# EXPERIMENTAL STUDIES OF NUCLEAR REACTIONS OF ASTROPHYSICAL RELEVANCE BY MEANS OF THE TROJAN HORSE METHOD APPLIED TO RESONANT REACTIONS*

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Nuclear astrophysics aims at bringing together experimental and theoretical data necessary to understand, model and predict astrophysical phenomena, from energy production in the Sun to more exotic scenarios, such as supernovae. Nuclear physics plays a major role as fusion reactions are responsible for the energy generation as well as for the synthesis of elements. In this context, resonant reactions have a particular interest since resonances might significantly change the nucleosynthesis path and the energy production rate, modifying stellar evolution, including the last evolutionary stages. Such experimental studies are extremely difficult, since nuclear reactions in most astrophysical environments occur at energies well below the Coulomb barrier among charged particles, as large as a few MeV. This is why indirect methods have been introduced in recent years. Among them, the Trojan Horse Method has proven very effective for reactions induced by charged particles and neutrons. In this work, we will discuss two recent indirect measurements of reactions of astrophysical interest, having a resonant behaviour: the $^{19}\mathrm{F}(p,\alpha)^{16}\mathrm{O}$ reaction that displays resonances at energies above the particle emission threshold and the $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$, dominated by a near-threshold resonance due to the 6.356 MeV $^{17}\mathrm{O}$ level.

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## 1. On the need for indirect methods in nuclear astrophysics

Nuclear reactions play a major role in astrophysics as they drive energy generation and production of chemical elements in many astrophysical sites such as stars or the early universe. Therefore, reaction cross sections $\sigma(E)$

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are among the most important input parameters of the models used to study these phenomena. Typical temperatures range from $\sim 10^6$ K in the case of quiescent stellar burning to $\sim 10^9$ K for explosive scenarios. However, since the Boltzmann constant equals $8.6\times10^{-5}$ eV/K, typical center-of-mass energies at which these reactions take place are $\leq 1$ MeV. Therefore, energies of astrophysical interest are usually so low that, for charged particles, the Coulomb barrier strongly hinders fusion.

Extrapolation is then often the only way to obtain the cross section at the energies of astrophysical interest (the so-called Gamow window [1, 2]), in the case of reactions among charged particles. Yet, the rapid exponential decrease of the cross section makes extrapolation very sensitive to the behaviour of the cross section close to the extrapolation region. To put extrapolation on a sounder basis, the astrophysical factor is often used [1, 2]

$$S(E) = \sigma(E)\,E\,\exp\left(2\pi\eta\right)\,, \tag{1}$$

where $\eta = Z_1 Z_2 e^2/\hbar v$ is the Sommerfeld parameter, $Z_1$ and $Z_2$ the atomic numbers of the interacting nuclei and $v$ their relative velocity. At odds with the cross section, the astrophysical factor is a slowly changing function of energy. Therefore, the use of $S(E)$ allows us to perform more accurate extrapolations, unless resonances appear in the cross section.

Therefore, unknown or unpredicted resonances due to excited states of the intermediate compound nuclei might introduce systematic errors in the extrapolation procedure, possibly enhancing the $S(E)$-factor and significantly influencing astrophysical models. These apply both to low-energy resonances and to broad states laying right below the particle decay threshold, whose high-energy tails may introduce a pronounced increase of $S(E)$. Moreover, the effect of interference between resonances is an additional source of uncertainty.

