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Bijaya Acharya, Lucas Platter, and Gautam Rupak

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Universal behavior of p-wave proton-proton fusion near threshold

Bijaya Acharya,[1,2,*] Lucas Platter,[2,3,†] and Gautam Rupak[4,‡]

[1] *Institut für Kernphysik and PRISMA+ Cluster of Excellence, Johannes Gutenberg-Universität Mainz, 55128 Mainz, Germany*
[2] *Department of Physics and Astronomy, University of Tennessee, Knoxville, TN 37996, USA*
[3] *Physics Division, Oak Ridge National Laboratory, Oak Ridge, TN 37831, USA*
[4] *Department of Physics & Astronomy and HPC² Center for Computational Sciences, Mississippi State University, Mississippi State, MS 39762, USA*

We calculate the p-wave contribution to the proton-proton fusion S-factor and its energy derivative in pionless effective field theory (EFT) up to next-to-leading order. The leading contributions are given by a recoil piece from the Gamow-Teller and Fermi operators, and from relativistic $1/m$ suppressed weak interaction operators. We obtain the value of $(2.5 \pm 0.3) \times 10^{-28}$ MeV fm^2 for the S-factor and $(2.2 \pm 0.2) \times 10^{-26}$ fm^2 for its energy derivative at threshold. These are smaller than the results of a prior study that employed chiral EFT by several orders of magnitude. We conclude that, contrary to what has been previously reported, the p-wave contribution does not need to be considered in a high-precision determination of the S-factor at astrophysical energies. Combined with the chiral EFT calculation of Acharya *et al.* (2016) [1] for the s-wave channel, this gives a total threshold S-factor of $S(0) = (4.047^{+0.024}_{-0.032}) \times 10^{-23}$ MeV fm^2.

The Sun is powered by nuclear burning of hydrogen, the most abundant element in the universe, into helium. The elementary proton-proton (pp) fusion process that results in a deuteron, a positron and a neutrino is the first step in the chain of reactions producing heavier elements in stellar environments [2]. Solar models for quantities such as core temperatures and neutrino flux are sensitive to the pp fusion cross section. At the relevant solar core temperatures ($T \sim 1.5 \times 10^7$ K), the Coulomb repulsion and the slow weak process result in a very small cross section. Thus, experimental measurements are prohibitive and non-existent. Theoretical calculations with well-justified uncertainty estimates are essential for providing critical input data for stellar models [3–5]. Inference of solar neutrino masses from terrestrial measurements depends crucially on the pp fusion rate. This reaction involves all the fundamental interactions except gravity. It is important in the field of astro, nuclear and particle physics, and there is an active effort to calculate the cross section with ever higher accuracy and precision [see Ref. [4] (SFII) for an extensive review of the existing literature].

The reaction cross section $\sigma(E)$ at center-of-mass (c.m.) kinetic energy E is conventionally expressed in terms of the S-factor $S(E) = E \exp(2\pi\eta_p)\sigma(E)$. The Sommerfeld parameter $\eta_p = \sqrt{m_p/E}\,\alpha/2$ with proton mass $m_p = 938.28$ MeV and fine structure constant $\alpha = 1/137$. SFII provides the best estimates of $S(0) = (4.01 \pm 0.04) \times 10^{-23}$ MeV fm^2 at threshold, and $S'(0)/S(0) = (11.2 \pm 0.1)$ MeV^{-1} for the logarithmic derivative. SFII also estimated the contribution of the $S''(0)$ term to be $\sim 1\%$ at the solar core temperature and recommended that a modern calculation be undertaken. The threshold S-factor and its energy derivatives have since been calculated in pionless [6] and chiral [1, 7] effective field theories (EFTs). Reference [7], the only study so far to have included capture from the p-wave, has claimed that this channel makes a significant contribution to $S(E)$, of roughly the same size as the s-wave $S''(0)$ term, in the astrophysically relevant $E \sim 10$ keV region [1]. An independent calculation of the p-wave contribution is therefore imperative, especially since the s-wave $S(E)$ has now been constrained to subpercentage precision [1].

