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ARTICLES

Reaction cross sections for $\nu\,^{13}\mathrm{C}\rightarrow e^-\,^{13}\mathrm{N}$ and $\nu\,^{13}\mathrm{C}\rightarrow\nu'\,^{13}\mathrm{C}^*$ for low energy neutrinos

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Cross sections for $\nu+{}^{13}\mathrm{C}$ reactions are calculated both for charged- and neutral-current reactions in order to estimate the efficiency of a ^{13}C target as a solar neutrino detector. The relevant transition matrix elements are obtained using the semiphenomenological effective-operator approach for p-shell nuclei.

I. INTRODUCTION

The solar neutrino problem is one of the outstanding issues in particle physics and astrophysics.[1-3] A number of ideas have been proposed to solve this problem[2] In order to distinguish among the possibilities, it is most desirable to carry out another solar neutrino experiment that has characteristics different from the original ^{37}Cl experiment.[1,2] In particular, if one could measure the neutral-current-induced reaction, it would unambiguously settle the long-standing issue of whether the deficit of the neutrino capture rate in the ^{37}Cl detector is due to the neutrino oscillation.[4-7]

In a recent communication[8] we suggested the possible advantage of a ^{13}C-enriched scintillation counter as a solar neutrino detector. The point is that the replacement of ^{12}C in a scintillation detector with its isotope ^{13}C not only makes the detector much more sensitive to charged-current reactions, but also turns it into a highly efficient detector for neutral-current reactions. It was also pointed[8] out that even the natural abundance of ^{13}C may be large enough to detect neutral-current reactions in a large volume scintillator. Reasonably reliable estimates of the cross sections for the relevant neutrino-nucleus reactions are a prerequisite for deriving meaningful physical conclusions. However, since its main purpose was to discuss the basic idea in the previous publication,[8] the neutrino capture cross sections presented therein were obtained in a rather crude method. In this paper we wish to improve upon these previous estimates. We will adopt here the semiphenomenological effective-operator approach,[9,10] which provides probably the most suitable framework for this purpose. In fact, this approach has proven to be successful in reproducing a great majority of the observed Gamow-Teller and $M1$ transition rates over a wide range in the periodic table.[9,10]

II. CROSS SECTIONS OF $\nu+{}^{13}\mathrm{C}$ REACTIONS

The reactions of our concern are the charged-current-induced reaction

$$^{13}\mathrm{C}+\nu_e\rightarrow{}^{13}\mathrm{N}+e^- \,, \tag{1}$$

and the neutral-current-induced reaction

$$^{13}\mathrm{C}+\nu_e\rightarrow{}^{13}\mathrm{C}^*+\nu_e \,. \tag{2}$$

Figure 1 illustrates the low-lying nuclear levels[11] relevant to this argument. We are excluding from our consideration negative-parity states with $J\geq\frac{5}{2}$ as well as positive-parity states. The contribution of forbidden transitions leading to these states is negligibly small for the solar neutrino energy region. The effects of forbidden-type transition operators that interfere with the allowed-type operator, as examined by Bahcall and Holstein,[12] are also quite small for a light nucleus, unless the allowed-type contribution is substantially suppressed for some reason. (The change caused by the inclusion of this effect is, in fact, negligible compared with other ambiguities contained in this treatment; see the following.)

Let us first discuss the charged-current reaction. The cross section σ for the reaction $\nu_e+|A\rangle\rightarrow e^-+|B\rangle$ is given by

$$\sigma=(2\pi)^4\int d^3p_e\delta(E_e+E_B-E_\nu-E_A)$$
$$\times|\langle e^-(p_e);B|H_{\mathrm{eff}}|\nu(p_\nu);A\rangle|^2 \,, \tag{3}$$

where E_ν (E_e) is the neutrino (electron) energy, and E_B (E_A) is the energy of the final (initial) nuclear state. As for the effective weak interaction Hamiltonian H_{eff}, we start with the effective single-particle Hamiltonian. The single-particle Hamiltonian in the allowed-type transition approximation is given by

