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EFFECTS OF THE AXIAL ISOSCALAR NEUTRAL CURRENT FOR SOLAR NEUTRINO DETECTION

J. Bernabéu[1,2], T.E.O. Ericson[2], E. Hernández[1] and J. Ros[1]

Abstract

An essential assumption in the analysis of all the large solar neutrino experiments sensitive to neutral currents has been that the axial transitions are purely isovector. The recent results on the spin structure of the proton suggest the presence of an axial isoscalar neutral-current interaction. This would modify the assumed transition strengths for the neutral-current detection of solar neutrinos. We demonstrate that in the long wavelength limit a deuterium target is insensitive to such a mechanism. Our results for the situation of the planned BOREX experiment show that the suggested isoscalar strength would increase the observed rate by 30–40%, depending on the transition.

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ERRATUM

Page 7, Eqs. (23) and (24)

$\frac{A_c^{(0)}}{2M}$ should be $\frac{A_c^{(0)}}{4M}$

$\frac{F_c^{(0)}}{2M}$ should be $\frac{F_c^{(0)}}{4M}$

Table 2, second column, third row: 0.0097 should be 0.097.

Table 3

second row	$\frac{F_N^{(0,1)}(I=1)}{2M}$	should be	$\frac{F_N^{(0,1)}(I=1)}{4M}$
third row	$\frac{F_N^{(0,1)}(I=0)}{2M}$	should be	$\frac{F_N^{(0,1)}(I=0)}{4M}$
last row	$\frac{F_c^{(0,1)}}{2M}$	should be	$\frac{F_c^{(0,1)}}{4M}$

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1 Introduction

The properties of the neutrino and its weak interaction phenomena play a central role in particle physics, astrophysics and cosmology. If the Sun is used as a source of neutrinos, many questions which are difficult or impossible to study with terrestrial sources are within reach. The rate of charged-current–solar-neutrino reactions in the ^{37}Cl detector [1] is about one third of that predicted by the standard solar model [2] and conventional weak interaction theory. For 15 years no other experiment confirmed the result of the ^{37}Cl experiment. Recently [3], however, the Kamiokande II group presented its results for the number of electrons with energy above 7.5 MeV. These confirm that the solar neutrino deficit is a real effect. The Kamiokande II experiment is sensitive to the direction of the neutrinos and demonstrates conclusively that they are emitted by the Sun. The two ^{71}Ga experiments [4], SAGE and GALLEX, are sensitive to the low-energy neutrinos produced in the Sun by the primary pp reaction.

As it stands now, it is impossible to decide whether the solar neutrino problem is due either to basic properties of the neutrino or to the shortcomings of the standard solar model. The neutrino phenomena most widely considered are neutrino flavour oscillations in a vacuum [5] and in matter, the MSW effect [6], or a large neutrino magnetic moment [7]. The last effect could be responsible for the time dependence of the yield suggested by the data [8] with an anticorrelation with the 11-year sunspot cycle. The reality of this effect is at present under intensive debate. High-rate experiments sensitive to different aspects of the problem are under way. For a discussion about their information content, see Ref. [9]. The Sudbury Neutrino Observatory experiment [10] and the BOREX detector [11] would test neutrino-flavour oscillations or transition magnetic moments of Majorana neutrinos, by comparing simultaneously the neutrino flux seen by neutral- and charged-current interactions in the same experiment. Neutral-current reactions proceed with the same cross-section for all neutrino flavours, whereas charged-current transitions are induced by left-handed electron neutrinos only. As a consequence, a difference in the measured neutrino flux for charged-current and neutral-current reactions would be a signal of one of these properties in the neutrino behaviour. In order to extract the flux from the solar neutrino measurements, one needs a good determination of the charged-current and neutral-current excitation strengths. Detailed calculations of the corresponding cross-sections for neutrino disintegration of the deuteron are available [12] in the literature, as well as for the ^{11}B–^{11}C neutrino reaction [11]. In all of them, however, the axial neutral-current transition strength has been taken up to now as being purely isovector for the nucleon.

The results on the spin distribution of the proton measured by the EMC experiment [13] suggest that there is a non-vanishing expectation value of the axial isoscalar neutral-current for the nucleon, which can affect neutrino-nuclear scattering. This component can modify the expected transition strengths associated with neutral-current detection of the solar neutrino flux. The study of this effect is the main objective of the present paper. In Section 2 we use the EMC data, together with other available sources, to estimate the axial neutral-current coupling constants for protons and neutrons. Section 3 provides the study of the neutral-current transition strength from the deuteron to the proton–neutron continuum state. We show that the axial isoscalar component cannot contribute to the transition amplitude in the long wavelength limit in the impulse approximation. Consequently, the heavy-water detector of the Sudbury Collaboration is not affected by the new mechanism. In Section 4 we analyse the BOREX neutral-current transitions, which involve a superposition of axial isovector and isoscalar components. We show that both the charged- and neutral-current interactions of the incoming neutrinos are described in terms of multipole form factors, most of which can be determined from other experiments. Section 5 presents the results of our study and the main conclusions.

2 Proton spin and neutral-currents

The crucial point of our paper is that there are indications of a neutral-current axial isoscalar coupling, which can lead to important modifications of the neutral-current neutrino cross-sections. This section explores the evidence for such a coupling from the EMC studies.

The EMC measurement [13] of the polarization-dependent structure function of the proton has led to the result

$$\int_0^1 dx\, g_1{}^{P}(x, Q^2) = 0.126 \pm 0.018 \quad \text{at} \quad \overline{Q}^2 \simeq 10\ \text{GeV}^2\ , \tag{1}$$

where g_1 (x,Q^2) is the structure function for polarized muon–polarized proton deep inelastic scattering. This result has given rise to a lively debate [14] about the composition and the spin structure of the proton.

