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The astrophysical reaction ^{8}Li(n, γ) ^{9}Li from measurements by reverse kinematics

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Abstract. We study the breakup of ^{9}Li projectiles in high-energy (28.5 MeV/u) collisions with heavy nuclear targets (^{208}Pb). The wavefunctions are calculated using a simple potential model for ^{9}Li. A good agreement with the measured data is obtained with insignificant E2 contribution.

The reaction ^{8}Li(n, γ)^{9}Li competes with the ^{8}Li(α, n)^{11}B reaction in the inhomogeneous Big Bang model, which assumes the existence of proton-rich and neutron-rich regions of space during the early Big Bang. This reaction also competes with the ^{8}Li β-decay in the r-process. Depending on the rate, this reaction may affect the primordial abundance and reduce the stellar production of $A > 12$ nuclei by as much as 50% [1]. The Coulomb dissociation method [2, 3] has proven to be a useful tool for extracting radiative capture reaction cross sections of relevance in nuclear astrophysics. In particular, it appears that the Coulomb dissociation of ^{9}Li is very useful [4] for elucidating the role of the inhomogeneous nucleosynthesis in the Big Bang model—the formation of ^{9}Li via the ^{8}Li(n, γ)^{9}Li reaction. However, a few lingering questions still need to be addressed, including the importance of E2 excitations for the MSU experiment [4], performed at approximately 28.5 MeV/u. We attempt to resolve this issue by using a relatively simple but still realistic nuclear model that, however, yields agreement with data and suggests that E2 excitations are negligible for the kinematical conditions of the MSU experiment.

For ^{9}Li we adopt a single-particle model in which the $J_0 = 3/2^-$ ground state can be described as a $j_0 = \mathrm{p}_{3/2}$ neutron interacting with the ^{8}Li core. The core is assumed to be inert with an intrinsic spin $I_c = 2^+$. A realistic value for the spectroscopic factor for this configuration is about 0.94 [5]. Since we are only interested in a rough estimate of the E2 contribution in the Coulomb breakup of ^{9}Li, we take $S = 1$.

The single particle states, $R_{E_x lj}(r)$, for the excitation energy $E_x = E_n + |E_0|$, where E_n is the neutron–^{8}Li relative energy, are found by solving the Schrödinger equation with a nuclear + spin–orbit potential, given by

$$V(r) = V_0 \left[1 - F_{s.o.}(\boldsymbol{l} \cdot \boldsymbol{s}) \frac{r_0}{r} \frac{\mathrm{d}}{\mathrm{d}r}\right] f(r), \qquad f(r) = \left[1 + \exp\left(\frac{r-R}{a}\right)\right]^{-1}, \tag{1}$$

with parameters $a = 0.52$ fm and $r_0 = 1.25$ fm. The radius of the ^{8}Li core, $R = 2.499$ fm, was chosen so that the radiative capture cross section would fall in the range of the experiment of [4]. The spin–orbit strength is set to $F_{s.o.} = 0.351$ fm. This is consistent with the choice

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for the ^{8}B model used in [6] but is slightly smaller than the value 0.38 fm which works well for the low-lying spectra of ^{17}O, ^{13}C, and ^{11}Be [7]. The well depth for the ground state $l_0(j_0 I_c)J_0 = (p_{3/2}, 2^+)3/2^-$, was adjusted to reproduce the one-neutron separation energy, $E_0 = -4.05$ MeV, and it is $V_0((p_{3/2}, 2^+)3/2^-) = -45.301$ MeV. A neutron-unbound state exists in ^{9}Li at 4.31 MeV [8, 9] above the neutron decay threshold. Its predicted spin parity is $5/2^-$. These are described as $p_{3/2}$ waves coupled to the ground state of the core, and the well depth for this channel, $V_0 = -33.814$ MeV has been adjusted to reproduce the resonance energy and its predicted [9] width, $\Gamma = 0.118$ MeV. For all other partial waves ($s_{1/2}$, $p_{1/2}$, $d_{3/2}$, etc) we choose identical well depths and set them equal to the value -45.301 MeV, as for the ground state.

