Photon-induced Nucleosynthesis: Current Problems and Experimental Approaches

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Photon-induced reactions play a key role in the nucleosynthesis of rare neutron-deficient p-nuclei. The paper focuses on (γ,α), (γ,p), and (γ,n) reactions which define the corresponding p-process path. The relation between stellar reaction rates and laboratory cross sections is analyzed for photon-induced reactions and their inverse capture reactions to evaluate various experimental approaches. An improved version $S_{C}(E)$ of the astrophysical S-factor is suggested which is based on the Coulomb wave functions. $S_{C}(E)$ avoids the apparent energy dependence which is otherwise obtained for capture reactions on heavy nuclei. It is found that a special type of synchrotron radiation available at SPring-8 that mimics stellar blackbody radiation at billions of Kelvin is a promising tool for future experiments. By using the blackbody synchrotron radiation, sufficient event rates for (γ,α) and (γ,p) reactions in the p-process path can be expected. These experiments will provide data to improve the nuclear parameters involved in the statistical model and thus reduce the uncertainties of nucleosynthesis calculations.

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I. INTRODUCTION

Nucleosynthesis of heavy nuclei proceeds mainly via neutron capture reactions and subsequent $\beta$-decays in the so-called s-process and r-process [1]. About 99% of the heavy nuclei above iron are synthesized in these processes. However, there are 35 nuclei on the neutron-deficient side of the chart of nuclides, the so-called p-nuclei, which cannot be made by neutron capture. The dominant reactions for the synthesis of p-nuclei are photodissociations induced by photons of a hot stellar environment with temperatures around $1.8 \leq T_{9} \leq 3.3$ (where $T_{9}$ is the temperature in $10^{9}$ K) or thermal energies $155 \text{ keV} \leq kT \leq 285 \text{ keV}$. The required seed nuclei must have been synthesized earlier in the s- and/or r-process [2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12]. The oxygen- and neon-rich layers of type II supernovae are good candidates for the astrophysical site of the p-process (note that “p” should be interpreted here as photodissociation, not as proton capture); however, there is no firm conclusion on the astrophysical site yet [11].

The full reaction network for the nucleosynthesis calculations in the p-process contains by far more than 1000 nuclei and more than 10000 reactions, involving many unstable nuclei. It is impossible to measure all reaction rates in the laboratory. Usually, the reaction rates are calculated within the framework of the statistical model [17]. It is the aim of this paper to discuss how the reliability of these calculations can be tested using new experimental data which can be obtained with synchrotron-based photon sources at SPring-8 [18] or at other facilities like S-DALINAC at TU Darmstadt [19], ELBE at Forschungszentrum Rossendorf [20], or HiγS at TUNL, Duke University [21].

The dominant reactions in the p-process path are (γ,n) and (γ,α) reactions leading to a path of the p-process which is located about 10 mass units from stability on the neutron-deficient side around $A \approx 200$ and close to stability around $A \approx 140$ [4, 11]. At lower masses around $A \approx 100$, the importance of (γ,p) and (p,γ) reactions increases. Recently it has been pointed out, that a fast expansion of high-entropy, proton-rich matter may produce a significant amount of $^{92,96}$Mo and $^{96,98}$Ru [22] which have been underestimated in most previous calculations; but unfortunately no astrophysical site could be assigned firmly to this scenario [22]. In addition, neutrino-induced nucleosynthesis has been suggested recently for the synthesis of light p-nuclei [23, 24].

Whereas a lot of work has been done for (γ,n) reactions in recent years, there are still only few experimental data for (γ,p) and especially for (γ,α) reactions at all, and only very few data at astrophysically relevant energies.

The paper is organized as follows. In Sect. II the present status for (γ,α), (γ,p), and (γ,n) reactions and indirect approaches are discussed. An improved treatment of the astrophysical S-factor is suggested which avoids an apparent energy dependence of the conventional S-factor based on the Gamow factor for capture reactions on heavy nuclei. Sect. III presents a detailed comparison between photodissociation reactions and capture reactions in the laboratory and in hot stellar environments with a focus on (γ,α) and (α,γ) reactions. In Sect. IV as an example the feasibility of (γ,α) experiments at SPring-8 is analyzed, and finally conclusions are given in Sect. V.

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II. PHOTODISSOCIATION FOR THE ASTROPHYSICAL P-PROCESS

In the following section we analyze the astrophysically relevant energy for photon-induced reactions. The most effective energy $E_{\text{eff}}$ and the width $\Delta$ of the Gamow-like energy window are derived from the thermal photon spectrum and from the energy dependence of the cross sections of the respective photon-induced ($\gamma,\alpha$), ($\gamma,p$), and ($\gamma,n$) reactions. This energy dependence can be estimated from the relation between photodissociation reactions and their inverse capture reactions. In addition, a brief overview of the available experimental data in this energy region is given.

A. ($\gamma,\alpha$) Reactions

Almost no experimental data exist for ($\gamma,\alpha$) reactions in the astrophysically relevant energy window that is defined in the following way. The photon density at temperature $T$ is given by the blackbody radiation:

$$n_{\gamma}(E,T) = \left(\frac{1}{\pi}\right)^{\frac{3}{2}} \frac{1}{E^3 \exp(E/kT) - 1}$$

(2.1)

The stellar reaction rate of a ($\gamma,\alpha$) reaction is obtained by

$$\lambda^*(T) = \int_0^\infty c \, n_{\gamma}(E,T) \, \sigma^*_\gamma(E) \, dE$$

(2.2)

with the cross section $\sigma^*_\gamma$ under stellar conditions, i.e., including thermal excitations of the target nucleus (see also Eqs. (1) and (2) in [17]). For completeness, we define here also the reaction rate in the laboratory $\lambda_{\text{lab}}$ with the target in the ground state (see also Sect. IV)

$$\lambda^*_{\text{lab}}(T) = \int_0^\infty c \, n_{\gamma}(E,T) \, \sigma^*_{\gamma,\alpha}(E) \, dE$$

(2.3)

where $\sigma^*_{\gamma,\alpha}$ is the ground state cross section. The rate $\lambda^*_{\text{lab}}$ can be measured in the laboratory using photon spectra with a (quasi-) thermal energy distribution which can be obtained by a superposition of bremsstrahlung spectra in a narrow energy window [13, 25] or by high-energy synchrotron radiation of GeV electron beams in a much broader energy window [26]. For simplicity of the following notation, we use the symbol $\lambda$ also for the reaction rate of the inverse ($\alpha,\gamma$) capture reaction:

$$\lambda^*(T) = \int_0^\infty c \, n_{\gamma}(E,T) \, \sigma^*_{\alpha,\gamma}(E) \, dE$$

(2.4)

The energy dependence of the ($\gamma,\alpha$) cross section can be derived from the inverse ($\alpha,\gamma$) cross section. The ($\alpha,\gamma$) capture reaction shows a roughly exponential energy dependence at low energies because of the tunneling probability through the Coulomb barrier. The nuclear part of the energy dependence can roughly be seen in the astrophysical S-factor defined by

$$\sigma(E) = \frac{1}{E} \exp(-2\pi\eta) \, S(E)$$

(2.5)

with the Sommerfeld parameter $\eta$. The S-factor has a much weaker energy dependence for capture reactions than the cross section.

