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Is it possible to determine the S–factor of the hep process from a laboratory experiment?

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Abstract

We discuss the problem of solar hep neutrinos originating from the reaction $p + {}^3\mathrm{He} \to {}^4\mathrm{He} + e^+ + \nu_e$ and obtain a relation between the astrophysical S–factor of the hep process and the cross section of the process $e^- + {}^4\mathrm{He} \to {}^3\mathrm{H} + n + \nu_e$ near threshold. The relation is based on the isotopic invariance of strong interactions. The measurement of the latter cross section would allow to obtain experimental information on $S(hep)$, the value of which, at the moment, is known only from theoretical calculations.

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I. INTRODUCTION

The reaction

$$p + {}^3\mathrm{He} \to {}^4\mathrm{He} + e^+ + \nu_e\,, \tag{1.1}$$

the so–called hep reaction, is the source of solar neutrinos with the highest energies (up to 18.8 MeV). It is expected that the cross section of this process at solar energies ($\lesssim$ 30 keV) is extremely small and cannot be measured in laboratory conditions [1]. The most detailed calculations of the cross section of the hep process were done in Refs. [2,3].

The process (1.1) proceeds via a Gamov–Teller transition (a change in isospin from 1 to 0 is involved). In ref. [2,3] there are two reasons for the strong suppression of the hep cross section. The first one is due to the fact that the matrix element of the process vanishes in the allowed approximation if only the main s–state components of ^{4}He and ^{3}He wave functions are taken into account [4]. The second reason lies in strong cancellations between different matrix elements of the weak nuclear current: it was found in Ref. [3] that the contribution to the matrix element of the hep process from the diagrams with π and ρ meson exchange currents (in which the transition of a nucleon into the Δ isobar is also taken into account) is comparable in modulus with the contribution of the one–body current, but has opposite sign. As a result there is a cancellation of the contributions of these two terms and the calculated cross section is about 5 times smaller than the cross section predicted by the one–nucleon term only.

The result of the calculation also depends on the two-body nuclear potential: by using two different, typical NN–potentials, for the S–factor of hep process in Ref. [3] the following values were obtained:

$$S_1(hep) = 1.44 \times 10^{-20}\ \mathrm{keV\ b}\,,$$
$$S_2(hep) = 3.14 \times 10^{-20}\ \mathrm{keV\ b}\,.$$

The average between these two values,

$$S_0(hep) = 2.3 \times 10^{-20}\ \mathrm{keV\ b}\,, \tag{1.2}$$

is used in the Standard Solar Model (SSM) [5]. If the value of the astrophysical S–factor of the hep process is given by Eq.(1.2) then the total flux of hep–neutrinos [6],

$$\Phi(hep) = 2.1 \times 10^3\ \mathrm{cm}^{-2}\,\mathrm{s}^{-1}, \tag{1.3}$$

is more than three orders of magnitude smaller than the flux of ^{8}B neutrinos,

$$\Phi({}^8\mathrm{B}) = 5.15 \times \left(1.00^{+0.19}_{-0.14}\right) \times 10^6\ \mathrm{cm}^{-2}\,\mathrm{s}^{-1},$$

and hep neutrinos give a negligible contribution to the event rates observed in solar neutrino experiments. Let us stress, however, that the value (1.2) of $S_0(hep)$ is the result of a very complicated and model-dependent calculations.

The recent interest in hep neutrinos was triggered by the results of the Super-Kamiokande experiment [7–9] in which the spectrum of recoil electrons in the solar-neutrino-induced process

$$\nu + e^- \to \nu + e^-$$

was measured. The spectrum of neutrinos from the decay ${}^8\mathrm{B} \to {}^8\mathrm{Be} + e^+ + \nu_e$ is determined by weak interactions and well known. The recoil electron spectrum measured in the Super-Kamiokande experiment is in agreement with the predicted spectrum in the whole energy range starting from 5.5 MeV, with the exception of the highest energy region, in which two data points with large errors (from the bins 13.5 – 14 MeV and 14 – 20 MeV) are above the prediction [7,8].

The Super-Kamiokande recoil electron spectrum can be fitted with the assumption that there is no distortion of the spectrum ($\chi^2 = 24.3$ at 17 d.o.f.) [8,9]. In Ref. [1] attention was payed, however, to the fact that the high energy points in the Super-Kamiokande data could be due to the contribution of hep neutrinos [10]. If one considers $S(hep)$ as a free parameter, then from the fit of the data (504 days of Super-Kamiokande) in Ref. [1] in the hypothesis of no oscillations $S(hep)/S_0(hep) = 26$ was obtained. In a more recent fit [8] (825 days of Super-Kamiokande) it was found $S(hep)/S_0(hep) = 16$ ($\chi^2 = 19.5$ at 16 d.o.f.).