The reason for this is seen as follows: Bjorken's result [15] for the first moment of $g_1(x,Q^2)$ in the limit $Q^2 \to \infty$, for a specific baryon B is the sum rule

$$\Gamma_B(Q^2) = \int_0^1 g_1{}^{B}(x, Q^2)\, dx = \langle B\uparrow| \sum_i \frac{1}{2}\, q_i^2\, \overline{\psi}_i\, \gamma_5\, \gamma_3\, \psi_i\, |B\uparrow\rangle\ , \tag{2}$$

where q_i^2 is the square of the electric charge of each quark (or antiquark) in units of the proton electric charge. When the quark charge matrix is expressed in terms of U(3) flavour generators, this gives the following expression

$$\Gamma_B(Q^2) = \frac{1}{2}\, \langle B\uparrow| \sum_i \overline{\psi}_i\, \gamma_5\, \gamma_3 \left\{ \lambda_3 + \frac{\lambda 8}{\sqrt{3}} + 2\sqrt{\frac{2}{3}}\, \lambda_0 \right\} \psi_i\, |B\uparrow\rangle\ , \tag{3}$$

where λ_a are the usual Gell-Mann matrices.

The matrix elements of the axial current operators with definite quark flavours for the proton can now be parametrized using the so-called axial and pseudoscalar form factors:

$$\langle p|\, \bar{\psi}_i\, \gamma_\mu\, \gamma_5\, \psi_i\, |p\rangle = G_A^{(i)}(Q^2)\ \bar{p}\, \gamma_\mu\, \gamma_5\, p + G_p^{(i)}(Q^2)\ q_\mu\, \bar{p}\, \gamma_5\, p\ . \tag{4}$$

In terms of definite U(3) flavour transformation properties, one can introduce the couplings for each of the terms of Eq. (3). One has

$$\begin{aligned} G_A^{(3)} &= G_A^{(u)} - G_A^{(d)} \\ G_A^{(8)} &= G_A^{(u)} + G_A^{(d)} - 2G_A^{(s)} \\ G_A^{(0)} &= G_A^{(u)} + G_A^{(d)} + G_A^{(s)}\ . \end{aligned} \tag{5}$$

The EMC result for the sum rule Γ_p determines experimentally the combination of these couplings indicated in Eq. (3). The charged weak currents transform, following

Cabibbo theory for the light quarks, as an octet under flavour SU(3). Then the two amplitudes $G_A^{(3)}$ and $G_A^{(8)}$ can be expressed through the amplitudes F and D known from semileptonic decays of baryons

$$G_A^{(3)} = F + D , \qquad G_A^{(8)} = 3F - D . \tag{6}$$

These two values, F and D, and the EMC value for Γ_p determine the U(3) flavour amplitudes

$$G_A^{(3)} = 1.254 \pm 0.006 , \qquad G_A^{(8)} = 0.68 \pm 0.04 , \qquad G_A^{(0)} = 0.12 \pm 0.17 . \tag{7}$$

It is remarkable that the singlet current coupling $G_A^{(0)}$ is compatible with zero, which seems to be in contradiction with naive expectations from models with constituent u, d quarks for the proton. In such models it would mean that the total helicity carried by all quarks (and antiquarks) in a polarized proton is small. However, as shown in Eq. (4), the quark flavours associated with $G_A^{(0)}$ refer to current operators. An (almost) vanishing value of the U(3) flavour singlet amplitude is expected [16] in a QCD picture in which there is a U(1) anomaly and the singlet axial current is not conserved even in the chiral limit. The singlet pseudoscalar η_0 meson will then decouple, and the η_0 is not a Nambu–Goldstone boson in this limit. There is thus only an octet, not a nonet, of Nambu–Goldstone bosons, and $G_A^{(0)}$ receives no contribution in a Goldberger–Treiman type of relation.

For our considerations we do not explicitly need constituent quark models of the proton, since we work instead with the effective nucleonic currents of Eq. (4). Using quark axial-current operators we can then separate out according to Eq. (5) the isoscalars built from the (u, d) operators and the s-quark operators in terms of the couplings $G_A^{(8)}$ and $G_A^{(0)}$ leading to

$$G_A^{(u)} + G_A^{(d)} = 0.31 \pm 0.12 , \qquad G_A^{(s)} = -0.19 \pm 0.06 . \tag{8}$$

As said before, the weak charged-currents of the unified electroweak theory reduce to Cabibbo components of the flavour SU(3) octet when limited to light (u, d, s)-quark operators. This ingredient has been used to determine the couplings (7) or (8). This means that there are no new implication of the EMC measurement for charged-current weak interactions.

For neutral-current weak interactions the hadronic current can be written [17] in the standard SU(2) × U(1) theory as

$$J_\mu^Z = J_\mu^{W^0} - 4 \sin^2 \Theta_w \, J_\mu^{em} , \tag{9}$$

where $J_\mu^{W^0}$ is the weak isospin-rotated current of the charged-current. In flavour space, the GIM mechanism, with the input of the charm quantum number, leads to diagonal currents. There is no flavour-changing neutral-current transitions, to lowest order. As the axial current components come only from $J_\mu^{W^0}$ in Eq. (9), we can write for the current operators

$$J_{\mu,A}^Z = J_{\mu,A}^{W^0} = \overline{\psi}_u \, \gamma_\mu \, \gamma_5 \, \psi_u - \overline{\psi}_d \, \gamma_\mu \, \gamma_5 \, \psi_d + \overline{\psi}_c \, \gamma_\mu \, \gamma_5 \, \psi_c - \overline{\psi}_s \, \gamma_\mu \, \gamma_5 \, \psi_s . \tag{10}$$

They are thus third components of the triplets associated with the first and second families of quarks under the weak isospin SU(2). For the first family, this coincides with strong isospin. Consequently, the first two terms of Eq. (10) correspond to the axial isovector current. Its matrix element for the proton reproduces $G_A^{(3)}$. In the second family one sees the mismatch between strong isospin and weak isospin properties. For light quarks, there is an axial isoscalar strangeness current in addition to the isovector current. Since the EMC result implies $G_A^{(s)} \neq 0$ for the proton according to Eq. (8), it can have important implications for neutral-current phenomena.