In the definition of the astrophysical S-factor only incoming s-waves are taken into account. However, the influence of the centrifugal barrier for higher partial waves with angular momenta $l > 0$ is small for reactions between heavy nuclei at energies around $5 - 10$ MeV. Thus, the energy dependence of the cross section for capture reactions, in particular, from the incoming $p$-wave, is close to s-wave capture.

In a very first approximation one may assume a constant S-factor of the inverse ($\alpha,\gamma$) capture reaction to the ground state. One finds a strict correspondence between the ($\alpha,\gamma_0$) and ($\gamma,\alpha_0$) cross sections where the $p$-wave $\alpha$ capture on even-even nuclei populates 1$^+$ states in the compound nuclei that decay directly to the ground state ($0^+$) by emitting $E1$ photons. By applying the reciprocity theorem which relates the ($\gamma,\alpha_0$) and ($\alpha,\gamma_0$) cross sections in the laboratory (see Sect. IV), one finds that the integrand of Eqs. (2.2) and (2.3) shows a maximum at

$$E_{\text{eff}}^{(\gamma,\alpha)} \approx 0.122 \times (Z_P^2/Z_T^2) A_{\text{red}} T_{\text{red}}^2)^{1/3} \text{ MeV} + Q_\alpha = E_{\text{eff}}^{(\alpha,\gamma)} + Q_\alpha$$

(2.6)

where $Z_P$, $Z_T$, and $A_{\text{red}}$ are the charge numbers of projectile and target and the reduced mass number of the inverse ($\alpha,\gamma$) reaction, and $Q_\alpha$ is the $Q$-value of the ($\alpha,\gamma$) capture reaction which corresponds to the binding energy of the $\alpha$ particle in the compound nucleus. Note that $Q_\alpha < 0$ for many $p$ nuclei. The integral of Eq. (2.6) and its maximum are shown in Fig. 10. The width of this maximum is approximately given by

$$\Delta^{(\gamma,\alpha)} \approx 0.237 \times (Z_P^2/Z_T^2) A_{\text{red}} T_{\text{red}}^2)^{1/6} \text{ MeV}$$

(2.7)

The Gamow-like window of the ($\gamma,\alpha$) reactions is approximately identical to the Gamow-window of the ($\alpha,\gamma$) capture reaction, but shifted by the binding energy $Q_\alpha$ [27, 28]. The reason for this similarity is that for both ($\gamma,\alpha$) and ($\alpha,\gamma$) reactions the most effective energy is governed by the tunneling through the Coulomb barrier in the entrance channel in ($\alpha,\gamma$) and in the exit channel in ($\gamma,\alpha$). Consequently, the position and width of the Gamow-like window for ($\gamma,\alpha$) reactions is temperature-dependent, and the width is much broader for ($\gamma,\alpha$) reactions compared to ($\gamma,n$) reactions. As an example, the position and width for the Gamow-like window of the reactions $^{148}\text{Gd}(\gamma,n)^{147}\text{Gd}$ ($S_n = 8984$ keV) and $^{148}\text{Gd}(\gamma,\alpha)^{144}\text{Sm}$ ($Q_\alpha = -3271$ keV) are listed in Table 11. This example has been chosen because the nucleus $^{148}\text{Gd}$ is a branching point in the astrophysical $p$-process which
TABLE I: Most effective energy $E_{\text{eff}}$ and width $\Delta$ of the Gamow-like window of the reactions $^{148}$Gd($\gamma$,n)$^{147}$Gd ($S_n = 8984$ keV) and $^{144}$Gd($\gamma,\alpha$)$^{141}$Sm ($Q_\alpha = -3271$ keV), derived under the assumption of a constant S-factor of the inverse ($\alpha,\gamma$) capture reaction. Note that the definitions for the widths are not fully consistent: for ($\gamma$,n) reactions $\Delta_{\text{FWHM}}$ is given [22], whereas for ($\gamma,\alpha$) reactions the usual definition for $\Delta_{1/e}$ is used [23]. Because of the asymmetry of the integrand in Eq. (2.2) there is no precise relation between $\Delta_{\text{FWHM}}$ and $\Delta_{1/e}$, but $\Delta_{\text{FWHM}} \approx 0.833 \Delta_{1/e}$ is a rough approximation for the less asymmetric ($\gamma,\alpha$) case. Additionally, the effective energy $E_{\text{eff}}^{(\alpha,\gamma)}$ for the ($\alpha,\gamma$) capture reaction is given. For the widths: $\Delta_{1/e}^{(\alpha,\gamma)} \approx \Delta_{1/e}^{(\gamma,\alpha)}$.

| $T_9$ | $kT$ (keV) | $E_{\text{eff}}^{(\gamma,\alpha)}$ (MeV) | $\Delta_{\text{FWHM}}^{(\gamma,\alpha)}$ (MeV) | $E_{\text{eff}}^{(\alpha,\gamma)}$ (MeV) | $\Delta_{1/e}^{(\alpha,\gamma)}$ (MeV) | $E_{\text{eff}}^{(\gamma,\alpha)}$ |
|-------|-----------|-----------------|-----------------|-----------------|-----------------|-----------------|
| 2.0   | 172       | 9.07            | 0.32            | 4.30            | 2.64            | 7.57            |
| 2.5   | 215       | 9.09            | 0.41            | 5.52            | 3.18            | 8.79            |
| 3.0   | 259       | 9.11            | 0.50            | 6.66            | 3.70            | 9.93            |

FIG. 1: Gamow window of selected ($\gamma,\alpha$) reactions around masses $A \approx 100, 150$, and 200: the integrand of Eq. (2.2) is shown (i) using a constant S-factor of the inverse ($\alpha,\gamma$) capture reaction (full lines) and (ii) using the theoretical ($\gamma,\alpha$) cross sections of [33] (dashed lines) which correspond to the energy-dependent S-factors of Fig. 2. The energy dependence of the S-factor leads to a shift of the Gamow window by typically a few hundred keV up to about 1500 keV. Note that the absolute scale of the data is arbitrary, though it reflects the S-factor dependence of the reaction.

The energy dependence of the astrophysical S-factor at low energies is not well determined. It cannot be taken directly from experimental data. In most cases the experimental S-factor decreases slightly with increasing energy [30, 34, 35, 36, 37, 38, 39, 40, 41]. However, the experimental data usually do not cover the full Gamow window but have been measured at somewhat higher energies. There is urgent need for improved experimental ($\alpha,\gamma$) data at low energies. It has to be pointed out that the experimental determination of ($\alpha,\gamma$) cross sections at astrophysically relevant energies is very difficult. The cross sections are very small around the effective energy $E_{\text{eff}}^{(\alpha,\gamma)}$ because of the rapidly decreasing tunneling probability below the Coulomb barrier. Nevertheless, a number of reactions has been measured in the last years at energies close or slightly above $E_{\text{eff}}^{(\alpha,\gamma)}$ including the following $p$-nuclei: $^{144}$Sm($\alpha,\gamma$)$^{148}$Gd [30], $^{112}$Sn($\alpha,\gamma$)$^{116}$Te [35, 36], $^{106}$Cd($\alpha,\gamma$)$^{110}$Sn [34], $^{96}$Ru($\alpha,\gamma$)$^{100}$Pd [37]. Further data for heavy nuclei exist for $^{76}$Ge($\alpha,\gamma$)$^{80}$Se [38], $^{106}$Cd($\alpha,\gamma$)$^{110}$Sn [39], and $^{139}$La($\alpha,\gamma$)$^{143}$Pr [40, 41], and extensive studies are in progress [42].