The great uncertainties in extrapolation and the importance of low-energy cross sections have led to the introduction of new facilities and methods to directly measure such cross sections at the energies of astrophysical interest. A very important improvement has been provided by the construction of underground laboratories (see, for instance, [3]). In this way, vanishingly small cross sections have been accessed down to astrophysical energies, for low-$Z$ nuclei. In such cases, a new challenge became evident. Indeed, at low energies, of the same order of the binding energies of electrons in atoms, their presence cannot be disregarded. At large distances of closest approach, the electron cloud shields the nuclear charges, since target nuclei are usually in the form of neutral molecules or atoms, while the projectiles are not always fully stripped of their electrons. The measured astrophysical factor is then exponentially enhanced with respect to the case of bare nuclei, by a factor $f(E) = \exp\left(\pi\eta\frac{U_e}{E}\right)$, where $U_e$ is the so-called electron screening potential (see [1, 2] for a general discussion and [4–6] for two examples). In principle, by calculating $U_e$ using atomic physics theory and dividing the measured $S(E)$ by the enhancement factor, the bare-nuclei astrophysical factor might be deduced, which is the parameter of interest in nuclear astrophysics. Nonetheless, our present understanding of the electron screening enhancement is far from exhaustive as experimental values of $U_e$ often exceed the theoretical upper limits [7], thus exponentially large systematic errors might be introduced into the extraction of $S(E)$. To sum up, even in those few cases where $S(E)$ is measured down to astrophysical energies, the electron screening prevents access to the bare-nuclei astrophysical factor, making extrapolation unavoidable.

Indirect approaches have then been developed attempting to bypass problems affecting direct measurements at low energies. For instance, the Trojan Horse Method (THM) [8, 9] is a well-known technique to measure $S(E)$ at astrophysical energies in the case of reactions involving charged particles and neutrons in the exit channel, with no Coulomb and centrifugal barrier suppression or electron screening effects. In the case of radiative capture reactions, the Asymptotic Normalization Coefficient (ANC) [10] method has allowed us to determine the direct-capture astrophysical factor with very high accuracy.

## 2. The Trojan Horse Method extension to resonant reactions

The THM was initially proposed by Baur [11] in 1986, and has undergone many upgrades along the years, especially to make it suitable for astrophysical applications [12]. An extensive review of the approach has been recently published [8]. In a simplified way, the astrophysical factor of the $A(x,b)B$ reaction, typically a reaction of astrophysical importance, is obtained by studying the $A(a,bB)s$ process performed at high energies (several tens of MeV), so neither Coulomb nor centrifugal barriers suppress the cross section and electron screening does not affect the $A$–$x$ interaction. To this purpose, the reaction mechanism sketched in Fig. 1 has to be selected, namely, the participant $(x)$–spectator $(s)$ mechanism leading to the population of excited states of the intermediate nucleus $F$, later decaying into $b+B$. Then, astrophysical energies are reached thanks to the energy shift introduced by the break up of $a$ and to the $x$–$s$ intercluster motion. Many efforts have been devoted to validate the THM approach. Among others, the influence of the choice of the $x$–$s$ momentum distributions entering the calculations has been investigated in [13], as well as the effect of distortions due to the low-momentum transfer [14], the independence of the deduced $A(x,b)B$ $S$-factor from the TH nucleus $a$ [15], and the use of the distorted wave Born approximation (DWBA) in the place of the plane wave impulse approximation (PWIA) [16]. In all the mentioned cases, deviations comparable or lower than the statistical error were extracted, especially because events in an $s$-momentum window $\leq 40$ MeV$/c$ are considered, making THM very robust. As a result, THM measurements have significantly impacted astrophysics, for instance, in the case of asymptotic giant branch star nucleosynthesis and meteorite composition [17–19] and in the case of Big Bang Nucleosynthesis [20]. Moreover, THM has been recently applied to reactions involving radioactive nuclei [21–23] and to neutron induced reactions [24].

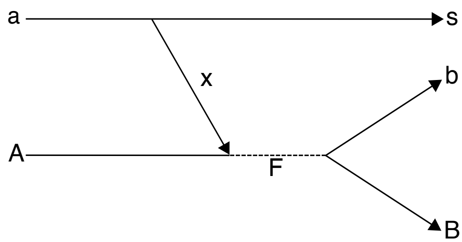


Fig. 1. Diagram illustrating the population of excited states of the intermediate system $F$ via a QF mechanism, proceeding through the direct break up of the TH nucleus $a$ and the capture of the participant cluster $x$ by the target nucleus $A$.

