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THE BARE ASTROPHYSICAL $S(E)$ FACTOR OF THE ^{7}Li(p, α)α REACTION

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ABSTRACT

The astrophysically important ^{7}Li(p, α)α reaction has been studied via the Trojan horse method in the energy range $E = 10$–400 keV. A new theoretical description, based on the distorted-wave Born approximation approach, allows one to extract information on the bare astrophysical S-factor, $S_b(E)$, with $S_b(0) = 55 \pm 3$ keV barns. The results are compared with direct experimental data leading to a model-independent value of the electron screening potential energy, $U_e = 330 \pm 40$ eV, much higher than the adiabatic limit $U_{\rm ad} = 175$ eV.

Subject heading: nuclear reactions, nucleosynthesis, abundances

1. INTRODUCTION

Nuclear reactions play a key role for the understanding of the early universe and of the evolution of stars (Fowler 1984). For this reason the cross section $\sigma(E_0)$ of a given reaction at the relevant thermal energy E_0 (the Gamow energy) must be known with a precision of 10% or better (all energies are given in the center-of-mass system except where quoted differently). In neutron-induced reactions, the lack of a Coulomb barrier and the typical $\sigma(E) \sim 1/E^{1/2}$ energy dependence led to a precise knowledge of $\sigma(E_0)$ for many cases. In comparison the Coulomb barrier of height E_c in charged-particle–induced reactions causes an exponential decrease of $\sigma(E)$ at $E < E_c$, $\sigma(E) \sim \exp(-2\pi\eta)$ (with η the Sommerfeld parameter), leading to a low-energy limit of direct $\sigma(E)$ measurements, which is typically much larger than E_0. Thus, one has to extrapolate the $\sigma(E)$ data determined at higher energies down to E_0 with the help of theory and other arguments. Of course, this extrapolation "into the unknown" may introduce a large uncertainty in $\sigma(E_0)$. The extrapolation is carried out usually for the astrophysical $S(E)$ factor defined as

$$S(E) = \sigma(E)E \exp(2\pi\eta) . \quad (1)$$

In recent years the availability of high-current low-energy accelerators, such as that at the underground Laboratori Nazionali del Gran Sasso (called LUNA; e.g., Bonetti et al. 1999), together with improved target and detection techniques have allowed us to perform $\sigma(E)$ measurements in some cases down to E_0 or at least close to E_0: e.g., ^{3}He(^{3}He, $2p$)^{4}He as low as $E = 16$ keV, while $E_0(\text{Sun}) = 21$ keV. Then, of course, no $\sigma(E)$ extrapolation is needed anymore for this reaction.

However, the measurements at ultralow energies suffer from the complication due to the effects of electron screening (Assenbaum, Langanke, & Rolfs 1987). This leads to an exponential increase of $\sigma(E)$ [or equivalently of $S(E)$] with decreasing energy relative to the case of bare nuclei. This can be described by an enhancement factor

$$f_{\rm lab}(E) = \frac{\sigma_s(E)}{\sigma_b(E)} = \exp\left(\frac{\pi\eta U_e}{E}\right) , \quad (2)$$

where $\sigma_s(E)$ and $\sigma_b(E)$ refer to the cross sections of electron-shielded and bare nuclei, respectively, and U_e is the electron screening potential energy. It should be pointed out that in the astrophysical environment the cross section under plasma conditions $\sigma_{\rm pl}(E)$ is related to the bare cross section by a similar enhancement factor

$$f_{\rm pl}(E) = \frac{\sigma_{\rm pl}(E)}{\sigma_b(E)} = \exp\left(\frac{\pi\eta U_{\rm pl}}{E}\right) . \quad (3)$$

This factor can be calculated if the plasma potential energy $U_{\rm pl}$ is known, which depends on detailed properties of the plasma, such as the Debye-Hückel radius. Clearly, a good understanding of U_e is needed in order to calculate $\sigma_b(E)$ from the experimental data $\sigma_s(E)$ using equation (2). In turn, the understanding of U_e may help to better understand $U_{\rm pl}$, needed to calculate $\sigma_{\rm pl}(E)$.