EFTs provide a description of interacting particles in terms of only those degrees of freedom that are relevant below a breakdown momentum scale, Λ. Low-energy observables are then calculated as an expansions in powers of Q/Λ, where Q is the characteristic momentum of the process under study. Such approaches have been widely used in nuclear physics. They provide a clear guidance on how to systematically construct the nuclear Hamiltonian and couplings to external electroweak sources as perturbation in Q/Λ. They also enable us to use the convergence of the expansion to estimate the uncertainty in theoretical calculations. The pp fusion process at solar energies $E \lesssim 100$ keV is peripheral, and thus can be accurately described in terms of the incoming pp s-wave phase shift and the outgoing deuteron bound state wave function to within 10% model-independently [2]. Thus the characteristic momentum scale $Q \sim p, \gamma, 1/a_{pp}, \alpha m_p \ll$

* acharya@uni-mainz.de
† lplatter@utk.edu
‡ grupak@ccs.msstate.edu

[1] After the authors of Ref. [7] were notified about this article, they revisited their calculation and published an Erratum [8] whose p-wave result, albeit much closer, is still not in agreement with this work within our uncertainty estimate. More importantly, as we will later discuss, the result of Ref. [8] for the total S-factor, with s- and p-waves included, does not agree with the value we quote here due to basis-truncation issues in their calculation.

m_π, where $p = \sqrt{m_p E} \lesssim 10$ MeV, $a_{pp} \sim 25$ MeV is the pp s-wave scattering length, $\gamma = 45.701$ MeV the deuteron binding momentum, $m_\pi \sim 140$ MeV the pion mass. It is therefore appropriate to employ Pionless EFT ($\not\pi$EFT) for the calculation of the pp fusion S-factor. This is an EFT with non-relativistic nucleons that interact through short-ranged forces without an explicit pion degree of freedom [9, 10]. Its breakdown scale is $\Lambda \sim m_\pi$ and the perturbative expansion is therefore in $Q/\Lambda \lesssim 1/3$. $\not\pi$EFT provides a simple description of pp fusion, to about 10% precision, in terms of nucleon-nucleon observables [11]. Calculation of the fusion rate to a few percent precision requires contribution from two-body currents that represent short-distance physics not constrained by elastic-channel nucleon-nucleon phase shifts [12].

In this Letter, we present the first calculation of the p-wave contributions to pp fusion in $\not\pi$EFT. The results are expressed in terms of model-independent parameters, and, therefore, universal. It provides an important constraint on the precise determination of the solar pp fusion rate and provides insights into further steps that are need to reduce uncertainties in the future.

a. **Pionless effective field theory:** The cross section calculation depends on the strong interaction, the Coulomb repulsion between the two protons, and the weak interaction. The dominant p-wave contribution requires the strong interaction only in the outgoing deuteron ($^3S_1$) channel, which is given by [9, 13–15]

$$\mathcal{L}_S = d_i^\dagger \left[\Delta - \left(i\partial_0 + \frac{\nabla^2}{4m}\right)\right] d_i + g_0 \left[d_i^\dagger (N^T P_i N) + \text{h.c}\right] , \tag{1}$$

where $m = 938.92$ MeV is the isospin-averaged nucleon mass, N represents a nucleon and the vector d_i represents the deuteron. $P_i = \sigma_2\sigma_i\tau_2/\sqrt{8}$, where the Pauli matrices $\boldsymbol{\sigma}$ and $\boldsymbol{\tau}$ respectively act on spins and isospins, projects the nucleons onto the spintriplet isosinglet $^3S_1$ deuteron channel. The two couplings Δ, g_0 are fixed by requiring that the deuteron bound state wave function has the correct exponential decay and normalization constant. In $\not\pi$EFT, this corresponds to ensuring the $^3S_1$ elastic scattering amplitude has a pole at $p^* = i\gamma$, and has the correct residue at the said pole. While these depend only on γ at leading order (LO), the contributions of the effective range $\rho = 1.764$ fm to the residue, which enter at next-to-leading order (NLO), can be expressed in terms of the deuteron wave function renormalization constant, Z_d, and treated exactly using the zed-parameterization [16].