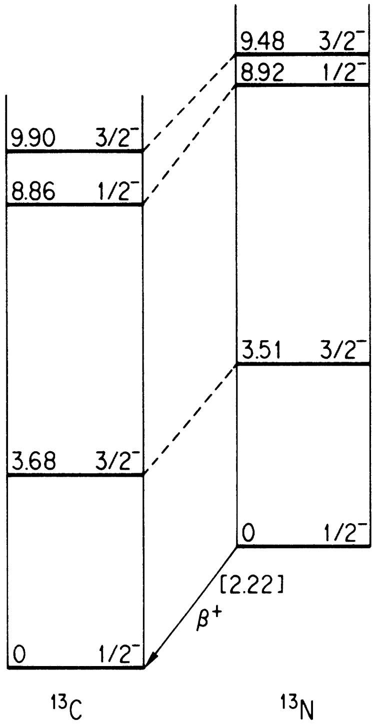

FIG. 1. Low-lying negative-parity levels in ^{13}C and ^{13}N that can be fed from the ground state via allowed-type transitions, $\Delta J=0$ or 1, $\Delta\pi=$no. All the listed levels (Ref. 11) have isospin $T=\frac{1}{2}$.

$$H_{\rm eff}=-\frac{G}{\sqrt{2}}\cos\theta_C t_+[f_V iL_4^* + f_A \boldsymbol{\sigma}\cdot\mathbf{L}^*] \ , \tag{4}$$

with

$$L_\mu^* = -i\bar{\psi}_e\gamma_\mu(1+\gamma_5)\psi_\nu \ . \tag{5}$$

Here, $G\cos\theta_C$ is the weak-coupling constant and $\boldsymbol{\sigma}$'s are the Pauli spin matrices. The nucleon form factors are[13] $f_V=1.0$ and $f_A=-1.262$. The standard calculation[14] yields

$$\sigma=\frac{G^2\cos^2\theta_C}{\pi}p_eE_eF(Z_B=7;E_e)\times\frac{1}{6(2J_A+1)}[f_V^2|\langle B|||\tau|||A\rangle|^2 + f_A^2|\langle B|||\tau\sigma|||A\rangle|^2] \ , \tag{6}$$

where the Fermi factor $F(Z_B;E_e)$ is included to take account of the Coulomb distortion effect; the transition matrix elements $\langle B|||\tau|||A\rangle$ and $\langle B|||\tau\sigma|||A\rangle$ are reduced in both ordinary and isospin spaces. We have used here the fact that all the states shown in Fig. 1 have isospin $T=\frac{1}{2}$. In terms of the reduced transition strengths Eq. (6) can also be written as

$$\sigma=\frac{G^2\cos^2\theta_C}{\pi}p_eE_eF(Z_B=7;E_e)\frac{1}{(2J_A+1)}\times[f_V^2B(F)+f_A^2B(\mathrm{GT})] \ , \tag{7}$$

where

$$B(F)\equiv\sum_{M_f,M_i}|\langle J_fM_f|t_\pm|J_iM_i\rangle|^2 = \tfrac{1}{6}|\langle B|||\tau|||A\rangle|^2 \ , \tag{8}$$

$$B(\mathrm{GT})\equiv\sum_{M_f,M_i}|\langle J_fM_f|t_\pm\boldsymbol{\sigma}|J_iM_i\rangle|^2 = \tfrac{1}{6}|\langle B|||\tau\sigma|||A\rangle|^2 \ . \tag{9}$$

On the other hand, for the β decay $|B\rangle\to|A\rangle$, there holds the relation[14]

$$ft=\frac{2\pi^3\ln 2}{m_e^5G^2\cos^2\theta_C}\times\frac{6(2J_B+1)}{f_V^2|\langle A|||\tau|||B\rangle|^2+f_A^2|\langle A|||\tau\sigma|||B\rangle|^2} \ , \tag{10}$$

which allows us to rewrite Eq. (6) as

$$\sigma=\frac{2\pi^2\ln 2}{m_e^5 ft}p_eE_eF(Z_B,E_e)\frac{2J_B+1}{2J_A+1} \ . \tag{11}$$

