



# Screening Effects in Stars and in the Laboratory

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Nuclear reactions are the driver of the evolution of many astrophysical objects. In the astrophysical environment their respective reaction rates are, however, modified due to the presence of other charges. The effects depend on the relative importance of Coulomb energy versus thermal energy and are distinguished between weak and strong screening. In the extreme case of pycnonuclear reactions, fusion reactions can be induced by the zero-point motion of nuclei in a Coulomb crystal. This paper reviews the various screening situations and discusses important applications. We also briefly review laboratory approaches to study screening effects.

**Keywords:** electron screening, stellar plasma, stellar burning, nucleosynthesis, pycno-nuclear reactions

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## 1 INTRODUCTION

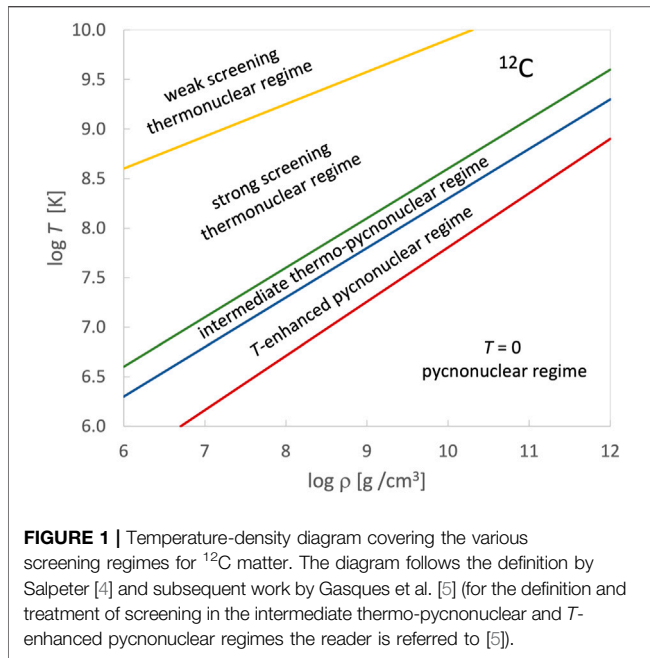
Willy Fowler liked to tease his colleagues by calling astrophysics and astronomy applied nuclear physics. With this witty remark he pointed to the fact that the dynamics, and hence the evolution, of many astrophysical objects is driven by nuclear reactions of charged nuclei, in this way also producing the elements in the Universe. During their long lives in hydrostatic equilibrium stars generate the necessary energy by fusing nuclei in their interior. When the nuclear energy source finally ceases, massive stars collapse under their own gravity resulting in a supernova explosion which finally liberates the produced elements to make them available for the next generation of stars or as the building blocks of life on Earth-like planets. During the collapse, it is nuclear reactions mediated by the weak interaction which fight against gravity [1]. Nuclear reactions are also essential to trigger thermonuclear supernovae where carbon burning is ignited as a runaway in a White Dwarf in a binary system [2].

In astrophysical environments, nuclear reactions take place in the presence of other charges that impact the rates at which such reactions occur. The conditions of density and temperature defining the stellar environment are usually such that atoms are fully ionized and hence the surrounding exists as a plasma made of nuclei of different kinds and charges, and of electrons. However, there are also important situations where atoms are only partially ionized, as is the case for heavier nuclei in the solar interior.

The effect of the surrounding charges on the reactions depend on the competition between Coulomb and thermal energies, which for a one-component plasma is defined by the parameter

$$\Gamma = \frac{(Ze)^2}{akT} \quad (1)$$

where  $k$  is the Boltzmann constant,  $T$  the temperature,  $Ze$  the charge of the nucleus and  $a$  an appropriate distance between the interacting nuclei. If  $\Gamma \ll 1$ , one speaks of a “weak screening regime” which is realized at high temperatures and/or low densities. In contrast, for high densities, i.e.,  $\Gamma \gg 1$ , one encounters the “strong screening regime”, if the most effective energy for nuclear



reactions (Gamow peak energy) is sufficiently larger than the thermal energy, or the “pycnonuclear regime”, where, even at  $T = 0$  K, fusion reactions are induced by density fluctuations. In between weak and strong regimes, one has a plasma with intermediate screening ( $\Gamma \approx 1$ ). The weak screening limit applies during stellar hydrostatic burning, while the ignition of the nuclear runaway in a type Ia supernova falls into the strong screening regime. Pycnonuclear reactions, instead, are mainly important for the crust evolution of neutron stars [3]. While screening regimes are not strictly defined, **Figure 1** sketches, as an example, the ranges of temperature and densities relevant to different screening in  $^{12}\text{C}$  matter.