In the original theory, based on the PWIA, a very simple relation between the cross section for the process depicted in Fig. 1 and the one of astrophysical interest was found (see [9] for instance). However, the $A(x,b)B$ $S$-factor is deduced in arbitrary units, and a normalisation constant is necessary for each wave contributing to it. Clearly, in the multi-resonant case this is a severe drawback, resulting in a very large total uncertainty. This is why extensive theoretical [10, 25] and experimental [26, 27] work was performed, to introduce an improved formalism called the modified R-matrix approach. Indeed, the participant cluster $x$ is virtual so the sub-reaction $A(x,b)B$ is half-off-energy-shell (HOES) and corrections are necessary to compare the THM cross section with the corresponding direct cross section, which is on-energy-shell (OES) [25].

In detail, in the plane wave approximation and assuming that the reaction yield is dominated by resonances, the plane wave amplitude of the process $a + A \to b + B + s$ (Fig. 1) is given by [8, 10]:

$$
\begin{aligned}
&M^{\mathrm{PWA(prior)}}(P,\boldsymbol{k}_{aA}) = (2\pi)^2\sqrt{\frac{1}{\mu_{bB}k_{bB}}}\varphi_a(\boldsymbol{p}_{sx})\\
&\times \sum_{J_F M_F j' l l' m_{j'} m_l m_{l'} M_x} i^{l+l'} \langle j m_j l m_l | J_F M_F\rangle \langle j' m_{j'} l' m_{l'} | J_F M_F\rangle\\
&\times \langle J_x M_x J_A M_A | j' m_{j'}\rangle \langle J_s M_s J_x M_x | J_a M_a\rangle\\
&\times \exp\left[-i\delta^{hs}_{bB\,l}\right] Y_{l m_l}(-\hat{\boldsymbol{k}}_{bB})
\end{aligned}
$$

$$\times \sum_{\nu\tau=1}^{N} \left[\Gamma_{\nu bBjlJ_F}\right]^{1/2} \left[\boldsymbol{A}^{-1}\right]_{\nu\tau} Y^{*}_{l'm'}(\hat{\boldsymbol{p}}_{xA})$$
$$\times \sqrt{\frac{R_{xA}}{\mu_{xA}}} \left[\Gamma_{\nu xAl'j'J_F}(E_{xA})\right]^{1/2} P_{l'}^{-1/2}(k_{xA}, R_{xA})$$
$$\times \Big( j_{l'}(p_{xA}R_{xA}) \left[(B_{xA\,l'}(k_{xA}, R_{xA}) - 1) - D_{xA\,l'}(p_{xA}, R_{xA})\right]$$
$$+2Z_xZ_Ae^2\mu_{xA} \int_{R_{xA}}^{\infty} \mathrm{d}r_{xA} \frac{O_{l'}(k_{xA}, r_{xA})}{O_{l'}(k_{xA}, R_{xA})} j_{l'}(p_{xA}r_{xA}) \Big) . \qquad (2)$$

Here, $F = b + B$, $P = (\boldsymbol{k}_{sF}, \boldsymbol{k}_{bB})$ is the six-dimensional momentum describing the three-body system $s$, $b$ and $B$, $\mu_{ij}$ is the $i - j$ reduced mass, $r_{ij}$ is the $i - j$ relative distance, $\boldsymbol{p}_{ij}$ 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 $\boldsymbol{k}_{ij}$ is calculated assuming the particles on-shell), $\delta^{hs}_{bB\,l}$ is the solid sphere scattering phase shift, $R_{xA}$ the $x + A$ channel radius, $B_{xA\,l'}(k_{xA}, R_{xA})$ of the outgoing spherical wave $O_{l'}(k_{xA}, R_{xA})$, $P_{l'}(k_{xA}, R_{xA})$ the $l'$-wave penetrability factor, $D_{xA\,l'}(p_{xA}, R_{xA})$ is the logarithmic derivative of the spherical Bessel function, $N$ the number of the levels included.