To compute the S-factors for the capture process $b + x \to a$ we have used the first-order perturbation theory. The matrix elements for electric multipole states (we couple angular momentum as $\boldsymbol{l} + \boldsymbol{s} = \boldsymbol{j}$, and $\boldsymbol{j} + \boldsymbol{I}_c = \boldsymbol{J}$) are given by

$$\langle JM|\mathcal{M}(E\lambda\mu)|J_0M_0\rangle = \frac{e_\lambda}{\sqrt{4\pi}}(-1)^{j_0+I_c+J_0+l_0+l+\lambda}\frac{\hat{\lambda}\hat{j}_0\hat{J}_0}{\hat{j}}$$
$$\times (J_0M_0\lambda\mu|JM)\left(j_0\lambda\frac{1}{2}0|j\frac{1}{2}\right)\begin{Bmatrix} j & J & I_c \\ J_0 & j_0 & \lambda \end{Bmatrix}\mathcal{O}(1\to 2;\lambda) \tag{2}$$

where $\mathcal{O}(1 \to 2; \lambda) = \int \phi^J_{lj}(r)\phi^{J_0}_{E_0 l_0 j_0}(r) r^\lambda \, \mathrm{d}r$ is the overlap integral, $\hat{m} \equiv \sqrt{2m+1}$, and $e_\lambda = Z_b e(-A_x/A_a)^\lambda + Z_x e(A_b/A_a)^\lambda$ is the effective electric charge ($b \equiv {}^8$Li, $x \equiv$ n). Here $R_{E_x lj}(r)$ is the radial wavefunction for the relative motion of the neutron and the core, normalized to $\sqrt{2m_{bx}/\pi\hbar^2 k}\,\sin(kr + \delta_{lj})$ at large r, where k is the relative momentum.

The response functions (multipole strengths) for the excitation of ^{9}Li are obtained by summing over all partial waves

$$\frac{\mathrm{d}B(E\lambda)}{\mathrm{d}E_x} = \sum_{lj}\frac{\mathrm{d}B(E\lambda; l_0 j_0 \to E_x lj)}{\mathrm{d}E_x}, \tag{3}$$

where

$$\frac{\mathrm{d}B(E\lambda; l_0 j_0 \to E_x lj)}{\mathrm{d}E_x} = \frac{m_{bx}}{\hbar^2 k}\sum_J (2J+1)\begin{Bmatrix} j & J & I_c \\ J_0 & j_0 & \lambda \end{Bmatrix}|\langle j\|\mathcal{M}(E\lambda)\|j_0\rangle|^2_J. \tag{4}$$

The sum over J in the last equality reduces to $|\langle j\|\mathcal{M}(E\lambda)\|j_0\rangle|^2/(2j_0+1)$ if the single-particle matrix elements are independent of the channel spin J.

The multipole strengths are presented in figure 1 for E1 and E2 excitations, as functions of the neutron energy, $E_n = \hbar^2 k^2/2m_{bx}$. They are dominated by s-wave components: the higher angular momentum waves become relevant for higher energies, as expected. This result has also been obtained in previous calculations [8, 10]. The quadrupole strength shows a resonance peak located at the excitation energy, $E_x = 0.247$ MeV. It is divided by ten in order to be displayed in the same figure.

The cross sections for direct capture are given by

$$\sigma^{(\lambda)}_{DC}(E_x) = \frac{(2I_c+1)}{(2J_0+1)}\frac{(2\pi)^3(\lambda+1)}{\lambda[(2\lambda+1)!!]^2}\frac{1}{k^2}\left(\frac{E_x}{\hbar c}\right)^{2\lambda+1}\frac{\mathrm{d}B(E\lambda)}{\mathrm{d}E_x}. \tag{5}$$

Figure 2 displays the direct capture cross section for the E1 transitions. The experimental data are from [4]. Due to the factor $(E_x/\hbar c)^{2\lambda+1}$ appearing in equation (5) (and since the peak of $\frac{\mathrm{d}B(E2)}{\mathrm{d}E}$ fm^{-4} is one order of magnitude greater than $\frac{\mathrm{d}B(E1)}{\mathrm{d}E}$ fm^{-2}, as seen from figure 1), the E2 contribution to the radiative capture cross section is about 3–4 orders of magnitude smaller than the E1 contribution.