Here, we review screening effects of nuclear reactions in these interesting and important astrophysical scenarios, and briefly discuss the extreme case where reactions are initiated by density fluctuations rather than by the thermal motion of the reaction partners. While we mainly focus on nuclear fusion reactions, screening is also relevant for processes mediated by the weak interaction, such as beta decay and electron capture. Important examples are the electron captures on  $^7\text{Be}$  during solar hydrogen burning and generally on nuclei during core-collapse of massive stars leading to a type II supernova. While most of the work related to screening has been based on progress in modelling, recent years witnessed also advances in experimental approaches to study screening effects in laboratory plasmas made possible by laser facilities, and in measurements of low-energy nuclear reactions. In the latter case the experimentally observed effects are somewhat larger than theoretically predicted.

## 2 STELLAR FUSION REACTIONS

Nuclear fusion reactions in stars occur at thermal energies  $E$  well below the Coulomb barrier between interacting nuclei. Thus the

energy dependence of the fusion cross section  $\sigma(E)$  is dominated by the tunnel probability and is usually written as [6]:

$$\sigma(E) = \frac{S(E)}{E} \exp[-2\pi\eta(E)] \quad (2)$$

Here,  $\eta = Z_1 Z_2 e^2 / \hbar v$  is the Sommerfeld parameter and  $v$  the relative velocity. The astrophysical S-factor  $S(E)$  accounts for nuclear effects and is a function usually mildly varying with energy. The astrophysical fusion rate is then given by folding the energy- (i.e. velocity-) dependent cross section with the stellar velocity distribution of the reaction partners, which is usually given by a Boltzmann distribution [6]. The product of the Boltzmann distribution and the tunnel probability has a pronounced peak at an energy (the Gamow peak energy) where the reaction occurs most effectively. The Gamow peak energy is noticeably larger than the thermal energy, but significantly smaller than the Coulomb barrier [6].

However, in the astrophysical environment the fusion reactions occur in the presence of other charges which modify the Coulomb barrier and hence have an impact on the astrophysical fusion rates. Starting with the pioneering work of Salpeter [4, 7], it has been customary and often possible to describe these screening effects by an enhancement factor  $F_{sc}$  which multiplies the bare nuclear reaction rate, yielding a screened rate:

$$R_{sc} = F_{sc} R_{bare} \quad (3)$$

While  $R_{bare}$  depends on the temperature via the Boltzmann distribution, the enhancement factor is also a function of plasma properties such as density, temperature, degeneracy and composition. The factorization of **Eq. (3)** allows to treat the problem as two distinct parts and is a very efficient way to incorporate the screening effects into astrophysical simulations. The screening factor  $F_{sc}$  is customarily given in the form of an exponential function for weak and strong screening as discussed below.

In the weak screening case, in which the average interaction energy between the reaction partners is much smaller than their kinetic energy, the nuclei are surrounded by a screening cloud, such that they are attracted by the electrons and repelled by the nuclei in the partner’s cloud. Due to the effect of the screening clouds, the reaction partners with charges  $Z_1, Z_2$  feel a screened Coulomb potential

$$U_{sc}(r) = \frac{Z_1 Z_2 e^2}{r} \exp(-r/R_D) \quad (4)$$

rather than the Coulomb repulsion of bare nuclei,  $U_b(r) = Z_1 Z_2 e^2 / r$ , where the parameter  $R_D$  (Debye length) is given by  $\zeta R_D = (\frac{kT}{4\pi\alpha\rho})^{1/2}$  with the fine-structure constant  $\alpha$  and the number density of nucleons  $\rho$  [8]. The parameter  $\zeta$  depends on the composition of the plasma and its degree of degeneracy [4, 8]. Salpeter’s model of weak screening is equivalent to the Debye-Hückel theory of dilute solutions of electrolytes [9].