Let us notice that the distortion of the recoil electron spectrum could also be due to the MSW effect or to vacuum oscillations (VO). The largest enhancement of the high energy part of the spectrum is expected for the VO solution. By fitting the data with $S(hep) = S_0(hep)$ in the case of VO solution it was found $\sin^2 2\theta = 0.79$, $\Delta m^2 = 4.3 \times 10^{-10}$ eV^2 ($\chi^2 = 44.1$ at 35 d.o.f.) [8] (see also Refs. [1,11,12]).

The investigation of the problem of hep neutrinos in the solar neutrino Super-Kamiokande experiment will be continued. The results of a new measurement of the spectrum of solar ν's in the region $E \geq 5$ MeV will be soon available from the SNO experiment [13]. In this experiment the spectrum of solar ν_e's will be determined from the measurement of the electron spectrum in the process $\nu_e + d \to e^- + p + p$.

One of the central theoretical problems connected with hep neutrinos is the astrophysical S–factor of the hep process [14]. "The most important unsolved problem in theoretical nuclear physics related to solar neutrinos is the range of values allowed by fundamental physics for the hep production cross section"(Bahcall [15]). In this letter we consider the possibility to determine $S(hep)$ from experimental data. We will obtain here a relation between $S(hep)$ and the total cross section of the process

$$e^- + {}^4\mathrm{He} \to {}^3\mathrm{H} + n + \nu_e \tag{1.4}$$

near threshold. The relation we obtain is based on the isotopic invariance of the strong interactions (we neglect the Coulomb interaction in the region of nuclear forces). From the existing nuclear data it follows that the violation of isotopic invariance for light nuclei cannot be larger than $10 - 20\%$.*

From the point of view of a possible investigation of the process (1.4) at small energies, we would like to stress two points:

1. The cross section of the process (1.4) does not contain the Coulomb penetration factor, which suppresses at small energies the cross sections of processes with initial particles having charges of equal sign.

*This estimate follows from nuclear mass differences, mirror nuclei spectra and so on.

2. There exist electron accelerators (microtrons) which allow to obtain high intensity electron beams in the range of energies which are appropriate for the investigation of the process (1.4).

II. THE RELATION BETWEEN $S(hep)$ AND $\sigma(e^- \, {}^4\mathrm{He} \to {}^3\mathrm{H}\, n\, \nu_e)$

Let us start by considering the process (1.1) at small solar energies ($\lesssim$ 30 keV). In Ref. [16] the general arguments are given that at small energies the cross sections of the reactions with charged initial particles have the form

$$\sigma(E) = \frac{1}{E} e^{-2\pi\eta} S(E) \,. \tag{2.1}$$

Here E is the kinetic energy of the initial particles in the C.M. system and

$$\eta = \frac{Z_1 Z_2 e^2}{v} \,, \tag{2.2}$$

where $Z_1 e$, $Z_2 e$ are the charges of the initial particles, $v = \sqrt{2E/\mu}$ is their relative velocity and μ is the reduced mass. In the expression (2.1),

$$P \equiv e^{-2\pi\eta} \tag{2.3}$$

is the probability of penetration of the incident particle through the Coulomb barrier [17] and the factor $S(E)$ is determined mostly by strong interactions. If there are no resonances at small energies, the function $S(E)$ depends very weakly on the energy E.

The relation (2.1) with $S \simeq$ const. allows to describe the existing low energy data and is used for the extrapolation of laboratory data to the energy region which is relevant for solar reactions (see, for example, Ref. [18]). Our further considerations will be based on this relation.

The standard weak interaction Hamiltonian density is given by

$$\mathcal{H}_I = \frac{G_F}{\sqrt{2}} \bar{\nu}_e \gamma^\alpha (1 - \gamma_5) e j_\alpha + \text{h.c.} \,, \tag{2.4}$$

where the hadronic $V - A$ current $j_\alpha = j^1_\alpha - i j^2_\alpha \equiv j^{1-i2}_\alpha$ is the "minus" component of the isovector j^a_α $(a = 1, 2, 3)$.