At low Q^2 the expectation value of the axial neutral-current for the proton is given by

$$\langle p| J^Z_{\mu,A} |p\rangle = \left[G_A^{(u)} - G_A^{(d)} - G_A^{(s)}\right] \bar{p}\, \gamma_\mu\, \gamma_5\, p\ . \tag{11}$$

We note that Eq. (11) involves a combination of the quark-flavour amplitudes different from that present in the amplitudes F and D of charged weak currents and in the EMC measurement. The analysis above gives for the axial neutral-current coupling of the proton

$$G_A^{(p)} = \left[G_A^{(u)} - G_A^{(d)} - G_A^{(s)}\right]_{proton} = 1.44 \pm 0.06\ . \tag{12}$$

To summarize: the weak neutral-current in the proton contains, besides the well-known axial isovector piece, an axial isoscalar contribution coming from a non-vanishing expectation value of the $\overline{\psi}_s\, \gamma_\mu\, \gamma_5\, \psi_s$ current operator.

The deviation of $G_A^{(p)}$ from the axial isovector value has been searched for in elastic neutrino–proton scattering [18]. The experimental results from Ahrens et al. [19] have been reanalysed including the axial isoscalar piece. However, this extraction of the $G_A^{(s)}$ amplitude from νp or $\bar{\nu}$p scattering is very sensitive to the nucleon form factors and their cut-off masses. Other tests of this new component come from neutrino scattering, from parity-violating electron scattering, and from parity violation in atomic physics, which has been discussed in Ref. [20].

In order to apply all this to nuclei we need the results (7) for the proton amplitudes, whereas for the neutron amplitudes the isovector coupling changes sign and the isoscalar couplings keep their signs. The axial neutral-current couplings of the proton and neutron are therefore

$$G_A^{(p)} = 1.44 \pm 0.06; \quad G_A^{(n)} = -1.06 \pm 0.06\ . \tag{13}$$

In terms of nucleonic axial neutral-currents there is therefore, in addition to the well-known isovector coupling g_A, an isoscalar coupling f_A, such that

$$g_A = \frac{1}{2}\left[G_A^{(p)} - G_A^{(n)}\right] = 1.254 \pm 0.006\ , \quad f_A = \frac{1}{2}\left[G_A^{(p)} + G_A^{(n)}\right] = 0.19 \pm 0.06\ . \tag{14}$$

This conclusion does not rely on any interpretation of the EMC result in terms of a constituent quark model for the proton. Furthermore, if we use the language of the parton model it is independent of quark and gluon densities contributing to the singlet axial current, whatever the combination [21]. We shall present our results as a function of f_A, and study its implications for neutral-current detection of the solar neutrino flux.

3 Neutrino disintegration of the deuteron

This problem has been discussed in detail by several authors [22] for low-energy neutrinos under the assumption that only the standard isovector currents contribute. They find that in this case the predominant transitions occur almost entirely as allowed Gamow–Teller transitions, i.e. by the axial current interaction to the isospin 1 singlet state

($^1S_0$; I = 1) of the two-nucleon system. Furthermore, it is established that the impulse approximation gives a correct description to better than 10% with calculable corrections from the pionic meson-exchange current. There is thus every reason to believe that the impulse approximation will give an excellent description of the isoscalar transitions, in particular since there will be no enhanced contributions of a pionic nature from meson-exchange terms in this case.

The unretarded nucleonic transition operators which appear in the long wavelength limit are

$$\sum_i 1^i \; ; \quad \sum_i \tau^i{}_{(\pm,3)} \tag{15}$$

for the isoscalar and isovector vector currents, and

$$\sum_i \vec{\sigma}^i \; ; \quad \sum_i \vec{\sigma}^i \, \tau^i{}_{(\pm,3)} \tag{16}$$

for the isoscalar and isovector axial currents.

The identity operator of Eq. (15) is incapable of producing inelastic transitions. The isovector vector operator is identical to a component of the isospin and can only produce transitions within an isospin multiplet. For an isospin 0 nucleus such as the deuteron, the isovector vector operator gives a null result in all cases. The only operators that concern us for the time being are therefore the two axial ones of Eq. (16). The isoscalar axial operator is proportional to the total spin carried by the nucleons. For the special case of the deuteron which is a nucleonic spin triplet state $^3S_1$ + $^3D_1$, spin is a good quantum number; the spin operator with no spatial component can then only reorient the deuteron but does not lead to inelastic transitions. Consequently we reach the important conclusion that the axial isoscalar neutral-current interaction is ineffective for solar neutrinos incident on a deuterium target such as the heavy-water detector of the Sudbury experiment. Finally the axial isovector current is the one that has so far been considered in the literature. The effect of this operator is to convert the deuteron S = 1, I = 0 state to the S = 0, I = 1 state without modifying the spatial wave function. The corresponding excitations are the ones of the broad singlet structure in the NN system, extending for several MeV near threshold, which nearly entirely exhausts the nucleonic strength in agreement with the detailed calculations [22].

4 Neutrino transitions in the ^{11}B–^{11}C system

The ^{11}B system is particularly suited to the method of comparing the neutral-current excitation of nuclear levels with the conventional charged-current transitions. The scheme of levels and relevant transitions in the ^{11}B–^{11}C system is presented in Fig. 1. This has led to a program of solar neutrino spectrometry using a massive calorimetric boron liquid-scintillation detector, the BOREX experiment [23], at the Gran Sasso Laboratory.

