Asymptotic normalization coefficient from the $^{12}$C($^7$Li,$^6$He)$^{13}$N reaction and the astrophysical $^{12}$C($p,\gamma$)$^{13}$N reaction rate

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Abstract

Angular distribution of the $^{12}$C($^7$Li,$^6$He)$^{13}$N reaction at $E(^7$Li) = 44.0 MeV was measured at the HI-13 tandem accelerator of Beijing, China. Asymptotic normalization coefficient (ANC) of $^{13}$N $\rightarrow$ $^{12}$C + $p$ was derived to be $1.64 \pm 0.11$ fm$^{-1/2}$ through distorted wave Born approximation (DWBA) analysis. The ANC was then used to deduce the astrophysical $S(E)$ factors and reaction rates for direct capture in $^{12}$C($p,\gamma$)$^{13}$N at energies of astrophysical relevance.

Key words: NUCLEAR REACTIONS $^{12}$C($^7$Li,$^6$He)$^{13}$N, $E(^7$Li) = 44.0 MeV, measured $\sigma(\theta)$, DWBA analysis, deduced asymptotic normalization coefficient, $^{12}$C($p,\gamma$)$^{13}$N $E$ = low, deduced astrophysical $S$-factor.

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1 Introduction

In stellar evolution, the $^{12}$C($p,\gamma$)$^{13}$N reaction plays a key role for the following reasons. Firstly, it is the first reaction in the Carbon-Nitrogen-Oxygen (CNO) cycle which dominates the energy production in stars with masses heavier than 1.5 M$_\odot$ [1,2]. Secondly, the $^{12}$C($p,\gamma$)$^{13}$N($\beta^+$)$^{12}$C reactions can enhance the $^{13}$C abundance [3], and thus influence the $^{12}$C/$^{13}$C ratio which is thought to be an important measure for stellar evolution and nucleosynthesis [4]. Thirdly, the supply of $^{13}$C by the $^{12}$C($p,\gamma$)$^{13}$N($\beta^+$)$^{13}$C reactions is also important for the $^{13}$C($\alpha,n$)$^{16}$O neutron source in the asymptotic giant branch (AGB) stars.
recent calculation with parametric one-zone nucleosynthesis showed that the more $^{13}\text{C}$ supply is needed at the end of the CNO cycle in solar-metallicity stars [5]. In view of the above mentioned significance, it is highly desired to carefully investigate the $^{12}\text{C}(p, \gamma)^{13}\text{N}$ cross section at energies below 1.0 MeV for the astrophysical interest.

The $^{12}\text{C}(p, \gamma)^{13}\text{N}$ reaction has been studied over a wide energy range down to about $E_{c.m.} = 70$ keV [6,7,8,9,10,11,12,13] since 1950. However, there exists obvious discrepancy in the experimental data at the lower energies. In the even lower energy range, no experimental data is available so far, and the astrophysical $S(E)$ factors can only be derived by the extrapolation which results in uncertainty inevitably. In the energy range of $E_{c.m.} \leq 200$ keV, the $^{12}\text{C}(p, \gamma)^{13}\text{N}$ reaction is dominated by the tail of the s-wave capture into the broad $1/2^+$ resonance at $E_r = 421$ keV. Although the contribution from direct capture is believed to be much smaller than that from the resonance tail, the interference between the two processes can lead to a considerable variation of $S(E)$ factors, since both of them proceed via s-wave and then decay by $E1$ transitions. As a result, the $S(E)$ factors either increase by the constructive interference or decrease by the destructive one. Thus the reliable experimental data on the direct capture are needed to derive the $S(E)$ factors, particularly at energies of $E_{c.m.} \leq 200$ keV. A practicable scheme to deduce the $S(E)$ factors for the direct capture in $^{12}\text{C}(p, \gamma)^{13}\text{N}$ is combining the asymptotic normalization coefficient (ANC) of $^{13}\text{N} \rightarrow ^{12}\text{C} + p$ and R-matrix approach [14,15], the ANC can be deduced from the angular distribution of one proton transfer reactions. The $(^7\text{Li}, ^6\text{He})$ reactions are considered to be a valuable spectroscopic tool because the shapes of their angular distributions can be well reproduced by the distorted wave Born approximation (DWBA) [16]. Thus, the $^{12}\text{C}(^7\text{Li}, ^6\text{He})^{13}\text{N}$ reaction is used to extract the nuclear ANC of $^{13}\text{N} \rightarrow ^{12}\text{C} + p$ in the present work.