It has been noticed especially for the case of the $^{144}$Sm($\alpha,\gamma$)$^{148}$Gd reaction that the theoretical prediction of the capture cross section slightly below 10 MeV (which is the astrophysically relevant energy, see Table I) varied by about two orders of magnitude, mainly depending on the $\alpha$-nucleus potential [30]. Consequently, $\alpha$-nucleus potentials for heavy $p$-nuclei were studied using elastic scattering at energies around the Coulomb barrier. Data for $^{144}$Sm [42, 92Mo [44], and $^{112}$Sn [43] are available, and a reasonable description of the scattering data and a variety of further experimental data including reaction cross sections and $\alpha$-decay half-lives have been obtained using systematic folding potentials [16, 47, 48, 49, 50]. However, the low-energy behavior of the potentials is not well-defined; this still leads to considerable uncertainties for the prediction of cross sections at astrophysically relevant energies. Further experimental data are strongly required, as e.g. pointed out in [48]. Another possibility for studying the $\alpha$-nucleus potential is using the sensitivity of ($n,\alpha$) reactions, for which details can be found in [51, 52].

There are also attempts for a comprehensive investigation of alpha-induced reactions on the same nucleus. For example, for the $^{106}$Cd nucleus ($\alpha,\gamma$), ($\alpha,p$) and ($\alpha,n$) reactions have been measured and compared to statistical model calculations with different input parameter sets [34]. It is found that there is no parameter set describ-
ing all three channels simultaneously. This reflects the limited knowledge on the optical potentials in the mass and energy region relevant to the astrophysical $p$-process. Similar study is in progress for the $^{112}$Sn isotope.

As already pointed out above, theoretical predictions of $(\alpha,\gamma)$ capture cross sections show a dramatic variation in their energy dependence (see e.g. Fig. 10 of [43] for $^{112}$Sn$(\alpha,\gamma)$116Te and Fig. 10 of [48] for $^{144}$Sm$(\alpha,\gamma)^{148}$Gd). To estimate the influence of an energy-dependent S-factor on the position of the Gamow window of $(\gamma,\alpha)$ photodissociation reactions, the following procedure can be applied. In a first step the statistical model $(\gamma,\alpha)$ cross sections compiled in [54] have been converted to the corresponding laboratory $(\alpha,\gamma_0)$ capture cross section to the ground state using time-reversal symmetry. This is a good approximation because of the dominating ground state contribution in the $(\gamma,\alpha)$ reaction (see Sect. 111). It is noted that since $(\gamma,\alpha)$ reactions are dominantly induced by E1 photons, the $(\alpha,\gamma_0)$ reaction of time-reversal symmetry is characterized by $p$-wave $\alpha$ capture on even-even nuclei. However, as far as the energy dependence is concerned, one expects essentially no difference between $s$-wave and $p$-wave capture because the effect of the centrifugal potential is very small in the energy region of interest.

Three examples have been chosen for the analysis of the energy dependence of the astrophysical S-factor $S(E)$ around $A \approx 100, 150,$ and 200: $^{92}$Mo$(\alpha,\gamma)^{96}$Ru, $^{144}$Sm$(\alpha,\gamma)^{148}$Gd, and $^{186}$Os$(\alpha,\gamma)^{190}$Pt and their inverse $(\gamma,\alpha)$ photodissociation reactions. The $(\alpha,\gamma_0)$ cross sections which are derived from the theoretical $(\gamma,\alpha)$ cross sections in [54] have been translated to astrophysical S-factors which are shown in Fig. 2. The astrophysical S-factor decreases with energy; the slope is about one order of magnitude per 2 MeV within the Gamow window.

The consequences of this energy dependence for the Gamow window of the $(\gamma,\alpha)$ cross section are shown in Fig. 1. For the three selected examples the integrand of Eq. (2.3) is shown using a constant S-factor (full lines) and using the model $(\gamma,\alpha)$ cross sections of [54] (dashed lines) which correspond to the energy-dependent S-factors of Fig. 2. The temperature for all figures is $T_0 = 2.5$ which is a typical $p$-process temperature. Using the theoretical energy dependence of the $(\gamma,\alpha)$ cross section [54], one finds a shift of the Gamow window to lower energies between a few hundred keV and up to about 1500 keV for the $^{190}$Pt$(\gamma,\alpha)^{186}$Os case. Although there is a significant change in the energy dependence of the cross section, the resulting shift of the energy of the Gamow window remains limited because the energy dependence of the thermal photon density is by far dominating; e.g., the exponential factor exp ($-E/kT$) changes by about 10 orders of magnitude between 5 MeV and 10 MeV whereas the variation of the cross section is about 2 orders of magnitude when one replaces the simple assumption of a constant S-factor with a realistic energy dependence (see Fig. 2).

FIG. 2: S-factor of the $^{92}$Mo$(\alpha,\gamma)^{96}$Ru, $^{144}$Sm$(\alpha,\gamma)^{148}$Gd, and $^{186}$Os$(\alpha,\gamma)^{190}$Pt calculated from the statistical model $(\gamma,\alpha)$ photodissociation cross sections of [54]. The thin lines correspond to the standard S-factor in Eq. (2.5), the thick lines use the new improved S-factor $S_C$ in Eq. (2.10). The new S-factor $S_C$ shows a much weaker energy dependence than the standard S-factor $S(E)$. The $S_C$ curves are shown in relative units and are normalized to $S(E)$ at $E = 10$ MeV. The chosen radius is $r_0 = 1.3$ fm. For $T_0 = 2.5$ the Gamow windows in Eq. (2.6) are located at $E^{\text{eff}}_{\text{cutoff}} = 6.75$ MeV, 8.79 MeV, and 10.09 MeV for the above capture reactions, respectively. The kinks in $S(E)$ and $S_C(E)$ at 9 MeV for the $^{92}$Mo$(\alpha,\gamma)^{96}$Ru reaction and at 12 MeV for the $^{144}$Sm$(\alpha,\gamma)^{148}$Gd and $^{186}$Os$(\alpha,\gamma)^{190}$Pt reactions indicate the opening of the $(\alpha,n)$ channel; here the competition between neutron emission and $\gamma$-ray emission reduces the $(\alpha,\gamma_0)$ cross section.

B. An improved treatment of the astrophysical S-factor

The Gamow factor exp ($-2\pi\eta$) in the conventional definition of the astrophysical S-factor in Eq. (2.3) presumably approximates the Coulomb-barrier tunneling probability for s-wave particles in the absence of the centrifugal potential. This probability, $P$, is expressed by $P = |\psi(R_N)|^2 / |\psi(R_c)|^2$, where the denominator is the probability of finding the particle at the classical turning point $R_c$ and the numerator is that at the nuclear radius $R_N$. For the one-dimensional Schrödinger equation for the Coulomb potential $V_C$ one finds

$$P = \exp \left[ - \frac{2}{h} \int_{R_c}^{R_N} \sqrt{2\mu \left( V_C(r) - E \right)} \, dr \right]. \quad (2.8)$$

The WKB solution of the three-dimensional Schrödinger equation gives the same equation with an extra proportional coefficient $\sqrt{(B_C - E)/E}$ with the Coulomb barrier height $B_C$ [55]. At low energies $E \ll B_C$ (typically by a factor of the order of 100) or, equivalently, when the classical turning point is much larger than the nuclear radius $R_c \gg R_N$, $P$ can be approximated by the Gamow factor exp ($-2\pi\eta$). In particular, Eq. (2.3) becomes $P = \exp (-2\pi\eta)$ for (the unrealistic case) $R_N = 0$ (point-like nuclei). In contrast, in the $\alpha$ capture reactions of current interest, the effec-
tive energy is no longer much smaller than the Coulomb barrier height. For example, as seen in Table [3] the effective energy $E_{\text{eff}}^{(\alpha,\gamma)}$ at $T_0 = 2 - 3$ is 7 - 10 MeV for $^{144}\text{Sm}$, while the Coulomb barrier height is around 20 MeV with the radius parameter $r_0 = 1.3$ fm in $R = r_0(A_1^{1/3} + A_2^{1/3})$.

