URTI
OF TH
OCUM
AR

CONF-820357--4

LA-UR -82-1627

MASTER

TITLE GAMOW-TELLER STRENGTH FUNCTIONS AND NEUTRINO PROBLEMS

AUTHOR(S) W. C. Haxton

DISCLAIMER

SUBMITTED TO Invited talk presented at the International Conference on Spin Excitations in Nuclei, Telluride, Colorado, March, 1982

Los Alamos Los Alamos National Laboratory
Los Alamos, New Mexico 87545

FORM NO 836 R4
ST NO 2629 5/81

GAMOW-TELLER STRENGTH FUNCTIONS AND NEUTRINO PROBLEMS

W. C. Haxton

Theoretical Division
Los Alamos National Laboratory
Los Alamos, New Mexico 87545
and
Department of Physics
Purdue University
West Lafayette, Indiana 47907

ABSTRACT

A quantitative understanding of spin strengths in nuclei is of vital importance in studies of nuclear double beta decay and in solar neutrino spectroscopy. The current status of these problems is outlined.

INTRODUCTION

Fifty years ago Pauli postulated the existence of a new particle, the neutrino or "little neutron," in order to conserve energy in beta decay, $n \to p + e^- + \bar{\nu}$. The name is suitable, as the neutrino has little or no mass and, like the neutron, spin 1/2 and no charge. Today we know that neutrinos come in three flavors associated respectively with three charged partners: the electron, the muon, and the tauon. Together with these charged partners they participate in a variety of charge-changing weak interactions, including the familiar beta decay.

The discovery a decade ago that a neutrino could scatter off a nucleon without being changed into its charged partner demonstrated the existence of a new class of weak interactions, those mediated by the neutral current. This helped substantiate what is now widely regarded as one of the major theoretical advances in modern physics, the unified description of the weak and electromagnetic interactions in the model of Glashow, Weinberg, and Salam (GWS) [1]. Within this model neutrinos are massless and certain quantities, such as lepton number, baryon number, and muon number, are exactly conserved.

Presently great theoretical effort is being expended in an even more ambitious endeavor, the search for a unified description of the electroweak and strong interactions [2]. Attempts to construct "grand unified" theories have already met with some success in that certain models predict the value of the Weinberg angle, a free parameter in the electroweak theory. The prejudice in grand unified theories for "naturality," the reluctance to postulate a priori global conservation laws, also suggests that many of the conservation laws of the GWS model are in fact only approximate, reflecting the enormous mass scale governing the strong-electroweak unification. The manner in which these conservation laws are broken should impose important constraints on formulations of grand unified theories.

How can we obtain these experimental constraints? The situation is quite different from that which prevailed a decade ago, when the electroweak unification mass ($\sim 10^2$ GeV) appeared directly accessible, stimulating the development of remarkable accelerator technology. The expected grand unification mass ($\sim 10^{15}$ GeV) may condemn us to investigation in the "low energy limit" for some time. Thus increasingly the future of particle physics will depend on the development of technologies to isolate rare events and to measure small branching ratios. Present experiments probing nucleon stability at the level of 10^{32} years may be in the vanguard of this effort [3].

The opportunities are present for nuclear physics to play an important role in this quest. The nucleus can serve as a filter for rare processes, isolating interactions according to spin, isospin, and parity. In Table 1 a few of the ongoing nuclear experiments that may have a profound impact on particle physics are listed. These and other possibilities for fruitful collaborations between the nuclear and particle physics communities remind one of the importance of the β-decay studies that tested the CVC hypothesis and the V-A theory of weak interactions 25 years ago [4,5]. In particular, today I will discuss two problems, ββ-decay and solar neutrino detection, which promise to constrain possible descriptions of the neutrino. These are particularly relevant to the discussions of this conference, as an understanding of spin excitations in nuclei is a prerequisite to their interpretation.

Table 1

	Symmetry/Interaction	Nuclear Test
1.	baryon number	decay of bound nucleon
2.	lepton number: masses and right-handed couplings of Majorana neutrinos	double beta decay
3.	time reversal	nuclear electric dipole moment; nuclear γ-decay
4.	flavor mixing; neutrino masses	neutrino oscillations (e.g., solar neutrino detection)
5.	separate lepton number	$\mu \rightarrow e$ conversion in muonic atoms
6.	$\Delta S = 0$ weak hadronic current	parity mixing of nuclear levels
7.	neutrino mass	tritium β-decay

DIRAC AND MAJORANA NEUTRINOS

One powerful proble of lepton number conservation, of the mass and charge conjugation properties of the electron neutrino, and of possible right-handed admixtures in the weak leptonic current is nuclear $\beta\beta$-decay, $(A,Z) \rightarrow (A,Z+2)$ [6]. This process, a fundamental nuclear decay mode, can be observed in a number of even-even nuclei where, due to the pairing interaction, the competing decay $(A,Z) \rightarrow (A,Z+1)$ is energetically inaccessible (see Fig. 1).

The question historically associated with $\beta\beta$-decay is whether the neutrino should be described by a Dirac or Majorana field. If the neutrino is a Dirac particle, it has a distinct antiparticle; if Majorana, the particle and antiparticle are indistinguishable. The neutrino is unique among the fermions in permitting these alternative descriptions: any fermion having a charge or measurable magnetic moment necessarily has a distinct antiparticle.

If we define the neutrino and antineutrino by

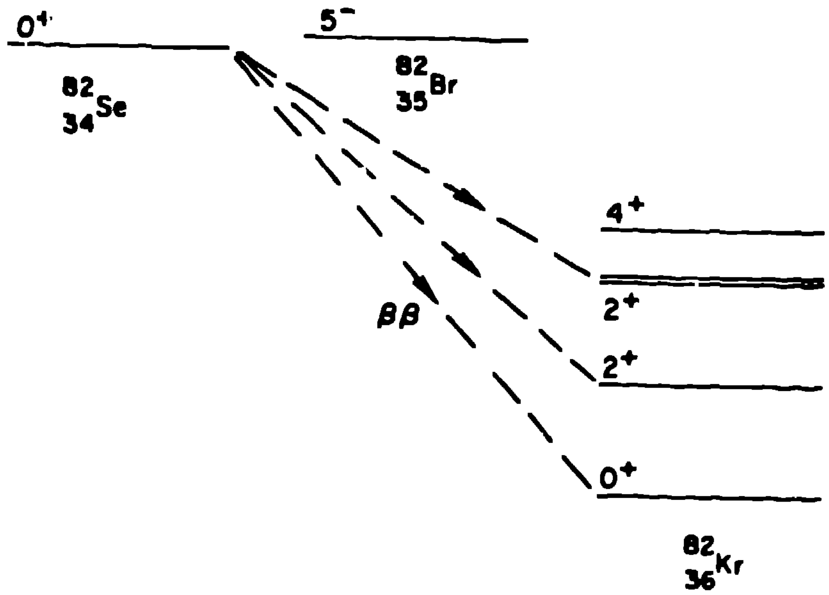

Fig. 1. Level scheme for the ββ-decay of ^{82}Se.

$$n \rightarrow p + e^- + \bar{\nu}_e$$

$$\nu_e + n \rightarrow p + e^- \tag{1}$$

the second order weak interaction

$$2n \rightarrow n + p + e^- + \bar{\nu}_e \rightarrow 2p + 2e^- + 2\bar{\nu}_e \tag{2}$$

will occur. This two-nucleon process will contribute to the decay $(A,Z) \rightarrow (A,Z+2)$, producing a final state with four leptons, regardless of the charge conjugation properties of the neutrino. If the neutrino is a Majorana field, a second decay mode is possible

$$2n \rightarrow n + p + e^- + \bar{\nu}_e \equiv n + p + e^- + \nu_e \rightarrow 2p + 2e^- \tag{3}$$

producing a neutrinoless final state. This second process enjoys a considerable phase space advantage over the reaction in Eq. (2) and, in the absence of the chirality suppression we will discuss momentarily, will dominate ββ-decay rates for a Majorana electron neutrino.