For the matrix element of the process (1.1) we have

$$\begin{aligned} \langle f|S|i\rangle = {} & -i \frac{G_F}{\sqrt{2}} \frac{1}{(2\pi)^3} \frac{1}{\sqrt{4 k_0 k'_0}} \ell^\alpha(k, k') \\ & \times \int d^4x \, e^{iq\cdot x} \langle {}^4\mathrm{He}| T\left(j_\alpha(x) e^{-i\int d^4y \mathcal{H}^0_I(y)} \right) |p\, {}^3\mathrm{He}\rangle \,. \end{aligned} \tag{2.5}$$

Here $\ell^\alpha(k, k') = \bar{u}(k')\gamma^\alpha(1 - \gamma_5) v(k)$ is the matrix element of the weak leptonic current, $q = k + k'$ (k and k' being the momenta of the e^+ and ν_e, respectively) and $\mathcal{H}^0_I = \mathcal{H}^h_I + \mathcal{H}^{em}_I$ is the Hamiltonian density of strong ($\mathcal{H}^h_I$) and electromagnetic ($\mathcal{H}^{em}_I$) interactions.

Let us first consider only that part of the matrix element of the process (1.1) which is determined by the strong interactions and gives the major contribution to the S–factor. Neglecting the Coulomb interaction in the region of nuclear forces, we have for the hadronic part of the matrix element (2.5)

$$\langle f|S|i\rangle = -i\frac{G_F}{\sqrt{2}}\frac{1}{(2\pi)^3}\frac{1}{\sqrt{4k_0k_0'}}\ell^{\alpha}(k,k')\langle{}^4\mathrm{He}|J_{\alpha}^{(-)}(0)|p\,{}^3\mathrm{He}\rangle(2\pi)^4\delta(P'-P)\,, \tag{2.6}$$

where $J_{\alpha}^{(-)}(x)\equiv J_{\alpha}^{1-i2}(x)$ is the hadronic weak $V-A$ current in the *Heisenberg representation* and P (P') is the total four-momentum of the initial (final) states. It is evident that $\langle{}^4\mathrm{He}|J_{\alpha}^{(-)}(0)|p\,{}^3\mathrm{He}\rangle$ includes *all* possible contributions coming from strong interactions (for example, in addition to the one–body nucleonic current, also two–body exchange currents, effects of the Δ isobar and so on).

Using the charge symmetry of strong interactions we have

$$\begin{aligned}&\langle{}^4\mathrm{He}|J_{\alpha}^{1-i2}|p\,{}^3\mathrm{He}\rangle = \langle{}^4\mathrm{He}|\mathcal{U}^{-1}\mathcal{U}J_{\alpha}^{1-i2}\mathcal{U}^{-1}\mathcal{U}|p\,{}^3\mathrm{He}\rangle\\ &= -\langle{}^4\mathrm{He}|J_{\alpha}^{1+i2}|n\,{}^3\mathrm{H}\rangle = -\langle n\,{}^3\mathrm{H}|J_{\alpha}^{1-i2}|{}^4\mathrm{He}\rangle^{*}.\end{aligned} \tag{2.7}$$

Here $\mathcal{U}=\exp\{i\pi T_2\}$ is the unitary operator of rotation by an angle π around the second axis in isospace. In Eq.(2.7) we took into account that

$$\begin{aligned}\mathcal{U}J_{\alpha}^{1-i2}\mathcal{U}^{-1} &= -J_{\alpha}^{1+i2}\,,\\ \mathcal{U}\,|p\,{}^3\mathrm{He}\rangle &= |n\,{}^3\mathrm{H}\rangle\,,\\ \mathcal{U}\,|{}^4\mathrm{He}\rangle &= |{}^4\mathrm{He}\rangle\,.\end{aligned} \tag{2.8}$$

Thus the hadronic part of the matrix elements of the process $e^-+{}^4\mathrm{He}\longrightarrow{}^3\mathrm{H}+n+\nu_e$ is connected with the matrix element of the hadronic weak current for the hep process by the simple charge symmetry relation (2.7). With the help of Eq.(2.7) we will obtain a relation which connects the S–factor of the hep process (1.1) with the cross section of the process (1.4).