The quantity U_e may also be calculated from an atomic model as the difference in electron binding energies between the atoms of the entrance channel and of the composite atom. This model corresponds to the adiabatic approximation to the electronic screening potential $U_{\rm ad}$, which should be well fulfilled at ultralow energies. The low-energy data of several fusion reactions involving light nuclides have indeed shown the exponential enhancement according to equation (2) (Costantini et al. 2000 and references therein). However, the deduced U_e values were much higher in all cases than the adiabatic limit (Fiorentini, Kavanagh, & Rolfs 1995): e.g., ^{3}He(d, p)^{4}He with $U_e = 218 \pm 18$ eV and $U_{\rm ad} = 120$ eV (Formicola et al. 2001). The results are disturbing: if one does not understand the effects of electron screening under laboratory conditions, one might also not understand fully the effects under astrophysical conditions.

A weak point in the laboratory approach—and thus in the deduced U_e value—is the need for an assumption about

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the energy dependence of $\sigma_b(E)$ at ultralow energies. This is usually based on theoretical calculations or on the trend of the data at higher energies, where the effects of electron screening are negligible (e.g., $E/U_e \geq 100$). On the other hand, a direct measurement of $\sigma_b(E)$ using bare nuclei (e.g., a crossed beam setup) appears difficult if not impossible due to luminosity problems in combination with ultralow $\sigma(E)$ values. Thus, an alternative experimental method for the determination of $\sigma_b(E)$ in charged-particle–induced reactions appears highly desirable. One such method is the Trojan horse method (THM), which will be applied here to the two-body reaction ^{7}Li$(p, \alpha)\alpha$, which is a key reaction in primordial nucleosynthesis (Pinsonneault et al. 1992). The values for the astrophysical S-factor from the indirect methods can be compared to direct measurements (Rolfs & Kavanagh 1986; Engstler et al. 1992) of $S_b(0) \approx 58$ keV barns and $U_e \approx 300$ eV ($U_{\rm ad} = 175$ eV).

2. THE TROJAN HORSE METHOD

The idea of the THM (Baur 1986) is to extract the cross section of an astrophysically relevant two-body reaction

$$A + x \rightarrow C + c \tag{4}$$

(here: $^7\mathrm{Li} + p \rightarrow \alpha + \alpha$) at low energies from a suitably chosen three-body reaction

$$A + a \rightarrow C + c + b \tag{5}$$

(here: $^7\mathrm{Li} + d \rightarrow \alpha + \alpha + n$). This is done with the help of direct nuclear reaction theory assuming that the Trojan horse a is predominantly composed of the clusters x and b. In the present application, this assumption is trivially fulfilled: $a =$ deuteron, $x =$ proton, $b =$ neutron. In the three-body process the wave function of the Trojan horse $a = b + x$ has a momentum distribution of finite width, i.e., the clusters b and x have a Fermi motion. In the THM one attempts to choose the kinematic condition such that the momentum transfer to b is small and, therefore, the nucleus b can be considered as a spectator during the interaction. This condition is called the quasi-free condition. It can be fulfilled, at least in a given region of the available three-body phase space, by choosing suitable kinematic conditions among the ejectiles (here: α-α coincidences at selected detection angles) such that the momentum transfer to the spectator (here: the undetected neutron) remains restricted (here: $\Delta p_n \leq 40$ MeV c^{-1}). In this selected phase space other reaction mechanisms, as e.g., sequential processes with ^{5}He or ^{8}Be intermediate states, can be expected to be suppressed to a negligible level.

For high energies in the entrance channel of reaction (eq. [5]) (here: $E_{^7\mathrm{Li}} = 19$–20 MeV), the three-body cross section is not suppressed by the $A + a$ Coulomb barrier (here: $E_c = 1.0$ MeV). However, ultralow effective energies in the reaction (eq. [4]) can be reached due to the Fermi motion of the particle x inside a. Thus, the two-body cross section will not be suppressed by the $A + x$ Coulomb barrier—the essential trick of the THM—in contrast to direct measurements of the reaction (eq. [4]). In addition, due to the high energy in the entrance channel of reaction (eq. [5]), the effects of electron screening are negligible and, thus, one can determine the bare cross section $\sigma_b(E)$ down to ultralow energies. This is another important feature of the THM.