At large distances ($r \geq R_N$), where the nuclear potential is absent, the solution of the radial Schrödinger equation is given by a linear combination of the regular and irregular Coulomb functions, $F(r)$ and $G(r)$. (The assumption of the absence of the nuclear potential holds for a potential well with a radius $R_N$, but obviously not for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.) Thus, as it is well known, a better approximation for the Wood-Saxon potential or other realistic potentials.)

Accordingly, a better astrophysical $S$-factor $S_C(E)$ is defined for the present case:

$$S_C(E) = \frac{1}{E} \left( \frac{1}{F^2_0(E, R_N)} + \frac{1}{G^2_0(E, R_N)} \right) S_C(E) \quad (2.10)$$

A minor disadvantage of the new $S_C$ is in the choice of the radius parameter $r_0$ though the $r_0$ dependence of $S_C$ is moderate (see Fig. 3). We have chosen here $r_0 = 1.3$ fm.

First we apply the new $S$-factor $S_C$ to the energy dependence of the $(\alpha,\gamma)$ cross section which is derived from the calculation of $^{54}$. The new $S$-factor is shown by the thick lines in Fig. 2 in comparison with the standard $S$-factor (thin lines). One finds a much weaker energy dependence of $S_C(E)$ compared to $S(E)$.

Additionally we apply the definition of the new $S$-factor $S_C$ to the experimental data of the $^{144}\text{Sm}(\alpha,\gamma)^{148}\text{Gd}$ reaction $^{33}$ in Fig. 3. Whereas the standard $S$-factor shows a noticeable energy dependence (upper part), the experimental data are almost constant when presented using the new $S_C(E)$. The influence of the chosen radius parameter $r_0$ on the energy dependence of the improved $S$-factor $S_C(E)$ is small.

The energy dependence of the $(\alpha,\gamma)$ cross section derived from a constant $S_C(E)$ is almost identical to that of the capture cross section converted from the statistical-model $(\gamma,\alpha)$ cross section of $^{54}$. Or, in turn, one finds an almost energy-independent $S$-factor $S_C(E)$ for the $(\alpha,\gamma)$ cross section derived from $^{54}$. As can be seen from Fig. 2, $S_C(E)$ varies typically by less than one order of magnitude within an energy range of about 5 MeV. (Note that the kinks in $S_C(E)$ at 9 MeV for the $^{92}\text{Mo}(\alpha,\gamma)^{96}\text{Ru}$ reaction and at 12 MeV for the $^{144}\text{Sm}(\alpha,\gamma)^{148}\text{Gd}$ and $^{186}\text{Os}(\alpha,\gamma)^{190}\text{Pt}$ reactions indicate the opening of the $(\alpha,n)$ channel.) As a consequence of the similar energy dependencies, a similar shift of the Gamow window to lower energies is found $(i)$ from a model prediction of the $(\gamma,\alpha)$ cross section and time-reversal symmetry, and $(ii)$ from an improved and almost energy-independent $S$-factor $S_C$ as defined in Eq. (2.10).

Eq. (2.6) can thus provide only a rough estimate for the position of the Gamow window for reactions between heavy nuclei. Taking a more realistic energy dependence of the capture cross section, e.g. from the improved and almost energy-independent $S_C(E)$ or from the calculations of $^{54}$, leads to a shift of the Gamow window to lower energies by several hundred keV.

For completeness it has to be pointed out that the shift of the Gamow window to lower energies applies not only to the $(\gamma,\alpha)$ photodissociation, but also to the $(\alpha,\gamma)$ capture reaction. In such cases it is important to calculate astrophysical reaction rates $\langle \sigma v \rangle$ by numerical integration of the energy-dependent cross section instead to use simple approximations which are based on the $S$-factor at the most effective energy $S(E_{\text{eff}})$ and listed in textbooks (e.g. $^{33}$). The slight shift of the Gamow window to lower energies does not affect the comparison between $(\alpha,\gamma)$ and $(\gamma,\alpha)$ reactions in Sect. III.

Summarizing the above, the observed noticeable energy dependence of the astrophysical $S$-factor for $(\alpha,\gamma)$ reactions discussed in the preceding subsection is a consequence of the definition of the $S$-factor with the Gamow factor $\exp(-2\pi\eta)$. The Gamow factor is valid only for
increases slightly with temperature from about 300 keV
reactions of current interest. The new and improved S-
should be preferentially performed using the new S-factor
Extrapolations of experimental data to lower energies
is much weaker compared to the conventional S-factor.
and the overall agreement between these
window has been discussed first by [25]; it is located close
Experimental data have been measured for a num-
A. \((\gamma,\alpha)\) Reactions
Most of the above arguments for \((\gamma,\alpha)\) reactions in
sect. \[\text{IIA}\] hold for \((\gamma,p)\) reactions, too. Recently, much
effort has been spent on \((p,\gamma)\) capture reactions for nuclei
with masses around \(A \approx 100\) [56, 57, 58, 59, 61, 62, 63, 64, 65]; the results are summarized in [15, 66]. Addition-
\(\text{D. \((\gamma,n)\) Reactions}\)
Recently, much effort has been spent on the analysis of
\((\gamma,n)\) reaction rates. The astrophysically relevant energy
width of less than 1 MeV. The Gamow-like window is
is thermally excited. A state at \(E_{\alpha} = 3\) [32]. The shift of \(E_{\alpha}\) to lower energies
\((\alpha,\gamma)\) and \((\gamma,\alpha)\) Reactions
For simplicity, the following discussion is restricted to even-even nuclei, the most interesting cases for \(p\)-
process nucleosynthesis [4]. The discussion is illustrated using the example of the photodissociation reac-
tion \(148\text{Gd}(\gamma,\alpha)144\text{Sm}\) and the inverse capture reaction
\(144\text{Sm}(\alpha,\gamma)148\text{Gd}\). At the end of this section, some re-
marks on odd nuclei are given.
The most effective energy \(E_{\alpha}\) for \((\gamma,\alpha)\) reactions has already been defined in Eq. (2.3). Compared to the \((\alpha,\gamma)\)
capture reaction, this energy is shifted by the binding energy \(Q_{\alpha}\). For simplicity, the shift of \(E_{\alpha}\) to lower energies
because of the energy dependence of the S-factor is ne-
glected in the following section (see Fig. 1 and Sect. \[\text{III}\]).
In laboratory experiments the target nucleus is in its
ground state. For \((\gamma,\alpha)\) reactions a photon with the en-
energy around \(E_{\alpha}^{(\gamma,\alpha)}\) is required, and the dominating E1
transition leads to the excitation of a 1− state in \(148\text{Gd}\) which decays dominantly to the ground state of \(144\text{Sm}\)
by p-wave \(\alpha\) emission (see Fig. 4 left). In general, \(\alpha\)
emission at energies below the Coulomb barrier favors
the ground state decay because of the largest tunneling
probability for the highest energy. For simplicity, the
energy and width of the Gamow window are calculated from
Eqs. (2.6) and (2.7). The slight shift to lower energies
as discussed in Sect. \[\text{III}\] is neglected in Figs. 3 and
4.
In a hot stellar environment the nucleus \(148\text{Gd}\) may be
thermally excited. A state at \(E_x\) with \(J^x\) is thermally
populated according to the Boltzmann statistics:
\[n(E_x, J^x) = n(0, 0^+) \exp (-E_x/kT)\] (3.1)
The situation is illustrated for the first two excited states of
\(148\text{Gd}\) in Fig. 4 (middle and right). The thermal exci-
tation reduces the required photon energy \(E_{\alpha}^{(\gamma,\alpha)}\) by \(E_x\) leading to the same window of excitation energies in
\(148\text{Gd}\) (Fig. 4 grey shaded). Now the dominant E1 tran-
sition from the \(2^+ (3^-)\) state leads to \(J^x = 1^-, 2^-, 3^-\)
(2\(^+\), 3\(^+\), 4\(^+\)). Most of these states decay by \(\alpha\) emission mainly to the ground state of \(^{144}\)Sm (full line) with the exception of the unnatural parity states with \(J^\pi = 2^-\), 3\(^+\) decaying to excited states in \(^{144}\)Sm (dashed lines).

