Neutrino-nucleus reactions and the structure of neutral currents

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Cross sections for neutrino scattering on ^{14}N to specific nuclear states are calculated. The ratio of the isovector neutral to charged current events is found to be independent of the nuclear models and sufficiently sensitive to permit discrimination among various fashionable models of neutral currents. Low energy neutrino elastic cross sections for p-shell nuclei are predicted to be of the order of 10^{-39} cm^2.

[RADIOACTIVITY Neutral and charged currents. CVC. Elastic and inelastic neutrino reactions. Gauge theory models. Shell model. Structure of ^{14}N.]

I. INTRODUCTION

The recent discovery[1,2] of neutral currents in high energy neutrino experiments has perhaps been one of the most exciting steps in our understanding of weak interactions. From the analysis of such experiments we have learned a lot about the space-time structure of the hadronic part of these currents. Thus we now know that theoretical models with pure axial or pure vector currents are ruled out. Models with pure pseudoscalar or pure tensor hadronic currents, as well as isoscalar models, are also ruled out.[3] Furthermore, from such experiments we have deduced that the overall strength of the neutral currents contribution is about 30% of the corresponding one for charged currents.

It is important to keep in mind that inclusive high energy experiments will not be able to determine completely the structure of neutral currents. They are plagued by ambiguities, such as the $V-A$ and the isovector-isoscalar ambiguities.[3] One therefore needs exclusive experiments. In this connection low energy measurements can supplement high energy ones and aid in the determination of the structure of the neutral currents. Most prominent among such experiments are the following:

(i) Experiments measuring the change in polarization of polarized laser beams[4] (atomic physics experiments) or the change of polarization of polarized electrons scattered by nuclear targets.[5]
(ii) Experiments measuring the inelastic neutrino cross sections to individual nuclear levels.[6]

By a judicious choice of the target nucleus and selection of the final states, one can eliminate both the isoscalar-isovector ambiguity and the $V-A$ ambiguity. In such experiments, which are admittedly difficult, one can isolate the various pieces of the hadronic current. The difficult task is the identification of the specific states expected to be excited. The simplest procedure seems to be to identify these levels by measuring the γ-ray energy, provided of course that they deexcite mainly by γ emission, i.e.,

$$\begin{aligned} &\nu + A \to \nu' + A^* \\ &\qquad\qquad \hookrightarrow A + \gamma, \\ &\bar{\nu} + A \to \bar{\nu}' + A^* \\ &\qquad\qquad \hookrightarrow A + \gamma. \end{aligned} \tag{1}$$

The extraction of the parameters characterizing the neutral currents cannot be made in a way which is completely independent of the nuclear models. A reasonable procedure would be to start from situations with the least amount of ambiguities regarding the nuclear structure. From elastic scattering, if the recoiling nucleus can be detected, or from atomic physics experiments one may deduce the isoscalar and isovector vector coupling constants. The next step then would be to look for allowed Gamow-Teller-like transitions ($\Delta J = 1$, no change of parity). These transitions will primarily determine the axial (Gamow-Teller) coupling constants, both isoscalar and isovector. They may also yield the vector coupling constants as small contributions arising from the magnetic moment term, provided that both neutrino and an antineutrino reactions are done for the same target. These results are particularly important for isoscalar currents. Indeed the mere observation of isoscalar Gamow-Teller-like transitions with cross sections $\gtrsim 10^{-42}$ cm^2 will automatically rule out most of the currently fashionable models of neutral currents,[6] including the standard Weinberg-Salam scheme.

In a previous work[6] we have shown that the above tasks can best be accomplished by choosing ^{14}N as the nuclear target and observing, in the reaction of Eqs. (1), the transitions from the ground state