The tunnel process occurs at radii which are significantly smaller than the Debye radius. Thus, the exponential in the screened Coulomb potential (**Eq. (4)**) can be expanded to first

order,  $\exp(-r/R_D) \approx 1 - r/R_D$ . ... Hence, the effect of screening is a reduction of the Coulomb barrier by a constant screening energy  $U_{sc} = (Z_1 Z_2 e^2)/R_D$ . Equivalently, the partners in the plasma fuse with a slightly increased energy ( $E + U_{sc}$ ) and the enhancement factor in Eq. (3) reduces to  $F_{sc} = \exp(-U_{sc}/kT)$ .

The screening enhancement factors are incorporated in the stellar evolution codes. For the solar plasma, the Debye radius is of the order  $10^{-11}$  m and the screening enhancement of nuclear reactions at the relevant solar energies is relatively small, about 5% for the important proton-proton reaction [10]. Special care has to be taken for resonances as the modification of the Coulomb potential lowers the resonance energy compared to the reaction threshold. Hence, the resonance lifetime is increased due to screening. If screening is sufficiently strong, the resonance might get bound in the astrophysical environment; we will mention an example below. The theory for resonant screening has been discussed for example in Ref. [11].

The Debye-Hückel theory of screening, as adopted by Salpeter, is based on a mean-field approach. Hence, it depends on the distance between the reaction partners, but not on their velocity. It has been argued that such a static approach would not be appropriate for the solar plasma [12]. In fact, studies of dynamical screening performed within the molecular dynamics approach for a proton plasma [13–15] indicate that fast moving protons which might not be accompanied by their full screening cloud [15] experience less screening than slow moving ones. These studies suggest that the static screening overestimates the enhancement of nuclear reactions in stellar plasmas, as fast moving particles with energies around the Gamow peak dominate the reaction rates. This view has, however, been disputed by Bahcall and collaborators who argued that Salpeter's screening approach is valid also at the Gamow peak energy due to the nearly perfect thermodynamic equilibrium present in the solar plasma [16].

At high densities, where  $\Gamma \gg 1$ , the bulk of the nuclei form a lattice (or quantum fluid). If the Gamow peak energy is significantly larger than the thermal energy, the main contribution to the reaction rate will come from fusing nuclei with energies around the Gamow peak. Again, Salpeter laid out the basis for the treatment of this situation of "strong screening" [4]. He showed that electrons can be treated as a uniform background and derived an enhancement factor assuming that the fusing nuclei move while being surrounded by a uniform cloud of electrons that neutralise the nuclear charge. In this approximation, the enhancement factor can again be written as  $F_{sc} = \exp(-U_{sc}/kT)$  and the screening energy  $U_{sc}$  is given by the difference of electrostatic energies between the fused system with combined ( $Z_1 + Z_2$ ) charge and that of the two individual nuclei with charges  $Z_1$  and  $Z_2$ .

An improved treatment of strong screening is obtained by determining the pair correlation function  $g(r)$  which is a measure for the probability for a given nucleus to find a reaction partner at distance  $r$  in the plasma. (In multicomponent plasmas  $g$  carries indices defining the partners.) The calculation of  $g(r)$  is identical to determining the mean potential of the plasma  $w(r)$ , with  $g(r) = \exp(-w(r)/kT)$ .

Defining  $g$  and  $g_0$  as the pair correlation functions with and without screening, the enhancement factor can be approximated

as  $F_{sc} = g(0)/g_0(0)$ , where one considers that the ratio is quite insensitive to  $r$  at nuclear distances. The enhancement factor is usually split into a "classical" and a "quantum" part [17, 18],  $F_{sc} = \exp(h_0 + h_1)$ . The classical part is approximately related to the difference of free energies  $\Delta F$  for the many-body system before and after the fusion reaction [19], leading to  $\exp(h_0) = \exp(\Delta F/kT)$ . The parameter  $h_0$  can also be calculated using Monte Carlo techniques. A parametrized form of  $h_0$  as a function of the Coulomb coupling parameter  $\Gamma$  and valid for  $1 < \Gamma < 170$  is given in [20]. The quantum contribution  $\exp(h_1)$  considers that the potential is not constant during the tunnel process. It has been calculated in different approximate ways (by path integrals [21], within the WKB approach [22, 23], and directly by Monte Carlo techniques [24, 25]) and found to be small compared to the classical contribution.

An important astrophysical application of strong screening occurs for carbon fusion reactions which are believed to trigger type Ia supernovae in degenerate White Dwarfs in the so-called single-degenerate scenario [2]. Under these degenerate conditions carbon burning occurs as a thermonuclear runaway where the reaction rate is enhanced by many orders of magnitude due to screening effects. The different approximations to the quantum effects lead to noticeable differences in the enhancement factors, which, however, are small compared to the total enhancement. Hence the differences are not expected to be too relevant [5]. A convenient parametrization for the screening of the carbon fusion rate can be found in [5].