We include the Coulomb interaction between the protons using the t-matrix $-it_C(E;\boldsymbol{q},\boldsymbol{p})$ for incoming (outgoing) momentum $\boldsymbol{p}$ ($\boldsymbol{q}$). It can be expressed in closed form using the momentum-space Coulomb wave function $\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q})$ as: $t_C(E;\boldsymbol{q},\boldsymbol{p}) = (E - q^2/m_p + i0^+)\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q})$. Coulomb amplitude t_C includes non-perturbative resummation of Coulomb photon exchanges.

The capture from pp p-wave initial state receives contribution from two sets of weak interactions. The first set constitutes the usual Fermi and Gamow-Teller interactions:

$$\mathcal{L}_W^{(\mathrm{FGT})} = -\frac{G_V}{\sqrt{2}} \left(l_+^0 N^\dagger \tau^- N + g_A \boldsymbol{l}_+ \cdot N^\dagger \boldsymbol{\sigma}\tau^- N\right) , \tag{2}$$

where G_V and g_A are the vector and axial coupling constants, for which we use the latest Particle Data Group [17] values of $1.1363(3) \times 10^{-11}$ MeV^{-2} and 1.2724(23), respectively. l_+^μ is the leptonic Dirac current, and $\tau^- = (\tau_1 - i\tau_2)/2$ is the isospin lowering operator.

The second set of interactions constitutes relativistic p/m effects:

$$\mathcal{L}_W^{(\mathrm{rel})} = \frac{G_V}{\sqrt{2}} \Big[g_A l_+^0 N^\dagger \boldsymbol{\sigma} \cdot \frac{i\overleftrightarrow{\boldsymbol{\nabla}}}{2m} \tau^- N + \boldsymbol{l}_+ \cdot N^\dagger \Big(\frac{i\overleftrightarrow{\boldsymbol{\nabla}}}{2m}\tau^- - \mu_V \boldsymbol{\sigma} \times \frac{\overline{\boldsymbol{\nabla}}}{2m}\tau^-\Big) N\Big] , \tag{3}$$

where $\mu_V = (\mu_p - \mu_n)/2$ denotes the isovector magnetic moment, $\overleftrightarrow{\boldsymbol{\nabla}} = \overleftarrow{\boldsymbol{\nabla}} - \overrightarrow{\boldsymbol{\nabla}}$ and $\overline{\boldsymbol{\nabla}} = \overleftarrow{\boldsymbol{\nabla}} + \overrightarrow{\boldsymbol{\nabla}}$.

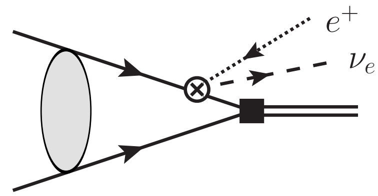

FIG. 1. Feynman diagram for pp fusion: solid lines nucleons, short-dashed line positron e^+, dashed line neutrino ν_e and double line deuteron. The blob represents Coulomb amplitude t_c, " ⊗ " a weak vertex, " ■ " a strong interaction vertex.