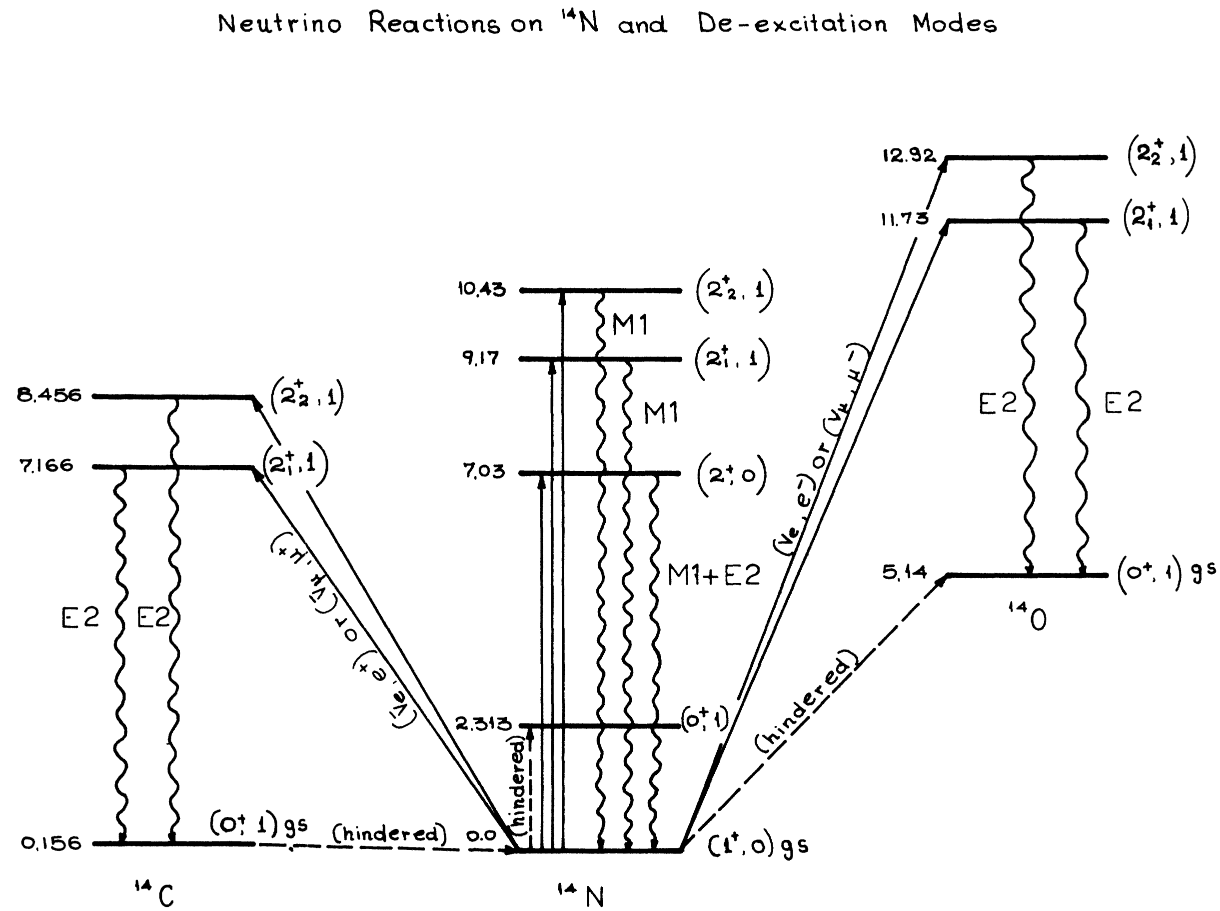

FIG. 1. Reactions to levels of interest in the present work for the $A=14$ system (taken from Ref. 33).

to the 2^+ $I=0$ and 2_1^+ $I=1$ states. In the present paper we will supplement the above-mentioned work by investigating the charged current reactions

$$\nu_e + {}^{14}\mathrm{N} \to {}^{14}\mathrm{O}(2^+) + e^- \quad \llcorner\!\to {}^{14}\mathrm{O}(\mathrm{g.s.}) + \gamma,$$

$$\bar{\nu}_e + {}^{14}\mathrm{N} \to {}^{14}\mathrm{C}(2^+) + e^+ \quad \llcorner\!\to {}^{14}\mathrm{C}(\mathrm{g.s.}) + \gamma. \tag{2}$$

These reactions may be very useful, in conjunction with reactions (1), in investigating the neutral current. The reason is that the ratio of the neutral to charged current events, for transitions to a given nuclear level, is much less dependent on the nuclear model than each of these two processes separately. Furthermore, knowledge of the cross sections for (2) may be useful in those cases in which the experimental resolution is not good enough to distinguish between charged and neutral current events and a background subtraction appears necessary. The ^{14}N target has the advantage that the charged current reaction will also predominantly excite the 2^+ $I=1$ states (the ground state transition has negligible cross section for the same reason that the inverse β decay is greatly hindered). Thus both the charged and neutral current processes will lead to states which deexcite by γ emission and can therefore be measured simultaneously (see Fig. 1).

Finally we investigate the possibility of measuring the elastic neutrino cross section by detecting the recoiling nucleus. Such experiments[7] can probably be done for targets which can simultaneously serve as detectors, e.g., C and Pb. We present estimates for such cross sections at various neutrino energies and nuclei recoiling with energies greater than a minimum (e.g., $E_N \geq 2$ MeV and $E_N \geq 2.5$ MeV).

In our work we have adopted a phenomenological point of view and we do not feel committed to any gauge theory model of neutral currents. We have, however, included in our calculations most of the currently fashionable of these models as a guide to the experimentalists about the expected cross sections. It is the task of the suggested experiments to settle the question of which, if any, of these is correct.