FIG. 4: Gamow window (shaded area) for the \(^{148}\)Gd(\(\gamma,\alpha\))\(^{144}\)Sm reaction for the ground state of \(^{148}\)Gd (0\(^+\); 0 keV; left) and for the first excited states (2\(^+\); 784 keV; middle; and 3\(^-\); 1273 keV; right) at a temperature of \(T_\text{lab} = 2.5\) \(\text{keV} \cdot \text{m}^{-1}\). Only the dominating E1 excitation is shown. See Sect. III for details.

The contributions of these thermally excited states on the stellar cross section \(\sigma^*\) can roughly be estimated. The smaller occupation probability which scales with \(\exp(-E_x/kT)\), see Eq. (3.1), is compensated by the higher photon density at the relevant effective energy \(E^\text{eff}_{\gamma,\alpha} - E_x\) which is required for the photodissociation of the excited state at \(E_x\). Consequently, all thermally populated states have comparable contributions to the stellar cross section \(\sigma^\text{lab}\) and rate \(\lambda^\text{lab}\) are usually only small contributions to the stellar cross section \(\sigma^*\) and rate \(\lambda^*\) (see also [73, 74]).

A precise estimate of the contributions of thermally excited states requires the E1 \(\gamma\)-ray strength function at the required energy \(E^\text{eff}_{\gamma,\alpha} - E_x\), i.e. at relatively low energies far below the giant dipole resonance (GDR). Based on the Brink-Axel hypothesis, the photon excitation cross section of an excited state is similar to the photon excitation of the ground state, but shifted by the energy \(E_x\) [75, 76, 77]. For increasing \(E_x\) the E1 \(\gamma\)-ray strength decreases at the requested energy \(E^\text{eff}_{\gamma,\alpha} - E_x\) because of the increasing distance to the dominating peak at the GDR. The properties of the low-energy tail of the GDR are discussed in [78], and its astrophysical relevance is analyzed e.g. in [79].

Typically one finds stellar enhancement factors of the order of 100 – 10000 for photon-induced reactions on heavy nuclei. The stellar enhancement factor depends on the low-energy behavior of the E1 \(\gamma\)-ray strength function and the level densities which enter into statistical model calculations. Stellar enhancement factors for heavy nuclei are listed in [13, 68, 80, 81]. Because of the high level density the influence of electromagnetic selection rules is relatively small for heavy nuclei compared to their dominating role in the photodissociation of light nuclei as discussed in [74].

On the other hand, the situation is different for the determination of \(E^\text{eff}_{\gamma,\alpha}\) of the \((\alpha,\gamma)\) capture reaction. Again, in the laboratory the target nucleus \((^{144}\text{Sm})\) is in its ground state. The effective energy \(E^\text{eff}_{\gamma,\alpha}\) is defined by the tunneling through the Coulomb barrier, e.g., for \(^{144}\)Sm\((\alpha,\gamma)^{148}\)Gd one obtains \(E^\text{eff}_{\gamma,\alpha} = 8.79\) MeV at \(T_\text{lab} = 2.5\). Levels with natural parity in \(^{148}\)Gd around \(E_x = E^\text{eff}_{\gamma,\alpha}\) may be populated and subsequently decay mainly by E1 transitions to all low-lying levels of \(^{148}\)Gd. This is illustrated in Fig. 5 left part.

FIG. 5: Gamow window (shaded area) for the \(^{144}\)Sm\((\alpha,\gamma)^{148}\)Gd reaction for the ground state of \(^{144}\)Sm (0\(^+\); 0 keV; left) and for the first excited states (2\(^+\); 1660 keV; middle; and 3\(^-\); 1810 keV; right) at a temperature of \(T_\text{lab} = 2.5\). See Sect. III for details.

If \(^{144}\)Sm is thermally excited, the situation is illustrated in Fig. 5 middle for \(E_x = 1660\) keV, \(J^\pi = 2^+\), and right for \(E_x = 1810\) keV, \(J^\pi = 3^-\). The required energy for the tunneling through the Coulomb barrier does not change for the thermally excited \(^{144}\)Sm: thus, one finds the same effective energy \(E^\text{eff}_{\gamma,\alpha} = 8.79\) MeV as for the ground state. Levels with many parities may be populated now, but shifted by \(E_x\) to higher energies. These excited states in \(^{148}\)Gd decay again to all low-lying states as indicated in Fig. 5.

Again, the contribution of thermally excited states to the stellar cross section \(\sigma^\text{lab}(\alpha,\gamma)\) can be estimated. The occupation probability again scales with \(\exp(-E_x/kT)\), see Eq. (3.1). But now there is no compensation because the same energy \(E_x\) is required for the tunneling, independent of the excitation energy of \(^{144}\)Sm. Hence, the contribution of excited states to the stellar cross section \(\sigma^*\) and reaction rate \(\lambda^*\) is much smaller for \((\alpha,\gamma)\) reactions compared to \((\gamma,\alpha)\) reactions.

Summarizing the above, a laboratory measurement of the \((\alpha,\gamma)\) capture cross section \(\sigma_{\text{lab}}^\text{(\alpha,\gamma)}\) is much closer to the stellar cross section \(\sigma^\text{lab}(\alpha,\gamma)\), compared with a laboratory measurement of the \((\gamma,\alpha)\) photodissociation cross section \(\sigma_{\text{lab}}^\gamma\), which is a tiny fraction of the stellar cross section \(\sigma^\gamma\). The stellar reaction rates \(\lambda^\text{\(\alpha,\gamma\)}\) of the \((\alpha,\gamma)\) reac-
tion and $\lambda^*_{(\gamma,\alpha)}$ of the $(\gamma,\alpha)$ reaction are linked together by the detailed balance theorem (e.g. [17]). Therefore, putting aside experimental difficulties for $(\alpha,\gamma)$ reactions around the effective energy, $(\alpha,\gamma)$ experiments seem to be better suited for the determination of astrophysical $(\gamma,\alpha)$ reaction rates than $(\gamma,\alpha)$ experiments. However, a careful consideration of Figs. 4 and 5 shows that the experimental $(\gamma,\alpha)$ cross section consists mainly of a well-defined transition from the ground state of the target nucleus via E1 excitation and $\alpha$ emission to the ground state of the daughter nucleus. This well-defined transition gives the chance for a precise test of theoretical predictions which cannot be performed in the same way for $(\alpha,\gamma)$ reactions (except that all $\gamma$-rays of the $(\alpha,\gamma)$ reaction could be resolved in the $(\alpha,\gamma)$ experiment - which is quite difficult for heavy nuclei).