Starting from the effective neutrino–quark Lagrangian for neutral-current interactions

$$\mathcal{L}(x) = -\frac{G}{2\sqrt{2}} \, \bar{\nu} \, \gamma^\mu \, (1-\gamma_5) \, \nu \, J^Z_\nu \; , \tag{17}$$

the hadronic neutral-current operator J^Z_ν is given by Eq. (9). Here G is the Fermi coupling constant. The nucleonic matrix elements of the axial current have a dependence on g_A and f_A as shown in Eq. (11) at low Q^2. In general, for nuclear transitions one can use an elementary particle approach [24] to express the matrix elements X^μ of the hadronic neutral-current J^μ in terms of form factors. In the Breit reference system, such a matrix element can be written

$$\begin{aligned} X^\mu(p_2,p_1)_{\lambda_1}^{\lambda_2} &= t^\mu \sum_L (-1)^{j_2+\lambda_2}\, C(j_1\, j_2\, L;\ \lambda_1 - \lambda_2,\ 0)\, A^{(L)}\,(Q^2) \\ &+ \sum_{L,1} n^\mu_{(\lambda_1-\lambda_2)}(-1)^{j_2+\lambda_1+1}\, C(j_1\, j_2\, L;\ \lambda_1,\ -\lambda_2,\ \lambda_1-\lambda_2) \\ &\times C(L1\, l; \lambda_1 - \lambda_2, -\lambda_1+\lambda_2, 0)\, F^{(l,L)}(Q^2)\,, \end{aligned} \tag{18}$$

where p_1, j_1, λ_1 (p_2, j_2, λ_2) are the initial (final) momentum, spin, and helicity of the nucleus, $Q^2 = -(p_1-p_2)^2$ and t^μ, $n^\mu_{(i)}$ is the usual tetrad of reference. For most situations, parity and angular momentum selection rules imply that only a few multipole form factors $A^{(L)}$ and $F^{(l,L)}$ contribute. Notations and general results using the elementary particle approach can be found in Ref. [24].

In the long wavelength approximation, the only non-vanishing [25] multipole form factors are $A^{(0)}(Q^2 \approx 0)$ and $F^{(0,1)}(Q^2 \approx 0)$. For *inelastic* transitions, the Conserved Vector Current (CVC) requires that the vector matrix elements vanish in the limit $Q^2 \to 0$. Therefore only the axial-current ones remain: of these $A^{(0)}$ induces parity change and $F^{(0,1)}$ no parity change. One notes in particular that the spin transitions are determined by the non-vanishing form factors $F^{(0,1)}$ ($Q^2 \approx 0$) and these are the dominant ones at low energies for transitions with spin–parity 1^+ quantum numbers. This is the case for ^{11}B, as seen in Fig. 1. The low-energy cross-section integrated over the angular distribution can be expressed in terms of $F_N^{(0,1)} \equiv F^{(0,1)}$ ($Q^2 \approx 0$) for the neutral-current transition, neglecting the contributions from all the other form factors

$$\sigma = \frac{(GE')^2}{4\pi}\, \frac{1}{2j_1+1} \left|\frac{F_N^{(0,1)}}{2M}\right|^2 . \tag{19}$$

Here $E' = E - \Delta$ is the outgoing neutrino energy, with Δ the nuclear excitation energy, j_1 is the initial nuclear spin and M the average nuclear mass. This neutral process can occur for neutrinos of any flavour $\nu + ^{11}B \to \nu + ^{11}B^*$.

For the charged-current transitions $\nu_e + {}^{11}B \to e^- + {}^{11}C^*$, the cross-section to the isospin analogue excited states has a similar structure to that of Eq. (19). Again there is only a dominant axial form factor $F_c^{(0,1)} = F_c^{(0,1)}$ ($Q^2 \approx 0$) for the charged-current transition to the final excited states of ^{11}C and we find

$$\sigma = \frac{G_c^2}{\pi}\, \frac{1}{2j_1+1} \left|\frac{F_c^{(0,1)}}{2M}\right|^2 E_e'\, |\vec{p}_e'|\, F(Z,E_e')\,. \tag{20}$$

Here G_c is the strangeness-conserving weak coupling constant, E_e' ($\vec{p}_e'$) is the energy (momentum) of the outgoing electron, $F(Z,E_e')$ is the Coulomb function for the final electron, and Z is the charge of the daughter nucleus.

As seen in Fig. 1, one must also consider the charged-current transition to the ^{11}C ground state, the isospin partner of the ^{11}B ground state. The unretarded vector form factor $A^{(0)}$ is simply proportional to the isospin flip operator connecting these two states. The cross-section for the ground state charged-current transition then gives

$$\sigma = \frac{G_c^2}{\pi}\, \frac{1}{2j_1+1} \left[\left|\frac{A_c^{(0)}}{2M}\right|^2 + \left|\frac{F_c^{(0,1)}}{2M}\right|^2\right] E_e'\, |\vec{p}_e'|\, F(Z,E_e')\,. \tag{21}$$

The two terms in the square bracket correspond to the well-known Fermi and Gamow–Teller transitions, respectively. The low-energy processes described by Eqs. (20) and (21) are open only to electron neutrinos.

The neutral-current transition form factors $F_N^{(0,1)}$ of Eq. (19) contain two contributions, associated with the isoscalar and isovector components of the nucleonic axial current. Their isovector part can be related, using isospin invariance for the nuclear states, to the charged-current transition form factors leading to the analogue excited levels of ^{11}C. With the normalization used to write the cross-sections as Eqs. (19) and (20), one has

$$F_N^{(0,1)} = F_N^{(0,1)}\ (I=0) + F_N^{(0,1)}\ \ (I=1)$$
$$|F_N^{(0,1)}\ (I=1)| = |F_c^{(0,1)}| , \qquad (22)$$

where $F_c^{(0,1)}$ are the form factors corresponding to the charged-current transition to the isospin analogue states described by $F_N^{(0,1)}$ $(I=1)$. The strict equality shown in Eq. (22) holds for transitions between states of isospin 1/2.