II. NEUTRAL CURRENT FORMALISM

It is generally assumed that the effective hadronic neutral current has a similar space-time structure to that of the charged current. Thus it can be described in terms of 12 form factors; i.e., the isovector and isoscalar vector, weak magnetism, scalar, axial vector, tensor, and pseudoscalar form factors. The scalar form factor vanishes because of conserved vector currents (CVC) (it is also second class). The pseudoscalar form factor also vanishes when the masses of the usual leptons (e, ν_e, μ, ν_μ) are neglected. The tensor form factor should, in principle, be included in a phenomenological theory like ours. Since, however, there is no definite evidence at present for the existence of second class charged currents, we decided to drop this term in order to keep the analysis as simple as possible. In the context of the quark model the interaction Lagrangian is then written as

$$\mathcal{L}_{\text{int}} = -\frac{G}{\sqrt{2}}\bar{\nu}\gamma^\mu(1-\gamma_5)\nu J_\mu^{NC}, \tag{3a}$$

$$J_\mu^{NC} = \tfrac{1}{2}[\bar{q}\gamma_\mu(\alpha_V^1 - \alpha_A^1\gamma_5)\tau_3 q + \bar{q}\gamma_\mu(\alpha_V^0 - \gamma_5\alpha_A^0)q + \cdots], \tag{3b}$$

where τ_3 is the third Pauli matrix, $q = \binom{u}{d}$, and the dots stand for contributions from other quarks. α_V^τ ($\tau = 0, 1$) are the amplitudes for the isoscalar and isovector vector currents, while α_A^τ are the corresponding quantities for the axial current. The matrix elements of (1b) between two nucleon states are written as

$$\langle N(p+k)|J_\mu^{NC}|N(p)\rangle = \bar{u}(p+k)(J_\mu^{NC})_{\text{eff}}u(p), \tag{4a}$$

with

$$(J_\mu^{NC})_{\text{eff}} = (\tfrac{1}{2}\tau_3 f_1^1 + f_1^0)\gamma_\mu + i\frac{\sigma_{\mu\nu}}{2m}k^\nu(\tfrac{1}{2}\tau_3 f_2^1 + f_2^0) - \gamma_\mu\gamma_5(\tfrac{1}{2}\tau_3 f_{\text{GT}}^1 + f_{\text{GT}}^0). \tag{4b}$$

(The pseudoscalar term has been neglected for ultrarelativistic electrons.) At small k^2 the quantities in (4b) are constants related to those of (3b) by

$$f_1^1 = \alpha_V^1 f_1^{cc}, \quad f_2^1 = \alpha_V^1 f_2^{cc}, \quad f_{\text{GT}}^1 = \alpha_A^1 f_{\text{GT}}^{cc}, \tag{5a}$$

$$f_1^0 = \tfrac{3}{2}\alpha^0, \quad f_2^0 = \tfrac{3}{2}\alpha_V^0(\mu_p + \mu_n), \quad f_{\text{GT}}^0 = \tfrac{3}{10}\alpha_V^0(g_A/g_V), \tag{5b}$$

where $f_1^{cc}, f_2^{cc}, f_{\text{GT}}^{cc}$ are the familiar quantities from the charged currents, namely

$$f_1^{cc} = 1, \quad f_2^{cc} = \mu_p - \mu_n, \quad f_{\text{GT}}^{cc} = g_A/g_V,$$

and $g_A/g_V = 1.24$, $\mu_p = 1.79$, $\mu_n = -1.91$. In determining $f_{\text{GT}}^0/f_{\text{GT}}^1$ we used SU(6). This is probably reasonable since SU(6) accurately predicts the F/D ratio for the axial charged current.[8,9] The quantities $f_1^\tau, f_2^\tau, f_{\text{GT}}^\tau$ ($\tau = 0.1$) play for the neutral currents the role played by the Fermi, magnetic, and Gamow-Teller coupling constants for the usual charged weak currents. There already exist important constraints upon these (or equivalently upon α_V^τ and α_A^τ) arising from fits to the available high energy data for elastic[1,2] as well as inclusive[10-12] reactions. Unfortunately these constraints do not uniquely determine[13-15] the basic quantities in (4b).

In the context of the (unified) gauge theories, there exist in the literature various models predicting the parameters $\alpha_V^\tau, \alpha_A^\tau$. As a guide to the experimentalists for the expected cross sections we give below these predictions. The currently fashionable models usually satisfy an $SU(2)\times U(1)$ algebra for the weak and electromagnetic interactions. The essential parameters in these are a, b, internally fixed within each model, as well as the Weinberg angle θ_w and the ratio $\rho = (m_w/m_z\cos\theta_w)^2$ which are determined by fitting to the (high energy) experimental data. In terms of these, the constants appearing in (3) are

$$\begin{aligned}\alpha_V^1 &= \tfrac{1}{2}\rho(2 + a + b - 4\sin^2\theta_w),\\ \alpha_A^1 &= \tfrac{1}{2}\rho(2 - a - b),\\ \alpha_V^0 &= \tfrac{1}{2}\rho(a - b - \tfrac{4}{3}\sin^2\theta_w),\\ \alpha_A^0 &= \tfrac{1}{2}\rho(b - a).\end{aligned} \tag{6a}$$

In Table I we give the values of a, b, for the various $SU(2)\times U(1)$ models considered,[13-17] as well as some best values for θ_w and ρ.