Let us continue considering the hep process. In the region of small energies we are interested in, only the contribution of the s–wave of the initial p–^{3}He system is relevant. Taking into account the Coulomb interaction between the initial p and ^{3}He, for the total cross section of the hep process in the center of mass system we have

$$\begin{aligned}\sigma(hep)=&\frac{(2\pi)^4}{v}\,\frac{G_F^2}{2}\,\frac{1}{4}\sum_{\mathrm{spins}}\int\frac{d^3k}{2k_0}\int\frac{d^3k'}{2k_0'}\\ &\times\left|\ell^{\alpha}\langle{}^4\mathrm{He}|J_{\alpha}^{(-)}|p\,{}^3\mathrm{He}\rangle\right|^2\delta(k_0+k_0'-\Delta)\,\frac{|\psi_{\vec p}^{(+)}(0)|^2}{|\psi_{\vec p}|^2}\,.\end{aligned} \tag{2.9}$$

Here $\Delta=m_p+m_{^3\mathrm{He}}-m_{^4\mathrm{He}}=19.284$ MeV (see, e.g., Ref. [19] for the values of the nuclear masses), $v=\sqrt{2E/\mu}$ is the relative velocity, E being the initial energy and μ the reduced mass of the $p-{}^3$He system, $\psi_{\vec p}^{(+)}(0)$ is the Coulomb wave function of the initial $p-{}^3$He system at $r=0$ and $\psi_{\vec p}(\vec r)=\exp(i\vec p\cdot\vec r)/(2\pi)^{3/2}$. Notice that in (2.9) we have neglected the small recoil energy of ^{4}He and that the factor 1/4 is due to averaging over the spin states of the initial particles.

For the hep process the non-relativistic expression

$$\frac{|\psi_{\vec{p}}^{(+)}(0)|^2}{|\psi_{\vec{p}}|^2} = \frac{2\pi\eta}{e^{2\pi\eta}-1} \tag{2.10}$$

holds, where

$$\eta = \frac{2e^2}{v} \simeq 8.66 \frac{1}{\sqrt{E[\mathrm{keV}]}} \, . \tag{2.11}$$

In the region of small energies $2\pi\eta \gg 1$, we have, for the Coulomb factor (2.10),

$$\frac{|\psi_{\vec{p}}^{(+)}(0)|^2}{|\psi_{\vec{p}}|^2} \simeq 2\pi\eta \, e^{-2\pi\eta} \, . \tag{2.12}$$

For the small energies we are interested in, we have $|\vec{q}\,|R \ll 1$, R being the radius of nuclear forces. Moreover, the parity of the initial and final nuclear states is the same and in the matrix element the linear term in the expansion of $\exp\{i\vec{q}\cdot\vec{x}\}$ vanishes. Thus, in Eq.(2.5) we can put $e^{-i\vec{q}\cdot\vec{x}} \simeq 1$ (allowed transition). In this approximation, the matrix element of the hadronic current does not depend on $\vec{q}$.† Moreover, in that region the matrix element $\langle {}^4\mathrm{He}|J_\alpha|p\,{}^3\mathrm{H}\rangle$ does not depend on the relative momentum $\vec{p}$ of the initial particles either. Since only the axial vector current contributes to the matrix element (Gamow–Teller transition) we obtain

$$\frac{1}{4k_0k_0'} \sum_{\mathrm{spins}} \left|\ell^\alpha \langle {}^4\mathrm{He}|J_\alpha^{(-)}|p\,{}^3\mathrm{He}\rangle\right|^2 =$$

$$\frac{1}{4k_0k_0'} \sum_{\mathrm{spins}} \ell^i \ell^{k*} \delta_{ik} \frac{1}{3} \left|\langle {}^4\mathrm{He}|\vec{J}^{(-)}|p\,{}^3\mathrm{He}\rangle\right|^2 = 2\left(1 - \frac{1}{3}\frac{\vec{k}\cdot\vec{k}'}{k_0k_0'}\right) \sum_{\mathrm{spins}} \left|\langle {}^4\mathrm{He}|\vec{J}^{(-)}|p\,{}^3\mathrm{He}\rangle\right|^2 \, . \tag{2.13}$$

It is obvious that the second term in the last equality does not give a contribution to the total cross section. Taking into account the Coulomb interaction of the final e^+ and ^{4}He, from Eqs.(2.9), (2.2) and (2.13) we obtain the total cross section of the hep process

$$\sigma(hep) = \frac{(2\pi)^7}{E} \, G_F^2 e^2 m_e^5 \mu \sum_{\mathrm{spins}} \left|\langle {}^4\mathrm{He}|\vec{J}^{(-)}|p\,{}^3\mathrm{He}\rangle\right|^2 f(\varepsilon_0)\, e^{-4\pi e^2/v} \, . \tag{2.14}$$