Early geochemical measurements of ββ-decay showed that the half lives for the decay $(A,Z) \rightarrow (A,Z+2)$ far exceeded the values expected for the process in Eq. (3), $\tau_{1/2} \sim 10^{12}$-10^{15} years. This was interpreted as a demonstration of the Dirac character of the electron neutrino, and prompted the introduction of lepton number ℓ to distinguish the neutrino from its antiparticle: the electron and neutrino are assigned $\ell = +1$, the positron and antineutrino $\ell = -1$. The assumption that additive lepton number is conserved then allows two-neutrino ββ-decay, but forbids neutrinoless ββ-decay, for which $\Delta(\Sigma\ell) = 2$.

Yet, with the discovery [5] in 1957 that the weak interaction violates parity conservation maximally (or nearly so), it became apparent that the Majorana/Dirac character of the electron neutrino was still in question. The particles which participate in the reactions of Eq. (1)

$$\begin{aligned} & n + p \rightarrow e^- + \bar{\nu}_e^{(+)} \\ & \nu_e^{(-)} + n \rightarrow p + e^- \end{aligned} \qquad (1a)$$

are the right-handed $\bar{\nu}_e$ and left-handed ν_e. Thus even if the neutrino is a Majorana particle

$$\begin{aligned} 2n & \rightarrow n + p + e^- + \bar{\nu}_e^{(+)} \\ & \equiv n + p + e^- + \nu_e^{(+)} \not\rightarrow 2p + 2e^- \end{aligned} \qquad (3a)$$

as the neutrino has the wrong helicity for absorption on a neutron. Therefore, if parity violation in the weak interaction is sufficiently close to maximal, the geochemical ββ-decay results imply neither a Dirac electron neutrino nor a conserved lepton number.

The great interest today in ββ-decay stems from the gauge theory prejudice [7] that a neutrino mass will break the γ_5-invariance of the weak current. Thus neutrinoless ββ-decay may occur, though at a rate suppressed by $(m_\nu/m_e)^2$, if the neutrino is a Majorana particle. A careful examination of ββ-decay rates then leads to the following conclusions:

(1) Present laboratory limits on neutrinoless ββ-decay place an upper bound on the neutrino mass of $\langle m_\nu^{Maj} \rangle \lesssim 10\text{-}50$ eV. This bound may impose a fundamental constraint on the charge conjugation properties of the neutrino if the tritium β-decay mass result $14 \text{ eV} \le m_\nu \le 52 \text{ eV}$ is correct [8].

(2) There is a hint, in the geochemical total ββ-decay rates for ^{128}Te and ^{130}Te, that no-neutrino decay is occurring. The rate, corresponding to $\langle m_\nu^{Maj} \rangle \sim 10$ eV, does not violate any experimental bound on lepton number violation.

(3) Systematic disagreement exists between the geochemical total $\beta\beta$-decay rates and the two-neutrino rates predicted by theory and measured in a single laboratory experiment. The origin of the discrepancy is unclear.

I would now like to summarize the experimental and theoretical work that leads to these results.

NUCLEAR DOUBLE BETA DECAY RATES

Our knowledge of $\beta\beta$-decay rates comes from two classes of experiments, geochemical and laboratory.

Geochemical measurements have determined the total $\beta\beta$-decay half lives for ^{130}Te, ^{128}Te, and ^{82}Se, as shown in Table 2. The noble gases are the rarest of the stable nuclides. Thus, over geologic times these reactions can produce significant elevations in the abundances of the daughter nuclei. The experimental procedure consists of outgasing by stepwise heating of Te- or Se-bearing ore samples, followed by high sensitivity mass spectrometry. The excess of the daughter isotope is determined by comparing the resulting noble gas isotopic distribution to that for the atmosphere. Once the ore age is fixed by geologic arguments or by K-Ar dating, this excess determines the total $\beta\beta$-decay rate.

Laboratory experiments have provided bounds on 2ν and 0ν $\beta\beta$-decay and, in one case, a 2ν half life. Of course, only the electrons are detected. A plot of the sum of the electron kinetic energies T is shown in Fig. 2. For 0ν $\beta\beta$-decay a spike is found at $T = T_o$, the total kinetic energy release; for 2ν decay, the distribution is continuous over the range from $T = 0$ to T_o.

Table 2: A summary of geochemical $\beta\beta$-decay results. The total kinetic energy carried off by leptons, T_o, is given in units of the electron mass.

Reaction	T_o ($m_e c^2$)	$\tau_{1/2}$ (years)	
^{130}Te → ^{130}Te	5.0	$(2.0 - 3.1)\cdot 10^{21}$	[9,10]
^{128}Te → ^{128}Xe	1.7	$(3.2 - 4.9)\cdot 10^{24}$	[11]
^{82}Se → ^{82}Kr	5.9	$2.76 \cdot 10^{20}$	[12]

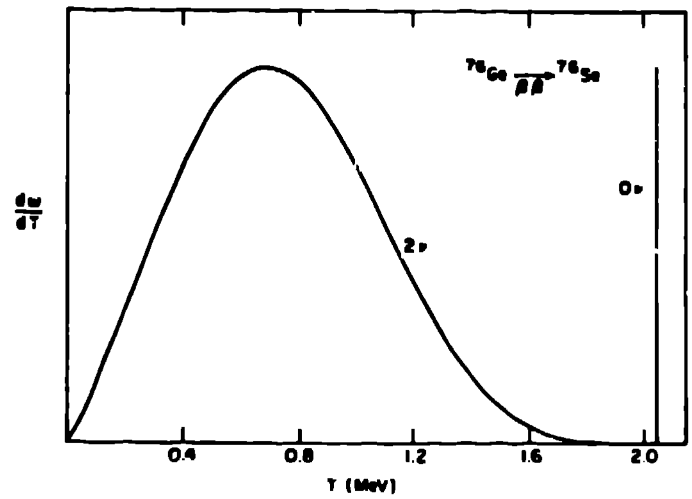

Fig. 2. Comparison of the differential decay rate dw/dT, where T is the sum of the kinetic energies carried off by the electrons, for 0ν and 2ν ββ-decay. The 0ν spectrum is a line at $T = T_o$.