In addition to these we also give results for two $SU(2)_L \times SU(2)_R \times U(1)$ models. The first has been suggested by Mohapatra and Sidhu (MS) and is characterized by no parity violation in atomic inter-interactions.[18] It implies

TABLE I. $SU(2)\times U(1)$ gauge theory models for neutral currents. See Eqs. (6a) in the text.

Model	a	b	$\sin^2\theta_w$	$\rho = \left(\frac{m_w}{m_z\cos\theta_w}\right)^2$
WS_1 (Ref. 16)	0	0	$\frac{3}{8}$	1
WS_2 (Refs. 15,16)	0	0	0.3	1
WP (Ref. 13)	1	0	$\frac{3}{8}$	1
GWP (Ref. 13)	$\frac{3}{4}$	0	$\frac{3}{8}$	1
F (Refs. 15–17)	0	1	0	0.6

$$\alpha_V^1 = \frac{\sin^2\theta}{1-\epsilon},$$

$$\alpha_A^1 = \frac{1}{1-\epsilon},$$

$$\alpha_V^0 = -\frac{\cos^2\theta}{3(1-\epsilon)}, \tag{6b}$$

$$\alpha_A^0 = 0,$$

where $\cos^2\theta \simeq 0.3$ and $\epsilon \simeq 0.1$ from the fit to the high energy data.[18] The second has been suggested by De Rújula, Georgi, and Glashow (DGG).[19] It implies

$$\alpha_V^1 = \frac{\cos^2\beta}{\cos^2\alpha}(1+\sin^2\alpha)(1-2\sin^2\vartheta),$$

$$\alpha_A^1 = \cos^2\beta, \quad \alpha_A^0 = 0, \tag{6c}$$

$$\alpha_V^0 = -\tfrac{2}{3}\sin^2\vartheta(1+\sin^2\alpha)\frac{\cos^2\beta}{\cos^2\alpha}.$$

In the DGG model we present below results for $\sin^2\vartheta \simeq 0.3$ and two possible solutions implied by the high energy data, i.e., $B = (\sin^2\beta = 0.5, \sin^2\alpha = 0.58)$ and $A = (\sin^2\alpha = 0.1, \sin^2\beta = 0.05)$.

III. FORMALISM FOR CHARGED CURRENTS

The effective operators J_μ^+ and J_μ^-, responsible for the neutrino (antineutrino) production of negatively (positively) charged leptons, are obtained from those of neutral currents by setting

$$\alpha_V^1 = \alpha_A^1 = 1, \quad \alpha_V^0 = \alpha_A^0 = 0, \quad \tfrac{1}{2}\tau_3 \to \tau^\pm.$$

In the case of charged current reactions we will consider only electron neutrinos. The final leptons will then be treated as ultrarelativistic, since

$$E_e = E_\nu - (B + E_x) \gg 2m_e c^2,$$

where B is the binding energy difference between the initial and final nuclei, and E_x the excitation energy of the final nuclear state (see Fig. 1). For ultrarelativistic electrons the Coulomb distortion of the final lepton wave function, in the field of a nucleus with $Z \sim 7$, is small.[20] Even for neutrino energies as low as $E_\nu \simeq 40$ MeV, the distortions of the wave functions were found[10] to be $\lesssim 5\%$. This means that the lepton wave functions may be treated as plane waves, without introducing errors larger than 10% in the cross sections. This further means that the Foldy-Wouthuysen reduction of the hadronic charged current is similar to that of the neutral current.[6]

The above considerations are not expected to be good for muonic neutrons at $E_\nu \lesssim 250$ MeV, which is the most interesting energy region for studying neutral currents with nuclear physics experiments.[6] Therefore muonic neutrino cross sections were not included in the present work.

IV. SENSITIVITY OF CROSS SECTIONS TO NUCLEAR MODELS

The conserved vector current (CVC) hypothesis has played a crucial role in the development of weak interactions. This theory relates the vector part of the weak charged current $V_\mu^\pm$ to that of the isovector electromagnetic current V_μ^3 as follows:

$$V_\mu^\pm = \mp[\tau^\pm, V_\mu^3]. \tag{7}$$

Thus $V_\mu^\pm$ is obtained from V_μ^3 by a mere rotation in isospin space. It is, of course, well known that the magnetic dipole part of the electromagnetic current contains a contribution from the orbital motion of the nucleon. If CVC is true, this implies that there should be a similar contribution in the weak vector current. This point seems to have been missed by the early authors concerned with CVC.[21,22] It has, however, been taken into account in the form factor treatment of Holstein.[23] It can most easily been seen by performing a nonrelativistic Foldy-Wouthuysen transformation of the interaction (4), keeping only terms which contribute up to linearly in the momentum transfer. We thus get

$$\mathcal{L}_{\text{int}} = \left(-f_{\text{GT}}^1\vec{\sigma} - \frac{f_1^1+f_2^1}{2m} i\vec{k}\times\vec{\sigma} + \frac{f_1^1}{m}\vec{p}\right)\cdot\vec{L}e^{+i\vec{k}\cdot\vec{r}}, \tag{8a}$$