The $^{12}\text{C}(^7\text{Li}, ^6\text{He})^{13}\text{N}$ angular distribution was measured at $E(^7\text{Li}) = 44.0$ MeV. The spectroscopic factor and ANC were derived based on DWBA analysis, and then used to calculate the astrophysical $S(E)$ factors and rates of $^{12}\text{C}(p, \gamma)^{13}\text{N}$ direct capture reaction at energies of astrophysical interest with the R-matrix approach. We have also computed the contribution from the resonant capture and the interference effect between resonant and direct captures.

2 Measurement of the $^{12}\text{C}(^7\text{Li}, ^6\text{He})^{13}\text{N}$ angular distribution

The experiment was performed at the HI-13 tandem accelerator of Beijing, China. A carbon target in thickness of 39.0 $\mu\text{g/cm}^2$ was bombarded with the 44.0 MeV $^7\text{Li}$ beam in intensity of about 100 pnA. The $^6\text{He}$ ions from the
Fig. 1. Angular distribution of the $^{12}$C($^{7}$Li, $^{6}$He)$^{13}$N reaction at $E(^{7}$Li) = 44.0 MeV.

$^{12}$C($^{7}$Li, $^{6}$He)$^{13}$N reaction were analyzed with the Q3D magnetic spectrograph. In order to attain a good angular resolution, the solid angle for the reaction products was set to be 0.23 msr. A two dimensional position sensitive silicon detector (PSSD) was placed at the focal plane of the spectrograph, which assured the full detection of $^{6}$He ions emitted within the solid angle. During the measurement, a Faraday cup placed at 0° in the reaction chamber was utilized to record the beam current which served as the normalization standard in determining the absolute cross section, while two independent silicon detectors (SSDs) at ±30° were used for both monitoring the beam balance and relative normalizing the measured cross sections. The normalization was checked for several angles from 23° to 40° with the elastic scattering data of 36.0 MeV $^{7}$Li on $^{12}$C [17].

The $^{12}$C($^{7}$Li, $^{6}$He)$^{13}$N differential cross sections were measured in the angular range of $7° \leq \theta_{lab} \leq 23°$, corresponding to $11° \leq \theta_{c.m.} \leq 37°$. The measured angular distribution is shown in Fig. 1. The experimental errors are from the uncertainties of statistics, target thickness (5%) and solid angle (6%).

3 Determination of the $^{13}$N Nuclear ANC

The spins and parities of $^{12}$C and $^{13}$N ground states are $0^+$ and $1/2^-$, respectively. The $^{12}$C($^{7}$Li, $^{6}$He)$^{13}$N cross section is dominated by the $(0^+,0) \rightarrow (1/2^-,1/2)$ transition, thus only $1p_{1/2}$ orbit in $^{13}$N can be populated. If the reaction is peripheral, the differential cross section can be expressed as

$$\frac{d\sigma}{d\Omega}_{exp} = \left(\frac{C_{7Li}}{b_{7Li}}\right)^2 \left(\frac{C^{13N}_{1,1/2}}{b^{13N}_{1,1/2}}\right)^2 \left(\frac{d\sigma}{d\Omega}\right)_{DWBA},$$

(1)
Table 1
Optical potential parameters used in the DWBA calculations, where $U$, $W$ are in MeV, $r$ and $a$ in fm.