Equation (2) has some important consequences:

— $\boldsymbol{A}_{\nu\tau}$ is the same level matrix as in the conventional R-matrix theory [28]. This is a crucial point as this implies that the same reduced widths $\gamma$ for the entrance and exit channels and the same resonance energies appear in the THM cross section and in the OES one. Therefore, reduced widths $\gamma$ and resonance energies, which altogether univocally fix the OES $S$-factor, can be deduced from the analysis of the measured cross section of the $a + A \to b + B + s$ process using Eq. (2). In this way, no extrapolation to astrophysical energies would be necessary, HOES effects are fully taken into account and, finally, normalisation (mandatory since we are using the plane wave approximation) can be extended to many resonances, notwithstanding that they are populated in different waves.

— Equation (2) justifies the THM unique feature, the possibility to reach zero energy with no Coulomb suppression. Indeed, it explicitly contains the factor $P_{l'}^{-1/2}(k_{xA}, R_{xA})$ fully compensating for the Coulomb and centrifugal barrier penetration factor in the $x + A$ entrance channel. Moreover, Eq. (2) shows that THM can be fruitfully used as a spectroscopic tool, making it possible to observe very weak resonances, populated with large $l$ waves, since they are not suppressed as they are in direct measurements. This can be useful, for instance, to investigate cluster states laying close to the particle emission thresholds. It is worth noting that in the case of radioactive beams, the use of indirect methods only allows us to cover energies as low as 0 MeV, as it has been shown [21–23], using experimental data.

— As remarked so far, the plane wave approximation makes normalisation to existing data mandatory, even if the modified R-matrix approach limits the need to a single constant over the whole measured energy range. However, Eq. (2) can be further upgraded within the distorted-wave Born approximation (DWBA) or continuum-discretised coupled channel (CDCC) formalisms [10]. Such upgrades might prove very useful in all cases where no direct data are available, or data with large uncertainties are solely present, especially in the case of reactions between radioactive nuclei and neutrons.

The use of reactions with three particles in the exit channel has many advantages. First, it is possible to calculate the $Q$-value of the reactions occurring in the target from the energies and angles of emissions of two out of three particles. Therefore, on the one hand, we can choose which particles are more convenient to detect so, for instance, in the case of reactions having a neutron in the exit channel we can detect the other two ejectiles. On the other hand, we can select to a great accuracy the reaction channel of interest, reducing essentially to zero the contribution of background. Once the reaction has been identified, the $Q$-value can be fixed at the theoretical one and one of the measured quantities can be dropped from calculations. By removing the variable (energy or angle) most affected by uncertainty, a significant improvement in the energy resolution can be achieved, that is of paramount importance in the case of multi-resonant reactions. Finally, Eq. (2) treats in the same way both positive and negative values of the $x-A$ relative energy $E_{x-A}$. Following [8], under the non-essential hypothesis that $a$ is at rest in the laboratory frame, the $x-A$ relative energy can be written as

$$E_{x-A} = \frac{m_x}{m_x + m_A} E_A - \frac{p_s^2}{2\mu_{sF}} + \frac{\boldsymbol{p}_s \cdot \boldsymbol{p}_A}{m_x + m_A} - \varepsilon_{sx} \,, \tag{3}$$

where $m_i$, $\boldsymbol{p}_i$ and $E_i$ are the mass, momentum and energy of the $i^{\text{th}}$ particle, $\mu_{sF}$ the $s-F$ reduced mass and $\varepsilon_{sx}$ the $x$–$s$ binding energy. As already discussed, $a$ binding energy and $x$–$s$ inter-cluster motion make it possible to reach astrophysical energies in the $A(x,b)B$ sub-reaction even if the THM reaction is induced at energies of many MeV per nucleon. Moreover, $E_{x-A}$ can assume negative values for a proper choice of the quantities appearing in Eq. (3). It has been proven that, in the case of resonances lying at negative $E_{x-A}$ values, Eq. (2) allows one to determine the ANCs of these states [29–32]. Therefore, Eq. (2) establishes a link between the THM and the ANC.