Therefore, the $(\alpha,\gamma)$ experiment, in principle, provides the best determination of an individual astrophysical $(\alpha,\gamma)$ and $(\gamma,\alpha)$ reaction rate, but $(\gamma,\alpha)$ experiments provide the most stringent test for the statistical model and are thus absolutely necessary because of the huge number of reaction rates to be determined. So both types of experiments are complementary and should be performed to learn more about stellar $(\alpha,\gamma)$ and $(\gamma,\alpha)$ reaction rates.

Many of the above arguments remain valid for the case of odd nuclei. However, in nuclei with odd proton or neutron number excited states are located at lower excitation energies compared to even-even nuclei. A significant population of these levels is found at typical $p$-process temperatures around $T_0 \approx 2 - 3$. These low-lying excited states have various quantum numbers $J^\pi$. The Gamow window for $(\gamma,\alpha)$ and $(\alpha,\gamma)$ reactions is the same as for even-even nuclei (see Figs. 4 and 5); however, it may be slightly modified by the selection rules for the electromagnetic excitation process and the additional centrifugal barrier for the $\alpha$ particle emission.

Contrary to even-even nuclei – because of the low-lying excited states – the stellar $(\alpha,\gamma)$ cross section $\sigma^*_{(\alpha,\gamma)}$ and capture rate $\lambda^*_{(\alpha,\gamma)}$ may now be significantly different from the laboratory measurements for $\sigma^{lab}_{(\alpha,\gamma)}$ and $\lambda^{lab}_{(\alpha,\gamma)}$. Similar to even-even nuclei, the stellar photodissociation rate $\lambda^*_{(\gamma,\alpha)}$ is much larger than the laboratory rate $\lambda^{lab}_{(\gamma,\alpha)}$. The influence of the electromagnetic selection rules and the additional centrifugal barrier is difficult to predict in general.

Measurements of laboratory cross sections for $(\alpha,\gamma)$ capture and $(\gamma,\alpha)$ photodissociation reactions are necessary to get the most stringent restrictions for the theoretical predictions of cross sections and reaction rates of odd nuclei. Neither a single $(\alpha,\gamma)$ experiment nor a $(\gamma,\alpha)$ experiment is able to provide stellar reaction rates for odd nuclei. In this case experimental data have to be completed by theoretical considerations to derive stellar reaction rates from experimental data.

### B. $(p,\gamma)$ and $(\gamma,p)$ Reactions

The above arguments for odd nuclei are also valid for the relation between $(\gamma,p)$ and $(p,\gamma)$ reactions which are also governed by the tunneling probability through the Coulomb barrier. Depending on spins and parities of target and residual nuclei, the influence of electromagnetic selection rules and the additional centrifugal barrier may be enhanced in the laboratory especially when the reaction cross section is strongly suppressed by a strong mismatch between the spins of the target ground state and the populated low-lying states in the residual nucleus. As soon as many levels are thermally populated, it is very unlikely that electromagnetic selection rules and the centrifugal barrier suppress all possible transitions for the stellar $(\gamma,p)$ or $(p,\gamma)$ cross section or reaction rate. Again, laboratory measurements of the $(p,\gamma)$ and $(\gamma,p)$ reactions are required although the measured cross sections in the lab may be only a small contribution of the stellar cross section. Instead, the experimental data have to provide stringent restrictions for theoretical calculations.

### C. $(n,\gamma)$ and $(\gamma,n)$ Reactions

Contrary to the previous cases, $(n,\gamma)$ and $(\gamma,n)$ reaction cross sections show a much weaker energy dependence because there is no repulsive Coulomb interaction. Consequently, the role of the centrifugal barrier is enhanced, and the ground state contributions to the stellar reaction rates of both $(n,\gamma)$ and $(\gamma,n)$ reactions may be small. Measurements of $(n,\gamma)$ cross sections in the laboratory and the derived $(n,\gamma)$ rates may be significantly different from stellar rates at high temperatures around $T_0 = 2 - 3$. However, at typical $s$-process temperatures ($kT \approx 25$ keV) the population of excited states remains small, and stellar $(n,\gamma)$ reaction rates can usually be determined from laboratory measurements with good accuracy. At $s$-process temperatures the stellar $(\gamma,n)$ reaction rate is negligible because of (i) the negligible number of thermal photons with energies of the order of the neutron separation energy and (ii) the negligible thermal population of excited states in most cases. Nevertheless, $(\gamma,n)$ experiments may be used to provide restrictions for theoretical predictions of the inverse $(n,\gamma)$ capture cross sections [69, 82, 83]. This is especially important for unstable $s$-process branching nuclei with short half-lives which cannot be studied by direct $(n,\gamma)$ capture experiments.

### IV. PHOTODISSOCIATION AT ELECTRON SYNCHROTRON BASED $\gamma$-RAY SOURCES

Measurements of photodissociation cross sections are currently limited to the neutron channel [23, 67, 68, 69]. Experimental techniques involved in these measurements are either photoactivation with continuous
bremsstrahlung \(25, 67\) or direct neutron counting with quasi-monochromatic \(\gamma\) rays from laser Compton backscattering \(68, 69\). It is a challenge to experimentalists to measure \((\gamma, \alpha)\) or \((\gamma, p)\) cross sections of astrophysical significance primarily because of a lack of appropriate \(\gamma\) sources that enables one to measure the Coulomb-suppressed cross sections. Although a pioneering attempt has been made of investigating the \(\alpha\) channel for \(^{92}\)Mo \(84\), bremsstrahlung may not ideally be suited to determining \((\gamma, \alpha)\) or \((\gamma, p)\) cross sections because unfolding integrated yield curves is not straightforward. Instead of relying on the unfolding procedure, a method of superposing bremsstrahlung spectra with different end-point energies was used to deduce the laboratory photoreaction rate in Eq. (2.3) at temperatures of billions of Kelvin and/or the excitation function of cross sections of s-wave nature \(25\). However, this method suffers from the fact that neither experiment nor simulation can determine the shape of the end-point portion of bremsstrahlung with high precision \(25, 67, 80\).