4.1 The charged-current transitions

The charged-current transition form factors can be obtained directly from other experimental sources. For the case of a charged (isovector) transition with an inverse beta decay, one can use the experimental (ft) value to obtain

$$G_c^2 \left[\left| \frac{A_c^{(0)}}{2M} \right|^2 + \left| \frac{F_c^{(0,1)}}{2M} \right|^2 \right] = \frac{(2j_2+1)}{2} \frac{\pi^3 \ \ln\ 2}{m_e^5 (ft)} , \qquad (23)$$

where j_2 is the initial (final) nuclear spin for beta decay (neutrino transition) and t is the β decay half-life. This is the case for the transition from ^{11}B to the ground state of ^{11}C. Using the experimental value [26] of (ft) for the β^+ decay $^{11}C \rightarrow\ ^{11}B + e^+\nu_e$, one obtains

$$\lambda_c \equiv \left| \frac{A_c^{(0)}}{2M} \right|^2 + \left| \frac{F_c^{(0,1)}}{2M} \right|^2 = 1.55 \ . \qquad (24)$$

Taking ^{11}C and ^{11}B as the components of an isospin doublet, the vector form factor is given by $|A_c^{(0)}/2M|^2 = 1$. This is due to the fact that, in the long wavelength approximation, the Fermi transition operator is given by the isospin lowering generator. So Eq. (24) can be regarded as the experimental determination of the axial form factor $F_c^{(0,1)}$ for the ground-state transition.

In the impulse approximation, in which the nuclear current is written in terms of the nucleonic degrees of fredom, the charged-current transition strengths to the levels of ^{11}C can be determined experimentally from forward (p,n) reactions on ^{11}B. These are sensitive to the spin–isospin-dependent component of the nuclear Hamiltonian. This spin–isospin operator is identical to the one present in the nucleonic axial isovector current in the allowed approximation $Q^2R^2 \ll 1$, where R is the nuclear radius. The reduced matrix element under rotations has the following form

$$\frac{F_c^{(0,1)}}{2M} = \sqrt{3}\, g_A\ [\sigma\ \tau_+]^{0,1} \ ;$$
$$[\sigma\ \tau_+]^{0,1} = \frac{(-1)^{j_2-m}}{C(j_1,j_2,1;-m,m,0)} \langle j_2, m; T', T_3'| \sum_{i=1}^{A} \tau_+^i\ \sigma_3^i\ |j_1, m; T, T_3\rangle \ , \qquad (25)$$

where $\tau_+ \equiv 1/2(\tau_1 + i\tau_2)$ converts a neutron into a proton, σ_3 is the third component of Pauli matrices and T, T_3 are the nuclear isospin and its third component.

The standard technique to determine these form factors experimentally is based on the observation that such matrix elements occur in forward proton–neutron charge-exchange reactions in the impulse approximation. So as to eliminate renormalization due to nuclear distorsions of the incident wave, one normalizes these transition rates to the accurately determined (ft) value for ground state axial β decay transition. The results of Taddeucci et al. [27] analysed in this fashion determine empirically the matrix elements $[\sigma\ \tau]^{0,1}$. For inelastic transitions, this leads to the values of $\lambda_c = |F_c^{(0,1)}/2M|^2$ given in Table 1. As shown there only the sum of the strengths of the 4.32 MeV and 4.80 MeV excited states of ^{11}C can be determined, although there is a suggestion in the data that the strengths are similar.

4.2 The neutral-current transitions

To continue with the programme of the determination of the nuclear form factors, we have to control the relative contribution of the isovector and isoscalar components of the neutral current. With a self-evident notation one gets for the isoscalar piece

$$\frac{F_N^{(0,1)}\ (I=0)}{2M} = \sqrt{3}\, f_A\ [\ \sigma\ ]^{0,1}\ ;$$

$$[\sigma]^{0,1} = \frac{(-1)^{j_2-m}}{C(j_1,j_2,1;-m,m,0)} \langle j_2, m; T, T_3 | \sum_{i=1}^{A} \sigma_3^i\ | j_1, m; T, T_3\rangle\ . \qquad (26)$$

The isovector piece is given by

$$\frac{F_N^{(0,1)}}{2M}\ (I=1) = \sqrt{3}\, g_A\ [\sigma\ \tau_3]^{0,1} \qquad (27)$$

with the last matrix element defined by Eq. (25) with the $\tau_+ \to \tau_3$ replacement. Actually $[\sigma\ \tau_+]^{0,1}$ and $[\sigma\ \tau_3]^{0,1}$ are equal, up to a sign, between states of isospin 1/2. The nucleonic axial isovector g_A and isoscalar f_A coupling constants have been discussed in Section 2 and are given in Eq. (14).

The operators $\vec{\sigma}$ and $\vec{\sigma}\tau_3$ are also the ones which appear in magnetic dipole transitions. So we can use a connection of the weak axial form factors $F_N^{(0,1)}$ with the electromagnetic vector form factors $F_{em}^{(1,1)}$. In the long wavelength limit, the nucleonic matrix element of the electromagnetic current operator responsible of isoscalar magnetic dipole transitions is

$$\begin{aligned}\frac{F_{em}^{(1,1)}\ (I=0)}{2M} &= -\frac{\sqrt{3}}{2}\,\frac{\Delta}{\sqrt{6}\, m_p}\,\frac{(-1)^{j_2-m}}{C(j_1,j_2,1;-m,m,0)} \\ &\quad \times \langle j_2, m; T, T_3 | \sum_{i=1}^{A} \left[(\mu_p + \mu_n)\sigma_3^i + \ell_3^i\right]\ | j_1, m; T, T_3\rangle\ ,\end{aligned} \qquad (28)$$

where μ is the total magnetic moment of the nucleon and ℓ_3 is the nucleonic (third component of the) orbital angular momentum operator. Rotational invariance imposes definite total angular momentum for the nuclear states, so that for *inelastic* transitions we have the following restriction

$$\langle | \sum_{i=1}^{A} \left(\vec{\ell}^{\,i} + \frac{1}{2}\,\vec{\sigma}^{\,i}\right) | \rangle = 0\ . \qquad (29)$$

Therefore, the isoscalar magnetic dipole amplitude is induced by the transition operator $(\mu_p+\mu_n-1/2)\ \sum_{i=1}^{A}\vec{\sigma}^i$, proportional to the operator responsible of the isoscalar weak neutral-current transition given by Eq. (26). We conclude that the experimental knowledge of the electromagnetic transition form factor $F_{em}^{(1,1)}$ (I = 0) determines the value of the axial isoscalar neutral-current form factor $F_N^{(0,1)}$ (I = 0) of interest, according to Eq. (26), up to the axial coupling f_A of the nucleon.