## 3. The $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$ reaction

The origin of nuclei heavier than iron is a very active field in nuclear astrophysics. More than $\sim 50\%$ of nuclei with $A \geq 56$ are produced by the s-process [33]. This consists in a succession of slow neutron captures, namely, the decay of the formed unstable nuclei occurs before a subsequent neutron capture. Therefore, a neutron source has to be available, releasing neutrons at a slow rate. Such a neutron source has been identified in a group of evolved stars, belonging to the asymptotic giant branch (AGB). In AGB stars, protons from the outer layers are mixed downward when the H-burning shell is extinguished, reaching to the intershell region, rich in carbon and helium. There, protons are captured by carbon nuclei, leading to the formation of a $^{13}$C enriched region (the so-called $^{13}$C pocket [34]). Then, $^{13}$C is destroyed through the $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$ reaction releasing neutrons. Typical temperatures range between $0.8\times10^{8}$ K and $1\times10^{8}$ K [33], corresponding to $^{13}$C–$\alpha$ center-of-mass energies of $\sim$140–230 keV (Gamow window [2]).

However, direct measurements stop at $\sim 280$ keV [35], since the cross section of the $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$ reaction at $\sim 300$ keV is already as low as $\sim 10^{-10}$ b. Extrapolation is therefore necessary to deduce the astrophysical factor at the energies of astrophysical interest. Unfortunately, many reasons make extrapolation very uncertain. Apart from the Coulomb barrier, exponentially damping the cross section, the trend of the cross section approaching the $^{13}$C–$\alpha$ threshold is complicated by the presence of a broad resonance formerly believed to lie at $-3$ keV (sub-threshold resonance), now known to be situated at 4.7 keV [36], and by the electron screening effect, already introducing a correction of about 20% at the lowest-energy data point [35]. Moreover, direct measurements suffer from large systematic errors also at larger energies, well within the MeV scale, determining a scatter in the absolute normalisation that is apparent in Fig. 2.

Indirect methods especially focused on the determination of the ANC of the near-threshold resonance, due to the 6.356 MeV state of $^{17}$O, using transfer reactions [37–40]. In contrast with direct measurements, these independent determinations were able to supply a very accurate and consistent values of the Coulomb modified ANCs, whose weighted average is $3.9\pm0.5\,\mathrm{fm}^{-1}$. On the other hand, the THM applied to the $^{13}\mathrm{C}(^{6}\mathrm{Li},\alpha^{16}\mathrm{O})n$ process [29, 30, 32] could deduce not only the ANC for this crucial state, but also the energy trend of the astrophysical factor of the $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$ reaction down to zero energy. Thanks to our approach, a very accurate result was obtained with no need of extrapolation, making it possible to derive also a very accurate reaction rate for astrophysical modelling. This has been possible since we could directly observe and fit the 6.356 MeV level in $^{17}$O, in an essentially background-free experiment.