The Feynman diagrams in Fig. 1 provide the dominant p-wave contribution to pp fusion. A straightforward calculation shows that p-wave capture from the weak interaction vertex in Eq. (2) comes from the deuteron recoil momentum $\boldsymbol{k}$. Thus this p-wave contribution scales as kp/Q^2 compared to the LO s-wave amplitude in $\not\pi$EFT [11, 18]. We name this recoil contribution T_{FGT}. The weak interaction vertex generated by the $\overleftrightarrow{\boldsymbol{\nabla}}$ terms in Eq. (3) contribute even in the zero-recoil limit. Relative to the LO s-wave amplitude, it is suppressed by a factor of p/m and we name this relativistic contribution $\boldsymbol{T}_{\mathrm{rel}}$. The contribution from the $\mu_V\overline{\boldsymbol{\nabla}}$ term is suppressed by k^2 and we do not include it. Compared to the LO s-wave amplitude, at momentum $p \sim \gamma \sim Q$, the recoil contribution kp/Q^2 and the relativistic contribution p/m are similar scaling as $Q^3/\Lambda^3 \sim 0.04$. We use this estimate for $T_{\mathrm{FGT}} \sim |\boldsymbol{T}_{\mathrm{rel}}|$ that holds up to $p \lesssim \gamma$ to keep the EFT analysis simple. Empirically, at solar energies $E \lesssim 100$ keV, p/m is small but pk/γ^2 is smaller. Thus $\boldsymbol{T}_{\mathrm{rel}}$ contribution is larger making $p/m \lesssim 0.01$ to be a better estimate for the relative contribution of the p-wave amplitude. Furthermore, the cross section (and therefore the S-factor) can be decomposed into a partial wave expansion as $\sigma(E) = \sigma_0(E)+\sigma_1(E)+\ldots$, where the subscript l refers to the l-th pp partial wave. We threfore anticipate the p-wave cross section (and S-factor) to be smaller by a factor of $p^2/m^2 \lesssim 10^{-4}$ compared to the s-wave value. We include the NLO correction from the effective range ρ. Initial state p-wave strong interactions are suppressed by relative powers of Q^3/Λ^3. Higher order corrections to the weak interactions are suppressed by at least $Q^2/\Lambda^2 \sim 0.01$, and constitute a 10% uncertainty in the p-wave cross section.

b. **The p-wave cross section:** The p-wave amplitude is

$$i\mathcal{M}_1 = i8\frac{G_v}{\sqrt{2}}\epsilon_i^{d*}\, u_N(-\boldsymbol{p})\mathbb{P}_j u_N(\boldsymbol{p}) \times\Big\{ \left(l_+^0 T_{\rm FGT} - \boldsymbol{l}_+\cdot\boldsymbol{T}_{\rm rel}\right) {\rm Tr}\left[P_i\tau^-\mathbb{P}_j^\dagger\right] + g_A\left(l_+^k T_{\rm FGT} - l_+^0 T_{\rm rel}^k\right){\rm Tr}\left[P_i\sigma_k\tau^-\mathbb{P}_j^\dagger\right]\Big\}, \quad (4)$$

where $\mathbb{P}_j$ is the spintriplet-isotriplet projector, $i\sigma_2\sigma_j(1+\tau_3)/4$. The non-relativistic two component nucleon spinor fields $u_N(\boldsymbol{p})$ are normalized as $[u_N(\boldsymbol{p})]_\alpha[u_N^*(\boldsymbol{p})]_\beta = \delta_{\alpha\beta}$ when summed over polarizations. The amplitudes from the loop integrals are

$$T_{\rm FGT} = g_0\sqrt{Z_d}m\int\frac{d^3q}{(2\pi)^3}\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q})\frac{\boldsymbol{q}\cdot\boldsymbol{k}}{(\gamma^2+q^2)^2}\,, \quad (5)$$

and

$$\boldsymbol{T}_{\rm rel} = g_0\sqrt{Z_d}m\int\frac{d^3q}{(2\pi)^3}\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q})\frac{\boldsymbol{q}}{m}\frac{1}{\gamma^2+q^2}\,. \quad (6)$$

The solid angle integral of $\boldsymbol{q}\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q})$ picks out the vector direction $\boldsymbol{p}$ and constitutes the $l=1$ partial wave contribution. The c.m. deuteron momentum $\boldsymbol{k}$ is related to the positron/neutrino pair momenta $\boldsymbol{p}_{e,\nu}$ from momentum conservation as $\boldsymbol{k} = -(\boldsymbol{p}_e+\boldsymbol{p}_\nu)$. The expressions for $T_{\rm FGT}\propto e^{i\delta_1}$ and $\boldsymbol{T}_{\rm rel}\propto e^{i\delta_1}$ are derived further below. Since $T_{\rm FGT}\boldsymbol{T}_{\rm rel}^* = T_{\rm FGT}^*\boldsymbol{T}_{\rm rel}$, Eq. (4) gives