Here $f(\varepsilon_0)$ is given by

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† This independence of the matrix element upon $\vec{q}$ is confirmed by the detailed calculations made in Refs. [2,3], in which not only the one–nucleon term but also two–body terms due to the exchange of π and ρ mesons were taken into account. Nevertheless we think that further investigation of the q–dependence of the nuclear matrix element is an important and interesting issue.

$$
f(\varepsilon_0) = \int_1^{\varepsilon_0} F(-2,\varepsilon)(\varepsilon_0-\varepsilon)^2\sqrt{\varepsilon^2-1}\,\varepsilon\, d\varepsilon\,, \tag{2.15}
$$

where $\varepsilon = k_0/m_e$, $\varepsilon_0 = \Delta/m_e$ and $F(-2,\varepsilon)$ is the Fermi function (ratio of the modulus squared of the positron wave function in the Coulomb field of the final nucleus, calculated at $r=R$, to the modulus squared of the plane wave). Tables of the Fermi function are given in Ref. [20]. For the hep reaction we have $\varepsilon_0 \simeq 37.7$. Using the approximation $F(-2,\varepsilon)\simeq 1$ valid for small Z and large positron energies, we obtain [20] $f(\varepsilon_0)\simeq \varepsilon_0^5/30 \simeq 2.55\times 10^6$.

Finally, from Eqs.(2.1) and (2.14), we derive the following expression for the S–factor for the hep process:

$$
S(hep) = (2\pi)^7 G_F^2\, e^2 \mu\, m_e^5 f(\varepsilon_0) \sum_{\text{spins}} \left| \langle {}^4\mathrm{He}|\vec{J}^{(-)}|p\,{}^3\mathrm{He}\rangle \right|^2 . \tag{2.16}
$$

Let us consider now the process (1.4). Neglecting the electromagnetic interaction of hadrons, for the matrix element of the process we get an expression similar to Eq.(2.6):

$$
\langle f|S|i\rangle = -i\frac{G_F}{\sqrt{2}}\frac{1}{(2\pi)^3}\frac{1}{\sqrt{4k_0k_0'}}\ell^\alpha(k,k')\langle n\,{}^3\mathrm{H}|J_\alpha^{(-)}(0)|{}^4\mathrm{He}\rangle(2\pi)^4\delta(P'-P)\,, \tag{2.17}
$$

where $\ell^\alpha(k,k') = \bar{u}(k')\gamma^\alpha(1-\gamma_5)u(k)$ and k and k' are the four–momenta of e^- and ν_e, respectively.

The threshold for the process (1.4) is given by (see, e.g., Ref. [19] for a table of nuclear masses)

$$
E_{\text{th}} = \frac{(m_n + m_{{}^3\mathrm{H}})^2 - m_{{}^4\mathrm{He}}^2}{2m_{{}^4\mathrm{He}}} \simeq 21.167 \text{ MeV}. \tag{2.18}
$$

We will consider the process (1.4) at electron energies close to the threshold energy E_{th}. For the total cross section we have

$$
\begin{aligned}
&\sigma(e^-\,{}^4\mathrm{He}\to{}^3\mathrm{H}\,n\,\nu_e) = \\
&\frac{1}{v_e}(2\pi)^4\frac{G_F^2}{2}F(2,E_e/m_e)\frac{1}{2E_e}\int\frac{d^3k'}{2k_0'}\int d^3p\int d^3p'\frac{1}{2}\sum_{\text{spins}}\left|\langle n\,{}^3\mathrm{H}|J_\alpha^{(-)}|{}^4\mathrm{He}\rangle\ell^\alpha\right|^2\delta(P'-P)\,,
\end{aligned} \tag{2.19}
$$

where $E_e \equiv k_0$ is the electron energy, v_e the velocity of the electron, p' and p are the total and relative momenta, respectively, of the $n-{}^3$H system and the Fermi function $F(2,E_e/m_e)$ takes into account the Coulomb interaction between the initial e^- and ^{4}He. Within the same approximations we have used in the derivation of the relation (2.16), the cross section for the process (1.4) turns out to be