In view of various approximations involved in the THM one cannot precisely extract absolute values of the two-body cross section. However, it is possible to obtain reliable information on the energy dependence of $\sigma_b(E)$. Then the relative $\sigma_b(E)$ results of the THM can be normalized to the absolute $\sigma_b(E) \equiv \sigma(E)$ values of the direct data at higher energies, where the effects of electron screening are negligible, i.e., at $E/U_e \geq 100$. The agreement in the energy dependence of both data sets at $E/U_e \geq 100$ represents a sensitive test of the THM and a necessary condition for the credibility of the deduced $\sigma_b(E)$ THM values at $E/U_e < 100$.

The three-body cross section of the reaction (eq. [5]) can be formulated in the distorted-wave Born approximation (DWBA) in postrepresentation (for details, see Typel & Wolter 2000):

$$\frac{d^3\sigma}{dE_C\, d\Omega_C\, d\Omega_c} = KF\,|\,T_{\rm fi}\,|^2 , \tag{6}$$

with a kinematic factor

$$KF = \frac{\mu_{Aa} m_c}{(2\pi)^5 \hbar^7} \frac{p_C p_c^3}{p_{Aa}} \left[\left(\frac{\boldsymbol{p}_{Bb}}{\mu_{Bb}} - \frac{\boldsymbol{p}_{Cc}}{m_c}\right) \cdot \frac{\boldsymbol{p}_c}{p_c}\right]^{-1} ,$$

in the usual notation for (relative) momenta and (reduced) masses and where B denotes the $C + c$ system. The T-matrix element

$$T_{\rm fi} = \langle \chi_{Bb}^{(-)} \Psi_{Cc}^{(-)} \phi_b \,|\, V_{xb} \,|\, \chi_{Aa}^{(+)} \phi_A\, \phi_a \rangle \tag{7}$$

contains distorted waves $\chi_{Bb}^{(-)}$ and $\chi_{Aa}^{(+)}$, wave functions ϕ_a, ϕ_A, and ϕ_b of the Trojan horse a, and the nuclei A and b, the two-body scattering wave function $\Psi_{Cc}^{(-)}$, and the interaction V_{xb} between the transferred particle x and the spectator b. In a so-called surface approximation (Baur 1986) the wave function $\Psi_{Cc}^{(-)}$ is replaced by its asymptotic form which contains the (on-shell) S-matrix elements $S_l(Cc \rightarrow Ax)$ of the (inverse) two-body reaction in the various partial waves, which are related in the standard way to the cross section of the two-body reaction (eq. [4]).

In a plane wave (instead of distorted wave) approximation the relation of the three-body cross section to the two-body cross section can be written more explicitly:

$$\frac{d^3\sigma}{dE_C\, d\Omega_C\, d\Omega_c} = KF\,|\,W(\boldsymbol{Q}_{Bb})\,|^2\, \frac{16\pi^2}{k_{Ax} Q_{Aa}} \frac{v_{Cc}}{v_{Ax}} \frac{d\sigma^{\rm THM}}{d\Omega_{Ax}} \tag{8}$$

with momenta

$$\hbar \boldsymbol{Q}_{Aa} = \boldsymbol{p}_{Aa} - \frac{m_A}{m_A + m_x} \boldsymbol{p}_{Bb} , \tag{9}$$

$$\hbar \boldsymbol{Q}_{Bb} = \boldsymbol{p}_{Bb} - \frac{m_b}{m_b + m_x} \boldsymbol{p}_{Aa} , \tag{10}$$

and the momentum distribution $W(\boldsymbol{Q}_{Bb})$ of the product of the interaction potential and the wave function of a, $V_{xb}\, \phi_a$. The two-body THM cross section in equation (8) is given as

$$\frac{d\sigma^{\rm THM}}{d\Omega_{Ax}} (Cc \rightarrow Ax) = \frac{1}{k_{Cc}^2} \,|\sum_l (2l + 1) \times P_l(\hat{Q}_{Aa} \cdot \hat{k}_{Cc}) [S_l J_l^{(+)} - \delta_{(Ax)(Cc)} J_l^{(-)}]\,|^2 . \tag{11}$$

