonant contribution to the reaction rate. The transition is a p–wave capture from the 1/2$^+$ ground state of $^{11}\text{Be}$ to the ground state of $^{12}\text{Be}$ and the 0$^+$ state at 2.7 MeV excitation energy, while the transition to the 2$^+$ state at 2.1 MeV is negligible.

4 Discussion

The new reaction rates could change the reaction flow in the inhomogeneous big bang nucleosynthesis. The smaller rate for $^{8}\text{Li}(n,\gamma)^{9}\text{Li}$ could mean that the main reaction flow will proceed through the reaction $^{8}\text{Li}(\alpha,n)^{11}\text{Be}$. The higher rate for $^{9}\text{Be}(n,\gamma)^{10}\text{Be}$ might give more importance to this reaction. Detailed network calculations with the new rates are planned for the near future.

Acknowledgments

We thank the Fonds zur Förderung der wissenschaftlichen Forschung in Österreich (project S7307–AST) and the Österreichische Nationalbank (project 5054) for their support.
References

1. T. Kajino and R.N. Boyd, *Astrophys. J.* **359**, 267 (1990).
2. R.G. Breit and E.P. Wigner, *Phys. Rev.* **49**, 519 (1936).
3. J.M. Blatt and V.F. Weisskopf, *Theoretical Nuclear Physics*, (Wiley & Sons, New York, 1962).
4. M. Wiescher and K.-U. Kettner, *Astrophys. J.* **263**, 891 (1982).
5. H. Herndl, J. Görres, M. Wiescher, B.A. Brown and L. van Wormer, *Phys. Rev. C* **52**, 1078 (1995).
6. K.H. Kim, M.H. Park and B.T. Kim, *Phys. Rev. C* **35**, 363 (1987).
7. H. Oberhummer and G. Staudt, in *Nuclei in the Cosmos*, ed. H. Oberhummer, (Springer–Verlag, Berlin, New York, 1991) p. 29.
8. P. Mohr, H. Abele, R. Zwiebel, G. Staudt, H. Krauss, H. Oberhummer, A. Denker, J.W. Hammer and G. Wolf, *Phys. Rev. C* **48**, 1420 (1993).
9. A.M. Kobos, B.A. Brown, R. Lindsay and G.R. Satchler, *Nucl. Phys. A* **425**, 205 (1984).
10. H. de Vries, C.W. de Jager and C. de Vries, *At. Data Nucl. Data Tables* **36**, 495 (1987).
11. B.A. Brown, A. Etchegoyen, W.D.M. Rae and N.S. Godwin, code OXBASH, 1984 (unpublished).
12. S. Cohen and D. Kurath, *Nucl. Phys.* **73**, 1 (1965).
13. E.K. Warburton and B.A. Brown, *Phys. Rev. C* **46**, 923 (1992).
14. J.C. Blackmon, A.E. Champagne, J.K. Dickens, J.A. Harvey, M.A. Hofstee, S. Kopecky, D.C. Larson, D.C. Powell, S. Raman and M.S. Smith, *Phys. Rev. C* **54**, 383 (1996).
15. S. Cohen and D. Kurath, *Nucl. Phys. A* **101**, 1 (1967).
16. V.F. Sears, *Neutron News*, Vol. 3, 26 (1992)
17. A.G.M. van Hees and P.M.W. Glaudemans, *Z. Phys. A* **315**, 223 (1984).
18. R.A. Malaney and W.A. Fowler, *Ap. J.* **333**, 14 (1988).
19. T. Rauscher, J.H. Applegate, J.J. Cowan, F.-K. Thielemann and M. Wiescher, *Ap. J.* **429**, 499 (1994).
20. S.F. Mughabghab, *Phys. Rev. Lett.* **54**, 986 (1985).
21. A. Mengoni, T. Otsuka, T. Nakamura and M. Ishihara, *Nucl. Phys. A* (Proc. Suppl.), in print.
22. F. Ajzenberg–Selove, *Nucl. Phys. A* **506**, 1 (1990).