

On the investigation of resonances above and below the threshold in nuclear reactions of astrophysical interest using the Trojan Horse Method.

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Abstract. The occurrence of resonances in reactions of astrophysical interest might significantly enhance the astrophysical factor with respect to the direct reaction contribution, divert nucleosynthesis path and change the energy production, with significant impact on astrophysics. Moreover, the determination of resonance parameters, that is, energy, spin-parity and partial widths, allows one to perform nuclear structure studies leading, for instance, to determine the cluster structure of the state under investigation. However, nuclear reactions in stars take place at energies well below ~ 1 MeV owing to the typical temperatures characterising these environments. Therefore, the Coulomb barrier exponentially suppressing the cross section and the electron screening effect, due to the shielding of nuclear charges by atomic electrons, make it very difficult to provide accurate astrophysical factors. The THM is an indirect method allowing to overcome such difficulties. It makes use of quasi-free reactions with three particles in the exit channel, $a + A \rightarrow c + C + s$, to deduce the cross section of the reaction of astrophysical interest, $a + x \rightarrow c + C$, under the hypothesis that A shows a strong $x + s$ cluster structure, right at astrophysical energies. By using a generalised R-matrix approach, the resonance parameters can be deduced from THM data allowing one to perform a full spectroscopic study of low-energy and sub-threshold resonances. In this work, we will discuss two examples of reactions of astrophysical interest, whose cross sections show a resonant behaviour: the $^{19}\text{F}(p, \alpha)^{16}\text{O}$ cross section that displays resonances at energies above the particle emission threshold and the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction, dominated by the -3 keV sub-threshold resonance due to the 6.356 MeV level in ^{17}O .

1. Introduction

Nuclear astrophysics deals with the investigation of nuclear physics phenomena influencing astrophysical sites such as stars or the early universe. Nuclear reactions power many astrophysical processes like energy production and synthesis of the chemical elements, therefore, cross sections $\sigma(E)$ are among the main input parameters of the codes used to model these scenarios. However, in the case of many astrophysical phenomena, such as quiescent stellar burning, energies of interest are so low that for charged particles the Coulomb barrier strongly



diminishes cross sections making the measurement of such cross sections at energies of interest for astrophysics not always possible.

These energies (the so-called Gamow window [1, 2]) usually vary between few keV and few hundreds of keV, therefore cross sections can be much smaller than 1 nb, making extrapolation from high energies the only way to estimate them at the energies of interest. Extrapolation is often performed by means of the astrophysical factor [1, 2]:

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

where $\eta = Z_1 Z_2 e^2 / \hbar v$ is the Sommerfeld parameter, Z_1 and Z_2 the atomic number of interacting nuclei and v their relative velocity. $S(E)$ is a smoothly varying function of energy because the gross of the Coulomb barrier penetration factor is compensated for owing to the presence of the inverse of the Gamow factor. In the case of resonant reactions, proceeding through an excited state of the intermediate compound nucleus, significant deviations from the smooth behaviour might be expected.

Therefore, extrapolation might introduce systematic errors since the occurrence of unknown or unpredicted resonances might dramatically enhance the astrophysical $S(E)$ -factor, strongly influencing astrophysical models. These considerations apply to both low-energy resonances and sub-threshold resonances as they may produce sizeable modifications of the S -factor due to, for instance, interference with another resonance, especially in the case of broad resonances [1, 2].

Great improvements in the experimental approach used to measure vanishingly small cross sections have been devised, making it possible to extend the measurements of some reactions down to astrophysical energies (see, for instance, [3]). However, approaching interaction energies comparable with the electron binding energies in atoms the presence of atomic electrons cannot be neglected. Electron clouds determine an enhancement of the astrophysical factor related to the shielding of the nuclear charges by the surrounding negatively-charged electrons and not to nuclear interaction (see [1, 2] for a general discussion and [4, 5] for two examples). Even in those few cases when the reaction cross sections were measured inside the Gamow window, electron screening prevented the access to the bare-nucleus cross section, making extrapolation unavoidable.

Indeed, electron screening behaves differently in the laboratory and in astrophysical environments, as in the former projectile and target are in the form of ions and atoms or molecules, while in stars, for instance, matter is in the form of plasma. However, our present understanding of the electron screening effect is rather imperfect as experimental values often exceed theoretical upper limits [6], thus potential systematic errors might be introduced in the evaluation of the bare-nucleus $S(E)$ and, as a consequence, in its extrapolation to low energies.