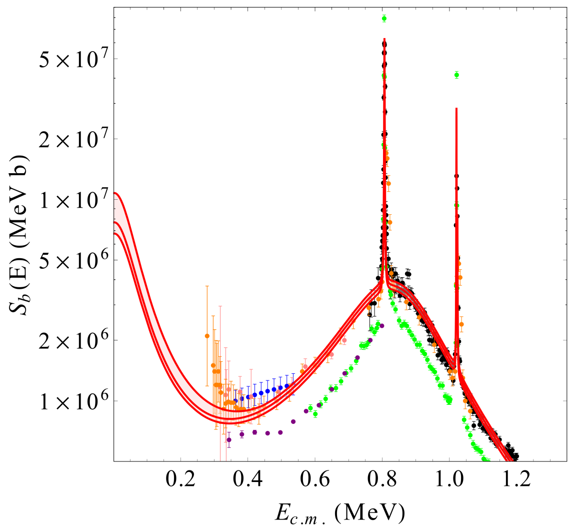


Fig. 2. (Colour on-line) R-matrix astrophysical factor (central/red curve) calculated adopting the resonance deduced from the THM cross section using Eq. (2). Normalisation and statistical uncertainties (including correlations) are given by the lower and upper red lines. Direct data, reported without any scaling, are shown as symbols [35, 41–45].

In early THM measurements [29, 30], THM data were normalised to the astrophysical factor recommended by Ref. [41]. In this work, new accurate data were used to renormalise existing $^{13}\mathrm{C}(\alpha,n)^{16}\mathrm{O}$ astrophysical factors, to perform an extensive R-matrix fit of available data in the literature. However, using this normalisation THM data supplied an inconsistently large ANC for the mentioned 6.356 MeV level in $^{17}$O, about a factor of two larger than previous ANC measurements. In a recent work [32], we decided to change the paradigm and normalise the THM $S$-factor to the 6.356 MeV state, to pick up the correct absolute value of direct data. Therefore, we concluded that only Refs. [35, 43] supply a direct data set compatible with the THM $S$-factor and the ANC of the threshold level measured in many experiments [37–40]. With this new normalisation, we were able to deduce the $S$-factor in Fig. 2 (shown as a grey/red band) and a consistent ANC of $3.6 \pm 0.7\,\mathrm{fm}^{-1}$. The corresponding THM $S$-factor at $E_{\mathrm{cm}} = 140$ keV is $1.80^{+0.50}_{-0.17} \times 10^6$ MeV b, to be compared with the astrophysical factor by Ref. [41] $S(140\,\mathrm{keV}) = 2.2^{+1.1}_{-0.8} \times 10^6$ MeV b. Interesting astrophysical consequences of the modified astrophysical were also found [32].

## 4. The $^{19}\mathrm{F}(p,\alpha)^{16}\mathrm{O}$ reaction

Most of the stellar observables originate in the outer layers. For instance, spectroscopic studies supply information on the composition, temperature and other physical conditions characterising the stellar surface. Conversely, it is very complicated to gather data about the internal structure, composition and physical properties of the inner areas, where the radiated energy is produced and nucleosynthesis takes place. Therefore, probes of stellar interiors have been proposed over the years, to constrain models of stellar internal structure and nucleosynthesis. Fluorine abundance can be used to restrict the stellar parameter space for the s-process, since its abundance is very sensitive to the physical conditions in the inner layers of AGB stars [19, 46]. Then, exhaustive knowledge of the destruction cross sections is mandatory to use the fluorine abundance as a sensitive probe.

Fluorine is mainly destroyed in $\alpha$-induced reactions in the intershell region of AGB stars [47], and at the bottom of the convective envelope of AGB stars [48] by $(p,\alpha)$ reactions at temperatures $\leq 4\times10^7$ K. In the case of the $^{19}\mathrm{F}(p,\alpha)^{16}\mathrm{O}$ reaction, only one set of data is available at the energies $E_{\mathrm{cm}} \leq 300$ keV, where fluorine burning is most effective at the bottom of the convective envelope, affected by large uncertainties [49], and limited to the $\alpha_0$ channel, corresponding to the emission of $\alpha$-particles by $^{20}\mathrm{Ne}$ leaving $^{16}\mathrm{O}$ in its ground state. This is presently considered the larger contributor to the total cross section [49], even if more data on the other channels are strongly needed. Before 2015, only the linear extrapolation by Ref. [50] was present below 500 keV, calling for indirect measurements to assess the trend of the astrophysical factor inside the Gamow window. Two THM measurements were performed at INFN — Laboratori Nazionali del Sud, Catania, Italy [51] and at INFN — Laboratori Nazionali di Legnaro, Legnaro, Italy [52].