The same analysis can be applied to the isovector magnetic dipole amplitude and gives the following electromagnetic form factor

$$\frac{F_{em}^{(1,1)}}{2M}(I=1) = -\frac{\sqrt{3}}{2}\frac{\Delta}{\sqrt{6}\,m_p}\frac{(-1)^{j_2-m}}{C(j_1,j_2,1;-m,m)} \langle j_2,m;T',T_3|\sum_{i=1}^{A}\left[(\mu_p-\mu_n)\sigma_3^i\,\tau_3^i+\ell_3^i\,\tau_3^i\right]|j_1,m;T,T_3\rangle\,. \tag{30}$$

The first term of the transition operator in Eq. (30), the spin–isospin current, is proportional to the isovector weak axial current form factor $F_N^{(0,1)}$ (I = 1), already determined from the isospin rotated (p,n) transition strengths of Table 1. So the empirical determination of both the electromagnetic $F_{em}^{(1,1)}$ (I = 1) and the weak neutral-current $F_N^{(0,1)}$ (I = 1) isovector form factor would allow the extraction of the orbital angular momentum term in Eq. (30), expected to be relatively small for these isovector magnetic dipole transitions.

From the formulae given above in this section one notices the connection of the weak neutral current form factors with the electromagnetic form factors. One obtains

$$\frac{F_N^{(0,1)}(I=0)}{F_N^{(0,1)}(I=1)} = \left(\frac{f_A}{g_A}\right)\frac{\mu_p-\mu_n}{\mu_p+\mu_n-\frac{1}{2}}\frac{F_{em}^{(1,1)}(I=0)}{F_{em}^{(1,1)}(I=1)}\left[1+\frac{\langle\|\ \ell_3\tau_3\ \|\rangle}{\langle\|\ (\mu_p-\mu_n)\sigma_3\tau_3\ \|\rangle}\right]\,. \tag{31}$$

The ratios shown in Eq. (31) are of interest since their sensitivity to nuclear physics details is small. This is important in case an insufficient phenomenological information is available.

The isoscalar and isovector form factors $F_{em}^{(1,1)}$ (I = 0) and $F_{em}^{(1,1)}$ (I = 1), as well as their relative sign, can be determined from the electromagnetic decay widths [26] of the excited levels of ^{11}B and ^{11}C. The problem is that only the width of the 2.00 MeV state of ^{11}C has been actually measured, contrary to the case of ^{11}B for which this information is available also for other states. For the other levels we shall use the Cohen–Kurath wave functions [28] obtained from effective interactions in the 1p-shell to predict the ratio between the two electromagnetic form factors. This ratio, together with the experimental width for ^{11}B*, determines $F_{em}^{(1,1)}$ (I = 0) and hence from Eq. (26) $F_N^{(0,1)}$ (I = 0) of interest.

There are at least three possible strategies to determine the relevant form factors F_C, F_N (I = 1) and F_N (I = 0) for the different transitions.

Method 1. When the (p,n) transition strength and the electromagnetic widths for the corresponding states in both ^{11}B and ^{11}C are known experimentally, one can extract all form factors without appealing to a particular nuclear model.

Method 2. When the electromagnetic width for a state in ^{11}B is experimentally known, but not that for the analogue one in ^{11}C, we can use the Cohen–Kurath wave functions to predict the ratio F_{em} (I = 0)/F_{em} (I = 1), which then determines F_N (I = 0). If the integrated strength of the (p,n) transitions to several levels is also known experimentally, we can use the nuclear model to predict its distribution among these levels. Then the F_C's and F_N (I = 1) follow.

Method 3. From the experimental width of the ^{11}B level and the two (theoretically) independent ratios of nuclear matrix elements shown in Eq. (31), we can determine the form factors separately for each level.

We shall now apply these methods to the different situations.

4.3 The transition to the first excited state

In this case we use Method 1. The isovector form factor is determined by the (p,n) reaction [27] and the isoscalar form factor is obtained from the combined information on the magnetic dipole transitions [26] in ^{11}B and ^{11}C. We obtain for the neutral-current form factors of ^{11}B the values

$$
\begin{aligned}
&3/2\ (\mathrm{p,n})\ 1/2^* \ \rightarrow\ \left|\frac{F_N^{(0,1)}\ (I=1)}{2M}\right| = 0.632\ g_A\ ; \\
&1/2^* \xrightarrow{\gamma} 3/2 \ \rightarrow\ \left|\frac{F_N^{(0,1)}\ (I=0)}{2M}\right| = 0.791\ f_A\ .
\end{aligned}
\tag{32}
$$

We emphasize that these predictions do not rely on a specific model for the nuclear structure of ^{11}B.

Just to demonstrate the possibilities offered by the Cohen–Kurath wave functions obtained from effective interactions in the 1p-shell, we compare in Table 2 the theoretical predictions and the phenomenological determinations of the individual ratios appearing in Eq. (31). The excellent agreement provides a guide for the predictions of other transitions with incomplete experimental information.

The experimental electromagnetic form factor $F_{em}^{(1,1)}$ (I = 1) is obtained from the combined information on $1/2^* \xrightarrow{\gamma} 3/2$ in ^{11}B and ^{11}C. The spin current contribution to it is deduced from $F_N^{(0,1)}$ (I = 1), so the phenomenological ratios of Table 2 follow. The empirically small result for the relative contribution of the orbital angular momentum term to the spin–isospin term is a typical result: the strength of isovector magnetic dipole transitions is dominated by the spin–isospin current. The isoscalar magnetic dipole strength is experimentally small, but this is an artefact of the nucleon magnetic moment combination appearing in Eq. (28). As seen from Table 2, the relative importance of the axial isoscalar neutral-current transition can be quite large.