where $-\vec{k}$ is the electron momentum transfer, $\vec{p}$ is the nucleon momentum operator, and $\vec{L}e^{+i\vec{k}\cdot\vec{r}}$ is the lepton current. Performing a multipole expansion of the exponential, keeping again terms linear in $\vec{k}$, Eq. (8a) becomes

$$\mathcal{L}_{\text{int}} = [-f_{\text{GT}}^1 j_0(kr)\vec{\sigma} + i\vec{\mu}\times\vec{k}]\cdot\vec{L}, \tag{8b}$$

where

$$\vec{\mu} = \frac{1}{2m}\left[(f_1^1+f_2^1)j_0(kr)\vec{\sigma} + f_1^1\frac{3j_1(kr)}{kr}\vec{l}\right] \tag{8c}$$

(the isospin operator $\vec{t}$ is understood) and $j_0(kr)$ and $j_1(kr)$ are the well known spherical Bessel functions. If now one makes the long-wavelength approximation $j_0(kr)\to 1, j_1(kr)\to \frac{1}{3}kr$, then $\vec{\mu}$ becomes the standard magnetic moment operator. Similar approximations hold also in the case of the standard β decay.[24]

The differential cross section for inelastic scattering of neutrinos (antineutrinos) to individual states, including terms linear in $\vec{k}$, takes the form

$$\frac{d\sigma}{d\Omega} = \left(\frac{G}{2\pi}\right)^2 \frac{E_f P_f}{2J_i+1}\Big[(f_{\text{GT}}^1)^2|F_A(k^2)|^2(1-\xi/3) \pm \tfrac{4}{3}f_{\text{GT}}^1\tfrac{1}{2}(f_1^1+f_2^1)(2E_\nu - \Delta E) \times F_A(k^2)F_M(k^2)(1-\xi)\Big] \tag{9}$$

(+ for neutrinos, − for antineutrinos). Here $\xi = \hat{\nu}\cdot\hat{P}_l'$ and $E_l', \vec{P}_l'$ are the final lepton energy and momentum, ΔE is the energy transfer, and $F_A(k^2), F_M(k^2)$ are the two relevant nuclear form factors defined by

$$F_A(k^2) = \langle J_F \| J_0(kr)\vec{\sigma}t_q \| J_i\rangle, \tag{10a}$$

$$F_M(k^2) = \left\langle J_F \left\| \left[J_0(kr)\vec{\sigma} + \frac{f_1}{f_1+f_2}\,\frac{3J_1(kr)}{kr}\,\vec{l}\right]t_q \right\| J_i \right\rangle, \tag{10b}$$

where t_q is the relevant isospin operator. The orbital contribution[25-27] for strong Gamow-Teller-like transitions is negligible. In our case it represents at most 15% of the spin magnetic moment contribution which in itself is a small correction at low energies. Thus we can write $F_M(k^2) \simeq F_A(k^2)$, i.e., the differential cross section contains only one form factor, namely $F_A(k^2)$. The form factors for the transitions to the members of the same isomultiplet are expected to be the same except for small isospin violating effects and the slightly different momentum transfers involved. Thus the ratio of the differential cross sections for neutral to charged events is independent of the nuclear model. In the experiments considered here, however, one is interested in total cross sections. Since the spin and magnetic moment form factors are multiplied by different angular functions, i.e., $1-\xi/3$ and $1-\xi$, respectively, the ratio of the integrated cross sections will not be completely independent of the nuclear model. Even in this case, however, the normalization of the form factor cancels out in the ratio and only a nuclear physics dependence through the shape of the form factor may survive. In this case the ratio $R = \sigma(\nu,\nu)/\sigma(\nu,e^-)$ takes the form:

$$R = \frac{(E_\nu')^2}{(E_e'P_e')}\,\frac{(\alpha_A^1)^2 f_{\rm GT}^{cc} + \frac{4}{3}\alpha_A^1\alpha_V^1(2E_\nu - \Delta E)\frac{1}{2}(f_1^{cc}+f_2^{cc})(J_M/J_A)}{f_{\rm GT}^{cc}(J_A^{cc}/J_A) + \frac{4}{3}(2E_\nu - \Delta E)\frac{1}{2}(f_1^{cc}+f_2^{cc})(J_m^{cc}/J_A)}, \tag{10c}$$

where

$$J_A = \int_{-1}^{1} (1-\xi/3)\,|F_A(\vec{k}^2)|^2 d\xi, \tag{10d}$$

$$J_M = \int_{-1}^{1} (1-\xi)\,|F_M(\vec{k}^2)|^2 d\xi \tag{10e}$$

and J_A^{cc} and J_M^{cc} are the corresponding quantities for the charged currents. The ratio of the shapes of the form factors for nucleons in the $1p$ shell of the harmonic oscillator takes the simple form

$$J_M/J_A = \frac{\int_{-1}^{1} (1-\xi)[1-\alpha^2/6 + f_1/(f_1+f_2)]^2 e^{-\alpha^2/2} d\xi}{\int_{-1}^{1} (1-\xi/3)(1-\alpha^2/6)^2 e^{-\alpha^2/2} d\xi}, \tag{10f}$$

where

$$\alpha^2 = (\vec{k}c)^2/(\hbar\omega m c^2), \tag{10g}$$

m is the nucleon mass, and $\hbar\omega$ is the parameter of the harmonic oscillator. The value $\hbar\omega = 14$ MeV was chosen to fit the electron scattering data.