| Channel | $U_V$ | $r_R$ | $a_R$ | $W_V$ | $r_I$ | $a_I$ | $r_c$ |
|---------|-------|-------|-------|-------|-------|-------|-------|
| $^7\text{Li} + ^{12}\text{C}$ | 194.55 | 0.50  | 0.82  | 7.66  | 1.31  | 0.75  | 1.30  |
| $^6\text{He} + ^{13}\text{N}$ | 170.50 | 0.79  | 0.67  | 10.70 | 1.31  | 0.75  | 1.30  |

where $(\frac{d\sigma}{d\Omega})_{\text{exp}}$ and $(\frac{d\sigma}{d\Omega})_{\text{DWBA}}$ denote the measured and calculated differential cross sections respectively. $C_{1,1/2}$ and $C_{7\text{Li}}$ are the nuclear ANCs of $^{13}\text{N} \rightarrow ^{12}\text{C} + p$ and $^{7}\text{Li} \rightarrow ^6\text{He} + p$, $b_{1,1/2}$ and $b_{7\text{Li}}$ being the single particle ANCs of the bound state protons in $^{13}\text{N}$ and $^{7}\text{Li}$, which can be calculated with the bound state single particle wave function and Whittaker function at larger radius.

The ratio of $C/b$ is so called spectroscopic factor. The proton spectroscopic factor of $^7\text{Li}$ was derived to be $0.41 \pm 0.05$ in our previous work [18]. The nuclear ANC of $^{13}\text{N}$ ground state can then be extracted by normalizing the theoretical differential cross sections to the experimental data via Eq. (1).

The angular distribution of the $^{12}\text{C}(^7\text{Li}, ^6\text{He})^{13}\text{N}$ reaction was calculated with the DWBA code FRESCO [19]. The bound state wave function was obtained by solving the Schrödinger equation using a Woods-Saxon potential with standard geometrical parameters ($r_0 = 1.25$ fm and $a = 0.65$ fm), the potential depth was adjusted so as to reproduce the observed binding energy of the valence proton. The optical potential parameters for both the entrance and exit channels were extracted by fitting the angular distributions of elastic scattering with the code SFRESCO [19]. The $^7\text{Li} + ^{12}\text{C}$ elastic scattering was measured in the present work. Since there is no experimental data for $^6\text{He}$ elastic scattering on $^{13}\text{N}$, the angular distribution of $^6\text{Li} + ^{13}\text{C}$ elastic scattering at $E(^6\text{Li}) = 28.0$ MeV [20] was used to extract the potential parameters of the exit channel. The extracted potentials are listed in Table 1.

Generally, the spectroscopic factor or nuclear ANC is determined by fitting the theoretical calculations to the experimental data at the first peak in the angular distribution for the forward angles [21], since the experimental differential cross sections for the backward angles are very sensitive to the inelastic coupling effects and other high-order ones, which can not be well described theoretically. In the DWBA calculation, the differential cross sections at three forward angles were used to extract the ANC. The normalized angular distribution is also presented in Fig. 1. The nuclear ANC and the spectroscopic factor for $^{13}\text{N} \rightarrow ^{12}\text{C} + p$ are deduced to be $1.64 \pm 0.11$ fm$^{-1/2}$ and $0.64 \pm 0.09$, respectively. The spectroscopic factors extracted by different experiments and the theoretical values are listed in Tab. 2. The spectroscopic factor obtained in our work agrees with the theoretical ones reported in Refs. [22,23] and the experimental results given in Refs. [26,28,29,37]. The Nuclear ANC of $^{13}\text{N} \rightarrow ^{12}\text{C} + p$ from this work is in good agreement with the value of $1.65 \pm 0.20$.
Table 2
The theoretical and experimental proton spectroscopic factors in the $^{13}$N ground state.