Table 3: Laboratory Limits on 2ν and 0ν ββ-decay.

Reaction	T_o $(m_e c^2)$	$\tau_{½}$ (years)	
$^{48}Ca \rightarrow {}^{48}Ti$	8.4	$\gtrsim 10^{21.3}$, 0ν	[13]
		$\gtrsim 10^{19.56}$, 2ν	[13]
$^{76}Ge \rightarrow {}^{76}Se$	4.0	$\gtrsim 10^{21.7}$, 0ν	[14]
$^{82}Se \rightarrow {}^{82}Kr$	5.9	$\gtrsim 10^{21.5}$, 0ν	[15]
		$10^{19.0 \pm 0.2}$, 2ν	[16]

The experimental task of discerning signal from background is thus simpler in the case of 0ν decay, accounting for the stringent limits shown in Table 3.

The remaining task is to compare these results with the theoretical predictions for ββ-decay mediated by Dirac and Majorana neutrinos. If the neutrino is Dira

$$\omega = \omega_{2\nu} \tag{4a}$$

and if Majorana

$$\omega = \omega_{2\nu} + \omega_{0\nu}(\eta, \langle m^{Maj}\rangle_\nu) \quad . \tag{4b}$$

(I have allowed breaking of the γ_5-invariance of the weak leptonic current by a mass $\langle m^{Maj}\rangle_\nu$, as discussed earlier, and by an explicit right-handed current of strength η.) Clearly the experimental limits on 0ν $\beta\beta$-decay constrain the mass and right-handed coupling of a Majorana electron neutrino provided one can calculate $\omega_{0\nu}$ as a function of η and $\langle m^{Maj}\rangle_\nu$. One can also use total geochemical rates to constrain these parameters.

The lepton number-conserving process of Eq. (2) gives rise to the nuclear decay shown in Fig. 3a. The corresponding decay rate can be evaluated in time-dependent perturbation theory. If the dependence of the energy denominator on the energy of the intermediate nuclear state is approximated by an average value [6]

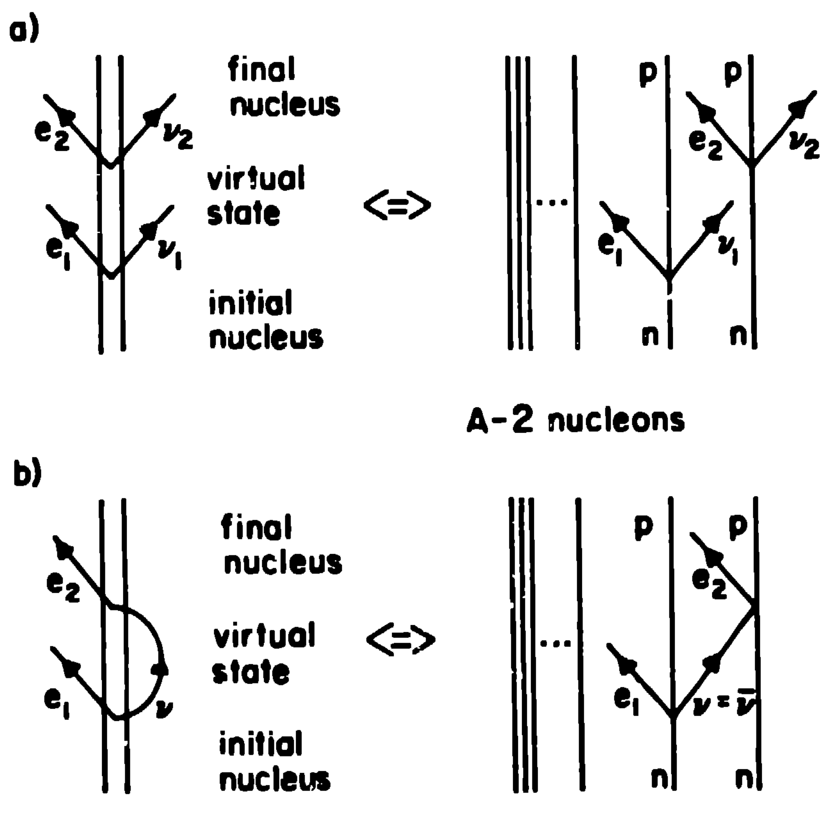

Fig. 3. Two-nucleon mechanisms for 2ν (a) and 0ν (b) $\beta\beta$-decay.

$$\frac{1}{E_N - E_I + \nu + \varepsilon} \approx \frac{1}{\langle E_N - E_I \rangle + \nu + \varepsilon} \tag{5}$$

the sum over virtual nuclear states can be compl ed by closure. [In Eq. (5) E_N and E_I represent the energies of the intermediate and initial nuclear states, while ν and ε are the energies of the neutrino and electron emitted in the first β-decay.] Evaluating each β-decay in the allowed approximation and specializing to transition between $J^\pi = 0^+$ nuclear states, one finds

$$\omega_{2\nu} = \xi_{2\nu} \frac{16G^4}{\pi^7} [F^{PR}(Z + 2)]^2 \frac{1}{(\langle E_N - E_I \rangle + T_o/2 + m_e)^2} \frac{m_e^{11}}{8!} f(T_o/m_e)$$

$$\times [F_1^4 |M_F|^2 + F_A^4 |M_{GT}|^2 - 2F_1^2 F_A^2 \mathrm{Re}(M_F \cdot M_{GT}^*)] \tag{6a}$$

where

$$F^{PR}(Z) = \frac{2\pi\ \alpha\ Z}{1 - \exp(-2\pi\alpha Z)} \tag{6b}$$

$$f(\varepsilon) = \varepsilon^7 [1 + \frac{\varepsilon}{2} + \frac{\varepsilon^2}{9} + \frac{\varepsilon^3}{90} + \frac{\varepsilon^4}{1980}] \tag{6c}$$

$$M_F = \langle F | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j) | I \rangle \tag{6d}$$

$$M_{GT} = \langle F | \tfrac{1}{2} \sum_{ij} \vec{\sigma}(i)\cdot\vec{\sigma}(j)\tau_+(i)\tau_+(j) | I \rangle \tag{6e}$$

The factor $\xi_{2\nu}$ is a phase space correction needed because Eq. (6a) has been written with the approximate Coulomb correction of Eq. (6b); it varies from 1.2 to 5, roughly, between Z = 22 and Z = 54 [17]. The double Fermi matrix element M_F is nonzero only through isospin impurities in the final state, and quite generally can be neglected.