We can determine $\sigma[^{13}\mathrm{C}\to{}^{13}\mathrm{N(g.s.)}]$ from the ft value[11] for ^{13}N(g.s.)$\to{}^{13}$C(g.s.) in a model-independent manner. The fact that the β decay ^{13}N(g.s.)$\to{}^{13}$C(g.s.) is a superallowed transition with $\log(ft)^{\rm exp}=3.667\pm0.001$ leads to a rather large cross section for the reaction $^{13}\mathrm{C}+\nu_e\to{}^{13}\mathrm{N(g.s)}+e^-$. The substitution of $(ft)^{\rm exp}$ into Eq. (11) yields

$$\langle\sigma[^{13}\mathrm{C}\to{}^{13}\mathrm{N(g.s.)}]\rangle=8.12\times10^{-43}\ \mathrm{cm}^2 \ .$$

Here and hereafter, the symbol $\langle\sigma\rangle$ stands for the cross section averaged over the neutrino spectrum from the ^{8}B β decay.[3] We remark that, due to the threshold energy $E_{\rm th}=2.2$ MeV, only solar neutrinos from the ^{8}B decay are relevant in this case.

For the transitions to the excited states in ^{13}N, no experimental ft values are available for direct determination of the cross section. Therefore, for the purpose of setting a general guideline of arguments, we first invoke a model calculation of the cross sections with reasonably realistic shell-model wave functions. We then try to improve the estimates using the semiphenomenological effective-operator method,[9,10] which allows us to correlate the transition matrix elements of our concern with experimentally known β-decay strengths in p-shell nuclei. As input for this method, we shall use the following three ground-ground β decays: ^{15}O$(\beta^+)^{15}$N, ^{13}N$(\beta^+)^{13}$C, and ^{11}C$(\beta^+)^{11}$B. In addition, information obtained from the ^{13}C$(p,n)^{13}$N reaction[15] will also be considered.

For a model estimation of the transition matrix elements, we use the Cohen-Kurath (CK) wave functions.[16] [In particular, the version called "(8−16)POT" is adopted here, but the essential feature of our argument does not change with other choices.] The CK wave functions are

TABLE I. Reduced Gamow-Teller strengths for $^{13}C(g.s) \leftrightarrow ^{13}N$ (upper part) and the beta decay of ^{11}C, ^{13}N, and ^{15}O (lower part). $B(GT)$ is defined by Eq. (9). $B(GT)_{theor}$ is the usual shell-model prediction based on the Cohen-Kurath wave functions (Ref. 16) $B(GT)^*_{theor}$ is the result of the effective transition operator method (Refs. 9 and 10); i.e., $B(GT)^*_{theor}$ has been calculated using the rescaled reduced matrix element $\langle B|||\tau\sigma|||A\rangle_{resc}$, Eq. (15). The experimental data are taken from Ref. 11 or Ref. 15.

States in ^{13}N	$B(GT)_{theor}$	$B(GT)^*_{theor}$	$B(GT)_{exp}$
$\frac{1}{2}^-$ (g.s)	0.646	0.420	0.411±0.004
			0.398±0.008[a]
$\frac{3}{2}^-$(3.51 MeV)	4.74	2.14	1.64±0.10[a]
$\frac{1}{2}^-$(8.92 MeV)	1.18	0.524	
$\frac{3}{2}^-$(9.48 MeV)	0.596	0.260	

	$B(GT)_{theor}$	$B(GT)^*_{theor}$	$B(GT)_{exp}$
$^{15}O(\beta^+)^{15}N$	0.667	0.448	0.522±0.007
$^{13}N(\beta^+)^{13}C$	0.646	0.420	0.415±0.006
$^{11}C(\beta^+)^{11}B$	2.49	1.37	1.44±0.04