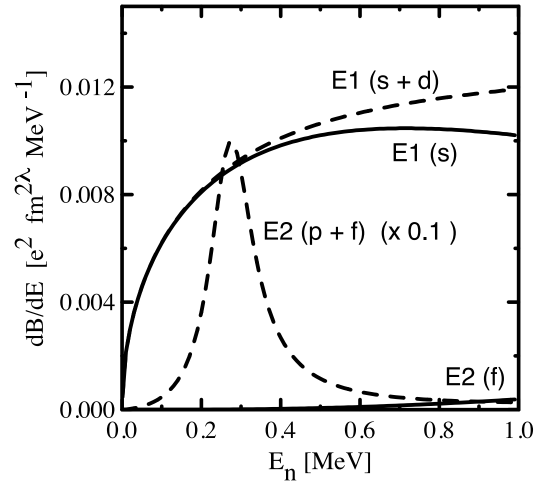

Figure 1. Response function (equation (3)) in units of e^2 fm^2 MeV^{-1} (e^2 fm^4 MeV^{-1}) for E1(E2) transitions in the reaction ^{9}Li(γ, n)^{8}Li, as a function of the neutron energy relative to the ^{8}Li core. The lower curves are the f (solid) and p + f (dashed) waves contribution, while the upper curves display the contribution of s (solid) and s+d (dashed) waves.

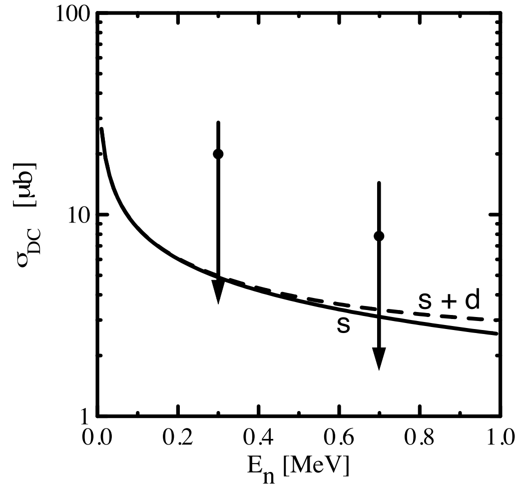

Figure 2. Radiative capture cross sections, in μb, for the reaction ^{8}Li(n, γ)^{9}Li in the direct capture model. The solid curve is obtained with s-waves, while the dashed curve includes d-wave transitions.

Since there are no data for the elastic scattering of ^{9}Li on Pb targets at this bombarding energy, we construct an optical potential using an effective interaction of the M3Y type [12,13] modified so as to reproduce the energy dependence of total reaction cross sections, i.e. [13],

$$t(E,s) = -\mathrm{i}\frac{\hbar v}{2t_0}\sigma_{NN}(E)[1-\mathrm{i}\alpha(E)]t(s), \tag{6}$$

where $t_0 = 421$ MeV fm^3 is the volume integral of the M3Y interaction $t(s)$, s is the nucleon–nucleon separation distance, v is the projectile velocity, σ_{NN} is the nucleon–nucleon cross section, and α is the real-to-imaginary ratio of the forward nucleon–nucleon scattering amplitude. At 28.5 MeV/nucleon, we use $\sigma_{NN} = 20$ fm^2 and $\alpha = 0.87$.

The optical potential is given by

$$U(E,\boldsymbol{R}) = \int \mathrm{d}^3r_1\,\mathrm{d}^3r_2\,\rho_P(\boldsymbol{r}_1)\rho_T(\boldsymbol{r}_2)t(E,s), \tag{7}$$

where $\boldsymbol{s} = \boldsymbol{R} + \boldsymbol{r}_2 - \boldsymbol{r}_1$, and ρ_T (ρ_P) is the ground state density of the target (projectile).

Following [14], the Coulomb amplitude is given by

$$f_C = \sum_{\lambda\mu} f_{\lambda\mu}^{(JM)}, \tag{8}$$

where

$$f_{\lambda\mu}^{(JM)} = \mathrm{i}^{1+\mu}\frac{Z_T e m_{PT}}{\hbar^2}\left(\frac{E_x}{\hbar c}\right)^{\lambda}\sqrt{2\lambda+1}\exp\{-\mathrm{i}\mu\phi\}\Omega_\mu(q)$$
$$\times G_{E\lambda\mu}\left(\frac{c}{v}\right)\langle JM|\mathcal{M}_{E\lambda,-\mu}|J_0M_0\rangle, \tag{9}$$

$$\Omega_\mu(q) = \int_0^\infty \mathrm{d}b\, b J_\mu(qb) K_\mu\left(\frac{E_x b}{\gamma\hbar v}\right)\exp(\mathrm{i}\chi(b)), \tag{10}$$

$q = 2k\sin(\theta/2)$, θ (ϕ) is the (azimuthal) scattering angle, and m_{PT} is the reduced mass of the target + projectile. $J_\mu(K_\mu)$ is the cylindrical (modified) Bessel function of order μ, and the functions $G_{\pi\lambda\mu}(c/v)$ are tabulated in [15]. The angular momentum algebra connects $\langle JM|\mathcal{M}_{E\lambda,-\mu}|J_0M_0\rangle$ with the reduced matrix elements of equation (2).