Suppose we use Method 3 for this transition. In this case, the experimental electromagnetic width of ^{11}B and the nuclear model estimate (Table 2) for the ratio $F_{em}^{(1,1)}(I=0)/F_{em}^{(1,1)}(I=1)$ would give rise to the prediction $\left|F_N^{(0,1)}(I=0)/2M\right| = 0.70\ f_A$, in fair agreement with the result 0.79 f_A of Eq. (32). The use of a second theoretical ratio of Table 2 would, in addition, determine the isovector form factor $\left|F_N^{(0,1)}(I=1)/2M\right| = 0.68\ g_A$, to be compared with the phenomenological value 0.63 g_A of Eq. (32). Although Method 1 cannot provide the relative sign between the isoscalar and isovector components of the neutral-current form factor, nuclear theory unequivocally and confidently determines the relative sign to be positive. A negative relative sign would not only mean a relative change of sign between the orbital and spin contributions in Table 2, but it would mean that the orbital matrix element would be larger than the spin one. This would be unacceptable from the viewpoint of nuclear theory for such well understood nuclei.

4.4 The other two inelastic transitions

As shown in Table 1, the (p,n) transition strengths to the 4.32 and 4.80 MeV levels of ^{11}C are not resolved experimentally. As dictated by our Method 2, we fix their summed strength to the observed one and take the distribution between the two states as predicted by the Cohen–Kurath wave functions. The model-dependent result is that the ratio is close to unity:

$$\left|\frac{F_N^{(0,1)}\ (I=1)[3/2 \to 5/2^*]}{F_N^{(0,1)}\ (I=1)[3/2 \to 3/2^*]}\right|_{th} = 0.944\ . \qquad (33)$$

Together with the experimental result of Table 1, this gives the individual isovector transition form factors shown in Table 3.

As for the isoscalar neutral-current form factors, we take for each of the two levels their measured electromagnetic lifetimes for the corresponding states in ^{11}B and the theoretical electromagnetic isoscalar–isovector ratio. We obtain the values shown in Table 3.

As by-products of our calculation with Method 2, we predict the electromagnetic lifetimes in ^{11}C to be 1.74 fs for the $5/2^*$ level and 0.61 fs for the $3/2^*$ level, which is consistent with the experimental [26] upper limits of 12 fs and 11 fs, respectively. Furthermore, the predicted (p,n) individual strengths are very similar, in agreement with the qualitative observation of Ref. [27].

We have also calculated the relevant form factors using Method 3 for the same states. Instead of using the theoretical input (33) and the summed (p,n) strength to these two levels, we have taken for each level the relative orbital angular momentum contribution to the magnetic dipole form factor as given by the theoretical nuclear model. Once the ratios of Eq. (31) are theoretically given one experimental quantity, the electromagnetic width of each level in ^{11}B, determines the magnitude of the form factors. They are shown in Table 3. As noted there, the predicted integrated (p,n) strength to the two levels turns out to be too high by about 40%, when compared with the experiment. Because of this discrepancy, we shall give our results for the absolute cross-sections based on Methods 1 and 2.

From the results shown in Table 3, we conclude on the potential importance of the axial isoscalar neutral-current form factor $F_N^{(0,1)}\ (I=0)$. For these transitions, the neutrino neutral-current strength offers an appreciable sensitivity to a non-vanishing axial isoscalar neutral-current coupling f_A of the nucleon.

5 Results and discussion

The presence of an axial isoscalar neutral-current interaction can modify the expected transition strengths in the detection of the solar neutrino flux. We have seen in Section 3 that a deuteron target is not affected by this mechanism in the long wavelength approximation. The BOREX detector, on the other hand, is sensitive in its neutral-current signals to the axial isoscalar component. We now quantify these effects for the different transitions involved.

We first emphasize that the effect under study only modifies the neutral-current transition strengths or, alternatively, their ratios to the analogue charged-current ones, in the set of tools available [11] in the ^{11}B experiment. From Eqs. (20) and (21) and the determination of the isovector form factors made in Section 4, we can calculate the charged-current cross-section averaged over the incident neutrino spectrum [29]. We obtain the results shown in Table 4.

The different transitions can be used as direct measurements of the isovector strengths. We note that the ground-state–ground-state transition strength is predicted in a model-independent way and that also the transition to the first excited state rests on a firm theoretical prediction (Method 1). These two low-lying transitions are thus the most adequate ones to extract the electron–neutrino flux arriving at the detector.

The neutral-current transition strengths induced by all neutrino flavours arriving at the detector are obtained from Eq. (19) and the form factors determined in Section 4. When averaged over the incident neutrino spectrum, the neutral-current cross-sections are as shown in Table 5.

The EMC experiment suggests, according to Eq. (14), a value of the ratio of isoscalar versus isovector couplings $(f_A/g_A) \simeq 0.15$. With this value one notes from Table 5 that the effect of the axial isoscalar neutral-current coupling increases the expected strength by an amount of 40% for the $3/2^- \to 1/2^{-*}$ transition, and somewhat less for the other transitions.

We urge the completion of complementary experiments measuring the electromagnetic widths of the ^{11}C levels to avoid the remaining uncertainties of the analysis. The nucleonic axial isoscalar neutral-current coupling f_A can be independently determined in other neutral-current experiments, particularly in dedicated elastic neutrino–proton scattering [30], and in the CERN SMC experiment on deep inelastic polarized-muon–polarized-deuteron scattering [31]. As shown here, the relevance of these ingredients to probe the properties of the neutrino in solar neutrino experiments is of great value.

In conclusion, the effects of the axial isoscalar neutral-current coupling, as suggested by the EMC experiment, are very important for the planned BOREX solar neutrino detector: the neutral-current transitions in ^{11}B increase their strengths to an appreciable level.

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Table 1

Empirical values for the quantity λ_c defined by Eq. (24) for transitions from ^{11}B(g.s.) to ^{11}C states

(J^π) E (MeV)	λ_c
$(3/2^-)$ 0	1.55
$(1/2^-)$ 2.00	0.63
$(5/2^-)$ 4.32 } $(3/2^-)$ 4.80 }	1.53

Table 2

The theoretical and phenomenological properties of ratios of transition matrix elements between the first excited and the ground state of ^{11}B. The sign of the phenomenological ratio of neutral current form factors has been theoretically fixed.