| $S_{13N}$ | Experiments or theory | Year | Reference |
|-----------|-----------------------|------|-----------|
| 0.61      | theory                | 1967 | [22]      |
| 0.56      | theory                | 1969 | [23]      |
| 0.78 - 1.35 | $^{12}$C($d, n$) | 1970 | [24]      |
| 0.74      | $^{12}$C($d, n$)      | 1971 | [25]      |
| 0.53 ± 0.12 | $^{12}$C($d, n$) | 1972 | [26]      |
| 0.70 - 1.48 | $^{12}$C($^3$He, $d$) | 1969 | [27]      |
| 0.56 - 0.78 | $^{12}$C($^3$He, $d$) | 1976 | [28]      |
| 0.68 ± 0.12 | $^{12}$C($^3$He, $d$) | 1976 | [29]      |
| 0.81 ± 0.12 | $^{12}$C($^3$He, $d$) | 1979 | [30]      |
| 0.48 ± 0.12 | $^{12}$C($^3$He, $d$) | 1980 | [31]      |
| 1.34      | $^{12}$C($a, t$)      | 1969 | [32]      |
| 0.91      | $^{12}$C($a, t$)      | 1972 | [33]      |
| 0.72      | $^{12}$C($^7$Li, $^6$He) | 1979 | [34]      |
| 0.38 ± 0.05 | $^{12}$C($^7$Li, $^6$He) | 1986 | [35]      |
| 0.64 ± 0.09 | $^{12}$C($^7$Li, $^6$He) | 2008 | present work |
| 0.25, 0.40 | $^{12}$C($^{10}$B, $^9$Be) | 1974 | [36]      |
| 0.62      | $^{12}$C($^{14}$N, $^{13}$C) | 1975 | [37]      |
| 0.29, 0.40 | $^{12}$C($^{16}$O, $^{15}$N) | 1979 | [38]      |

fm$^{-1/2}$ extracted from the $^{12}$C($^{10}$B, $^9$Be)$^{13}$N reaction [39].

4 Astrophysical $S(E)$ factors of the $^{12}$C($p, \gamma$)$^{13}$N reaction

Following the approach used in our previous work [15], the astrophysical $S(E)$ factors of the $^{12}$C($p, \gamma$)$^{13}$N reaction were calculated by the R-matrix method. For the radiative capture reaction $B + b \rightarrow A + \gamma$, the cross section to the state of nucleus A with the spin $J_f$ can be written as [14,15]

$$\sigma_{J_f} = \sum_{J_i} \sigma_{J_iJ_f}, \quad (2)$$
\[ \sigma_{J_iJ_f} = \frac{\pi}{k^2} \frac{2J_i+1}{(2J_b+1)(2J_B+1)} \sum_{I_l} \left| U_{I_lJ_iJ_f} \right|^2, \]

where \( J_i \) denotes the total angular momentum of the colliding nuclei \( B \) and \( b \) in the initial state, \( J_b \) and \( J_B \) are the spins of nuclei \( b \) and \( B \), and \( I, k \) and \( l_i \) are their channel spin, wave number and orbital angular momentum in the initial state, respectively. \( U_{I_lJ_iJ_f} \) is the transition amplitude from the initial continuum state \(( J_i, I, l_i)\) to the final bound state \(( J_f, I)\). In the single-level, single-channel approximation, the resonant amplitude for the capture into the resonance with energy \( E_{Rn} \) and spin \( J_i \), and subsequent decay into the bound state with the spin \( J_f \) can be expressed as

\[ U_{R_{I_lJ_iJ_f}} = -ie^{i(\omega_{l_i}-\phi_{l_i})} \frac{\left[ I_{l_i}^{J_i}(E)\Gamma_{\gamma J_f}(E) \right]^{1/2}}{E - E_{Rn} + i\Gamma_{J_i}/2}. \]