We conclude therefore that indeed the ratio of neutral to charged current cross sections is more independent of the nuclear model than either one of them separately.

V. ELASTIC SCATTERING

Up to now the elastic neutrino scattering on a nucleus has been considered out of the experimental feasibility. However, it has recently been pointed to us by Deutsch[7] that it is perhaps possible to detect the recoiling nucleus with a target, like ^{12}C and Pb, which can simultaneously serve as a detector, provided that the recoiling energy is $E_N \gtrsim 2$ MeV. Since the cross section for so high a momentum transfer is very small, only its coherent part, which is proportional to A^2, is going to be significant. Thus from such measurements one can extract the isoscalar vector coupling constant α_V^0 from light nuclei ($N \simeq Z$), and perhaps also the isovector vector coupling constant α_V^1 from nuclei with a large neutron excess. In fact these processes provide the best way to extract the values of α_V^0, α_V^1 from neutrino reactions. Other possibilities to measure α_V^0, α_V^1 consist in using the parity changing nuclear transitions in inelastic neutrino scattering, or the small vector current contribution in the Gamow-Teller processes.[6] However, in both of these latter cases the extraction of α_V^T is not so reliable as in the case of elastic neutrino scattering. Notice also that α_V^T can also be extracted from atomic physics experiments,[4] and it is interesting to see whether these two methods agree with each other.

Below we estimate the elastic cross section induced by neutrinos with energy E_ν leading to a recoiling nucleus with energy E_N greater than a given value E_0. This cross section is a rapidly decreasing function of E_0 for fixed E_ν, and an increasing function of E_ν for fixed E_0. For p-shell nuclei with $N=Z$ it takes the form

$$\sigma(E_\nu,E_0)=(\alpha_V^0)^2\frac{G^2}{2\pi}A^2(E_\nu-E_0)^2I(A_s,A,E_\nu,E_0) \tag{11a}$$

with

$$I(A_s,A,E_\nu,E_0)=\frac{9}{4}\int_{-1}^{\xi_1}(1+\xi)\left[1-\left(1-\frac{A_s}{A}\right)\frac{\alpha^2}{6}\right]^2 \times e^{-\alpha^2/2}d\xi, \tag{11b}$$

where $A_s=4$ is the number of s-shell nucleons, $\xi=\hat{\nu}\cdot\hat{\nu}'=\cos\varphi$ (φ is the angle between incoming and outgoing netrinos) and α^2 is given in (10g), and

$$\xi_1=\frac{E_\nu{}^2+(E_\nu-E_0)^2-2Amc^2E_0}{2E_\nu(E_\nu-E_0)}.$$

For a fixed E_0 there is a threshold neutrino energy given by $(E_\nu)_{th}=[\frac{1}{2}(E_0Amc^2)]^{1/2}$, from which for ^{12}C we get $(E_\nu)_{th}=106$ and 118.5 MeV, for $E_0=2$ and 2.5 MeV, respectively. From (11) we then obtain $\sigma(140,2)=6.75$ and $\sigma(140,2.5)=1.53$ in units of $(\alpha_V^0)^2\times10^{-40}$ cm^2 for $E_\nu=140$ MeV. The same quantities for $E_\nu=250$ MeV become $\sigma(250,2)=19.8$ and $\sigma(250,2.5)=8.1$. When the values of α_V^0 for the various models are introduced, these cross sections turn out to be two orders of magnitude smaller than the already measured high energy νp and $\bar{\nu}p$ elastic scattering processes. This is partly due to the smallness of the parameter α_V^0, but mainly due to the phase space restrictions imposed by $E_N \geq E_0$. In fact due to this restriction the above cross sections represent (written in the same order) only 1.6%, 0.35%, 4.5%, and 1.9% of the respective total cross sections. Thus the total elastic neutrino cross sections for ^{12}C are 10^{-39} cm^2, i.e., of the same order of magnitude as the high energy elastic νp, $\bar{\nu}p$ results. Of course, they are expected to be even larger for heavier nuclei.