[a]Deduced from the (p,n) reaction data (Ref. 15).

the most commonly used shell-model wave functions for the normal-parity states in p-shell nuclei. They reproduce the magnetic moments, the $M1$ decay rates, and the ft values of allowed β decays in the p-shell nuclei reasonably well. Table I gives the reduced Gamow-Teller strengths $B(GT)$ calculated with the CK wave functions for the transition relevant to this discussion. The table also gives the available experimental values. A comparison of the calculated and experimental values indicates that the CK prediction overestimates $B(GT)$ by a factor of 1.3–3, but the $B(GT)$'s calculated with the CK wave functions reproduce at least the observed systematics reasonably well.

We give in Table II the cross section for the $^{13}C+\nu_e \rightarrow ^{13}N+e^-$ reaction for ^8B solar neutrinos calculated with the CK wave functions. The table shows that, apart from the ground state, only the lowest excited state ($\frac{3}{2}^-$; $E_x=3.51$ MeV) is important for the solar neutrino detection, and the contributions from other excited levels are negligible. The cross section for this lowest excited state amounts to as high as 70% of that for the ground state.

Next, we discuss the neutral-current reaction (2). The neutral current in the standard theory is given by

$$J^N_\mu = V^3_\mu + A^3_\mu - 2\sin^2\theta_W J^{em}_\mu \ . \tag{12}$$

For solar neutrino energies it is not necessary to retain the induced or recoil terms. Then the vector current does not give rise to any inelastic processes, and therefore we can simply take

$$J^N_\mu = f_A \bar{\psi}\gamma_\mu\gamma^5\tau_3/2\psi \ .$$

The resulting effective Hamiltonian is

$$H_{eff} = -\frac{G}{2\sqrt{2}} f_A \sigma\cdot\mathbf{L}^*\tau_3 \ . \tag{13}$$

The corresponding cross section is given by

$$\sigma = \frac{G^2}{\pi} f_A^2 (E_\nu - \Delta m)^2 \tfrac{1}{48} |\langle B|||\sigma\tau|||A\rangle|^2 \ , \tag{14}$$

where Δm is the mass difference between the initial nuclear level $|A\rangle$ and final nuclear level $|B\rangle$. Using the reduced matrix element obtained with the CK wave functions, we obtain the cross sections for neutral-current reactions as shown in Table II. We see from the table that the reaction leading to the first excited state ($\frac{3}{2}^-$; $E_x=3.68$ MeV) is by far the most important, $\langle\sigma[^{13}C \rightarrow ^{13}C(3.68 \text{ MeV})]\rangle = 2.58\times10^{-43}$ cm^2, and that the contributions of the other final states are smaller by two orders of magnitude and practically negligible.

TABLE II. Cross sections for $\nu+^{13}C$ reactions (in units of cm^2) averaged over the ^8B-decay neutrino spectrum. The row labeled "CK" corresponds to the result of the usual shell-model estimate based on the Cohen-Kurath wave functions. The row labeled "rescaled" gives the estimates obtained with the use of the rescaled reduced matrix element $\langle B|||\tau\sigma|||A\rangle_{resc}$ in Eq. (15). The row labeled "direct" means the result directly obtained from Eq. (11). The underlined entries represent the best available estimates.