### A. Blackbody synchrotron radiation at SPring-8

Recently, it was found \(26\) that the high energy part of the intense synchrotron radiation produced by a ten-Tesla superconducting wiggler at SPring-8 well coincides with the stellar blackbody radiation at temperatures in the range \(T_B = 1.9 - 4.4\) depending on the magnetic field strength from 4 to 10 T. This temperature range of the synchrotron radiation overlaps with that expected in the \(p\)-process nucleosynthesis. The synchrotron radiation at an equivalent temperature \(T_B = 4.4\) has the highest \(\gamma\) flux that is experimentally most favorable. The slight difference of the experimentally most favorable synchrotron radiation from the upper limit of the \(p\)-process temperature does not hamper the value of experimental data that constrain the laboratory cross section as discussed below. It is noted that the integrand \(n_{\gamma}(E, T) \sigma_{\gamma,\alpha}(E)\) in Eq. (2.3) after replacing the blackbody radiation \(n_{\gamma}(E, T)\) with the synchrotron radiation can be directly obtained from the experimental yield (see details below). One of the most important applications is to determine the photodissociation rate of \(^{180}\)Ta\(^m(\gamma,\alpha)\)\(^{179}\)Ta which is of direct relevance to the \(p\)-process nucleosynthesis of \(^{180}\)Ta\(^m\). It was also shown in \(26\) that the “blackbody synchrotron radiation” can be used to study the \(^{181}\)Ta\((\gamma,\alpha)\)\(^{177}\)Lu reaction by photodissociation. In this section, we perform a comprehensive study of the experimental feasibility of determining the laboratory \((\gamma, \alpha)\) or \((\gamma, p)\) reaction rates.

We remark again that the laboratory reaction rate is a small fraction of the stellar photoreaction rate, where photoreactions on nuclei in excited states thermally populated under standard conditions dominate over the laboratory rate by a factor of 100 - 10000 (see Sect. III). The laboratory rate can cast new light into the \((\gamma, \alpha)\) cross section \(\sigma_{\gamma,\alpha}(E)\) in Eq. (2.3), especially the \(\gamma\) transmission coefficient involved in the Hauser-Feshbach model cross section. Photomolecular reactions best probe the \(E_1\) \(\gamma\)-strength function which constitutes the main term of the \(\gamma\) transmission coefficient. Although the laboratory cross section has a threshold at the particle separation energy, the laboratory rate can be used to constrain the model \(E_1\) \(\gamma\) strength function not only above but also below the reaction threshold. The \(E_1\) \(\gamma\) strength function below the reaction threshold has direct impact on the stellar reaction rate since \(E_1\) giant resonance is built on individual excited states according to the Brink hypothesis \(72\).

There is an important feature that helps an experimental study of the \(\alpha\) channel. As pointed out in Sec. II A, the reaction becomes exothermic with a positive \(Q\)-value in the \(\alpha\) channel of photoreactions on heavy nuclei: \(Q(\gamma, \alpha) > 0\). Therefore, the most effective \(\gamma\)-ray energy in Eq. (2.3) for heavy nuclei decreases by the amount of \(Q_\alpha = -Q(\gamma, \alpha)\). As a result, the required \(\gamma\) energy is comparable to or even smaller than that for the neutron channel (see Fig. 1), while the energy of \(\alpha\) particles is boosted by the exothermic reaction. The \(\gamma\) flux of the synchrotron radiation is very high in the low energy region, which may overcome small \((\gamma, \alpha)\) cross sections governed by the Coulomb barrier. On the other hand, most of \((\gamma, \alpha)\) reactions on stable nuclei respectively boost the exothermic reaction. The \(\gamma\) flux of the synchrotron radiation is very high in the low energy region, which may overcome small \((\gamma, \alpha)\) cross sections governed by the Coulomb barrier. On the other hand, most of \((\gamma, \alpha)\) reactions on stable nuclei so that an experimental technique of direct particle counting needs to be applied. In this respect, the exothermic reaction helps to detect the promptly emitted \(\alpha\) particles.

### B. Expected \((\gamma, \alpha)\) yields at SPring-8

We discuss below measurements of the laboratory \((\gamma, \alpha)\) rate with the synchrotron radiation produced by the 10-T superconducting wiggler at SPring-8. It is foreseen to propose a dedicated beamline for experiments with the blackbody synchrotron radiation (BSR).

The number of \(\alpha\) particles \(N\) produced per second by irradiating a target sample with BSR is expressed by

\[
N(t) = n_T \int_{S_\alpha} n_{\gamma, BSR}(E, T) \sigma_{\gamma,\alpha}(E) dE, \tag{4.1}
\]

where \(n_T\) is the number of target nuclei per unit area included in the target sample and \(n_{\gamma, BSR}(E, T)\) is the flux of BSR \([s^{-1} \text{MeV}^{-1}]\) with an equivalent blackbody temperature \(T\). One can see that it is straightforward to convert this experimental quantity to the laboratory reaction rate in Eq. (2.3).

By using the \(\gamma\)-ray flux (Fig. 3 of Ref. \(26\)) and the \((\gamma, \alpha)\) cross section of \(54\), the event rate of Eq. (4.1) was calculated for all the 233 reactions compiled in \(54\). Fig. 3 shows the integrand of Eq. (4.1) for \((\gamma, \alpha)\) reactions on \(^{195}\)Pt, \(^{144}\)Nd, and \(^{99}\)Ru for \(T_B = 4.4\) corresponding to the maximum magnetic field of 10 T. (The same examples have been chosen as in the previous Sects. III and...
except that the unstable $^{148}$Gd has been replaced by the experimentally most favorable case $^{144}$Nd.) The integrand shows a peak at 7.50 MeV for $^{190}$Pt, at 7.20 MeV for $^{144}$Nd, and at 9.49 MeV for $^{96}$Ru. The most probable $\alpha$ energy is 10.75 MeV for $^{190}$Pt ($Q_\alpha = -3.25$ MeV), at 9.10 MeV for $^{144}$Nd ($Q_\alpha = -1.90$ MeV), and at 7.80 MeV for $^{96}$Ru ($Q_\alpha = 1.69$ MeV). Note that $Q_\alpha = -Q(\gamma,\alpha)$.

The integrand in Eq. (4.1) and Fig. 6 exhibits a sudden decrease at an energy corresponding to the neutron separation energy, i.e., at 8.9 MeV for $^{190}$Pt, 7.8 MeV for $^{144}$Nd, and at 10.7 MeV for $^{96}$Ru; as soon as the neutron channel opens, $\alpha$ emission competes with neutron emission which is not suppressed by the Coulomb barrier. A similar behavior is also found for the $(\alpha,\gamma)$ cross sections, see Fig. 2.

The results assure an experimental feasibility when both the number of target nuclei for the event rate and the kinetic energy for the $\alpha$ detection are taken into account as follows. When a target foil of a thickness corresponding to half the range of $\alpha$ particles with the most probable energy is used, the event rate is 44 cph (counts/hour) for $^{190}$Pt, 2200 cph for $^{144}$Nd, and 540 cph for $^{96}$Ru. These event rates are sufficient from the viewpoint of the direct detection of $\alpha$ particles. Several foils can be mounted inside a vacuum chamber for a simultaneous irradiation with BSR and $\alpha$ particles emitted can be detected with suitable detectors, for example, arrays of silicon detectors each surrounding the individual target foil. Table I lists experimental conditions for $(\gamma,\alpha)$ reactions with the event rate at the level of 10 cph or more. The count rate estimates are based on 100% enriched targets. Highly enriched targets are necessary for the suggested experiments because the detected $\alpha$ particle energy does not allow to distinguish among $(\gamma,\alpha)$ reactions on different nuclei in the target. One sees that the synchrotron radiation from the 10-T superconducting wiggler at SPring-8 can provide an unprecedented experimental opportunity to measure $(\gamma,\alpha)$ cross sections on many nuclei in the $p$-process path. It is to be noted that besides the event rate, the experimental feasibility, as a matter of course, depends also on the availability, the enrichment, and the cost of target foils. Finally, the experimental background conditions have to be analyzed carefully because of the high intensity of low-energy photons in the BSR spectrum.