It is similar in form to a usual two-body cross section except for the functions $J_l^{(\pm)}$, which are given explicitly in Spitaleri et al. (2001). These can be well approximated at low relative energies by

$$J_l^{(\pm)} = D_l k_{Ax} Q_{Aa} R^2 e^{\mp i\sigma_l} \times [G_l(k_{Ax} R) \pm iF_l(k_{Ax} R)] , \quad (12)$$

where D_l is a proportionality constant, F_l and G_l are the regular and irregular Coulomb wave functions, σ_l is the Coulomb phase shift in partial wave l, and R is a cutoff radius originating from the plane-wave and surface approximations. In the present case we have used for R the sum of the nuclear radii, assuming for each nucleus $r = r_0 A^{1/3}$, with $r_0 = 1.4$ fm. The argument of the Legendre polynomial P_l is the cosine of the center-of-mass scattering angle for the two-body reaction. The expression (eq. [8]) is similar to the result of the plane wave impulse approximation (PWIA), which was applied in earlier applications of the THM to the ^{7}Li(p, α)α and ^{6}Li(d, α)α reactions (Cherubini et al. 1996; Calvi et al. 1997; Spitaleri et al. 1999; Aliotta et al. 2000). In PWIA the momentum distribution of the Trojan horse a and the off-shell two-body cross section appear as factors, but the off-shell cross section can only be related to the on-shell cross section of the two-body reaction by a heuristic approximation for the Coulomb penetrability effect. In our present DWBA approach for the two-body process, Coulomb effects are fully included, despite the use of a plane wave approximation in the three-body T-matrix element and the relation between on-shell and off-shell cross sections is evident.

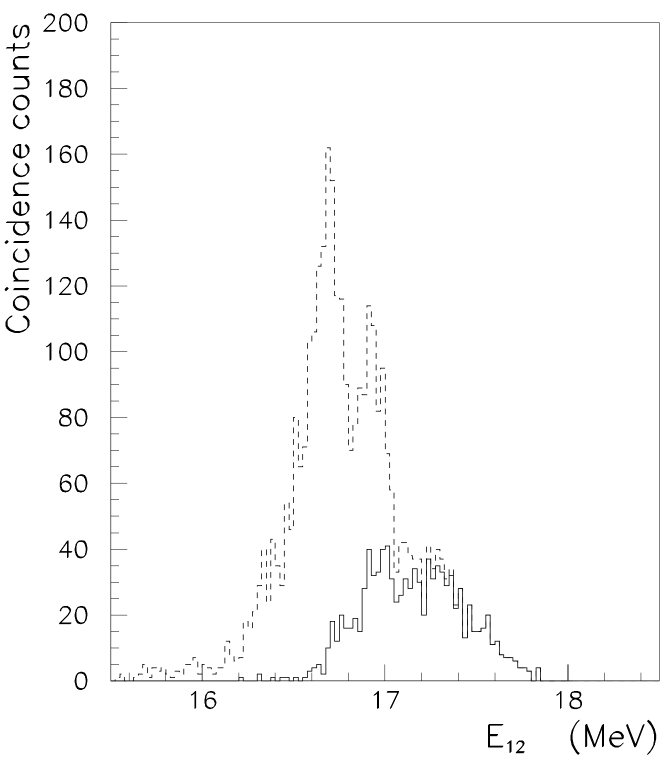

FIG. 1.—Coincidence events between the two α particles in the THM reaction $d(^7\mathrm{Li}, \alpha\alpha)n$ at $E_{^7\mathrm{Li}} = 20$ MeV for detectors at $\theta_{lab} = 45° \pm 0\overset{\circ}{.}75$ and $-45° \pm 0\overset{\circ}{.}75$, as a function of relative energy E_{12} between the two α particles. The solid histogram shows the results restricted by the condition of a low spectator momentum ($p_n \leq 40$ MeV c^{-1}) and represents predominantly the case of the quasi-free process. Without this restriction, the energy region of the quasi-free process overlaps with the energy region of the sequential decay via the 16.6 and 16.9 MeV compound states in ^{8}Be (*dashed histogram*).