The lepton number-violating decay of Fig. 3b can occur via the two-nucleon process of Eq. (3). The rate can be calculated for the following choice of the leptonic current

$$j_\mu^{lep}(x) = \bar{\psi}_e(x)\gamma_\mu[(1 - \gamma_5) + \eta(1 + \gamma_5)]\psi_\nu^{Maj}(x) \quad . \tag{7}$$

The γ_5-invariance is broken by an explicit right-handed current admixture η and by a Majorana neutrino mass $\langle m^{Maj} \rangle_\nu$. The general result for the decay rate for transitions between $J^\pi = 0^+$ states is somewhat complicated [17,18]. We given only the rate for $\eta = 0$.

$$\omega_{0\nu}^{\eta=0} = \xi_{0\nu} \frac{G^4}{8\pi^5} [F^{PR}(Z+2)]^2 m_e^7 \left(\frac{\langle m^{Maj}_\nu\rangle}{m_e}\right)^2$$

$$f_1(T_o/m_e)|F_A^2 M_2'' - F_1^2 M_1''|^2 \tag{8a}$$

where

$$f_1(\varepsilon) = \varepsilon[1 + 2\varepsilon + \frac{4\varepsilon^2}{3} + \frac{\varepsilon^2}{3} + \frac{\varepsilon^4}{30}] \tag{8b}$$

$$M_2'' = \langle F|\tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\vec{\sigma}(i)\cdot\vec{\sigma}(j) \frac{g(r_{ij})}{r_{ij}}|I\rangle \tag{8c}$$

$$M_1'' = \langle F|\tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j) \frac{g(r_{ij})}{r_{ij}}|I\rangle \tag{8d}$$

with $g(r_{ij}) \simeq 1$ a slowly-varying function of $r_{ij} = |\vec{r}_i - \vec{r}_j|$. In the general result for $\eta \neq 0$ two additional matrix elements

$$M_3 = \langle F|\tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\hat{r}_{ij}\cdot\vec{\sigma}(i)\hat{r}_{ij}\cdot\vec{\sigma}(j) \frac{g(r_{ij})}{r_{ij}} |I\rangle \tag{9a}$$

$$M_4 = \langle F|\tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\hat{R}_{ij}\cdot(\hat{r}_{ij} \times (\vec{\sigma}(i) - \vec{\sigma}(j))) \frac{R_{ij}}{r_{ij}^2} g(r_{ij})|I\rangle$$

with $R_{ij} = |\vec{r}_i + \vec{r}_j|$, also appear. Note that the the matrix elements in Eq. (8a) differ from the 2ν operators M_F and M_{GT} only by their radial dependence, $g(r_{ij})/r_{ij}$.

To evaluate the expressions in Eqs. (6) and (8) one must calculate the two-body nuclear matrix elements. Recently the group at Los Alamos (Haxton, G. J. Stephenson, Jr., and D. Strottman) has tackled this structure problem with state-of-the-art shell model techniques [17,19,20]. Although the length of this talk precludes a detailed description of this work, I will provide a brief summary.

The $\beta\beta$-decay transition $^{48}Ca \rightarrow {}^{48}Ti$ appears extremely favorable: the structure is thought to be relatively simple, and the large kinetic energy release $T_o = 4.3$ MeV promises considerable phase space enhancement of the rate. However, Lawson showed in a simple Nilsson model that $M_{GT} = 0$ as the result of a K selection rule [21]. This constraint is relaxed somewhat in more realistic intermediate coupling models. Our shell model diagonalization was performed with the Kuo-Brown full g-matrix [22] in the 2p1f model space. A closed ^{40}Ca core is assumed, and all configurations of eight valence nucleons for which the $1f_{7/2}$ occupation is at least four are allowed.

The treatment of the decays $^{76}Ge \rightarrow {}^{76}Se$ and $^{82}Se \rightarrow {}^{82}Kr$ is more complicated. A direct shell model calculation in the canonical model space, involving the orbitals $1f_{5/2}$, $2p_{3/2}$, $2p_{1/2}$, and $1g_{9/2}$ lying between the magic numbers 28 and 50, is not feasible. We instead employ a weak coupling approximation in which full shell model calculations are performed separately for the valence protons and neutrons [19]. The proton-neutron interaction is then diagonalized in a basis formed from the 50 proton and 50 neutron wave functions lowest in energy, yielding wave functions of the form

$$\psi^{pn}_{J=0} = \sum_{i,j=1}^{50} \xi_{ij} [\psi^{p}_{J_i \gamma_i} \otimes \psi^{n}_{J_j \gamma_j}]_{J=0} \quad . \tag{10}$$

We again employ the Kuo g-matrix. It should be noted that certain spin partners, $1f_{7/2}$ and $1g_{7/2}$, are outside the model space. Naively one expects the influence of these subshells to be small. Furthermore, inclusion of these orbitals would introduce spurious center-of-mass wave function components that could have serious effects, as the g-matrix is not translationally invariant.

The calculations for the decays $^{130}Te \rightarrow {}^{130}Xe$ and $^{128}Te \rightarrow {}^{128}Xe$ were also performed in a weak coupling basis. The model space includes the orbitals $1g_{7/2}$, $2d_{5/2}$, $2d_{3/2}$, $3s_{1/2}$, and $1h_{11/2}$ lying between the magic numbers 50 and 82, and the interaction employed is that of Baldridge and Vary [23]. The separate proton and neutron calculations involve sufficiently many configurations that some restrictions must be imposed on the occupation of the less favored orbitals.

The reliability of our limits on η and $\langle m^{Maj}_{\nu} \rangle$ will depend on the quality of these wave functions. There is one obvious check suggested by the similarity of the 0ν matrix elements to those governing 2ν $\beta\beta$-decay: do these wave functions properly reproduce the 2ν decay rates?

The results shown in Table 4 are surprising. Theoretical and laboratory rates for 2ν $\beta\beta$-decay in ^{48}Ca and ^{82}Se are in good agreement. However, the upper bounds that can be placed on $|M_{GT}|$ from total geochemical rates are consistently much smaller than values predicted by theory. The laboratory and geochemical rates are also in sharp disagreement for the one case permitting a comparison $^{82}Se \rightarrow {}^{82}Kr$. The large theoretical matrix elements for ^{76}Ge, ^{82}Se, ^{128}Te, and ^{130}Te come about through a coherent addition of many amplitudes in the two-body density matrix. Very recently Zamick and Auerbach [24] have obtained similar theoretical values for M_{GT} in ^{48}Ca and ^{76}Ge using a Nilsson model with pairing. Importantly, they attribute the coherence found in the shell model to pairing, and demonstrate that large matrix elements result in their treatment for any reasonable choice of the pairing strength.

What is the reason for this disagreement? Perhaps the most troublesome aspect of the theoretical treatment is the replacement of the E^{-1}-weighted sum over intermediate nuclear states by the non-energy-weighted sum. Yet both the coherence described above and tests involving explicit summation over low-lying intermediate states [19] indicate that no significant bias is introduced by the closure approximation. A possibility of great interest in view of the discussions at this conference, the coupling to delta-hole excitations, had little effect on M_{GT} in the calculations of Zamick and Auerbach [24].

Table 4: Calculated and Experimental Double Gamow-Teller Matrix Elements M_{GT}.