At high densities and sufficiently low temperatures the nuclei settle into a Coulombic lattice, for which usually a body-centred-cubic (*bcc*) structure is assumed [26]. At very high densities the lattice is destroyed due to the zero-point motion of the nuclei and the system rather becomes a quantum fluid. However, due to its strong dependence on charge, this transition occurs at such high densities that it is usually irrelevant for astrophysical purposes [27]. An exception is  ${}^4\text{He}$ , as we will see below. The energy scale of the zero-point motion can be estimated by the ion plasma frequency [7] and allows to tunnel the Coulomb barrier, which is significantly modified due to the interaction with other ions and the neutralising electron background. Such density-induced reactions are called pycnonuclear reactions [7] and they are the reason why no Coulomb crystal exists at arbitrarily large densities.

The pycnonuclear reaction rate  $R_{\text{pyc}}$  for a one-component crystal can be expressed (in  $\text{s}^{-1}\text{cm}^{-3}$ ) as [5]:

$$R_{\text{pyc}} = \rho X_i A Z^4 S(E_{\text{pk}}) C_{\text{pyc}} 10^{46} \lambda^{3-C_{\text{pl}}} \exp(-C_{\text{exp}}/\sqrt{\lambda}) \quad (5)$$

where  $\rho$  is the density in  $\text{g}/\text{cm}^3$ ,  $X_i$ ,  $A$ ,  $Z$  are the mass fraction, mass number and charge of the fusing nuclei and the astrophysical S-factor  $S$  is evaluated at the energy  $E_{\text{pk}}$  of the plasma frequency (see [5] for details). Above (below) the neutron drip density at  $4 \times 10^{11} \text{g}/\text{cm}^3$ ,  $X_i < 1$  ( $X_i = 1$ ) due to the presence of free neutrons. The parameter  $\lambda$  is defined as  $\lambda = \frac{\hbar^2}{mZ^2e^2} (n_i/2)^{1/3}$  with nucleon mass  $m$  and number density  $n_i$ . The three parameters ( $C_{\text{pyc}}$ ,  $C_{\text{pl}}$ ,  $C_{\text{exp}}$ ) in Eq. (5) have to be calculated using a model for the tunnel process. Various values are presented in [7, 24, 27].

Beyond the static approximation to the rate given in Eq. (5), effects related to the polarization of the surrounding Coulomb lattice were discussed by Schramm and Koonin [27]. They found that these dynamical effects, which are non-negligible per se, cancel each other. It is estimated that pycnonuclear reaction rates can be uncertain by a few orders of magnitude [26].

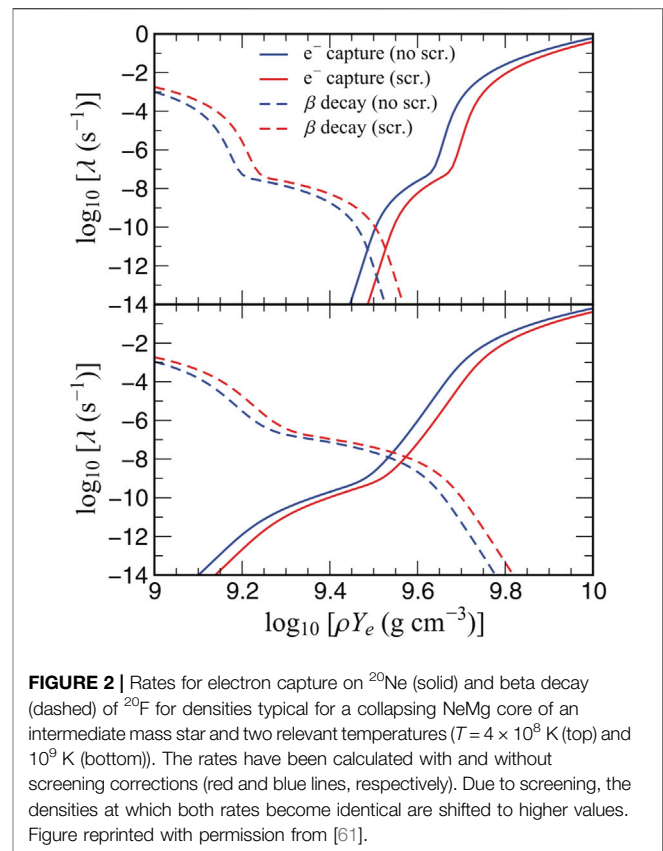
Pycnonuclear reactions play a role in the structure of isolated neutron stars which accrete matter from the interstellar medium [28]. Due to the accretion, the original matter of the neutron star crust is successively pushed to higher densities and hence undergoes a series of double electron captures which make the matter more neutron-rich. These captures produce free neutrons above the neutron drip density, until the system finally reaches densities at which nuclei become able to fuse by pycnonuclear reactions. In the pioneering work of [3, 29], where the original composition of the crust was assumed to consist purely of  $^{56}\text{Fe}$ , pycnonuclear fusion is expected to occur for  $^{34}\text{Ne}$  nuclei at densities slightly exceeding  $10^{12}\text{ g/cm}^3$ . Dynamical studies of the neutron star crust evolution encompassing a large reaction network are presented in [30]. These authors also find that the uncertainty in the pycnonuclear reaction rates have only a modest effect on the crust evolution.