In our application of the THM the ^{7}Li(p, α)α reaction is dominated at energies $E \leq 1$ MeV by the contribution from the $l = 1$ partial wave in the entrance channel. Neglecting the contributions from other partial waves we can write

$$\frac{d^3\sigma}{dE_{\alpha_1}\, d\Omega_{\alpha_1}\, d\Omega_{\alpha_2}} = KF\,|\,W(\boldsymbol{Q}_{(\alpha_1\,\alpha_2)n})\,|^2 \times \frac{C_1}{\boldsymbol{T}_1}\,\frac{v_{\alpha_1\,\alpha_2}}{v_{\mathrm{Li}-p}}\,\frac{d\sigma_1}{d\Omega_{\mathrm{Li}-p}}\,(\alpha_1 + \alpha_2 \to {}^7\mathrm{Li} + p) \quad (13)$$

with a constant C_1 and the P-wave Coulomb penetrability $\boldsymbol{T}_1 = [G_1^2(k_{\mathrm{Li}-p} R) + F_1^2(k_{\mathrm{Li}-p} R)]^{-1}$. Introducing the astrophysical S-factor

$$S(E_{\mathrm{Li}-p}) = E_{\mathrm{Li}-p}\,\sigma(E_{\mathrm{Li}-p})\,\exp\,(2\pi\eta_{\mathrm{Li}-p})$$

for the ^{7}Li(p, α) α reaction, we obtain the form

$$\frac{d^3\sigma}{dE_{\alpha_1}\, d\Omega_{\alpha_1}\, d\Omega_{\alpha_2}} = KF\,|\,W[\boldsymbol{Q}_{(\alpha_1\,\alpha_2)n}]\,|^2 \times \frac{\mu_{\mathrm{Li}-p}\, k_{\mathrm{Li}-p}}{\mu_{\alpha_1\,\alpha_2}\, k_{\alpha_1\,\alpha_2}}\,\frac{\tilde{C}_1\,\exp\,(-2\pi\eta_{\mathrm{Li}-p})}{E_{\mathrm{Li}-p}\,\boldsymbol{T}_1(k_{\mathrm{Li}-p} R)} \times S(E_{\mathrm{Li}-p}) P_1^2(\hat{Q}_{\mathrm{Li}-d} \cdot \hat{k}_{\alpha_1\,\alpha_2}) , \quad (14)$$

which is the desired relation between the three-body cross section and the S-factor. The angular distribution of the ^{7}Li(p, α)α reaction is contained in the Legendre polynomials $P_1(\hat{Q}_{\mathrm{Li}-d} \cdot \hat{k}_{\alpha_1\,\alpha_2})$. We see that at low energies the suppression of the cross section by $\exp\,(-2\pi\eta_{\mathrm{Li}-p})$ is cancelled by the Coulomb penetrability factor $\boldsymbol{T}_1(k_{\mathrm{Li}-p} R)$. The momentum distribution of the $|\,W\,|^2$ is related to the Fourier transform of the ground state wave function Φ_a of the deuteron by

$$W[\boldsymbol{Q}_{(\alpha_1\,\alpha_2)n}] = \left(E_d - \frac{\hbar^2 Q^2_{(\alpha_1\alpha_2)n}}{2\mu_{p-n}}\right)\Phi_a[\boldsymbol{Q}_{(\alpha_1\,\alpha_2)n}] ,$$

where $E_d = -2.2$ MeV is the binding energy of the deuteron. We assume a Hulthén wave function for the deuteron

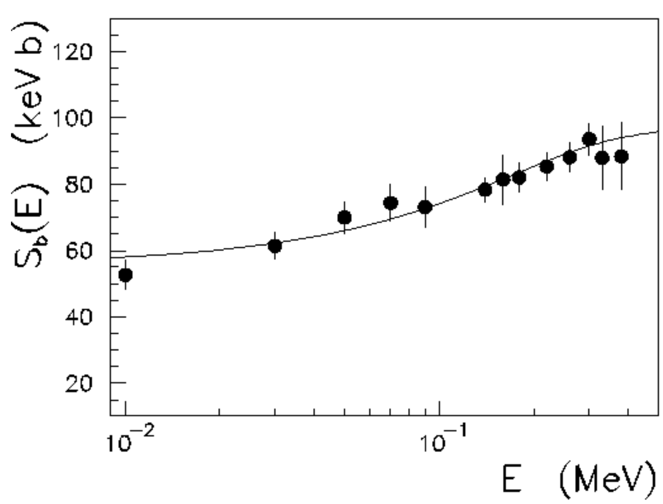

FIG. 2.—Bare astrophysical $S_b(E)$ factor of ^{7}Li(p, α)α deduced from the THM and normalized to the direct data (Rolfs & Kavanagh 1986; Engstler et al. 1992) at $E = 200$–400 keV. The solid curve represents a prediction (Aliotta et al. 2000).