Nucleus	$\|M_{GT}\|_{theory}$		$\|M_{GT}\|_{exp}$	
^{130}Te	1.48	[17]	0.10-0.13*	[9,10]
^{128}Te	1.47	[17]	0.18-0.23*	[11]
^{82}Se	0.94	[19]	1.43	[16]
			0.27*	[12]
^{76}Ge	1.28	[19]		
^{48}Ca	0.22	[20]	$\lesssim$ 0.19	[13]

*Maximum values determined from total geochemical rates.

Alternatively, one can question the geochemical assumptions. Are noble gases retained in the ore over geologic times? The consistency between geologic estimates of the ore age and the results of K-Ar dating demonstrates that this ligher noble gas does remain in the ore [25]. Furthermore, there is reasonable consistency between the geochemical half lives determined from different ores by different geochemists [25]. In summary, we can find no obvious flaw in the geochemistry, in the theoretical treatment, or in the ^{82}Se laboratory experiment of the Moe and Lowenthal. The origin of the inconsistencies in Table IV is unknown.

Observing that the 0ν matrix elements differ from their 2ν counterparts only by the gentle radial dependence $g(r_{ij})/r_{ij}$, Primakoff and Rosen [6] suggested in their early work on $\beta\beta$-decay that a scaling relation might exist between 0ν and 2ν matrix elements, $M_2''/M_{GT} \simeq 1/R_o$ with $R_o = 1.2\,A^{1/3}$ the nuclear radius. Our calculations for ^{76}Ge, ^{82}Se, ^{128}Te, and ^{130}Te demonstrate that this scaling holds remarkably well, though with a somewhat different strength, $M_2''/M_{GT} = (0.57 \pm 0.03/R_o$. One then expects the discrepancies in estimates of 2ν and total $\beta\beta$-decay rates to carry over to 0ν $\beta\beta$-decay. The resulting uncertainties in the bounds on $\langle m^{Maj}\rangle_\nu$ and η for the $\beta\beta$-decay of ^{82}Se are apparent in Fig. 4.

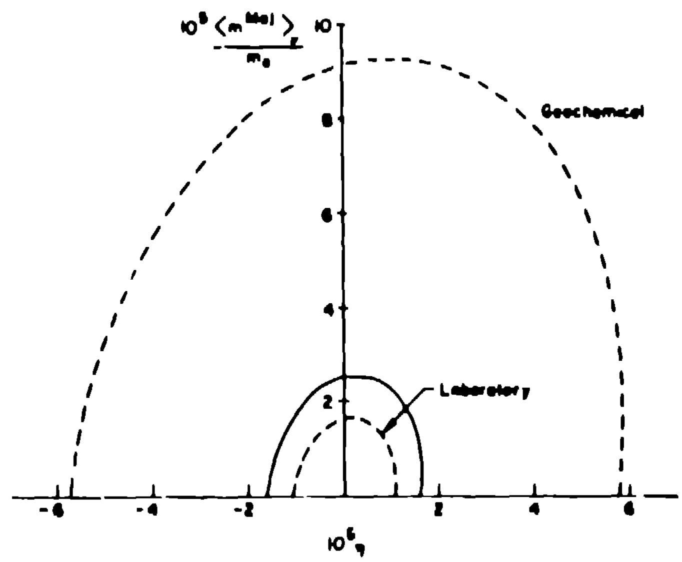

Fig. 4. Dashed lines show boundaries of allowed regions in η - $\langle m^{Maj}\rangle_\nu$ plane which result if all ^{82}Se $\beta\beta$-decay matrix elements are normalized to reproduce the total geochemical rate [12] and the laboratory 2ν rate of Moe and Lowenthal [16]. Solid line employs theoretical matrix elements.

Pontecorvo suggested that the question of matrix element normalization might be circumvented by comparing ββ-decay rates for two different nuclei, ^{128}Te and ^{130}Te. As the structure of these isotopes differs only by a neutron pair, Pontecorvo assumed that the ββ-decay nuclear matrix elements would be identical. The Los Alamos calculations (see Table IV) and those of Vergados [26] support this assumption. The ratio of the total ββ-decay rates for these isotopes should then be determined by phase space. Because the energy releases for these decays are quite different (see Table 2), one finds

$$\frac{\tau_{1/2}^{2\nu}(128)}{\tau_{1/2}^{2\nu}(130)} = 5100$$

and

$$\frac{\tau_{1/2}^{0\nu}(128)}{\tau_{1/2}^{0\nu}(130)} = 25$$

so that the ratio of half lives tests sensitively the decay mechanism. The experimental result [11]

$$\frac{\tau_{1/2}(128)}{\tau_{1/2}(130)} = 1590$$

suggests that both 2ν and 0ν ββ-decay may be contributing. The values for η and $\langle m^{\mathrm{Maj}}_{\nu}\rangle$ that are consistent with the experimental ratio can be derived given only the relative strength of the 2ν and 0ν matrix elements. The solution is shown in Fig. 5.

I believe that this purported evidence for lepton number violation is really quite weak. It is difficult to accept the theoretical demonstration of matrix element equality in view of the alarming discrepancy between theory and geochemistry in the absolute rates. Furthermore, Kirsten has recently reported [25] a measurement of the half life ratio which is inconsistent with the older Missouri group value [11] given above and consistent with lepton number conservation. While Kirsten's measurement was made with relatively young ore, and thus his statistical error exceeds that of the Missouri group, certainly his result indicates that the experimental situation is unsettled.

In summary, the laboratory limits on 0ν ββ-decay yield $\langle m^{\mathrm{Maj}}_{\nu}\rangle \lesssim (9\text{-}52)$ eV, with the range reflecting the discrepancy between the

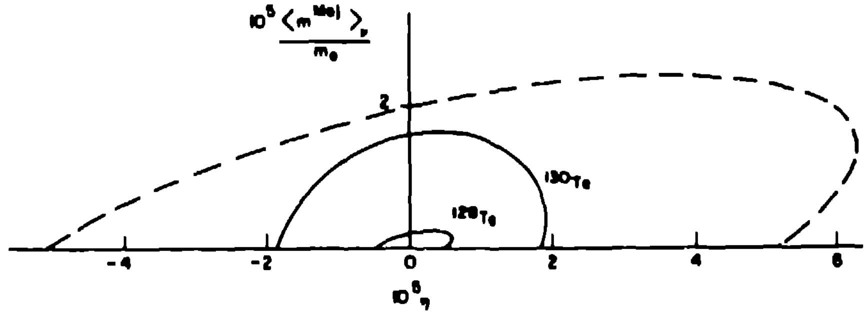

Fig. 5. Values of η and $\langle m^{Maj}\rangle_\nu$ lying on the dashed line give, under the assumption of equal matrix elements in ^{130}Te and ^{128}Te, the ratio of geochemical rates measured by the Missouri group [11]. The solid lines show the boundaries of the allowed region derived by using theoretical matrix elements and attributing the entire decay rate to 0ν $\beta\beta$-decay. The inconsistency between solid and dashed lines reflects the disagreement between theory [17,19] and the geochemical results [9,10].

laboratory and geochemical measurements for ^{82}Se. If the magnitudes of matrix elements are taken from the Los Alamos calculations, $\langle m^{Maj}\rangle_\nu \lesssim 13$ eV; in this case one finds that the limits on 0ν decay in ^{48}Ca and ^{76}Ge impose somewhat less stringent constraints on $\langle m^{Maj}\rangle_\nu$. The geochemical rates in ^{130}Te and ^{128}Te give, under the Pontecorvo assumption of equal matrix elements, $\langle m^{Maj}\rangle_\nu \lesssim 10$ eV. (We choose the inequality because of the conflict between the results of Kirsten and the Missouri group.)