Here it is assumed that the boundary parameter is equal to the shift function at resonance energy, and \( \phi_{l_i} \) is the hard-sphere phase shift in the \( l_i \)th partial wave,

\[ \phi_{l_i} = \arctan \left[ \frac{F_{l_i}(k, r_c)}{G_{l_i}(k, r_c)} \right], \]

where \( F_{l_i}^2 \) and \( G_{l_i}^2 \) are the regular and irregular Coulomb functions, \( r_c \) is the channel radius. The Coulomb phase factor \( \omega_{l_i} \) is given by

\[ \omega_{l_i} = \sum_{n=1}^{l_i} \arctan \left( \frac{\eta_i}{n} \right), \]

where \( \eta_i \) is the Sommerfeld parameter. \( \Gamma_{bI_i}^{J_i}(E) \) is the observable partial width of the resonance in the channel \( B + b \), \( \Gamma_{\gamma J_f}^{J_i}(E) \) is the observable radiative width for the decay of the given resonance into the bound state with the spin \( J_f \), and \( \Gamma_{J_i} \approx \sum_I \Gamma_{bI_i}^{J_i} \) is the observable total width of the resonance level. The energy dependence of the partial widths is determined by

\[ \Gamma_{bI_i}^{J_i}(E) = \frac{P_i(E)}{P_i(E_{R_n})} \Gamma_{bI_i}^{J_i}(E_{R_n}) \]

and

\[ \Gamma_{\gamma J_f}^{J_i}(E) = \left( \frac{E + \varepsilon_f}{E_{R_n} + \varepsilon_f} \right)^{2L+1} \Gamma_{\gamma J_f}(E_{R_n}), \]
where $\Gamma_{\beta I J}^f(E_{R_f})$ and $\Gamma_{\gamma I J}^i(E_{R_i})$ are the experimental partial and radiative widths, $\varepsilon_f$ is the proton binding energy of the bound state in nucleus $A$, and $L$ is the multipolarity of the gamma transition. The penetrability $P_l(E)$ is expressed as

$$P_l(E) = \frac{kr_c}{F_l^2(k, r_c) + G_l^2(k, r_c)}. \quad (9)$$

The nonresonant amplitude can be calculated by

$$U_{\ell_i J_i}^{\Lambda R} = -(2)^{3/2} \ell_i + L - \ell_f + 1 \epsilon(i\omega_i - \phi_i) \frac{\mu_{Bb}^{L+1/2}}{m_b^{1/2}} \left[ Z_{\ell_i} c + (-1)^{L} \frac{Z_{Bb} e}{m_b^{1/2}} \right] (k, r_c)^{L+1/2} \times \frac{(L+1)(2L+1)}{L} \frac{1}{(2L+1)!} C_{J_i, J_f} F_{\ell_i}(k, r_c) G_{\ell_i}(k, r_c) \times W_{\ell_f}(2kr_c) \sqrt{F_{\ell_i}(l_0 L 0 | l_f 0) U(L l_f J_i; l_i J_f) J'_L(l_i l_f)}, \quad (10)$$

where $k = (E + \varepsilon_f)/hc$ is the wave number of the emitted photon. $C_{J_i, J_f}$ is the nuclear ANC of $^{13}N \rightarrow ^{12}C + p$, and $l_f$ are the wave number and relative orbital angular momentum of the bound state. $W_l(2kr)$ is the Whittaker hypergeometric function with $\kappa = \sqrt{2\mu_{Bb} \varepsilon_f}$. $(l_0 L 0 | l_f 0)$ and $U(L l_f J_i; l_i J_f)$ are the Clebsch-Gordan and Racha coefficients, respectively. $J'_L(l_i l_f)$ is the integral expression defined as

$$J'_L(l_i l_f) = \frac{1}{r_c^{L+1}} \int_{r_c}^{\infty} dr' r'^{L+1} \frac{W_{\ell_f}(2kr)}{W_{\ell_f}(2kr_c)} \left[ \frac{F_{\ell_i}(k, r)}{G_{\ell_i}(k, r_c)} - \frac{G_{\ell_i}(k, r_c)}{F_{\ell_i}(k, r_c)} \right]. \quad (11)$$

The non-resonant amplitude contains the radial integral ranging only from the channel radius $r_c$ to infinity since the internal contribution is contained within the resonant part. Furthermore, the R-matrix boundary condition at the channel radius $r_c$ implies that the scattering of particles in the initial state is given by the hard sphere phase. Hence, the problems related to the interior contribution and the choice of incident channel optical parameters do not occur. Therefore, the direct capture cross section only depends on the ANC and the channel radius $r_c$. Using the experimental ANC $1.64 \pm 0.11$ fm$^{-1/2}$ from the present work, the non-resonant $^{12}C(p, \gamma)^{13}N$ cross sections vs. $E_{c.m.}$ were calculated, as shown in Fig 2. In the calculation, $r_c$ was taken to be 5.0 fm following the previous works in Ref. [15] and [40].