	Final states in ^{13}N			
Charged-current reaction	$\frac{1}{2}^-$(g.s)	$\frac{3}{2}^-$(3.51 MeV)	$\frac{1}{2}^-$(8.92 MeV)	$\frac{3}{2}^-$(9.48 MeV)
CK	9.10×10^{-43}	5.83×10^{-43}	2.2×10^{-45}	5.3×10^{-46}
Rescaled	8.02×10^{-43}	$\underline{2.62\times10^{-43}}$	$\underline{9.9\times10^{-46}}$	$\underline{2.3\times10^{-46}}$
Direct	$\underline{8.12\times10^{-43}}$			
	Final states in ^{13}C			
Neutral-current reaction	$\frac{1}{2}^-$(g.s)	$\frac{3}{2}^-$(3.68 MeV)	$\frac{1}{2}^-$(8.86 MeV)	$\frac{3}{2}^-$(9.90 MeV)
CK		2.58×10^{-43}	3.3×10^{-45}	6.2×10^{-46}
Rescaled		$\underline{1.16\times10^{-43}}$	$\underline{1.4\times10^{-45}}$	$\underline{2.7\times10^{-46}}$

III. CROSS SECTIONS OF $\nu+{}^{13}C$ CALCULATED WITH THE EFFECTIVE-OPERATOR METHOD

We expect that the preceding arguments based on the shell-model estimate are reliable enough for the purpose of selecting out transitions that are important for the solar neutrino detection. It would be safe to conclude that only the lowest two states in ^{13}N and the first excited state in ^{13}C are important for this purpose. At a more quantitative level, however, the reliability of the shell-model result should be taken with caution. It is well known[9,10] that observed values of B(GT) are in general significantly quenched compared with the corresponding shell-model values. Core polarization and exchange currents are two known important mechanisms that decrease B(GT), and a great deal of work went into the study of these effects.[9,10] The calculations so far done, however, still fall short of quantitatively reproducing the detailed systematics of B(GT). A similar situation also exists for $M1$ transitions. To quote an example that seems to have some relevance with this discussion, the observed value[17] of $B(M1)$ for the reaction $^{13}C(e,e')^{13}C(\frac{3}{2}^-;$ 3.68 MeV) is quenched by a factor of ~ 2 as compared with the CK prediction. A detailed study[18] of second-order core polarization, isobar-current and meson-exchange-current processes indicates that these processes account for a part of the $B(M1)$ quenching, but that the effects are not large enough to remove the discrepancy between the theory and experiment.

In this situation we take here a semiphenomenological effective-operator method.[9,10] In this approach one assumes that the modification of $\langle B|||\tau\sigma|||A\rangle$ due to the core polarization, meson-exchange correction, and other effects are taken into account by replacing $\tau\sigma$ by an effective operator, while keeping the coefficients of fractional parentage for $|A\rangle$ and $|B\rangle$ fixed as given by the shell-model wave functions. We denote by $\langle B|||\tau\sigma|||A\rangle_{\rm resc}$ the GT-transition matrix element rescaled by the introduction of the effective GT-transition operator, i.e.,

$$\langle B|||\tau\sigma|||A\rangle_{\rm resc} \equiv \langle B|||g_{LA}^{\rm eff}\tau \boldsymbol{l}+g_A^{\rm eff}\tau\sigma+g_{PA}^{\rm eff}[Y_2\times\sigma]^{(1)}|||A\rangle \ . \quad (15)$$

Once the effective coupling constants, $g_{LA}^{\rm eff}$, $g_A^{\rm eff}$, and $g_{PA}^{\rm eff}$, are determined in such a manner that a selected set of B(GT)'s for nuclides belonging to a common major shell is reproduced, all other B(GT)'s involving the same major shell can be predicted. This effective transition operator method is known to be highly successful in correlating a large number of observed B(GT)'s.[9,10] In this study the effective coupling constants appropriate for the $A=13$ system are determined using the experimental information on the β decay in ^{15}O, ^{13}N, and ^{11}C. To constrain further, we also invoke some theoretical arguments whose validity has been tested in the previous extensive studies.[9,10]