Thus a cautious interpretation of these results indicates $\langle m^{Maj}\rangle_\nu \lesssim 50$ eV. If, in addition, one chooses to believe the Pontecorvo assumption, the result of Moe and Lowenthal, or the theoretical matrix element calculation, then the more stringent constraint $\langle m^{Maj}\rangle_\nu \lesssim 10$ eV follows. Recently, a measurement of the endpoint spectrum in the β-decay of the triton yielded $14 < m_\nu \lesssim 46$ eV [8]. If this is correct, then the more stringent $\beta\beta$-decay limit demonstrates that the electron neutrino cannot be a

Majorana mass eigenstate! This exciting result underscores the importance of extending current 0ν $\beta\beta$-decay limits one to two orders of magnitude as, under the most cautious interpretation, such results could test the charge conjugation properties of the neutrino.

Finally, I would like to mention a few topics that may convey some of the flavor of present studies in $\beta\beta$-decay. Rosen [27] and Doi et al. [18] recently pointed out that $0^+ \to 2^+$ $\beta\beta$-decay transitions are of particular interest in that the 0ν mechanism can occur only via the right-handed current. Unfortunately, according to our Los Alamos work, the matrix elements that govern such transitions are quite weak, so that $0^+ \to 0^+$ transitions provide much more stringent constraints on η. If, however, 0ν decay is observed, $0^+ \to 2^+$ transitions may prove a valuable tool for separating mass effects from those of the right-handed current.

There has also been considerable discussion of mechanisms involving the $\beta\beta$-decay transition $\Delta_{33} \to n$ within the nucleus [18,28, 29]. It can be shown, in the allowed approximation, that this amplitude vanishes for $0^+ \to 0^+$ 2ν and, in the SU(6) limit, 0ν $\beta\beta$-decay. The possibility of strong $\Delta_{33} \to n$ $0^+ \to 2^+$ 0ν transitions is presently under study. Other lepton number-violating $\beta\beta$-decay mechanisms that have been discussed include Majoron production [30] and Higgs exchange [31] (Fig. 6). Recent work indicates that the

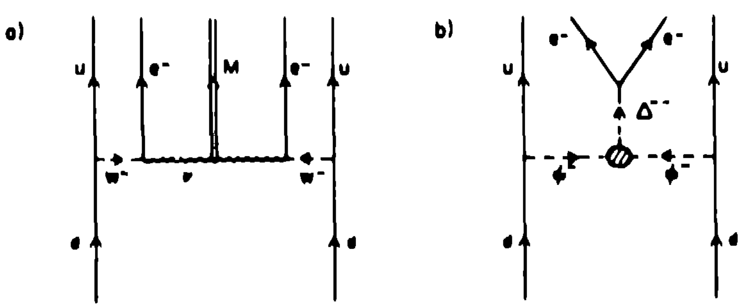

Fig. 6. Mechanisms for $\beta\beta$-decay involving (a) Majoron production [30] and (b) Higgs exchange [20,31,32].

Higgs exchange mechanism is much less important than originally believed [20,32]. Majoron production poses a difficult problem experimentally, as this light scalar would carry off kinetic energy, leaving an electron energy distribution that would be difficult to distinguish from that for 2ν decay.

We have not discussed the definition of $\langle m^{Maj}\rangle_\nu$ except in the case that the neutrino is a mass eigenstate. Wolfenstein [33] has shown that in CP-invariant theories where multiple Majorana neutrinos couple to the electron

$$\langle m^{Maj}\rangle_\nu = \sum_i C_i^2 N_i m_i^{Maj}$$

where N_i is the CP eigenvalue and $\sum_i C_i \leq 1$. Thus it may be a mass difference that is constrained in ββ-decay, and this quantity then is not simply related to that mass measured at the tritium ββ-decay endpoint. Doi et al. [18] have considered more general mass matrices arising in CP-noninvariant theories.

Finally, in view of the general concerns of this conference, there is an interesting possibility that rigorous upper bounds can be placed on $|M_{GT}|$ by measuring the GT strength distributions in the intermediate nucleus from both parent and daughter. Thus ββ-decay nuclei may be attractive candidates for °0 (p,n) and (n,p) studies [34]. Also, such GT strength distributions may provide important tests of the nuclear wave functions presently employed in ββ-decay studies.

SOLAR NEUTRINOS

I will now briefly discuss the solar neutrino puzzle and the importance of GT strength measurements to future plans for neutrino spectroscopy.

To date only a single solar neutrino experiment, the ^{37}Cl experiment of Ray Davis, Jr., and collaborators, [35] has been mounted. The resulting capture rate, 1.95 ± 0.3 SNU, is in serious disagreement with the predictions of the standard solar and weak interaction models, 8.0 SNU [36] (1 SNU = 10^{-36} captures/^{37}Cl atom/s). If this discrepancy is due to a misunderstanding of the physics of the solar interior, the implications for present theories of stellar evolution could be profound [37]. Alternatively, if the sun does produce the expected neutrino flux, then some mechanism must be altering the character of those neutrinos before they reach earth. This suggestion now seems particularly interesting in view of recent evidence for massive neutrinos [8] and neutrino oscillations [38].

These two classes of solutions to the ^{37}Cl puzzle can be distinguished. Proposed modifications of the standard model to accommodate the ^{37}Cl capture rate result primarily in a reduced flux of high energy (14 Mev endpoint) ^{8}B neutrinos, whose production depends most critically on the central temperature of the sun. Neutrino oscillations or decay would, except under unusual conditions, affect all components on the solar neutrino flux equally. Thus there has been great interest in mounting new experiments to complete the spectroscopy of the neutrino sources shown in Table 5. Today I would like to describe three possibilities for new experiments that I find particularly interesting. I will emphasize the importance of °0 (p,n) Gamow-Teller (GT) strength measurements in eliminating uncertainties in the neutrino capture cross sections estimates for each of these experiments.

Kuzmin [40] suggested a solar neutrino experiment based on the reaction $^{71}Ga(\nu,e)^{71}Ge$. The calculations of Bahcall [41] and others [42] indicate that the ^{71} Ga capture rate in the standard model is primarily (70%) determined by the flux of neutrinos from the driving reaction of the pp-chain, $p + p \to {}^{2}H + e^{+} + \nu$. The pp neutrino flux is effectively fixed, provided only that hydrogen burning is the sun's energy source, by the observed solar luminosity. Thus, if a ^{71}Ga experiment yields a capture rate much

Table 5: Reactions (1) - (4) produce solar neutrinos with continuous distributions, while (5) and (6) are line sources. E_{ν}^{max} is the maximum energy of the neutrinos for all reactions except (4), where it has been computed with respect to the center of the broad 2.9 MeV ^{8}Be resonance populated in the β-decay of ^{8}B. Fluxes are taken from the standard solar model calculation of Bahcall et al. [39].