$$\overline{|\mathcal{M}_1|^2} = 8\left(\frac{G_v}{\sqrt{2}}\right)^2\Big\{(E_\nu E_e + \boldsymbol{p}_\nu\cdot\boldsymbol{p}_e)\times\left(3|T_{\rm FGT}|^2 + 2g_A^2\boldsymbol{T}_{\rm rel}\cdot\boldsymbol{T}_{\rm rel}^*\right) + 6\left(\boldsymbol{p}_\nu\cdot\boldsymbol{T}_{\rm rel}\right)\left(\boldsymbol{p}_e\cdot\boldsymbol{T}_{\rm rel}^*\right) + 3\left(E_\nu E_e - \boldsymbol{p}_\nu\cdot\boldsymbol{p}_e\right)\boldsymbol{T}_{\rm rel}\cdot\boldsymbol{T}_{\rm rel}^* - \left(6+4g_A^2\right)\left(E_\nu\boldsymbol{p}_e + E_e\boldsymbol{p}_\nu\right)\cdot\boldsymbol{T}_{\rm rel}T_{\rm FGT}^* + 2g_A^2\left(3E_\nu E_e - \boldsymbol{p}_\nu\cdot\boldsymbol{p}_e\right)|T_{\rm FGT}|^2\Big\}, \quad (7)$$

where we used the polarization sum over the leptonic currents $l_+^\mu l_+^{\nu\,\dagger}$.

The spin averaged cross section for non-relativistic fields is given by Fermi's Golden Rule as

$$\sigma_1(E) = \int\frac{{\rm d}^3p_e}{(2\pi)^3}\frac{{\rm d}^3p_\nu}{(2\pi)^3}\frac{1}{4E_eE_\nu}\frac{1}{v_{rel}}\overline{|\mathcal{M}_1|^2} \times 2\pi\,\delta\left(\delta m + E - \frac{k^2}{2M_d} - E_e - E_\nu\right)\,, \quad (8)$$

where $\delta m = 2m_p - M_d = m_p - m_n + \gamma^2 m/(m_pm_n)$. The integral can be reduced to 4-dimensions. The magnitude p_ν is constrained from the Dirac δ-function. We are free to choose the spin quantization axis ($\hat{z}$ axis) along $\boldsymbol{p}$ direction. Azimuthal symmetry of the total lepton momentum $\boldsymbol{p}_e+\boldsymbol{p}_\nu = -\boldsymbol{k}$ implies dependence only on the difference in the azimuthal angle $\phi = \phi_e - \phi_\nu$ of the pair $\boldsymbol{p}_{e,\nu}$. The integral in Eq. (8) can then be written as

$$\sigma_1(E) = \frac{1}{(2\pi)^4}\int_0^{p_e^{\rm max}}{\rm d}p_e p_e^2\int_{-1}^1{\rm d}x_e\int_{-1}^1{\rm d}x_\nu \times\int_0^{2\pi}{\rm d}\phi\frac{p_\nu^2}{|1+\frac{p_\nu}{M_d}+\frac{p_e}{M_d}x_{e\nu}|}\frac{1}{4E_eE_\nu}\frac{1}{v_{rel}}\overline{|\mathcal{M}_1|^2}\,, \quad (9)$$

where $x_{e,\nu} = \hat{\boldsymbol{p}}\cdot\hat{\boldsymbol{p}}_{e,\nu}$ and $x_{e\nu} = \hat{\boldsymbol{p}}_e\cdot\hat{\boldsymbol{p}}_\nu = x_ex_\nu + \sqrt{1-x_e^2}\sqrt{1-x_\nu^2}\cos\phi$. The neutrino momentum magnitude is given by

$$p_\nu = -M_d - p_ex_{e\nu} + [(M_d+p_ex_{e\nu})^2 + 2M_d(2m_p - M_d + E - E_e) - p_e^2]^{1/2}\,, \quad (10)$$