$$
\begin{aligned}
&\sigma(e^-\,{}^4\mathrm{He}\to{}^3\mathrm{H}\,n\,\nu_e) = \\
&\quad\frac{32}{105}(2\pi)^6 G_F^2\sum_{\text{spins}}\left|\langle n\,{}^3\mathrm{H}|\vec{J}^{(-)}|{}^4\mathrm{He}\rangle\right|^2(E_e-E_{\text{th}})^{7/2}\mu\sqrt{2\mu}\,F(2,E_e/m_e)\,,
\end{aligned} \tag{2.20}
$$

where μ is the reduced mass of the $n-{}^3$H system (704.1 MeV) which we identify numerically with that of the $p-{}^3$He system (703.3 MeV), in agreement with our assumption of isotopic invariance of the strong interactions. Using now the isotopic relation (2.7), which connects the hadronic parts of the S–matrix elements of the processes (1.1) and (1.4), we obtain the following relation between the total cross section of the process (1.4) and the astrophysical S–factor of the hep–process:

$$\sigma(e^- \, {}^4\mathrm{He} \to {}^3\mathrm{H}\, n\, \nu_e) = \frac{32}{105}\frac{1}{(2\pi)e^2}\sqrt{\frac{2\mu}{m_e}}\,\frac{F(2,E_e/m_e)}{f(\varepsilon_0)}\left(\frac{E_e - E_{\rm th}}{m_e}\right)^{7/2}\frac{S(hep)}{m_e}\,. \tag{2.21}$$

For the case of electron energies $E_e > 20$ MeV, we can set the Fermi function equal to one. Then we get for the cross section of the reaction (1.4)

$$\sigma(e^- \, {}^4\mathrm{He} \to {}^3\mathrm{H}\, n\, \nu_e) \simeq 0.62\times 10^{-50}\ \mathrm{cm}^2 \times \left(\frac{E_e - E_{\rm th}}{m_e}\right)^{7/2}\frac{S(hep)}{S_0(hep)}\,. \tag{2.22}$$

The relation (2.21) allows to determine the astrophysical S–factor of the hep process $p + {}^3\mathrm{He} \to {}^4\mathrm{He} + e^+ + \nu_e$ directly from the measurements of the cross section of the process $e^- + {}^4\mathrm{He} \to {}^3\mathrm{H} + n + \nu_e$ near threshold.‡ If, for example, $S(hep) = 20 S_0(hep)$, a representative value among the ones indicated by the analysis of the Super–Kamiokande data [1,9,11,12], then at $E_e - E_{\rm th} \simeq 10$ MeV, for the cross section of the process (1.4) we get the value

$$\sigma(e^- \, {}^4\mathrm{He} \to {}^3\mathrm{H}\, n\, \nu_e) \simeq 4.0\times 10^{-45}\ \mathrm{cm}^2.$$

Obviously, the cross section of the process (1.4) is small (weak interactions and small energies). However, taking into account the importance of obtaining direct experimental information on $S(hep)$, it is worthwhile from our point of view to consider the possibility of performing such a measurement. One can use the advantage of high intensity beams of low energy electrons from microtrons and the possibility to detect the process (1.4) by radiochemical or mass-spectrometric methods.

III. CONCLUSIONS

The problem of hep neutrinos is one of the most important issues in solar neutrino physics and in solar neutrino oscillations. The direct measurement of the cross section of the process (1.1) at solar energies does not seem to be possible with the present techniques and the calculation of the cross section of the hep process is a very complicated problem.

In this paper we have obtained a relation between the astrophysical S–factor of the hep process and the total cross section of the process (1.4) near the threshold ($E_{\rm th} \simeq$ 21.167 MeV). The relation is based on the isotopic invariance of strong interactions. The measurement of the cross section of this process near threshold would allow to determine the major hadronic part of $S(hep)$. Though the smallness of the cross section (2.22) precludes

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‡Note that the kinetic energy of the $n - {}^3\mathrm{H}$ system in its C.M. system, which corresponds to the kinetic C.M. system energy of $p - {}^3\mathrm{He}$ in the hep reaction, is given by $E(n\,{}^3\mathrm{H}) \simeq E_e - E_{\rm th} - k'_0$, where k'_0 is the final neutrino energy.

going so close to the threshold such that $E_e - E_{\rm th}$ is of the order of the temperature in the solar core, measuring the process (1.4) at the more realistic energies $E_e - E_{\rm th} \sim 10$ MeV might allow to extrapolate $S(hep)$ to smaller energies. Therefore, we believe that it is worthwhile to consider the possibility of measuring the cross section (2.22) at microtron facilities.

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