Reaction	E_{ν}^{max} (MeV)	Flux ($10^{10}/cm^2 s$)
(1) $p + p \to {}^{2}H + e^{+} + \nu$	0.420	6.1
(2) $^{13}N \to {}^{13}C + e^{+} + \nu$	1.199	4.6×10^{-2}
(3) $^{15}O \to {}^{15}N + e^{+} + \nu$	1.732	3.7×10^{-2}
(4) $^{8}B \to {}^{8}Be + e^{+} + \nu$	14.02	5.85×10^{-4}
(5) $^{7}Be + e^{-} \to {}^{7}Li + \nu$	0.862 (89.6%) 0.384 (10.4%)	4.1×10^{-1}
(6) $p + e^{-} + p \to {}^{2}H + \nu$	1.442	1.5×10^{-2}

reduced from standard model predictions, our particle physics must be at fault. In this sense the ^{71}Ga experiment will provide a test of neutrino oscillations for small Δm^2 and large mixing angles far beyond that possible with terrestrial neutrino sources.

The primary obstacle to the Brookhaven proposal for a ^{71}Ga experiment is the cost of the requisite quantity of gallium, estimated to exceed $25 million. However, there are in addition nagging uncertainties in the capture cross section. Gamow-Teller transitions to two excited states in ^{71}Ge, the $5/2^-$ (175 keV) and $3/2^-$ (500 keV) states, can be excited by ^7Be neutrinos. The ^7Be neutrino flux, like the ^8B neutrino flux, depends sensitively on the sun's central temperature. If the GT strengths for these transitions prove to be unusually strong, one could no longer argue that the ^{71}Ga capture cross was insensitive to solar model assumptions. Bahcall has argued from nuclear systematics that upper limits on the relevant transition strengths are log(ft) = 6.0 ($5/2^-$) and 5.0 ($3/2^-$) [43]. Yet, in view of the importance of this experiment, the need for a definitive <u>measurement</u> of these transition strengths is clear. Presumably $0°$ (p,n) mappings with 175 keV resolution would settle this matter. (Note: In the discussion following this talk Dr. Orihara announced results of (p,n) measurements performed at Tohoku University showing that each of these transitions is quite weak. This substantiates the theoretical work of Bahcall and further strengthens the argument for doing the ^{71}G experiment.)

Davis, Sam Hurst, and collaborators [44] have considered designing a solar neutrino experiment based on the reaction $^{81}Br(\nu,e)^{81}Kr$ originally discussed by R. D. Scott [45]. The techniques developed by Davis to isolate ^{37}Ar could also be used to collect ^{81}Kr, and noble gas resonance ionization spectrometry might permit ^{81}Kr counting at the necessary sensitivity. The β-decay measurement of Bennett et al. [46] and the calculations of Bahcall [47] and Haxton [48] indicate that the capture rate is predominantly determined by the ^7Be neutrino flux. Thus the ^{81}Br and ^{37}Cl experiments would together determine the $^7Be/^8B$ neutrino flux ratio, a quantity that also serves to distinguish flaws in solar physics from those in particle physics: the flux ratio is likely to be unaffected by solar oscillations, but is highly sensitive to modifications in the standard solar model affecting the sun's central temperature.

The log(ft) value for the first GT transition, to the $1/2^-$ (190 keV) state in ^{81}Br, has been measured by Bennett et al. [46]. The strength of a second transition that can be excited by ^7Be neutrinos, to the $5/2^-$ (457 keV) state, is unknown. Thus a $0°$ (p,n) measurement with 270 keV resolution would be important. Unidentified levels exist at 549 keV and 608 keV in ^{81}Kr; if these states can be excited by the GT operator, an improvement in resolution would be required. The high energy ^{8}B capture rate should also be appreciable, so a complete mapping of the GT strength distribution below particle breakup in ^{81}Kr will be needed.

An experiment of quite a different kind has been recently proposed by G. A. Cowan and Haxton: a measurement of the concentrations of ^{97}Tc and ^{98}Tc produced by neutrino absorption in a deeply buried molybdenite ore body over the past several million years [49]. This experiment would test the long-term stability of the sun and, in particular, the suggestion that the solar neutrino puzzle and the recent Pleistocene glacial epoch are both the result of sudden mixing in the solar core four million years ago.

Only the ^{8}B neutrinos can induce the reaction $^{98}Mo(\nu,e)^{98}Tc$; as the ^{8}B phase space varies slowly as a function of nuclear excitation energy, no strong restrictions on resolution in GT mappings exist in this case. In addition to ^{8}B neutrinos, the ^{7}Be neutrinos may contribute importantly to the capture rate for $^{97}Mo(\nu,e)^{97}Tc$. A resolution of 250 keV should permit an accurate estimate of this capture rate.

I would like to acknowledge my collaborators in studies of ββ-decay and solar neutrino detection, G. A. Cowan, S. P Rosen, G. J. Stephenson, Jr., and D. Strottman. In addition to these, I thank N. Auerbach, S. Austin, F. Avignone, J. N. Bahcall, the Kent State group, J. Rapaport, and L. Zamick for helpful conversations and communications. This work was supported in part by the National Science Foundation (grant no. PHY-8021272) and by the Department of Energy.

REFERENCES

[1] See, for example, S. Weinberg, Nobel Prize Address, Science 210, 1212 (1980).

[2] See, for example, T. Goldman, talk presented to the APS Division of Particles and Fields, Santa Cruz (1981).

[3] M. Goldhaber, talk presented at the Workshop on Low Energy Tests of High Energy Physics, Santa Barbara (1982).

[4] Y. K. Lee, L. W. Mo, and C. S. Wu, Phys. Rev. Lett. 10, 253 (1963).

[5] C. S. Wu, E. Ambler, R. Hayward, D. Hoppes, and R. Hudson, Phys. Rev. 105, 1413 (1957).

[6] For general discussion see H. Primakoff and S. P. Rosen, Rep. Prog. Phys. 22, 121 (1950); E. J. Konopinski, "Theory of Beta Radioactivity" (Oxford University Press, London); talks presented by S. P. Rosen, G. J. Stephenson, Jr., and C. S. Wu, Neutrino Mass Mini-Conference, Telemark, Wisconsin (1980); S. P. Rosen, talk presented at Neutrino '81, Maui, Hawaii (1981); W. C. Haxton and S. P. Rosen, talks presented at Workshop on Low Energy Tests of High Energy Physics, Santa Barbara (1982); E. Fiorini, Revista del Nuovo Cim. 2, 1 (1971); D. Bryman and C. Picciotto, Rev. Mod. Phys. 50, 11 (1978).