According to Ref. 9, the $g_{LA}^{\rm eff}\tau\boldsymbol{l}$ term arising from the core polarization and the exchange current is very small, and there is little ambiguity in this conclusion. We therefore fix the value of $g_{LA}^{\rm eff}$ at a commonly accepted value,[9] $g_{LA}^{\rm eff}=0.011$. We emphasize that the results described in the following are not sensitive to this particular choice. The coupling constant $g_{PA}^{\rm eff}$ is due to the tensor force, and therefore its theoretical estimate[9] is expected to be fairly reliable. To obtain our best parameters, we first vary $g_{PA}^{\rm eff}$ within the range 0.19–0.25, which corresponds to the existing theoretical uncertainty,[9] and determines the relation between $g_{PA}^{\rm eff}$ and $g_A^{\rm eff}$ so that it reproduces the experimental B(GT) for $^{13}N(\beta^+)^{13}C$ within the experimental error. We then search the optimal value of $g_{PA}^{\rm eff}$ (and the corresponding value of $g_A^{\rm eff}$) that gives the best fit to the B(GT) for $^{15}O(\beta^+)^{15}N$ and $^{11}C(\beta^+)^{11}B$. The results are

$$g_A^{\rm eff}=0.69, \quad g_{PA}^{\rm eff}=0.19 \ . \quad (16)$$

Using these values and $g_{LA}^{\rm eff}=0.11$, we calculate $\langle B|||\tau\sigma|||A\rangle_{\rm resc}$ and the corresponding B(GT), to be denoted by $B({\rm GT})^*_{\rm theor}$. The results are given in Table I. One sees that the agreement with the data is now acceptable for our purpose.

Another useful piece of information comes from the knowledge of the empirical proportionality between the (p,n) reaction cross section in the forward direction and the Gamow-Teller transition strength.[15,19,20] According to the latest study by Tadeucci *et al.*,[15] the proportionality between the transitions involving different pairs of nuclides may be uncertain by as much as $\sim 50\%$. By contrast, for transitions with a given fixed pair of initial and final nuclides, they claimed that the proportionality holds much better, typically to an accuracy of $\sim 5\%$. In Table I, the values of $B({\rm GT})_{\rm exp}$'s indexed with "a" have been obtained by assuming the validity of this proportionality. Looking at the B(GT) for

$$^{13}{\rm C(g.s.)}\leftrightarrow{}^{13}{\rm N}(\tfrac{3}{2}^-;\ 3.51\ {\rm MeV})\ ,$$

we note that $B({\rm GT})_{\rm exp}=1.64\pm0.10$ is somewhat smaller than $B({\rm GT})^*_{\rm theor}=2.14$. This discrepancy may be attributed to either or both of the following two points: (1) the uncertainty in deriving $B({\rm GT})_{\rm exp}$ from the (p,n) reaction may be larger than claimed in Ref. 15; (2) the effective transition operator method is valid only up to a 20–30 % level. In the subsequent discussion we use $B({\rm GT})^*_{\rm theor}=2.14$ (with 20–30 % ambiguity attached) for

$$^{13}{\rm C(g.s.)}\leftrightarrow{}^{13}{\rm N}(\tfrac{3}{2}^-;\ 3.51\ {\rm MeV})\ .$$

We can improve estimates of the cross section for $^{13}C+\nu_e\rightarrow{}^{13}N+e^-$ and $^{13}C+\nu_l\rightarrow{}^{13}C^*+\nu_l$ by adopting $B({\rm GT})^*_{\rm theor}$ in Eq. (7). The cross sections averaged over the neutrino spectrum for the ^8B β decay are presented in Table II. Compared with the results obtained in the ordinary shell model, the improved estimates are smaller by a factor of ~ 2 except for the ground-state transition. We use for the best estimate the cross section obtained directly from the observed ft value for the ground-state transition $^{13}C+\nu_e\rightarrow{}^{13}N({\rm g.s.})+e^-$, and for the other transitions we use the σ's calculated with the use of $B({\rm GT})^*_{\rm theor}$. Figure 2 shows the σ's for the major transitions as functions of the incident neutrino energy. One should assign typical uncertainties of $\sim 30\%$ to

$$\sigma[^{13}{\rm C}+\nu_e\rightarrow{}^{13}{\rm N}(\tfrac{3}{2}^-;\ 3.51\ {\rm MeV})+e^-]$$