Because of the problems affecting direct measurements at astrophysical energies, indirect techniques have been developed to bypass them and attain the astrophysical factor at low energies. For instance, the Trojan Horse Method (THM) [7] is a valid technique to get information on the $S(E)$ at astrophysical energies for reactions having charged particles and neutrons in the exit channel, with no Coulomb and centrifugal barrier suppression neither electron screening. In the case of radiative capture reactions, the Asymptotic Normalization Coefficient (ANC) [8] approach has allowed to obtain the direct-capture zero-energy $S(E)$ factor with very high accuracy. The Coulomb dissociation (CD) was also developed [9] to determine the low-energy cross section of charged-particle induced reactions having a γ -ray in the exit channel.

2. The THM for resonant reactions

The THM has been established as a tool to investigate low-energy nuclear reactions by C. Spitaleri [10, 11] following the pioneering work of G. Baur [12]. In more recent years, A.M. Mukhamedzhanov has strongly improved the theoretical formalism [13, 14], introducing

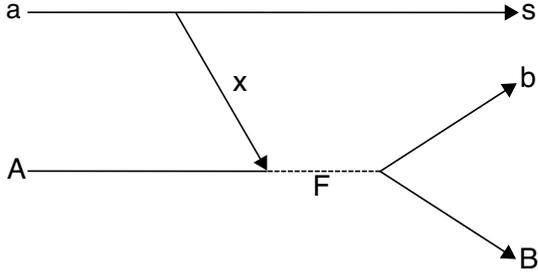


Figure 1. Sketch of the participant (x)-spectator (s) mechanism leading to the population of excited states of the intermediate nucleus F , necessary condition for the application of the THM formalism.

the modified R-matrix approach, allowing for the analysis of multi-resonance reactions with unprecedented accuracy (see, for instance, [15, 16, 17, 18]). These upgrade in the theory has impacted astrophysics, as it is shown in the case of evolved star nucleosynthesis and presolar grains [19, 20]. In the THM general framework, the cross section of the $A(x, b)B$ reaction is obtained through the $A(a, bB)s$ reaction performed at high energies (several tens of MeV), so it is not hindered by the Coulomb and centrifugal barriers of the target-projectile system and no electron screening enhancement alters the trend of the cross section. However, particle x is virtual so the $A(x, b)B$ THM is half-off-energy-shell (HOES) and cannot be simply compared with the corresponding direct (on-energy-shell, OES) cross section [13]. The modified R-matrix approach has been introduced to extract the astrophysical S-factor of interest from the quasi-free (QF) reaction yield, accounting for HOES effects.

Adopting the plane wave approximation (PWA) in the prior form and neglecting, for sake of simplicity, spins of the involved particles, the amplitude of the process $a + A \rightarrow b + B + s$ (Fig.1) takes the form:

$$M^{PWA(prior)}(P, \mathbf{k}_{aA}) = \left\langle \chi_{sF}^{(0)} \Psi_{bB}^{(-)} | V_{xA} | \varphi_a \varphi_A \chi_{aA}^{(0)} \right\rangle, \quad (2)$$

where $P = (\mathbf{k}_{sF}, \mathbf{k}_{bB})$ is the six-dimensional momentum describing the three-body system s , b and B . $\chi_{aA}^{(0)} = \exp(i\mathbf{k}_{aA} \cdot \mathbf{r}_{aA})$, $\chi_{sF}^{(0)} = \exp(i\mathbf{k}_{sF} \cdot \mathbf{r}_{sF})$, \mathbf{r}_{ij} and \mathbf{k}_{ij} are the relative coordinate and relative momentum of i and j nuclei, $\Psi_{bB}^{(-)}$ is the wave function of the fragments b and B in the exit channel, $F = b+B$, V_{xA} is the interaction potential of x and the target nucleus A , φ_a and φ_A are the bound state wave function of nuclei a and A , respectively. If we assume that resonant reaction mechanism is dominant in the explored energy region, taking into account spins of the interacting particles, considering only the s-wave bound state $a = s + x$ and neglecting the internal degrees of freedom of the transferred particle x , we get the prior PWA amplitude of the THM cross section in the form [14, 7]:

$$\begin{aligned} M^{PWA(prior)}(P, \mathbf{k}_{aA}) &= (2\pi)^2 \sqrt{\frac{1}{\mu_{bB} k_{bB}}} \varphi_a(\mathbf{p}_{sx}) \\ &\times \sum_{J_F M_F j' l' m_j' m_l' m_l' M_x} i^{l+l'} \langle j m_j l m_l | J_F M_F \rangle \langle j' m_j' l' m_l' | J_F M_F \rangle \\ &\quad \times \langle J_x M_x J_A M_A | j' m_j' \rangle \langle J_s M_s J_x M_x | J_a M_a \rangle \\ &\quad \times \exp[-i\delta_{bB}^{hs}] Y_{l m_l}(-\hat{\mathbf{k}}_{bB}) \\ &\quad \times \sum_{\nu\tau=1}^N [\Gamma_{\nu b B j l J_F}]^{1/2} [\mathbf{A}^{-1}]_{\nu\tau} Y_{l' m'}^*(\hat{\mathbf{p}}_{xA}) \\ &\quad \times \sqrt{\frac{R_{xA}}{\mu_{xA}}} [\Gamma_{\nu x A l' j' J_F}(E_{xA})]^{1/2} P_{l'}^{-1/2}(k_{xA}, R_{xA}) \\ &\quad \times (j_{l'}(p_{xA} R_{xA}) [(B_{xA l'}(k_{xA}, R_{xA}) - 1) - D_{xA l'}(p_{xA}, R_{xA})] \\ &\quad + 2Z_x Z_A e^2 \mu_{xA} \int_{R_{xA}}^{\infty} dr_{xA} \frac{O_{l'}(k_{xA}, r_{xA})}{O_{l'}(k_{xA}, R_{xA})} j_{l'}(p_{xA} r_{xA})) . \end{aligned} \quad (3)$$

Here, \mathbf{p}_{ij} is the $i - j$ relative momentum in the case of off-energy-shell particles, thus $E_{ij} \neq p_{ij}^2/2\mu_{ij}$ (while \mathbf{k}_{ij} is calculated assuming the particles on-shell), δ_{bBl}^{hs} is the solid sphere scattering phase shift, R_{xA} the $x + A$ channel radius,

$$B_{xA}l'(k_{xA}, R_{xA}) = R_{xA} \frac{\partial O_{l'}(k_{xA}, R_{xA})}{\partial r_{xA}} \Big|_{r_{xA}=R_{xA}} \quad (4)$$

is the logarithmic derivative as in the R-matrix method,

$$O_{l'}(k_{xA}, R_{xA}) = \sqrt{\frac{k_{xA}R_{xA}}{P_{l'}(k_{xA}, R_{xA})}} \exp \left[-i\delta_{xA}^{hs} \right] \quad (5)$$

is the outgoing spherical wave, $P_{l'}(k_{xA}, R_{xA})$ the l' -wave penetrability factor,

$$D_{xA}l'(p_{xA}, R_{xA}) = R_{xA} \frac{\partial j_{l'}(p_{xA}, R_{xA})}{\partial r_{xA}} \Big|_{r_{xA}=R_{xA}} \quad (6)$$

the logarithmic derivative and $j_{l'}(p_{xA}, R_{xA})$ the spherical Bessel function, N the number of the levels included. This is a generalization of the R-matrix approach because we consider reactions with three particles in the exit channel, where the TH-nucleus a in the initial states carries the transferred particle x , which is off-energy-shell.

Eq.3 has four important consequences:

- $\mathbf{A}_{\nu\tau}$ is the same level matrix as in the conventional R-matrix theory [21]. Therefore, it depends on the entry and exit channels reduced width amplitudes γ , energy levels and energy shifts. All of them can be extracted by fitting the experimental THM cross section and then can be used to deduce the $A(x, b)B$ astrophysical factor. In this way, we have an exact parameterisation of the astrophysical factor with no need of extrapolation. HOES effects can affect the phases determining interference and the relative heights of the resonances, but the reduced widths γ , containing the nuclear structure effects, appear in the same way in THM and direct data.
- The presence of the factor $P_{l'}^{-1/2}(k_{xA}, R_{xA})$ eliminates the penetration of Coulomb and centrifugal barriers in the $x + A$ channel, which is the entry channel of the reaction of astrophysical importance. The compensation of this penetrability factor is the main advantage of the THM, since it allows one to measure the astrophysical factor of the reaction of astrophysical interest down to zero energy. Moreover, resonances that can be populated with large l only are not suppressed by the centrifugal barrier, thus they can be observed even in those cases they are very weak in direct measurements, making the THM a powerful spectroscopic tool.
- The use of the PWA implies that normalisation is accomplished by extending the indirect measurement to an energy region where directly measured data are available and scaling the deduced γ -widths to match the values in the literature. This is because in PWA absolute values are greatly overestimated. However, Eq.3 can be modified as the $a - A$ and the $s - F$ interactions can be treated within the more advanced distorted waves (DWBA) and the CDCC formalisms [14]. This is a very important point as it opens the possibility to make unnecessary normalisation to direct data, at present a major drawback of THM. This is especially important in the investigation of reactions induced by radioactive ion beams, where direct data might be absent or scarce.
- In a THM experiments three particles are emitted in the exit channel. From the measurement of the energies and the angles of emissions of two out of three emitted

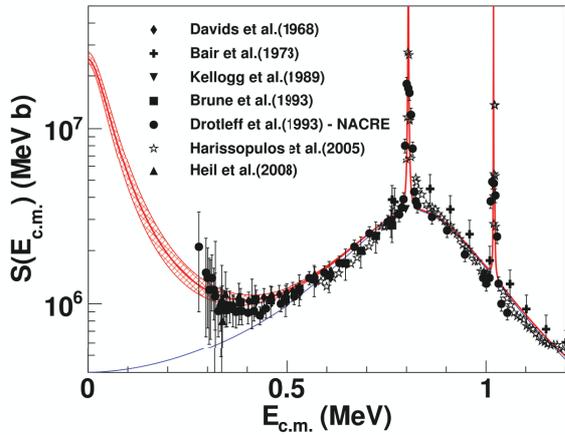


Figure 2. R-matrix calculated $S(E)$ -factor of the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction (red band) from Ref. [23]. The R-matrix $S(E)$ -factor not including the threshold resonance is displayed by the blue line. Black symbols are used for direct data normalized as in [36]. Different marks are used for each data set, as specified in the inset. See Ref. [23] for more details.

particles, all the kinematic variables can be calculated. The most important parameter for astrophysical applications is the $x - A$ relative energy. Following [7], under the non essential hypothesis that the nucleus a undergoing breakup is at rest in the laboratory system (similar formula can be found for breakup of the projectile), the $x - A$ relative energy can be written as:

$$E_{x-A} = \frac{m_x}{m_x + m_A} E_A - \frac{p_s^2}{2\mu_{sF}} + \frac{\mathbf{p}_s \cdot \mathbf{p}_A}{m_x + m_A} - \varepsilon_{sx}, \quad (7)$$

where m_i , \mathbf{p}_i and E_i are the mass, momentum and energy of the i -th particle, μ_{sF} the $s - F$ reduced mass and ε_{sx} the $x - s$ binding energy. Part of the projectile energy is spent to break the impinging nucleus a and thanks to the $x - s$ inter cluster motion, astrophysical energies can be achieved in the $x - A$ channel of the TH reaction using beam energies of few tens of MeV, bypassing Coulomb barrier. Moreover, since high beam energies are used, the electron screening enhancement does not occur. Furthermore, negative E_{x-A} energies can be explored by choosing a suitable combination of beam energy, spectator momentum and target nucleus a . From the THM measurement of sub-threshold states, yielding the reduced widths γ , the ANC can be deduced, making it possible to connect the two indirect approaches [22, 23, 24].

In the past years, a number of studies have been accomplished to validate the THM approach, making the method very robust. For instance, the effect of momentum distribution variations on the deduced astrophysical factors was considered. The use of realistic distributions in the place of the simple Hulthén function in momentum space (in the case of deuteron) has been investigated in [25], while the influence of the experimental momentum distribution in the place of the theoretical one was evaluated in [26]. In both cases, changes comparable or lower than the statistical uncertainty were retrieved, since events in a momentum region of particle s below about 40 MeV/c are usually considered in the data analysis [27]. The use of DWBA in the place of the PWA was tested in [28], showing again significant changes in the deduced astrophysical factors. The method has been extended to measurements of neutron induced reactions [29] and to reactions involving radioactive nuclei [30, 31].