[7] Talks presented by G. Senjanović, D. Wu, J. Schechter, Y. Tomozana, and L. Wolfenstein, Neutrino Mass Mini-Conference, Telemark, Wisconsin (1980).

[8] V. P. Lubimov, E. G. Novikov, V. Z. Nozik, E. F. Tretykov, and V. S. Kosik, Yad. Fiz. 32, 301 (1980) and Phys. Lett. 94B, 266 (1980).

[9] T. Kirsten, W. Gentner, and O. A. Schaeffer, Z. Physik 202, 273 (1967); T. Kirsten, W. Gentner, and O. Müller, Z. Naturforsch 22a, 1703 (1967); T. Kirsten, A. O. Schaeffer, E. Norton, and R. W. Stoenner, Phys. Rev. Lett. 20, 1300 (1968).

[10] E. C. Alexander, Jr., B. Srinivasan, and O. K. Manuel, Earth and Planet. Sci. Lett. 5, 478 (1969); B. Srinivasan, E. C. Alexander, Jr., and O. K. Manuel, Econ. Geol. 67, 592 (1972) and J. Inorg. Nucl. Chem. 34, 2381 (1972).

[11] E. W. Hennecke, O. K. Manuel, and D. D. Sabu, Phys. Rev. C11, 1378 (1972).

[12] B. Srinivasan, E. C. Alexander, Jr., R. D. Beatty, D. E. Sinclair, and O. K. Manuel, Econ. Geol. 68, 252 (1973).

[13] R. Bardin, P. Gollon, J. Ullman, and C. S. Wu, Phys. Lett. 26B, 112 (1967) and Nucl. Phys. A158, 337 (1970).

[14] E. Fiorini, A. Pullia, G. Bertolini, F. Capellani, and G. Restelli, Nuovo Cim. A13, 747 (1973).

[15] B. T. Cleveland, W. R. Leo, C. S. Wu, L. R. Kasday, A. M. Rushton, P. J. Gollon, and J. O. Ullman, Phys. Rev. Lett. 35, 737 (1975).

[16] M. K. Moe and D. Lowenthal, Phys. Rev. C22, 2186 (1980).

[17] W. C. Haxton, G. J. Stephenson, Jr., and D. Strottman, Phys. Rev. D25, 2360 (1982).

[18] M. Doi, T. Kotani, H. Nishiura, K. Okuda, and E. Takasugi, Phys. Lett. 102B, 219 (1981); Prog. Theor. Phys. 66, 1739 (1981); Prog. Theor. Phys. 66, 1765 (1981).

[19] W. C. Haxton, G. J. Stephenson, Jr., and D. Strottman, Phys. Rev. Lett. 47, 153 (1981).

[20] W. C. Haxton, S. P. Rosen, and G. J. Stephenson, Jr., Los Alamos preprint and submitted to Phys. Rev. Lett.

[21] R. D. Lawson, Phys. Rev. C24, 1500 (1981).
[22] T. T. S. Kuo, private communication; T. T. S. Kuo and G. E. Brown, Nucl. Phys. A114, 241 (1968).
[23] W. T. Baldridge and J. P. Vary, Phys. Rev. C14, 2246 (1976).
[24] L. Zamick and N. Auerbach, Rutgers University preprint.
[25] T. Kirsten, talk presented at the Workshop on Low Energy Tests of High Energy Physics, Santa Barbara (1982).
[26] J. D. Vergados, Phys. Rev. C12, 865 (1976).
[27] S. P. Rosen, talk presented to Orbis Scientiae 1981, Center for Theoretical Studies, University of Miami.
[28] A. Halprin, P. Minkowski, H. Primakoff, and S. P. Rosen, Phys. Rev. D13, 2567 (1976).
[29] C. Picciotto, Can. J. of Phys. 56, 399 (1978).
[30] H. M. Georgi, S. L. Glashow, and S. Nussinov, Harvard Univ. preprint HUTP-81/A026 (1981); G. B. Gelmini and M. Roncadelli, Phys. Lett. 99B, 411 (1981).
[31] R. N. Mohapatra and J. D. Vergados, Phys. Rev. Lett. 47, 1713 (1981).
[32] J. Schechter and J. W. F. Valle, Syracuse University preprint SU-4217-213 (to be published in Phys. Rev. D); L. Wolfenstein, private communication.
[33] L. Wolfenstein, Phys. Lett. 107B, 77 (1981).
[34] A series of such (p,n) measurements are being made at the Indiana University Cyclotron Facility.
[35] R. Davis, Jr., D. S. Harmer, and K. C. Hoffman, Phys. Rev. Lett. 20, 1205 (1968); R. Davis, Jr., Proc. of the International Conf. on Neutrino Physics and Astrophysics, Moscow, Vol. 2, 99 (1969); R. Davis, Jr., and J. M. Evans, Proc. of the 13th International Cosmic Ray Conf., Vol. 3, 2001 (1973); R. Davis, Jr., Proc. Informal Conf. on the Status and Future of Solar Neutrino Research, ed. G. Friedlander, BNL Report 50879, 1 (1970).
[36] J. N. Bahcall, talk presented at Neutrino '81, Maui, Hawaii (1981).
[37] R. T. Rood, in Proc. Informal Conf. on the Status and Future of Solar Neutrino Research, ed., G. Friedlander, BNL Report 50879, 175 (1978); J. N. Bahcall and R. L. Sears, Ann. Rev. Astron. and Astrophys. 10, 25 (1972).
[38] F. Reines, H. W. Sobel, and E. Pasierb, Phys. Rev. Lett. 45, 1307 (1980).
[39] J. N. Bahcall, S. H. Lubow, W. F. Huebner, N. H. Magee, Jr., A. L. Merts, M. F. Argo, P. D. Parker, B. Rozsnyai, and R. K. Ulrich, Phys. Rev. Lett. 45, 945 (1980).
[40] V. A. Kuzmin, Sov. Phys. JETP 22, 1051 (1966).
[41] J. N. Bahcall, B. Cleveland, R. Davis, Jr., I. Dostrovsky, J. C. Evans, Jr., W. Frati, G. Friedlander, K. Lande, K. Rowley, and J. Weneser, Phys. Rev. Lett. 40, 1351 (1978).

[42] N. Itoh and Y. Kohyama, Astrophys. J. 246, 989 (1981).

[43] J. N. Bahcall, Rev. Mod. Phys. 50, 681 (1978).

[44] G. S. Hurst, M. G. Payne, S. Kramer, and C. H. Chen, Phys. Today 33, 24; R. Davis, Jr., private communication.

[45] R. D. Scott, Nature 264, 729 (1976).

[46] C. L. Bennett, M. M. Lowry, R. A. Naumann, F. Loeser, and W. Moore, Phys. Rev. C22, 2245 (1980).

[47] J. N. Bahcall, Phys. Rev. C24, 2216 (1981).

[48] W. C. Haxton, Nucl. Phys. A367, 517 (1981).

[49] G. A. Cowan and W. C. Haxton, Science 216, 51 (1982).