The eikonal phase, $\chi(b)$, is given by

$$\chi(b) = 2\eta\ln(kb) - \frac{1}{\hbar v}\int_{-\infty}^{\infty}\mathrm{d}z\, U_{opt}(R), \tag{11}$$

where $\eta = Z_P Z_T e^2/\hbar v$, $\hbar k$ is the projectile momentum, and $R = \sqrt{b^2+z^2}$. The optical potential, U_{opt}, in the above equation is given by equation (7).

The cross section for Coulomb excitation of a state with angular momentum J and excitation energy E_x is obtained by an average (and a sum) over the initial (final) angular momentum projections:

$$\frac{\mathrm{d}^2\sigma_{E\lambda}}{\mathrm{d}\Omega\,\mathrm{d}E_x} = \frac{1}{2J_0+1}\sum_{M_0,M}|f_{\lambda\mu}^{(JM)}|^2. \tag{12}$$

As explained in detail in [2, 3, 14], the above cross section can be factorized in terms of a product of virtual photon numbers and breakup cross sections by real photons, which by detailed balance is directly related to the radiative capture cross sections. Thus a measurement of $\mathrm{d}\sigma_{E\lambda}/\mathrm{d}\Omega\,\mathrm{d}E_x$ can be used to obtain radiative capture cross sections of astrophysical interest. Experimentally, the nuclear contribution to the breakup cross section can be separated by repeating the measurement on light targets (see, e.g., [4]).

At the bombarding energies of tens of MeV/nucleon, the E2 virtual photon number is much larger than that of E1. As a consequence, even when the E2 contribution to the radiative capture cross section is small, it may be amplified in Coulomb breakup experiments [3]. A known example is the Coulomb breakup of ^{8}B [16] at 50 MeV/nucleon, which has been used to extract the S-factor for the radiative capture cross section ^{7}Be(p, γ)^{8}B at low energies. It has been claimed [17] that the E2 contribution accounts for as much as 20% of the differential cross sections for the experimental conditions. This poses an additional experimental problem, since one needs to separate the wanted E1 transition matrix elements from the unwanted E2 transition matrix elements. Nonetheless, it was shown in [18] that the E2 contribution in the Coulomb breakup of ^{8}B is much less than that predicted in [17].

In order to infer the relevance of E2 for the breakup of ^{9}Li (28.5 MeV) on lead targets, we plot in figure 3 the angle integrated cross section $\mathrm{d}\sigma_\lambda/\mathrm{d}E_x$, as a function of the neutron energy relative to ^{9}Li. The quadrupole contribution has been multiplied by ten in order to be more visible. We see that the E2 contribution to the breakup cross section is at least one order of magnitude smaller than the E1 contribution, even at the resonance region. We thus conclude that E2 transitions are not relevant in the experiment of [4]. Although we obtained this result by means of a single-particle model for the ^{8}Li(n, γ)^{9}Li reaction, we do not expect that it would change appreciably with more sophisticated models.

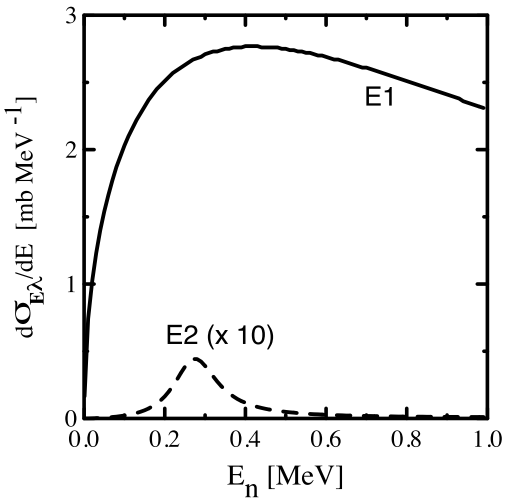

Figure 3. Coulomb breakup cross section $d\sigma_{E\lambda}/dE$ (in mb MeV^{-1}) for the reaction ^{9}Li (28.5 MeV/nucleon) + Pb $\longrightarrow$ ^{8}Li + n + Pb, as a function of the neutron–^{9}Li relative energy, in MeV. The E2 breakup contribution multiplied by ten is also shown.

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