In both experiments, the QF $^2\mathrm{H}(^{19}\mathrm{F},\alpha^{16}\mathrm{O})n$ reaction was measured using deuterons to transfer protons and induce the $^{19}\mathrm{F}(p,\alpha)^{16}\mathrm{O}$ reaction. Equation (2) was then used to correct the THM cross section for HOES and energy resolution effects. In the earlier experiment [51], energy resolution was not good enough to separate all the $^{20}\mathrm{Ne}$ states contributing to the $S$-factor. This work, however, for the first time pointed out the occurrence of a 113 keV peak sitting right inside the Gamow window, thus bearing a major astrophysical prominence. After this seminal work, two direct measurements [49, 53] were performed, allowing for a more accurate normalisation of the THM $S$-factor (though not influencing much the astrophysical factor at the Gamow energy [46]) and to extend the reach of the direct data almost to the upper tail of the Gamow peak [49]. The direct $S(E)$ showed a distinctive trend at the lowest energies, indicating the presence of a broad $2^+$ state at 251 keV, which was misidentified in Ref. [51] owing to the interplay between the poor energy resolution and its width (162 keV). The recent work in Ref. [52] allowed the problem of the ambiguous determination of such a resonance to be solved thanks to the improved energy resolution. Therefore, very good agreement was found between the indirect measurement (grey/red band in Fig. 3) and direct data from Refs. [49, 53] (solid symbols). Figure 3 shows the $S(E)$-factor calculated with the resonance parameters from the fitting of THM data below 600 keV using Eq. (2). The middle (red) curve marks the recommended $S(E)$-factor while the upper and lower limits are obtained from the combined statistical and systematic errors.

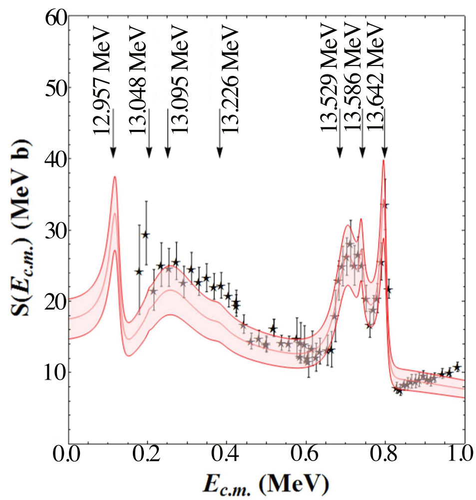


Fig. 3. (Colour on-line) Comparison of the THM $S$-factor with direct data. The THM result is shown as a grey/red band, highlighting the upper and lower limits allowed for by the normalisation and statistical errors. The black stars represent the experimental data from Refs. [49, 53]. Finally, the arrows indicate the $^{20}$Ne levels contributing to the $S(E)$-factor.

## 5. Concluding remarks

Resonant reactions play a key role in nuclear physics and nuclear astrophysics. In nuclear physics, they allow us to perform spectroscopic studies and, especially close to the decay thresholds, to investigate cluster structures. In nuclear astrophysics, they can divert the nucleosynthesis flow and change the energy production rate and elemental yield. For these reasons, they have been subject of many studies over the years. We have discussed an indirect method for exploring the near-threshold region of relevance for nuclear astrophysics studies. In particular, we have examined the THM measurements of the $^{13}$C$(\alpha, n)^{16}$O and $^{19}$F$(p, \alpha)^{16}$O reactions, of great importance for the synthesis of the elements heavier than iron and to constrain the scenario where their production is taking place. Unprecedented information about the trend of the astrophysical factor inside the Gamow window is reported, possibly having important astrophysical consequences.

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