and the maximal positron momentum is

$$p_e^{\rm max} = \left\{\left(2 - \frac{2m_p+E}{M_d}\right)\times\left([2m_p - M_d + E]^2 - m_e^2\right)\right\}^{1/2}\,. \quad (11)$$

c. **Results:** The cross section $\sigma_1(E)$ in Eq. (9) is evaluated by numerical integration using analytic expressions for $T_{\rm FGT}$ and $\boldsymbol{T}_{\rm rel}$. These can be derived from the coordinate space wavefunction

$$\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{r}) \equiv \int\frac{d^3q}{(2\pi)^3}e^{i\mathbf{q}\cdot\boldsymbol{r}}\chi_{\boldsymbol{p}}^{(+)}(\boldsymbol{q}) = \sum_{l=0}^\infty(2l+1)i^le^{i\delta_l}P_l(\hat{\boldsymbol{r}}\cdot\hat{\boldsymbol{p}})\frac{F_l(\eta_p;pr)}{pr}\,, \quad (12)$$

where $\delta_l = \arg\Gamma(l+1+i\eta_p)$ is the Coulomb phase shift and

$$F_l(\eta_p;\rho) = \frac{2^le^{-\pi\eta_p/2}|\Gamma(l+1+i\eta_p)|}{\Gamma(2l+2)}\rho^{l+1}e^{-i\rho}\times M(l+1-i\eta_p, 2l+2; 2i\rho)\,, \quad (13)$$

is the regular Coulomb wave function expressed in terms of the conventionally defined Kummer's function

$M(a,b,z)$. Equation (5) can then be written as

$$\begin{aligned} T_{\rm FGT} =& \frac{1}{6} g_0 \sqrt{Z_d} m(\boldsymbol{p}\cdot\boldsymbol{k}) e^{i\delta_1} e^{-\pi\eta_p/2} |\Gamma(2+i\eta_p)| \\ & \times \int_0^\infty dr r^3 e^{-(\gamma+ip)r} M(2-i\eta_p, 4, i2pr) \\ =& -\sqrt{\frac{8\pi\gamma}{1-\rho\gamma}} e^{i\delta_1} e^{-\pi\eta_p/2} (\boldsymbol{p}_\nu + \boldsymbol{p}_e)\cdot\boldsymbol{p} \\ & \times |\Gamma(2+i\eta_p)| \frac{1}{(\gamma^2+p^2)^2} e^{2\eta_p \arctan p/\gamma}, \end{aligned} \quad (14)$$

where we have used the NLO relation $g_0\sqrt{Z_d}m = \sqrt{8\pi\gamma/(1-\rho\gamma)}$. Similarly, Eq. (6) can be written as

$$\begin{aligned} \boldsymbol{T}_{\rm rel} =& \frac{1}{3} g_0 \sqrt{Z_d} m e^{i\delta_1} e^{-\pi\eta_p/2} |\Gamma(2+i\eta_p)| \frac{\boldsymbol{p}}{m} \\ & \times \int_0^\infty dr r(1+\gamma r) e^{-(\gamma+ip)r} M(2-i\eta_p, 4, i2pr) \\ =& \sqrt{\frac{8\pi\gamma}{1-\rho\gamma}} e^{i\delta_1} e^{-\pi\eta_p/2} \frac{\mathbf{p}}{m} |\Gamma(2+i\eta_p)| \\ & \times \frac{1}{2p^2+2p^2\eta_p^2} \\ & \left[1 + \frac{p^2+2p\eta\gamma-\gamma^2}{\gamma^2+p^2} e^{2\eta_p \arctan p/\gamma}\right]. \end{aligned} \quad (15)$$

In Fig. 2 we show the result for the S-factor $S_1(E)$. We perform a polynomial fit to the results shown in Fig. 2 and use it to extrapolate the S-factor and its derivative to zero energy. We obtain

$$\begin{aligned} S_1(0) &= (2.47 \pm 0.25 \pm 0.01) \times 10^{-28} \text{ MeV fm}^2, \\ S_1'(0) &= (2.16 \pm 0.22 \pm 0.01) \times 10^{-26} \text{ fm}^2, \end{aligned} \quad (16)$$

where the first errors indicate EFT uncertainties and the second ones are numerical errors from polynomial fits to $S(E)$.