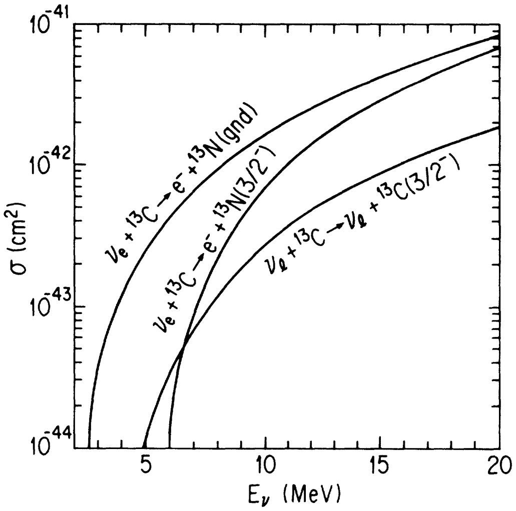

FIG. 2. Cross sections for the charged-current reaction ${}^{13}C+\nu_e \to {}^{13}N+e^-$ and the neutral-current reaction ${}^{13}C+\nu_l \to {}^{13}C^* + \nu_l$ as functions of the incident neutrino energy E_ν. In calculating the cross sections leading to the excited states, the rescaled transition matrix element $\langle B|||\tau\sigma|||A\rangle_{\rm resc}$, Eq. (15) was used.

and

$$\sigma[{}^{13}C+\nu_l \to {}^{13}C(\tfrac{3}{2}^-;\ 3.68\ {\rm MeV})+\nu_l]\ ,$$

whereas the uncertainty in $\sigma({}^{13}C+\nu_e \to {}^{13}N(g.s.)+e^-)$ is negligible in this context.

As was emphasized in Ref. 8, the fact that the number of final states that need to be taken into account for the ${}^{13}C$ detector is very limited offers the advantage that one can probably make a direct calibration measurement for these matrix elements using neutrino beams from stopped muons. The usual limitation with empirical calibrations of neutrino-nucleus reaction cross sections with the use of μ^+-decay neutrinos comes from the fact that one has to deal with many final nuclear states that can be excited. Since the number of relevant levels is only a few with the ${}^{13}C$ target, however, this difficulty can be avoided, especially when one selects particular final states through observing their characteristic decay modes; the cascade feeding of these particular final states is also negligible due to the fact that all higher-lying states are single-particle unbound levels. The cross section averaged over the μ^+-decay neutrino spectrum is estimated to be

$$\langle\sigma\rangle_{\mu\ {\rm decay}} = 2.22\times10^{-41}\ {\rm cm}^2 \qquad (17a)$$

for $\nu_e + {}^{13}C \to e^- + {}^{13}N(g.s.)$,

$$\langle\sigma\rangle_{\mu\ {\rm decay}} = 2.27\times10^{-41}\ {\rm cm}^2 \qquad (17b)$$

for $\nu_e + {}^{13}C \to e^- + {}^{13}N(\tfrac{3}{2}^-;\ 3.51\ {\rm MeV})$,

$$\langle\sigma\rangle_{\mu\ {\rm decay}} = 0.56\times10^{-41}\ {\rm cm}^2 \qquad (17c)$$

for $\nu_e + {}^{13}C \to \nu_e + {}^{13}C(\tfrac{3}{2}^-;\ 3.68\ {\rm MeV})$,

$$\langle\sigma\rangle_{\mu\ {\rm decay}} = 0.74\times10^{-41}\ {\rm cm}^2 \qquad (17d)$$

for $\bar\nu_\mu + {}^{13}C \to \bar\nu_\mu + {}^{13}C(\tfrac{3}{2}^-;\ 3.68\ {\rm MeV})$. In calculating $\langle\sigma\rangle_{\mu\ {\rm decay}}$, since the energy of the neutrinos produced in μ decay can reach 50 MeV, we have included the effect of finite momentum transfers in the following manner. For $B(F)$ we multiply the transition operator τ by $j_0(qr)$. For B(GT), we use Eq. (15), in which, however, $\tau\sigma$ is replaced by $\tau\sigma j_0(qr)$. These modifications[12] made $\langle\sigma\rangle_{\mu\ {\rm decay}} \sim 15\%$ smaller, whereas this type of correction is negligible ($<1\%$) for solar neutrino energies.