The elastic cross section on Pb offers also the possibility of measuring the isovector coupling constant α_V^1, once the isoscalar α_V^0 is determined. Since all nucleons participate in the reaction, the relevant form factor is not given reliably by the shell model. It is therefore advantageous to use the charged form factor $F_{ch}(k^2)$, deduced from electron scattering experiments, to calculate this process. Thus Eq. (11b) becomes

$$I(A,Z,E,E_0)=\int_{-1}^{\xi_1}(1+\xi)\left[\tfrac{3}{2}\alpha_V^0-\left(\frac{A-2Z}{2A}\right)\alpha_V^1\right]^2 \times\{F_{ch}(k^2)\}^2d\xi. \tag{12}$$

Since the nucleus is heavy the threshold for the incoming neutrino energy is high; e.g. $(E_\nu)_{th}=442$ and 495 MeV for $E_0=2$ and 2.5 MeV, respectively. The cross section is expected to rise steeply with E_ν and will take large values because the mass number is large.

Noncoherent cross sections arising from all inelastic processes, for the above recoiling nucleus energies, are expected to be A^2 times smaller; i.e., negligible even for light nuclei. Hence, if such events are measured, the extraction of α_V^0 and α_V^1 from (12) should be unambiguous and fairly independent of the nuclear models.

VI. RESULTS FROM INELASTIC NEUTRINO SCATTERING

Even though the ratio of neutral to charged current events is not crucially dependent on the nuclear model adopted, in the calculations we used the best nuclear wave functions one could construct in the framework of present day nuclear shell models.[6,28] These are the wave functions describing the states 1^+, 2_1^+, 2_2^+ and have been discussed previously.[6] For the sake of completeness we repeat their main features here:

$$|1^+I=0(\text{g.s.})\rangle=0.980\,|\bar{p}^2 1^+ I=0\rangle+\cdots,$$
$$|2_1^+I=1\rangle=0.714\,|\bar{p}^2 2^+, I=1\rangle+\cdots,$$
$$|2_2^+I=1\rangle=0.663\,|\bar{p}^2 2^+ I=1\rangle+\cdots,$$

where

$$|\bar{p}^2 1^+ I=0\rangle=0.956\,|L=2s=1\rangle+0.163\,|L=0s=1\rangle +0.247\,|L=1s=0\rangle,$$
$$|\bar{p}^2 2^+ I=1\rangle=0.907\,|L=2s=0\rangle+0.422\,|L=1s=1\rangle,$$

TABLE II. Cross sections $\sigma(\nu,e^-)$ and $\sigma(\bar{\nu},e^+)$ on ^{14}N leading to the 2^+ $I=1$ levels in units of 10^{-40} cm^2 for various neutrino energies in MeV.

State \ E_ν	(ν,e^-) 40	100	140	250	$(\bar{\nu},e^+)$ 40	100	140	250
2_1^+	0.2414	1.0237	1.0549	1.0584	0.2122	0.6056	0.6459	0.6581
2_2^+	0.1971	0.8709	0.8956	0.9588	0.1733	0.5619	0.5791	0.5951
Total	0.4385	1.8946	1.9505	2.0172	0.3855	1.1674	1.2250	1.2531