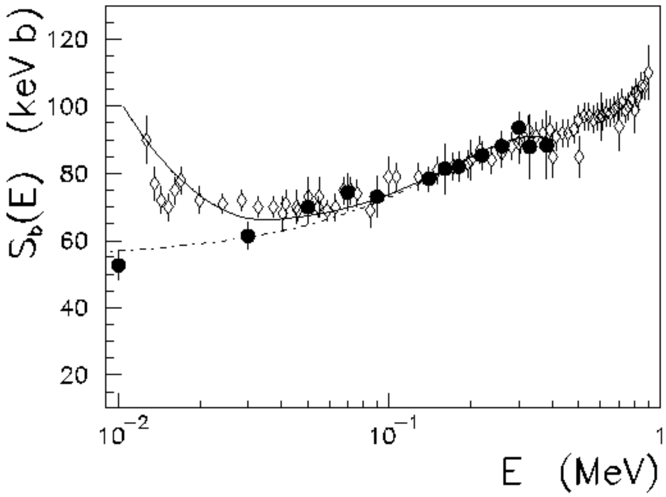

FIG. 3.—Bare $S_b(E)$ factor of ^{7}Li(p, α)α from the THM is compared with the screened $S_s(E)$ factor from direct data (Engstler et al. 1992), where the dashed curve represents a polynomial fitted to the $S_b(E)$ data, and the solid curve includes the effects of electron screening with $U_e = 330$ eV.

ground state

$$\Phi_a(\boldsymbol{r}_{p-n}) = \sqrt{\frac{ab(a+b)}{2\pi(a-b)^2}}\,\frac{1}{r_{p-n}}\,(e^{-ar_{p-n}} - e^{-br_{p-n}})\ ,$$

with parameters $a = 0.2317$ fm^{-1} and $b = 1.202$ fm^{-1} (Zadro et al. 1989) and find for the momentum wave function the result

$$\Phi_a[\boldsymbol{Q}_{(\alpha_1\,\alpha_2)n}] = \frac{1}{\pi}\sqrt{\frac{ab(a+b)}{(a-b)^2}} \times \left[\frac{1}{a^2 + Q^2_{(\alpha_1\alpha_2)n}} - \frac{1}{b^2 + Q^2_{(\alpha_1\alpha_2)n}}\right]. \quad (15)$$

The parametrization leads to a full width at half-maximum of 57.4 MeV for $|\Phi_a|^2$. This has also been used in earlier applications of the THM (Cherubini et al. 1996; Spitaleri et al. 1999, 2000; Aliotta et al. 2000).

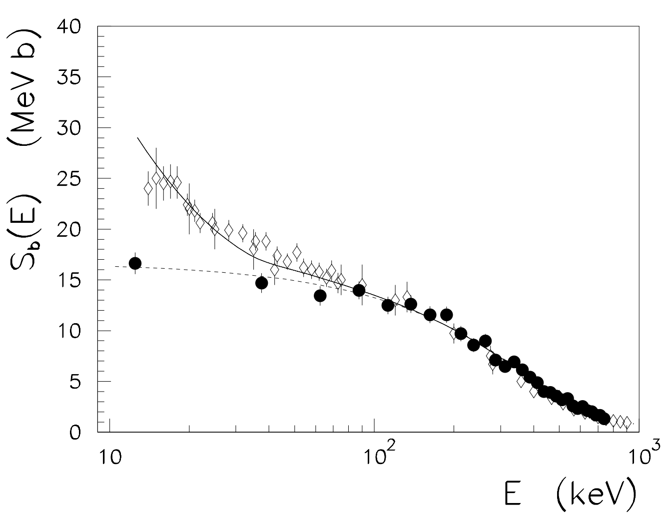

FIG. 4.—Bare $S_b(E)$ factor data (*filled circles*) for the reaction ^{6}Li(d, α)α from the THM (Spitaleri et al. 2001) and screened $S_s(E)$ factor data (*open diamonds*) from direct experiments (Engstler et al. 1992). The dashed curve is a polynomial fitted to $S_b(E)$ and the solid curve includes the effects of electron screening with $U_e = 340$ eV.