Special care has to be taken for  $^4\text{He}$  matter at high densities. Of particular interest is the triple-alpha rate, by which  $^4\text{He}$  is transformed into  $^{12}\text{C}$ . During helium burning in normal stars, temperatures are sufficiently high that the process occurs in two steps via the formation of an intermediate  $^8\text{Be}$  resonance (the ground state) at 92.2 keV above the  $\alpha + \alpha$  threshold [31]. In Wolf-Rayet stars, temperatures are lower and the triple-alpha reaction is enabled via the low-energy wing of the  $^8\text{Be}$  resonance [32, 33]. Finally, in accreting neutron stars, temperatures might be so low that the triple-alpha process occurs via pycnonuclear fusion [34]. At the relevant densities ( $\rho > 3 \times 10^8\text{ g/cm}^3$ ), zero-point motion transforms  $^4\text{He}$  matter into a quantum fluid. Under these conditions, the triple-alpha fusion rate was calculated in [35]. However, because of the low energy of the  $^8\text{Be}$  resonance, the nuclear  $\alpha$ - $\alpha$  potential must also be included in addition to the Coulomb interaction. As it was pointed out in [36], screening in  $^4\text{He}$  matter becomes strong enough around  $3 \times 10^9\text{ g/cm}^3$  to transform it to bound  $^8\text{Be}$  matter. This phase transition was confirmed in a study which described the  $^4\text{He}$  quantum fluid within the hypernetted chain approach taking higher-order correlations into account [37]. The study also showed that these correlations are strong enough to transform  $^4\text{He}$  matter into  $^{12}\text{C}$  at densities lower than the critical density for the phase transition to  $^8\text{Be}$  matter. The latter is then probably not realised in nature.

For screening effects of very strong magnetic fields on thermonuclear reactions the reader is referred to Refs. [38, 39] and references therein.

### 3 REACTION MEDIATED BY THE WEAK INTERACTION

Although reactions mediated by the weak force are usually slower than competing time scales, they can be significantly altered by



their astrophysical surroundings often with interesting consequences. For example, temperature can change the effective half lives of nuclei due to thermal population of excited nuclear states. In s-process nucleosynthesis this occurs at certain branching points making the rates for beta decays and neutron captures competitive. As a consequence, the matter flow branches at these nuclei allowing to determine the neutron density or the temperature of the astrophysical environment from the observed abundances [40, 41]. An extreme change in half life occurs for  $^{180}\text{Ta}$  where a long-lived excited state with half live comparable to the age of the Universe is thermally mixed with the ground state which lives only a few hours [42, 43]. Other than by s-process nucleosynthesis [44, 45],  $^{180}\text{Ta}$  can also be made by neutrino nucleosynthesis [46–48]. Finite-temperature effects also modify inelastic neutrino-nucleus cross sections, in particular at low and modest neutrino energies [49]. Inelastic neutrino scattering off nuclei had been discussed as a potential source for thermalization of neutrinos with matter during the collapse of massive stars [50, 51], but turns out to be rather unimportant [52].