The astrophysical S-factor is related to the cross section by

$$S(E) = E \sigma(E) \exp(E_G/E)^{1/2} \quad (12)$$
Fig. 2. The non-resonant cross sections of the $^{12}\text{C}(p, \gamma)^{13}\text{N}$ reaction computed with the ANC derived from the present experiment.

Fig. 3. Astrophysical $S(E)$ factors as a function of $E_{\text{c.m.}}$ for the $^{12}\text{C}(p, \gamma)^{13}\text{N}$ reaction. The dashed line and the solid line are the contributions from the direct proton capture and the total $S(E)$ factors, respectively. The experimental data are taken from Refs. [6,7,8,9,12]

where the Gamow energy $E_G = 0.978Z_1^2Z_2^2\mu$ MeV, $\mu$ is the reduced mass of the system. Using the resonance parameters ($E_R$=421 keV, $\Gamma_{\text{tot}}(E_R) = 36.5$ keV, and $\Gamma_\gamma(E_R) = 0.67$ eV) from Ref. [41], the S-factors for direct and resonant captures can be then derived.

Since the incoming angular momentum ($s$-wave) and the multipolarity ($E1$) for the direct and resonant capture $\gamma$-radiation are identical, there is an in-
terference between the two capture processes. Following the Ref. [9], the total S-factor is calculated by

\[ S_{\text{tot}}(E) = S_{\text{dc}}(E) + S_{\text{res}}(E) \pm 2[S_{\text{dc}}(E)S_{\text{res}}(E)]^{1/2} \cos(\delta_r), \]  

where \( \delta_r \) is the resonance phase shift, given by

\[ \delta_r = \arctan \left( \frac{\Gamma_p(E)}{2(E - E_r)} \right). \]

The experimental results from the direct measurement of \(^{12}\text{C}(p, \gamma)^{13}\text{N} \) reaction [9] show that the interference between the resonant and direct captures is constructive below the resonance energy, and destructive above it. Based on this interference pattern, the present total S-factors are then obtained, as shown in Fig. 3. One can see that the present S-factors of \(^{12}\text{C}(p, \gamma)^{13}\text{N} \) are in good agreement with the experimental data from Ref. [42] and the references therein. Very recently, N. Burtebaev et al. [12] measured the cross sections of the \(^{12}\text{C}(p, \gamma)^{13}\text{N} \) reaction at beam energies \( E_p = 354, 390, 460, 463, 565, 750, \) and 1061 keV. They obtained \( \text{ANC}^{(13}\text{N}) = 1.72 \text{ fm}^{-1/2}, \Gamma_\gamma = 0.65 \pm 0.07 \text{ eV} \) and \( \Gamma_p = 35.0 \pm 1.0 \text{ keV} \) by fitting their experimental data with R-matrix approach. Their results are in good agreement with ours.

5 Astrophysical \(^{12}\text{C}(p, \gamma)^{13}\text{N} \) reaction rates

The astrophysical \(^{12}\text{C}(p, \gamma)^{13}\text{N} \) reaction rate is calculated with

\[ N_A\langle \sigma v \rangle = N_A \left( \frac{8}{\pi \mu} \right)^{1/2} \frac{1}{(k_B T)^{3/2}} \int_0^\infty S(E) \exp \left[-\left(\frac{E_G}{E}\right)^{1/2} - \frac{E}{k_B T}\right] dE, \]  

where \( \nu = \sqrt{2E/\mu} \) and \( N_A \) and \( k_B \) are Avogadro and Boltzmann constants respectively. The updated reaction rates are shown in Fig. 4, together with the previous ones from NACRE and CF88 compilations. The present reaction rates are larger than that of NACRE and CF88 by about 20% in the low temperature range of \( T_9 < 0.5 \), and very close to their average value in the temperature range of \( 0.5 \leq T_9 \leq 10.0 \).