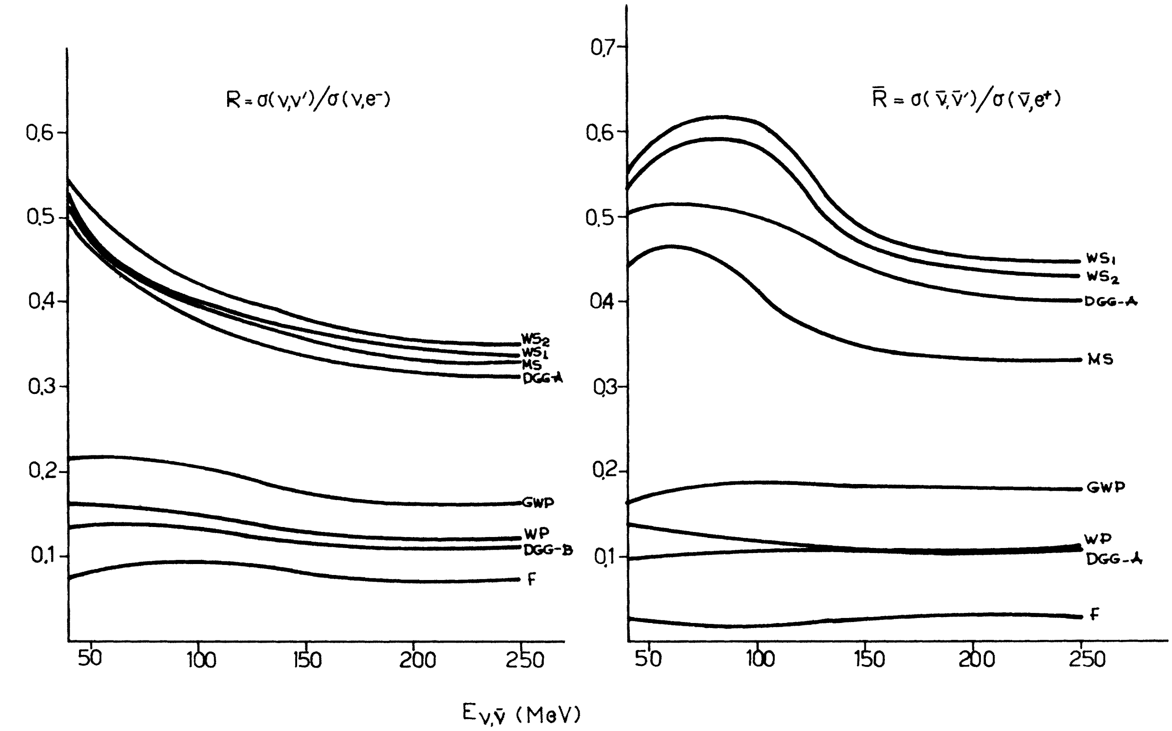

FIG. 2. The ratio $R(\bar{R})$ of neutral to charged current events on ^{14}N leading to the 2_1^+ $I=1$ level for neutrino (antineutrino) scattering. The results for the 2_2^+ $I=1$ state are almost the same, except for small differences near threshold.

and $+\cdots$ stands for core-excited components (4h-2p, 6h-4p, etc.; the 3h-1p excitations are unimportant). Since the g.s. is 96% a two-hole state, the structure of the components indicated by $\cdots$ is immaterial so long as our transition operator is a one-body operator. The above wave functions reproduce quite well all existing experimental data for the relevant states.[6,29] Using them we calculated the isovector transitions indicated in Fig. 1. The results for charged currents are presented in Table II. In Fig. 2 we also present the ratio of the isovector neutral to charged currents sections, for various gauge theory models. At low neutrino energies this ratio is expected to be $\simeq\frac{1}{2}$. For models with $\alpha_A^1 \simeq 1$ we expect this ratio to be close to $\sim\frac{1}{2}$ even at higher energies, with possible deviations arising from the vector current and being such that $R=\sigma(\nu,\nu')/\sigma(\nu,e^-) \lesssim \frac{1}{2}$ and $\bar{R}=\sigma(\bar{\nu},\bar{\nu}')/\sigma(\bar{\nu},e^+) \gtrsim \frac{1}{2}$. The situation is more complicated whenever α_A^1 is substantially less than unity (see Fig. 2).

VII. CONCLUSIONS

We have seen that the ratio of neutral to charged current cross sections in inelastic neutrino scattering is less dependent on the nuclear model than each one of them separately. We have also shown that, by choosing ^{14}N as the target nucleus, charged and neutral current events can be measured simultaneously. Our results indicate that, for isovector transitions, $R\equiv\sigma(\nu,\nu')/\sigma(\nu,e^-)$ and $\bar{R}\equiv\sigma(\bar{\nu},\bar{\nu}')/\sigma(\bar{\nu},e^+)$ are sensitive functions of the parameters α_V^1, α_A^1. So nuclear physics experiments can be very helpful in discriminating among the various gauge theory models. We particularly note in Fig. 2 that there are two clearly distinct groups of models; namely (WS_1, MS, WS_2, DGG-A) and (WP, F, GWP, DGG-B), which can definitely be discriminated by the suggested experiments.[6] It is amusing to remark that a characteristic feature of the models in the second group is that they include a $(u,b)_R$ doublet which provides a simple explanation of the high y anomaly[30] and the rise[31] of $\sigma(\bar{\nu}N\to\mu^+\ldots)/\sigma(\nu N\to\mu^-\ldots)$ at high energies. No such right hand currents are included in the models of the first group, which have therefore some difficulty with the aforementioned data.[30,31] Of course, even if these experimental results are not confirmed,[32] the second group of models is still acceptable as far as high energy data are concerned, provided that the predicted b quark is sufficiently heavy. From this point of view it seems very interesting that nuclear physics experiments can easily discriminate between these two groups of models, irrespective of the mass of the quark b.

In addition, with somewhat less confidence, the experiments suggested here can discriminate between F[15,17] and (WP, GWP, DGG-B) models.[13,19] Even a distinction between (DGG-B, WP) and GWP appears feasible. On the other hand inelastic neutrino experiments on nuclear targets (for isovector transition) do not seem to be able to distinguish among WS_1, WS_2, MS, DGG-A.

If it turns out to be feasible to measure elastic neutrino scattering by detecting the recoiling nucleus, this experiment will complement those measuring parity violation in atomic physics and will provide an independent determination of α_V^0, α_V^1.

The charged current cross sections presented in Table II may be used for background subtraction in neutral current experiments involving electron neutrinos (antineutrinos), when the resolutions are not good enough. However, these background problems due to the charged currents can be avoided by staying at low energies and using muonic neutrinos. Thus using for example $E_{\nu_\mu} \leq 130$ MeV, the isoscalar transition[6] to the 2^+ $I=0$ level at 7.03 MeV cannot be confused with charged events leading to the 2^+ $I=1$ $E_x = 7.01$ MeV state[32] in ^{14}C (see Fig. 1).

We conclude by saying that neutrino reactions leading to definite nuclear states seem to have measurable cross sections and that when they are done, they will shed light on the structure of the hadronic neutral current.

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