The plasma density has important effects also on beta decays, as the presence of electrons blocks final states for the decay. This Pauli blocking effect at finite density reduces the rate and has to be accounted for in the derivation of stellar beta decay rates [53, 54]. As a consequence of this reduction, beta decays are basically prohibited during final collapse of a massive star, but play an interesting role during silicon burning where beta decay and

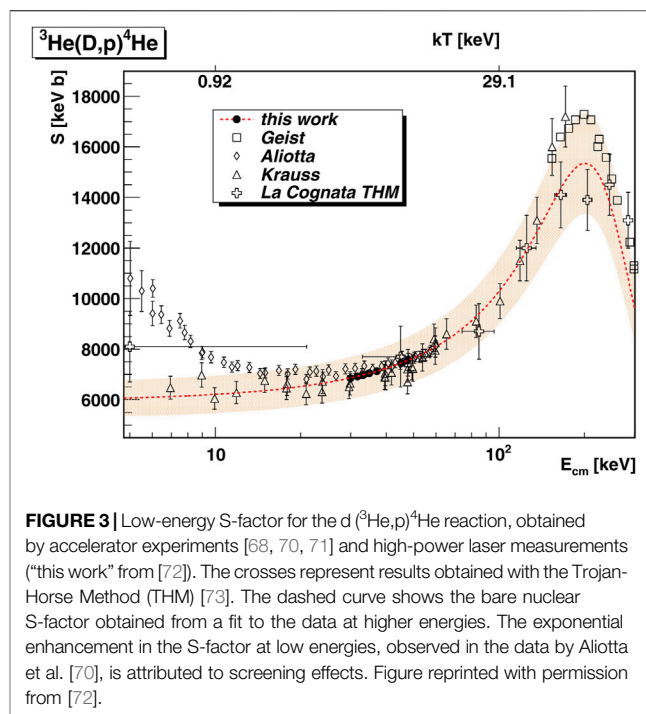
electron capture rates are comparable and contribute significantly to the cooling of the stellar core without changing the electron-to-nucleon ratio [55]. The blocking of beta decays allows the collapsing core to get successively more neutron-rich by continuous electron captures on nuclei [1, 56, 57]. These electron captures accelerate the collapse as they reduce both the pressure of the relativistic electron gas and the entropy of the core matter. The latter effect is responsible for the fact that nuclei survive the collapse. Coulomb corrections are important for the thermodynamic properties of the high density plasma [58], but they also modify the weak interaction rates [59]. Screening contributes to these rates in two different ways. First, it changes the threshold energy between initial and final states. Second, the electron energy is reduced by the presence of the background compared to the unscreened case. Refs. [60, 61] present a formalism on how screening corrections can be incorporated into the stellar electron capture and beta decay rates. In general, electron capture rates are slightly reduced due to screening, while beta decay rates are enhanced [61], as shown in **Figure 2**.

This result has important consequences for the URCA process which is essential for the cooling of NeMg cores in the final evolutionary stages of intermediate mass stars as it shifts the operation of the URCA process to higher densities [61–63]. Significant experimental and theoretical progress in determining stellar weak interaction rates has recently been reviewed in [64].

Electron captures on selected nuclei occur also during hydrostatic stellar burning. At such conditions screening cannot be treated as caused by a degenerate relativistic electron gas like during core collapse, but rather as due to an ideal electromagnetic plasma. A particularly interesting and important case is the electron capture on  ${}^7\text{Be}$  which, in competition with proton capture, determines the rate of high-energy neutrinos produced in the Sun. In the laboratory the half life of  ${}^7\text{Be}$  is about 53 days. As the electron density at the nucleus in the solar plasma is reduced compared to the atomic case, the half life is longer in the Sun. Under solar conditions,  ${}^7\text{Be}$  has a non-negligible probability, which also depends on screening effects, of having a bound K-shell electron. Therefore, a consistent treatment of  ${}^7\text{Be}$  ions embedded in the solar surroundings is required. Such studies have been performed on the basis of two quite distinct methods, a mean-field approach [65] and a density matrix formalism [66], yielding, however, nearly identical results for the solar  ${}^7\text{Be}$  half life (about 81 days).