3. EXPERIMENTS

For the THM application to $d(^7\text{Li}, \alpha\alpha)n$ one needs to measure the angles and energies of both α particles, which determine—together with the ^{7}Li projectile energy—all kinematic quantities, such as the relative energy E_{12} between the two α particles. In the postcollision description, the center-of-mass interaction energy E between the ^{7}Li and p nuclei is then given by $E = E_{12} - Q$, where $Q = 17.35$ MeV is the Q value of the ^{7}Li(p, α)α reaction.

New data were obtained at the Laboratori Nazionali del Sud, Catania, using ^{7}Li beams with energies of $E_{^7\text{Li}} = 19.0$, 19.5, and 20.0 MeV (with a typical current of 10 pnA) provided by the 15 MV tandem accelerator. The targets consisted of deuterated polyethylene films (typical thickness = 250 μg cm^{-2}). The reaction products were observed using a set of six position sensitive silicon detectors (PSDs) placed on each side of the beam axis with a distance of 40 and 75 cm from the target. The PSDs were placed at mean laboratory angles $\theta_{\text{lab}} = +45°$ ($\pm 2°$), $+34°$ ($\pm 5°$), 23° ($\pm 5°$), $-45°$ ($\pm 2°$), $-55°$ ($\pm 5°$) and $-65°$ ($\pm 5°$), where the number in parentheses indicates the angular range covered. The angular resolution was typically 0°.1 for each PSD, which is needed for an improved analysis of the THM. Coincidence events were recorded between all pairs of detectors placed on opposite sides of the beam axis.

In an earlier experiment (Calvi et al. 1997) at $E_{^7\text{Li}} = 20$ MeV ionization chambers were placed as transparent detectors in front of the PSDs at $\theta_{\text{lab}} = +45°(\pm 2°)$ and $-45°(\pm 2°)$. However, due to the high Q-value of the ^{7}Li(p, α)α reaction, the condition of α-α coincidences alone already provided sufficiently clean spectra, even without particle identification. Thus, all subsequent experiments used only PSDs, which allowed in turn for an improved energy resolution.

The new and earlier data were analyzed using the same procedure with the above improved DWBA method in relating the three-body cross section into the relevant ^{7}Li(p, α)α cross section. A previous analysis of the earlier data (Aliotta et al. 2000) used the simpler PWIA formalism.

4. DATA ANALYSIS

A crucial aspect of the data analysis is to assure the dominant contribution of the quasi-free process in the three-body reaction. Thus, one must select from the overall α-α coincidence events those corresponding to the quasi-free condition. This selection was performed (Spitaleri et al. 1999) by the requirement of low momenta of the spectator neutron in the laboratory system, $p_n \leq 40$ MeV c^{-1}, which amplifies the presence of the quasi-free process and in turn suppresses the contributions of other mechanisms, such as sequential processes leading to the same final state in the three-body reaction. In this respect, the angle settings of the PSDs were chosen such that the energy region of the quasi-free process was sufficiently separated from peaks involving e.g., the excitation of the 16.6 and 16.9 MeV states in the intermediate nucleus ^{8}Be. These states are located below the ^{7}Li(p, α)α threshold and thus cannot be excited in the quasi-free process. However, they can originate from a standard sequential reaction with a three-body final channel and can generate an undesirable background for the quasi-free process. With the restriction of the spectator momentum to low values, the contribution of these states is strongly

TABLE 1

FIT PARAMETERS OF THE POLYNOMIAL EXPRESSION $S(E) = S(0) + S_1 \times E + S_2 \times E^2$ APPLIED TO THE THM DATA[a]

	CALVI ET AL. 1997[b]	PRESENT WORK[c]			
S(E) PARAMETERS	$E_{^7Li} = 20$ MeV	$E_{^7Li} = 19$ MeV	$E_{^7Li} = 19.5$ MeV[c]	$E_{^7Li} = 20$ MeV[c]	MEAN VALUE
S(0) (keV barns)	53 ± 13	57 ± 8	55 ± 6	55 ± 3	55 ± 3
S_1 (barns)	210	214	214	209	210
S_2 (keV^{-1} barns)......	-386	-422	-356	-304	-310