3. The sub-threshold resonance case: the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction

In a class of evolved stars, those belonging to the so-called asymptotic giant branch (AGB), very peculiar conditions are present making it possible to synthesise nuclei heavier than iron

through a succession of slow neutron captures, followed by decay of the formed unstable nuclei. This path to heavy nuclei is named s-process and it is responsible of the production of $\sim 50\%$ of nuclei with $A \gtrsim 56$. In AGB stars, protons from the outer layers are mixed downward and quickly captured by carbon nuclei, eventually leading to the formation of a ^{13}C pocket [33]. Then, ^{13}C nuclei emits neutrons through the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction, at temperatures varying between 0.8×10^8 K and 1×10^8 K [34], which can be captured by seed nuclei to build up heavier nuclei, later transported to the stellar surface.

At 0.9×10^8 K, the energy range where the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction is most effective, the Gamow window [2], is $\sim 140 - 230$ keV. In such region, its direct measurement is very challenging because of the Coulomb barrier, exponentially suppressing the cross section, and the interplay between a threshold resonance determined by the population of the 6.356 MeV level in ^{17}O and atomic electron screening. Indeed, at ~ 300 keV the cross section of the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction is already as low as $\sim 10^{-10}$ b, thus direct measurements stopped at ~ 280 keV [35]. Moreover, direct measurements show contradicting results also at MeV energies, owing to uncertainties in the absolute normalisation. Therefore, extrapolation is necessary at present to assess the astrophysical factor of the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ reaction at astrophysical energy, but large errors might be introduced owing to the systematic uncertainty affecting high energy data. Therefore, [36] employed a broad data set including renormalised $^{13}\text{C}(\alpha, n)^{16}\text{O}$ astrophysical factors to perform an extensive R-matrix fit; a 140 keV astrophysical factor $S(140 \text{ keV}) = 2.2_{-0.8}^{+1.1} \times 10^6$ MeVb was then obtained. We consider 140 keV as this is the smallest energy of interest for astrophysics, the hardest to reach by conventional approaches and the most subject to systematic effects in the case of extrapolation. The renormalised data-set is shown in Fig.2 as black symbols.

The THM S-factor is shown instead as a red band in Fig.2. Thanks to our approach, a very accurate result is obtained with no need of extrapolation down to zero energy, making it possible to supply a very accurate reaction rate for astrophysical modelling. This is possible since we could directly access the 6.356 MeV level in ^{17}O , sitting at -3 keV in the $^{13}\text{C} - \alpha$ relative energy spectrum. The recommended THM S-factor at $E_{c.m.} = 140$ keV is $3.2 \pm 0.5 \times 10^6$ MeVb, about 45% larger than the value provided by [36]. The most important result is a significant reduction of the uncertainty on the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ S-factor at the Gamow peak, which decreased from about 43% to about 16%. Recently, two works [37, 38] have determined new interesting spectroscopic information about the 6.356 MeV. The consequences on the $^{13}\text{C}(\alpha, n)^{16}\text{O}$ S-factor are presently under investigation [39].

4. Resonances above the threshold: the $^{19}\text{F}(p, \alpha_0)^{16}\text{O}$ reaction

Since the s-process is quite complicated, involving large nuclear physics input as well as peculiar astrophysical conditions, the possibility to constrain such boundary conditions is very appealing. Fluorine might represent a strong constraint of stellar internal structure since its abundance is very sensitive to the physical conditions in the inner layers of AGB stars [40]. This entails the understanding of fluorine nucleosynthesis, which is incomplete to date. A possible explanation is the poor accuracy of our present picture of the fluorine destruction due to extra-mixing processes [40], where fluorine is exposed to protons at temperatures $\lesssim 4 \times 10^7$ K.