1/2*	Theory	Phenomenology
$\frac{F_N^{(0,1)}(I=0)}{F_N^{(0,1)}(I=1)}$	$1.043\left(\frac{f_A}{g_A}\right)$	$1.252\left(\frac{f_A}{g_A}\right)$
$\frac{F_{em}^{(1,1)}(I=0)}{F_{em}^{(1,1)}(I=1)}$	0.0097	0.111
$\frac{\langle\|\ell_3\iota_3\|\rangle}{\langle\|(\mu_p-\mu_n)\sigma_3\iota_3\|\rangle}$	−0.135	−0.088

Table 3

The values of transition matrix elements between the ground state and the two excited states 5/2* and 3/2*. The results provided by Methods 2 and 3 are given for comparison.

	Method 2		Method 3	
	5/2*	3/2*	5/2*	3/2*
$\left\|\frac{F_N^{(0,1)}\ (I=1)}{2M}\right\|$	0.673 g_A	0.713 g_A	0.856 g_A	0.833 g_A
$\left\|\frac{F_N^{(0,1)}\ (I=0)}{2M}\right\|$	0.689 f_A	0.820 f_A	0.689 f_A	0.820 f_A
$\frac{F_{em}^{(1,1)}\ (I=0)}{F_{em}^{(1,1)}\ (I=1)}$	0.076	0.077	0.076	0.077
$\frac{\langle \| \ell_3 \tau_3 \| \rangle}{\langle \| (\mu_p - \mu_n)\sigma_3\tau_3 \| \rangle}$	+0.093	+0.216	-0.142	+0.041
$\tau(^{11}C)$	1.74 fs	0.61 fs	1.74 fs	0.61 fs
$\lambda_c = \left\|\frac{F_c^{(0,1)}}{2M}\right\|^2$	0.718	0.807	1.164	1.102

Table 4

Charged-current cross-sections for $^{11}B \rightarrow {}^{11}C$, averaged over the incident solar neutrino spectrum. The transition strengths to the ground and first excited states of ^{11}C have been predicted using Method 1, those of the other two transitions using Method 2.

$J^{\pi}(^{11}C)$	$\sigma_c(10^{-43}\ cm^2)$
$3/2^-$	9.82
$1/2^{-*}$	2.09
$5/2^{-*}$	0.84
$3/2^{-*}$	0.72

Table 5

Neutral-current cross-sections for neutrino excitation of $^{11}B^*$ averaged over the incident solar neutrino spectrum. The transition strength to the first excited state has been predicted using Mehod 1, those to the other two states using Method 2.

$J^{\pi}(^{11}B)$	$\sigma_N(10^{-44} cm^2)$
$1/2^{-*}$	$7.52\ [1 + 1.25(f_A/g_A)]^2$
$5/2^{-*}$	$3.58\ [1 + 1.02(f_A/g_A)]^2$
$3/2^{-*}$	$3.10\ [1 + 1.15(f_A/g_A)]^2$

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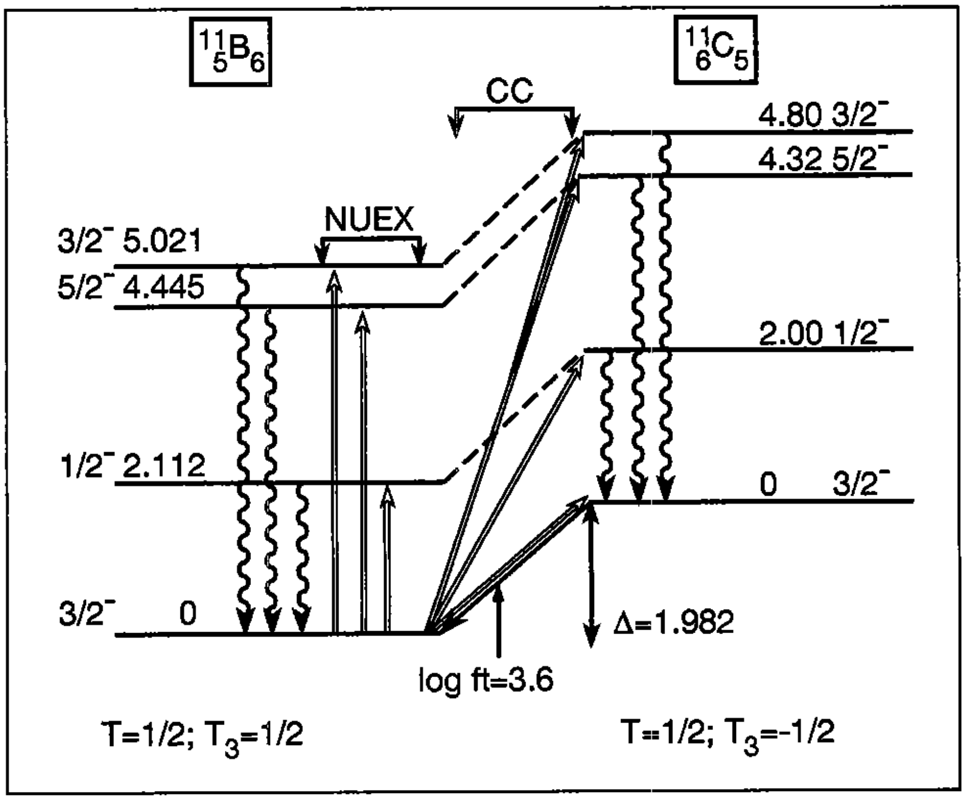
Fig. 1 The ^{11}B–^{11}C mirror nuclear system. The energies are in MeV. The excitations of ^{11}C from ^{11}B (g.s.) are denoted by CC (charged current) and the excitations of ^{11}B from ^{11}B (g.s.) by NUEX (neutral excitations).



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