[a] E in units of MeV.
[b] Coincidences between one detector pair.
[c] Weighted average of coincidences between nine detector pairs.

reduced. This is demonstrated in Figure 1, where coincidence α-α spectra as functions of the relative energy are shown with and without the restriction to low neutron momenta. Even with this condition counts due to sequential processes are clearly present in the low energy part of the spectrum. However, such events were not taken into account by the THM analysis that was performed only on events corresponding to positive interaction energy ($E_{12} \geq 17.35$ MeV).

5. DETERMINATION OF $\sigma_b(E)$ FOR ^{7}Li(p, α)α

The observed energy dependence of the three-body differential cross section with the condition $p_n \leq 40$ MeV c^{-1}, determined at each ^{7}Li energy and each detection pair, was converted into the two-body cross section $\sigma_b(E)$—or equivalently $S_b(E)$—using the improved DWBA procedures described in § 2, equation (14). The deduced $S_b(E)$ values (weighted average of all results), normalized to the direct data in the energy range $E = 200$–400 keV, are shown in Figure 2. The figure also shows a theoretical curve, which was obtained by Aliotta et al. (2000) from a fit with a superposition of direct and resonance contributions to the direct experimental data for $E \geq 100$ keV. It is in excellent agreement with the data. (For a comparison of the THM and the direct data in the energy range $E = 1$–7 MeV, see Zadro et al. 1989.) The apparent discrepancy between the predictions and the previous THM data at ultralow energies reported by Aliotta et al. (2000) is due mainly to the approximate parametrization of the penetration factor $\boldsymbol{T}_1$ adopted there.

Using the polynomial expression

$$S(E) = S(0) + S_1 \times E + S_2 \times E^2$$

to fit the individual $S_b(E)$ data sets at different $E_{^7Li}$ energies, results were obtained (Table 1) consistent within statistical error. Thus, the application of the THM leads to stable and reliable information on $S_b(E)$ with $S_b(0) = 55 \pm 3$ keV barns, where the quoted error is only the statistical one. The present data set also suffer from a systematic error of $\sim 10\%$ arising from the normalization procedure of the indirect data to the direct ones. The uncertainty on the energy is mainly due to the relative energy resolution of the experiment, which is not constant but depends on kinematics. Typical values are as small as 20 keV, thanks to the well known "magnifying glass" effect (Baur, Bertulani, & Rebel 1986). Previous work (Engstler et al. 1992) estimated $S_b(0) \approx 58$ keV barns and the same value was suggested by a recent R-matrix fit (Barker 2000).

From a comparison of the $S_b(E)$ energy dependence from the application of the THM with the direct $S_s(E)$ data (Fig. 3) one can deduce a model-independent value of the screening potential energy, $U_e = 330 \pm 40$ eV (Fig. 3, *solid curve*). This is significantly higher than the adiabatic limit $U_{ad} = 175$ eV. A similar result for the electron screening potential has been obtained recently for the ^{6}Li(d, α)α reaction using the THM (see Fig. 4): $U_e = 340 \pm 50$ eV (Spitaleri et al. 2001). The results confirm a significant discrepancy between the screening potentials U_e and U_{ad}, which is presently not understood. On the other hand it is in agreement with the hypothesis (Engstler et al. 1992) of an isotopic independence of the U_e values.

6. CONCLUSIONS

Improvements in experiment and theory have provided an increased credibility of the THM in the determination of the bare astrophysical $S_b(E)$ factor for charged-particle–induced reactions down to ultralow energies. They lead to $S_b(0) = 55 \pm 3$ keV barns consistent with fits to direct measurements and the NACRE compilation (Angulo et al. 1999) and, thus, to essentially unchanged astrophysical conclusions.

In order to elucidate further the applicability and usefulness of the THM, an experiment is in progress to study ^{7}Li(p, α)α with the THM reaction ^{7}Li(^{3}He, $\alpha\alpha$)d. Similarly, the ^{3}He(d, p)^{4}He reaction (§ 1) will be investigated via the THM reaction ^{6}Li(^{3}He, αp)^{4}He.

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