## 4 EXPERIMENTAL APPROACHES TO STUDY SCREENING EFFECTS

The strategy to determine nuclear reaction rates for astrophysical applications is a two-step process: at first one derives the cross section (astrophysical S-factor) for bare nuclei, which is then modified to account for environmental effects. The measurement of the bare cross section is, however, an experimental challenge, as the astrophysically most important energy region (around the



**FIGURE 3** | Low-energy S-factor for the  $d({}^3\text{He},p){}^4\text{He}$  reaction, obtained by accelerator experiments [68, 70, 71] and high-power laser measurements (“this work” from [72]). The crosses represent results obtained with the Trojan-Horse Method (THM) [73]. The dashed curve shows the bare nuclear S-factor obtained from a fit to the data at higher energies. The exponential enhancement in the S-factor at low energies, observed in the data by Aliotta et al. [70], is attributed to screening effects. Figure reprinted with permission from [72].

Gamow peak energy) lies far below the Coulomb barrier, where the cross sections are extremely small and their direct measurement in the laboratory is difficult with presently available techniques. Therefore, one measures the cross sections down to energies as low as possible and then extrapolates the data to the relevant energy region. It was generally believed that the uncertainty in the extrapolation procedure could be reduced by steadily lowering the energies at which data are taken in the laboratory. This approach, however, might be problematic as at very low energies the experimentally measured cross section does not represent the required case for bare nuclei: the cross section is enhanced due to screening effects connected with the electrons present in the target (and possibly in the projectile) [67].

The  ${}^3\text{He}(d,p){}^4\text{He}$  reaction is probably the best studied case for the experimental observation of laboratory screening, as the available low-energy data [68–71, 73] are the most definitive in this case.

In the laboratory the cross section was measured by shooting an ion beam of (bare) deuterons onto an atomic  ${}^3\text{He}$  gas (with two electrons per atom). At the lowest accessible energies, the separation of target and projectile during tunnelling is much smaller than the atomic radius; thus, the two electrons, originally bound to  ${}^3\text{He}$ , are now attracted by the joint charge of projectile and target, corresponding to  ${}^5\text{Li}$ . The gain in electronic binding energy (about  $\Delta E \approx 120$  eV) between the fused  ${}^5\text{Li}$  and the well-separated  $(d+{}^3\text{He})$  systems should be transferred to the relative motion so that the fusion process occurs with a slightly enhanced energy compared to the case of bare nuclei. Analogue to **Eq. (3)** the measured cross section should be increased by an exponential factor which can be expressed as  $f_{sc} \approx \exp(\pi\eta(E)\Delta E/E)$ , considering that  $\Delta E$  is much smaller than the relative energy  $E$

of the fusing nuclei [67]. Such an exponential enhancement is indeed experimentally observed [69, 70], as shown in **Figure 3**. However, the screening energy  $U_{sc}$  obtained by fitting the data to the exponential parametrization for  $f_{sc}$  is larger than the adiabatic limit  $\Delta E$ , i.e. the difference in binding energies. Yet, the applicability of the adiabatic limit was confirmed within a Time Dependent Hartree-Fock (TDHF) calculation of the electron dynamics in the fusion reaction [74]. In addition, a noticeably smaller experimental screening energy was observed in an experiment performed in inverse kinematics using a molecular deuteron target [75]. A TDHF study of screening effects for proton scattering on a molecular hydrogen target shows a strong sensitivity of the enhancement factor to the orientation of the molecule [76]. It is smaller if the proton has to pass the spectator nucleus before the fusion process occurs. No isotopic dependence of the screening effect has been found, as expected [77, 78]. The use of possibly incorrect stopping powers, which the fusion energies have to be corrected for, has also been suggested in [79] as a possible solution of the difference between experimental results and adiabatic limit. Indeed, theoretical calculations [80, 81] of stopping power for the same systems mentioned above using dynamic calculations do not agree with extrapolations to lower energies, thus corroborating the findings in [79].

Differences between experimental and adiabatic screening energies have been found also for other nuclear reactions [82]. In some cases, clustering configurations (e.g.,  $d \otimes \alpha$  in  ${}^6\text{Li}$ ) in the interacting nuclei could potentially explain the large screening potentials observed experimentally, as suggested in [83].

In general, an obvious source of uncertainty are the bare nuclear cross sections required to determine the experimental screening value. To eliminate this uncertainty experimental approaches like the Trojan Horse Method (THM) [84, 85] have been developed, which allow for an indirect determination of the bare nuclear cross sections also at very low energies. **Figure 3** includes bare nuclear cross sections obtained via the THM [73] that are consistent with the extrapolation of the cross sections from higher energies. The S-factor has also been determined by an experiment where a molecular deuteron target was bombarded with intense and ultrafast laser pulses [72]. The discrepancy between observed and expected screening enhancement for gaseous targets is still not explained.