IV. DISCUSSION

We now discuss consequences of the new estimates on ${}^8B$ solar neutrino detections. We use the latest estimate of the total flux of ${}^8B$ solar neutrinos given by Bahcall and Ulrich,[3]

$$\phi_{\nu_e}({}^8B) = 5.8\times10^6\ {\rm cm}^{-2}\ {\rm s}^{-1}\ .$$

We first consider the charged-current-induced reactions. The transition to the ${}^{13}N$ ground state can be monitored by the β decay back to ${}^{13}C$(g.s). On the other hand, the final state in the ${}^{13}C(g.s.) \to {}^{13}N(\tfrac{3}{2}^-;\ E_x=3.51$ MeV) transition is unstable against proton emission ($E_p=1.57$ MeV), which can also be detected. Since the contributions of other states are negligible, we consider the sum of the transitions leading to the lowest two states in ${}^{13}N$ as available signals for the charged-current reaction due to solar neutrinos. We recall here that the cross section for the ground-state transition has been obtained directly from the β-decay data and has little ambiguity, whereas the cross section for ${}^{13}C(g.s) \to {}^{13}N(\tfrac{3}{2}^-;\ E_x=3.51$ MeV) has a typical ambiguity of $\sim 30\%$. Combining the charged-current events leading to the lowest two states in ${}^{13}N$, we expect, corresponding to $\langle\sigma\rangle = 1.07\times10^{-42}$ cm², 6.22 SNU or 7870 events per kiloton yr for a ${}^{13}CH_2$ detector [1 SNU equals 10^{-36} captures per target atom per second]. The uncertainty in this event rate is expected to be $\sim 10\%$. Comparing this number with 1680 (1060) electron recoil events per kiloton yr from $\nu_e e^- \to \nu_e e^-$ for a cutoff recoil electron energy 3.5 (5.0) MeV, we note that one has 5–8 times more events with a ${}^{13}CH_2$ detector than with the usual ${}^{12}CH_2$ detector, where ${}^{12}C$ does not play a role as an active target.

Regarding the neutral-current reaction, this calculation gives for the transition to the ${}^{13}C(\tfrac{3}{2}^-;\ E_x=3.68$ MeV) state $\langle\sigma\rangle = 1.16\times10^{-43}$ cm², which yields 0.68 SNU or 860 events per kiloton yr in a ${}^{13}CH_2$ detector. The uncertainty in this estimate is $\sim 30\%$. The transition can be identified by the monochromatic γ ray emitted in the transition back to the ground state. Only the $\tfrac{3}{2}^-$ level needs to be considered here, since the excitation of other levels are negligible. Thus the ${}^{13}CH_2$ detector can be highly promising as a neutral-current detector. The new estimates of event rates, although smaller than the previous values[8] by a facator of ~ 2, are still encouragingly large.

It was remarked in Ref. 8 that the natural abundance ($\sim 1\%$) of ${}^{13}C$ is rather high, so that ${}^{13}C$'s contained in a currently proposed large CH_2 scintillator may provide useful targets for measuring the neutral current reactions once the detection threshold is sufficiently lowered; 30 events/yr are expected in a 3000 ton ordinary CH_2 scintillator for the neutral-current-induced reaction for the standard flux of ^8B solar neutrinos. For further discussion concerning characteristics and merits of the ^{13}C detector, we refer the reader to Ref. 8.

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