The total reaction rates as a function of temperature obtained in our work are parameterized with an expression used in the astrophysical reaction rate library REACLIB [43],
Fig. 4. (Color online) The $^{12}$C$(p, \gamma)^{13}$N reaction rates from present work as well as those from NACRE and CF88 compilations.

Table 3
The fitting parameters of the $^{12}$C$(p, \gamma)^{13}$N reaction rates.

| parameters | value   | parameters | value   |
|------------|---------|------------|---------|
| $a_0$      | 1.1397E+01 | $b_0$    | 4.0200E+01 |
| $a_1$      | -2.9913E-02 | $b_1$    | 2.1886E+00 |
| $a_2$      | -1.1507E+01 | $b_2$    | -1.5188E+02 |
| $a_3$      | 7.6988E+00  | $b_3$    | 1.1973E+02 |
| $a_4$      | -3.4841E+00 | $b_4$    | -3.7889E+00 |
| $a_5$      | 3.8122E-01  | $b_5$    | 1.2086E-01 |
| $a_6$      | 4.5641E-02  | $b_6$    | -8.1517E+01 |

\[
N_A\langle \sigma v \rangle = \exp(a_0 + a_1 T_9^{-1} + a_2 T_9^{-1/3} + a_3 T_9^{1/3} + a_4 T_9 \\
+ a_5 T_9^{5/3} + a_6 \ln T_9) + \exp(b_0 + b_1 T_9^{-1} + b_2 T_9^{-1/3} \\
+ b_3 T_9^{1/3} + b_4 T_9 + b_5 T_9^{5/3} + b_6 \ln T_9).
\] (16)

The value of fit parameters $a_{0-6}$ and $b_{0-6}$ are listed in Tab. 3, and the fitting errors are less than 1% in the temperature range of $0.01 \leq T_9 \leq 10$. 

10
6 Conclusion and discussion

The $^{12}\text{C}(p, \gamma)^{13}\text{N}$ reaction is of considerably astrophysical interest. The cross sections in the energy range of astrophysical relevance is very small and difficult to be measured directly. In the energy range of $E_{\text{c.m.}} \leq 200$ keV, the $^{12}\text{C}(p, \gamma)^{13}\text{N}$ cross section depends on the tail of the s-wave capture into the broad $1/2^+$ resonance at $E_\text{r} = 421$ keV, the direct capture to the ground state and their interference. The determination of the resonant parameters for the $1/2^+$ state and the ANC of $^{13}\text{N}$ ground state is helpful in extrapolating the experimental data at high energies down to the energies of astrophysical interest (around 25 keV).

We measured the angular distribution of the $^{12}\text{C}(^7\text{Li, } ^6\text{He})^{13}\text{N}$ reaction at $E_{^7\text{Li}} = 44.0$ MeV, and deduced the nuclear ANC and spectroscopic factor for the $^{13}\text{N}$ ground state. The astrophysical S-factors and reaction rates of $^{12}\text{C}(p, \gamma)^{13}\text{N}$ are then extracted with the R-matrix approach. The S-factor at 25.0 keV is found to be $1.87 \pm 0.13$ keV·b. The result is consistent with that of $1.75 \pm 0.22$ keV·b obtained by Burteaev et al. [12], and slightly higher than the earlier values of $1.45 \pm 0.20$ keV·b, $1.54 \pm 0.08$ keV·b and $1.33 \pm 0.15$ keV·b reported in Refs. [9,10,11]. In order to clarify the deviation, the further study of the resonant parameters of the $1/2^+$ state is needed.

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