Large screening enhancements were also found for the low-energy  $d + d$  fusion reaction in deuterated metals [86–88]. Here the low-energy S-factor shows an exponential enhancement, but the deduced screening energy for metals can exceed the value obtained for a gas-target experiment by more than an order of magnitude [88, 89]. Importantly, if the deuterons are implanted into semiconductor or insulator materials, the screening energy is consistent with the gas-target result. This observation led to the proposal that the classical Debye theory might be applicable to describe the screening enhancement in metals [90]. The Debye model predicts a characteristic temperature dependence of the screening proportional to  $T^{1/2}$  which was tentatively observed in [91]. Considering the ion-momentum distribution and inhomogeneous screening effects beyond the Debye model reduced the gap between prediction and observation [92, 93]. The Debye model was, however, questioned in [89]. These authors argued that the screening in

deuterated metals requires a dynamical treatment of the deuteron density which indeed reproduced the dependence of the screening energy on the charge of the host material. But also this improved model predicts screening energies which are about a factor of 2 smaller than the observed values [89, 94]. A quantum-mechanical model for the screening effects was presented in [95], but has not yet been fully applied due to computer limitations.

The unresolved discrepancy between theoretical and experimental screening energies for atomic or molecular targets is an obstacle to reducing existing uncertainties for nuclear reactions with light ions. The problem would, of course, become obsolete if it was possible to measure the relevant cross sections directly for the appropriate astrophysical plasmas. At least for interactions between light ions in the Sun, in early stellar burning stages, or in Big Bang nucleosynthesis, this perspective becomes possible using high-intensity lasers. For example, a measurement of the low-energy  $d({}^3\text{He}, p){}^4\text{He}$  S-factor is included in **Figure 3** [72]. Other laser-induced measurements of astrophysical S-factors dealt with the  $d(p, \gamma){}^3\text{He}$  [96] and the  $t+{}^3\text{He}$  [97] reactions, both of interest for Big Bang nucleosynthesis. These pioneering works were, however, performed under plasma conditions at which screening is expected to be negligible. Such dedicated studies of plasma screening effects are envisioned for future laser facilities, including the Extreme Light Infrastructure ELI-NP under construction in Bucharest [98] making use of two laser-induced colliding plasmas [99]. An alternative approach, i.e. to measure directly the bare reaction cross section, exploits a unique combination of experimental devices at the CRYRING storage ring at GSI. Here, a stored beam of ions can be made to collide with a transverse beam and the reaction products be detected with a newly developed detection array [100]. If successful, such measurements have the potential for significant breakthroughs in our understanding of the electron screening effect.

## 5 CONCLUSION

In astrophysical environments nuclear reactions are affected by the presence of other charges in the stellar plasma. Starting with the pioneering work of Salpeter [4, 7] such effects have been included in astrophysical simulations taking also into account that the relative importance of Coulomb versus thermal energy, defined by the Coulomb plasma parameter  $\Gamma$ , requires different screening treatments under different physical conditions. The progress of these treatments is mainly theoretical. In the case of weak screening, as it applies for example to the solar plasma, the enhancement of the fusion reactions is quite small (typically, a few percent) and appears not to be a considerable source of uncertainty [8]. Similar conclusions also apply to the electron capture reaction on  ${}^7\text{Be}$ , which has drawn specific attention in [101]. The uncertainty in the enhancement factors for the strong screening case and in particular in the pyconuclear regime remains noticeably larger.

Experimental approaches to screening are so far restricted to studies of light-ion fusion reactions where (at the lowest energies reachable in accelerator experiments) the presence of electrons in the nuclear target (and possibly also in the projectile) screens the Coulomb barrier, effectively enhancing the relative energy of the collision partners during the tunnel process. The energy shift

obtained from experimental data is often larger than the value expected if the collision proceeds adiabatically. This discrepancy is not yet fully understood and introduces some uncertainty into the determination of astrophysical cross sections for light ion reactions. For heavier nuclei the energies reachable in accelerator experiments are too high for screening effects to be relevant. In the future it might be possible to circumvent this uncertainty source by measuring the relevant cross sections directly in a laser-induced plasma.

## AUTHOR CONTRIBUTIONS

MA and KL have contributed, respectively, to some of the experimental and theoretical results presented in this review and have co-written the manuscript.

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