The $^{19}\text{F}(p, \alpha)^{16}\text{O}$ reaction is the main destruction channel of fluorine in this scenario. However, only one set of direct data is available so far at the energies where fluorine burning is most effective ($E_{c.m.} \leq 300$ keV, the Gamow energy [2]), still with quite large errors [41]. Furthermore, only the α_0 channel, corresponding to the emission of α -particles off ^{20}Ne compound system leaving ^{16}O in its ground state, has been investigated, being considered the larger contributor to the total cross section [41]. Before, only extrapolations were available [42] below about 500 keV, showing a non resonant behaviour, sharply contradicting the trend of the astrophysical factor at higher energies. This very simple recommended extrapolation to astrophysical energies has triggered the reassessment of the nuclear reaction rates involved in

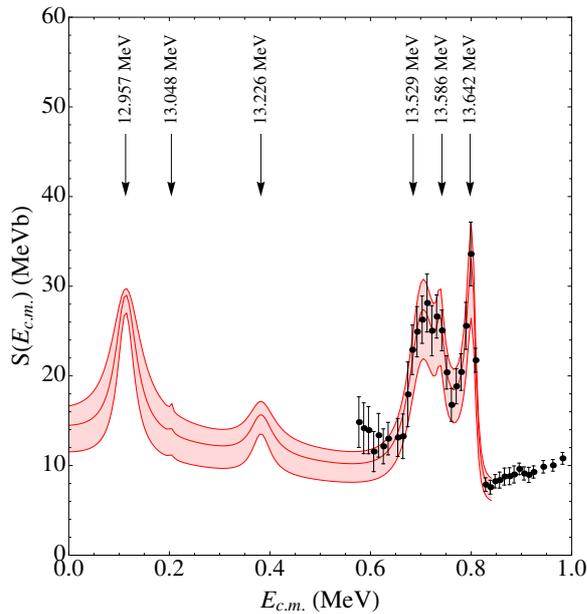


Figure 3. THM astrophysical factor of the $^{19}\text{F}(p, \alpha_0)^{16}\text{O}$ reaction, normalized to the data from [45] above 600 keV, as shown in [44]. The middle red line marks the recommended S-factor, while the upper and lower red lines stand for upper and lower limits set by combined statistical, normalization and energy shift error. The solid symbols represent the direct astrophysical factor in [45]. Finally, the arrows mark the ^{20}Ne states contributing to the S(E)-factor.

fluorine production and destruction, in particular by using the THM in its version modified to deal with resonant reactions [43, 44].

To this purpose, the QF $^{2}\text{H}(^{19}\text{F}, \alpha)^{16}\text{O}n$ reaction at 50 MeV beam energy was measured by means of a ^{19}F beam impinging onto deuterated polyethylene targets, thus using deuterons to transfer protons and induce the $^{19}\text{F}(p, \alpha)^{16}\text{O}$ QF reaction. More details on the experiment are given in [43]. Here we underscore that by using Eq.3 we derived the reduced widths of a number of resonances and, in particular, of a 113 keV peak sitting right inside the Gamow window, which was later confirmed by [41] and might have important consequences for astrophysics. The $p - ^{19}\text{F}$ relative energy spectrum spanned an energy interval from 0 to about 1 MeV, making it possible to normalise the THM astrophysical factor to the existing direct data. In the original work [43], THM data were normalised to a weighed average of direct data in the energy window 0.6 – 0.8 MeV [42]. Later, new direct data were made available in the normalisation energy region [45], leading to a reanalysis of the THM data [44].

Fig. 3 shows the S(E)-factor calculated with the resonance parameters from the fitting of THM data below 600 keV. The middle red curve marks the S(E)-factor computed using the parameters from the best fit, while the red band arises from the uncertainties on the resonance parameters, due to the combined statistical and systematic error. An average error of 20% is obtained. At present, the main source of uncertainty is due to the non resonant contribution to the astrophysical factor, since the one given in [42] is based on a very simple calculation. New direct measurements are of utmost importance to have a more realistic non-resonant contribution at low energies.

A new THM experiment has been recently performed to improve the energy resolution affecting THM data, which is a second important source of indetermination on the THM S-factor. Indeed, owing to the occurrence of many resonances below 600 keV, suitable energy resolution is at order to disentangle the occurring resonances. The poor energy resolution in the original experiment [43] prevented us to correctly identify resonances around 200 keV, where a small peak is apparent in Fig.3. Direct data [41] showed a peculiar behaviour at the lowest energies, demonstrating the occurrence of a broad 2^+ state at 251 keV, which was misidentified in the pioneering work [43] owing to the interplay between the poor energy resolution and its width (162 keV). Therefore, a direct comparison between the THM data [43, 44] and those in

Ref.[41] is presently impossible. This last work has triggered the new improved THM experiment mentioned above, aimed at passing the resolution of the pioneering work [43]. Preliminary results suggest an agreement between direct and indirect data and the manuscript is in preparation [46].

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