Our result for $S_1(0)$ agrees with the tentative estimates we made earlier based on the power counting, but does not agree with the value of $S_1(0) = 2.0\times10^{-25}$ MeV fm^2 claimed in Ref. [7]. In fact, the p-wave contribution is much smaller than the $\sim$ 1% contribution obtained by Ref. [7] in the entire $0-100$ keV energy region in which they perform their calculations. We, therefore, disagree with the findings of Marcucci *et al.* in Ref. [7] and claim that the p-wave contributions need not be considered in the calculation of the *pp* S-factor at astrophysically relevant energies since these are much smaller than the precision of the s-wave calculation [see Ref. [1] for a state-of-the-art uncertainty analysis]. Furthermore, Refs. [1, 19] have found that basis truncation errors accounted for a reduction in Ref. [7]'s s-wave S-factor by about 0.7 %. Since Marcucci *et al.* only addressed the error in the p-wave calculation in their Erratum [8], their value for the corrected S-factor, with combined s- and p-waves, is still incorrect, and does not agree with value $S(0) = (4.047^{+0.024}_{-0.032})\times10^{-23}$ MeV fm^2 MeV fm^2 calulated by Ref. [1] within the uncertainty band, which re-

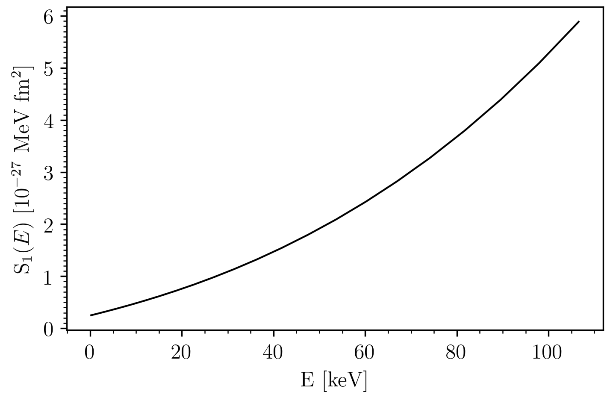

FIG. 2. The p-wave S-factor $S_1(E)$.

mains unmodified upon inclusion of the p-wave contribution calculated in this work [2]. Finally, we emphasize that, even though the p-wave numbers we calculated are negligible given the large uncertainty in the s-wave value, the correction we make to Ref. [7]'s results is at least as important as all sources of uncertainty combined.

d. **Conclusion:** We calculated for the first time the contribution of p-wave *pp* configuration to the fusion rate in $\not\pi$EFT. This analysis was motivated by a recent calculation with chiral potentials that suggested that the leading p-wave contributions are comparable to the next-to-next-to-leading s-wave contributions.

We determined the dominant Feynman diagrams contributing to the p-wave S-factor and calculated their contribution at low energies. The NLO calculation, include the recoil contributions from the Gamow-Teller and Fermi operators, as well as the relativistic $1/m$ suppressed weak interaction operators. We found that the p-wave contribution to the *pp* fusion S-factor is smaller than the value obtained by Ref. [7] by several orders of magnitude, and that the effect of p-wave fusion is therefore negligible for a high-precision determination.

Our analytic results for the p-wave fusion matrix element can serve as a benchmark for numerical calculations of chiral effective theory matrix elements. They are expressed in terms of the weak couplings constants g_A and G_F and two observables from the strong sector, the deuteron binding energy and wave function renormalization and therefore do not suffer from any short-distance model ambiguities. We note that the input observables needed to predict the p-wave fusion rate are

[2] The relationship between the chiral EFT counterterms c_D and d_R have since been updated [20]. This correction makes a negligible modification in the S-factor value compared to the uncertainty band.

s-wave observables. Our results illustrate furthemore the importance of calculating astrophysically relevant three-nucleon capture reactions in pionless effective field theory to reduce currently accepted uncertainties and to explore the importance of recoil corrections in these nuclei.

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