Similar calculations for $(\gamma,p)$ reactions show that there is an experimental opportunity for $^{11}$ $(\gamma,p)$ reactions on nuclei from $^{59}$Co to $^{89}$Y at SPring-8 (not listed here).

V. CONCLUSION AND RECOMMENDATIONS

The current status of astrophysically relevant photon-induced cross sections and reaction rates was discussed in detail with a focus on $(\gamma,\alpha)$ reactions where almost no experimental data are available in the energy range of the astrophysical $p$-process. It is shown that experimental $(\alpha,\gamma)$ data, in principle, provide the best estimate for stellar $(\gamma,\alpha)$ reaction rates though measurements of $(\alpha,\gamma)$ reactions around the effective energy below the Coulomb barrier themselves are very difficult. Although experimental $(\gamma,\alpha)$ data in the laboratory are not perfectly suited for a direct determination of stellar $(\gamma,\alpha)$ reaction rates, such data nevertheless provide best insight into particular ingredients of the required statistical model calculations, namely the $\gamma$-ray strength function, the level density, and the $\alpha$-nucleus potential.

It is found that the energy dependence of the widely used astrophysical S-factor $S(E)$ hampers the extrapolation of cross sections to low energies especially for the heavy nuclei under study in this work. As a result, the most effective energy $E_{\text{eff}}$ is slightly lower than estimated from the usual formula in Eq. (2.4). An improved S-factor $S_C(E)$ is suggested which is based on the Coulomb wave functions instead of the simplistic Gamow factor $\exp(-2\pi\eta)$. The energy dependence of $S_C(E)$ is much smaller than for $S(E)$; thus, extrapolations using $S_C(E)$ should be more reliable.

The high-energy part of synchrotron radiation produced by a 10-T superconducting wiggler at SPring-8 is an excellent tool to mimic the blackbody radiation of temperatures around $T_0 = 1.5 - 4.4$ which are typical or slightly higher than $p$-process conditions. Based on the statistical model cross sections of [54] and the photon intensity of [26], it is shown that a noticeable number of $(\gamma,\alpha)$ cross sections can be measured with reasonable count rates. These new experimental results are expected to lead to a significant improvement of nuclear parameters involved in the statistical model, in particular, $\gamma$-ray strength functions and $\alpha$-nucleus potentials, and thus may help to reduce the uncertainties of $p$-process nucleosynthesis calculations.
TABLE II: Conditions for \((\gamma,\alpha)\) reactions that can be experimentally studied with the synchrotron radiation from the 10-T superconducting wiggler at SPring-8. The event rate is given for a single target-foil of the thickness corresponding to a half the range of \(\alpha\) particles in the target foil. The table is separated into three subtables with rates below 100 cph, above 100 cph, and above 1000 cph.

| Target nucleus | Reaction | Q-value | Gamow peak | Integral | Target thickness | Event rate |
|----------------|----------|---------|------------|----------|------------------|------------|
| \(^{69}\text{Ga}\) | -4.5     | 9.9     | 1.0E+2     | 18       | 35               |
| \(^{72}\text{Ge}\) | -4.5     | 10.2    | 4.9E+1     | 22       | 17               |
| \(^{93}\text{Nb}\) | -1.9     | 8.1     | 2.5E+1     | 8.0      | 8                |
| \(^{96}\text{Mo}\) | -2.8     | 9.1     | 2.0E+1     | 8.5      | 9                |
| \(^{98}\text{Ru}\) | -2.2     | 10.0    | 1.5E+2     | 7.8      | 62               |
| \(^{102}\text{Pd}\) | -2.2     | 10.2    | 9.7E+1     | 9.0      | 44               |
| \(^{106}\text{Cd}\) | -1.6     | 10.0    | 1.8E+2     | 13       | 87               |
| \(^{112}\text{Sn}\) | -1.8     | 10.2    | 3.0E+1     | 16       | 14               |
| \(^{122}\text{Te}\) | -1.1     | 9.5     | 2.2E+1     | 19       | 9                |
| \(^{130}\text{Ba}\) | -0.52    | 9.5     | 3.6E+1     | 38       | 16               |
| \(^{145}\text{Nd}\) | +1.6     | 6.3     | 9.3E+1     | 16       | 32               |
| \(^{149}\text{Sm}\) | -1.2     | 9.4     | 1.1E+2     | 16       | 40               |
| \(^{151}\text{Eu}\) | +2.0     | 6.0     | 1.4E+2     | 15       | 32               |
| \(^{154}\text{Gd}\) | +0.92    | 8.7     | 1.57E+1    | 21       | 7                |
| \(^{156}\text{Dy}\) | +1.8     | 7.8     | 1.4E+2     | 20       | 66               |
| \(^{162}\text{Er}\) | +1.6     | 8.1     | 4.8E+1     | 20       | 23               |
| \(^{168}\text{Yb}\) | +2.0     | 8.6     | 5.0E+1     | 30       | 27               |
| \(^{170}\text{Yb}\) | +1.7     | 7.4     | 1.9E+1     | 24       | 8                |
| \(^{174}\text{Hf}\) | +2.5     | 7.5     | 1.3E+2     | 15       | 65               |
| \(^{176}\text{Hf}\) | +2.3     | 8.3     | 4.1E+1     | 16       | 22               |
| \(^{178}\text{Hf}\) | +2.1     | 7.5     | 1.6E+1     | 14       | 7                |
| \(^{180}\text{Ta}\) | +2.0     | 6.9     | 2.3E+1     | 10       | 9                |
| \(^{180}\text{W}\) | +2.5     | 7.8     | 3.8E+1     | 11       | 20               |
| \(^{184}\text{Os}\) | +3.0     | 7.5     | 8.7E+1     | 10       | 46               |
| \(^{186}\text{Os}\) | +2.8     | 7.5     | 4.1E+1     | 10       | 21               |
| \(^{190}\text{Pt}\) | +3.3     | 7.5     | 8.2E+1     | 11       | 44               |

| \(^{66}\text{Zn}\) | -4.6     | 10.0    | 8.7E+2     | 7        | 283              |
| \(^{70}\text{Ge}\) | -4.1     | 9.8     | 1.3E+3     | 11       | 459              |
| \(^{74}\text{Se}\) | -4.1     | 10.1    | 3.7E+2     | 13       | 134              |
| \(^{96}\text{Mo}\) | -2.1     | 9.6     | 7.2E+2     | 5        | 170              |
| \(^{106}\text{Ru}\) | -1.7     | 9.5     | 1.2E+3     | 8        | 543              |
| \(^{120}\text{Te}\) | -0.3     | 8.7     | 3.7E+2     | 19       | 156              |
| \(^{142}\text{Ce}\) | +1.3     | 7.2     | 1.5E+3     | 19       | 566              |
| \(^{147}\text{Sm}\) | +2.3     | 6.3     | 8.7E+2     | 18       | 346              |
| \(^{148}\text{Sm}\) | +2.0     | 7.2     | 1.6E+3     | 20       | 697              |
| \(^{150}\text{Sm}\) | +1.4     | 7.8     | 3.2E+2     | 20       | 140              |
| \(^{152}\text{Cd}\) | +2.2     | 7.2     | 1.7E+3     | 20       | 763              |

| \(^{67}\text{Zn}\) | -4.0     | 9.4     | 3.1E+3     | 7        | 1046             |
| \(^{144}\text{Nd}\) | +1.9     | 7.2     | 5.1E+3     | 20       | 2227             |

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