New Contribution to Neutrinoless Double Beta Decay in Gauge Models

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It is pointed out that in a general class of gauge models, there exist new contributions to neutrinoless double β transitions, that do not involve a Majorana neutrino but the decay of a doubly charged Higgs boson to electrons. Explicit calculations for the case $^{48}Ca \to {}^{48}Ti$ indicate that for reasonable choice of parameters, this new contribution may dominate over that involving light or heavy Majorana neutrinos.

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All observed weak-interaction processes to date appear to respect a $U(1)_L$ global symmetry associated with lepton number.[1] Lacking any dynamical reason for the existence of $U(1)_L$ it is important to ask whether lepton number is indeed a good or an approximate quantum number? In the context of modern gauge theories, if neutrinos are taken massive, satisfactory understanding of the smallness of their mass[2] requires that the leptonic $U(1)_L$ symmetry (local[2] or global[3]) be spontaneously broken. Phenomenological tests of this idea, of course, depend on the scale of the $U(1)_L$-symmetry breakdown. If this scale is low[4] ($\approx$TeV), an outstanding signature for this breaking is the observation of neutrinoless double beta decay[5] $[(\beta\beta)_{0\nu}]$

$$(N, Z) \to (N-2, Z+2) + e^- + e^-$$

with a lifetime $\tau_{(\beta\beta)_{0\nu}} \approx 10^{22}$–$10^{24}$ yr, not much higher than the existing experimental lower limits[5] on $\tau_{(\beta\beta)_{0\nu}} \gtrsim 10^{21}$ yr.

The traditional theoretical analysis of the $(\beta\beta)_{0\nu}$ process had to suppose the existence of a Majorana neutrino ν_m, which could be ν_e itself or could couple to the electron, and one had to make one (or both) of the following assumptions: (i) There exists a small right-handed component in the predominantly left-handed current. (ii) The Majorana neutrino ν_m has a mass not much less than a few eV.

These mechanisms have been analyzed in the past[5-9] and limits on the neutrino mass have been extracted.[7-9] In this note, we point out that a new contribution to $(\beta\beta)_{0\nu}$ involving no neutrinos, but the decay of doubly charged Higgs boson to electrons, can exist in gauge models where neutrinos are Majorana particles. For reasonable values of parameters, we estimate this contribution for the transition $^{48}Ca \to {}^{48}Ti$ and find that it may dominate over the neutrino-exchange graphs.

We first point out the origin of the new Higgs exchange contribution to $(\beta\beta)_{0\nu}$, in a gauge model based on the group $SU(2) \otimes U(1)$ with only left-handed Majorana neutrinos. It is only relevant for our purpose to display the leptonic and Higgs content of our model,[10]

$$\psi_L \equiv \begin{pmatrix} \nu_L \\ e_L^- \end{pmatrix} \quad (\tfrac{1}{2}, -1), \qquad e_R^- \quad (0, -2); \tag{1}$$

$$\phi \equiv \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix} \quad (\tfrac{1}{2}, 1), \qquad \vec{\Delta} \equiv \begin{pmatrix} \Delta^{++} \\ \Delta^+ \\ \Delta^0 \end{pmatrix} \quad (1, 2). \tag{2}$$

Gauge invariance allows the following Yukawa-type couplings among the above fields:

$$\mathcal{L}_Y = h_1 \bar{\psi}_L \phi e_R^- + i h_2 \psi_L^T \tau_2 \vec{\tau} \psi_L \vec{\Delta} + \text{H.c.} \tag{3}$$

As in the standard Glashow-Weinberg-Salam model, we assume that the ϕ also couples to the quarks. Spontaneous breakdown of the $SU(2)_L \otimes U(1)$ gauge symmetry is caused by a $\langle \phi \rangle = \binom{0}{v}$.

The existence of a trilinear Higgs coupling of the form $\mathcal{L}_3 \equiv \lambda M \phi^T \tau_2 \vec{\tau} \phi \cdot \vec{\Delta}$ implies that $\langle \Delta^0 \rangle = \kappa \neq 0$ leading to a Majorana mass term for the left-handed neutrino, with $m_\nu = h_2 \kappa$. If we assume that $h_2 \approx 1$, this implies $\kappa \lesssim 3.5 \times 10^{-8}$ GeV. Since $v \simeq 10^3$ GeV and we choose $\lambda \approx 1$ and $M \approx v$, we need a fine tuning of the mass term of $\vec{\Delta}$ to one part in 10^{-11}. This is obviously unsatisfactory but no worse than in other models with vastly disparate mass scales such as grand unified ones.

The presence of the $\mathcal{L}_3$ term breaks lepton number by two units and contributes to $(\beta\beta)_{0\nu}$ decay without exchange of Majorana neutrinos as shown in Fig. 1. It leads to an effective double β-decay Hamiltonian of the form

$$\mathcal{H}_{\Delta L=2} \simeq G_2 \bar{d}_L u_R \bar{d}_L u_R e_L{}^T c^{-1} e_L + \text{H.c.} \tag{4}$$

with

$$G_2 \simeq \frac{h_q{}^2 h_2 \lambda M}{m_\varphi{}^4 - m_\Delta{}^2}. \tag{5}$$

The strength of this interaction is estimated to be $\simeq G_F{}^2 \times 10^{-8}\ \text{GeV}^{-1}$, if we assume $h_q{}^2 \simeq 10^{-9}$ (since it is related to the quark masses) and $m_\varphi \approx m_\Delta \simeq 100$ GeV. This value is of the same order as the present limits on the interaction.[11]

We consider next the $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ model.[4] We remind the reader about the leptonic multiplets of the model: left-handed doublet

$$\psi_L \equiv (\nu_L, e_L{}^-):\ (\tfrac{1}{2}, 0, -1);$$

right-handed doublet

$$\psi_R \equiv (N_R, e_R{}^-):\ (0, \tfrac{1}{2}, -1);$$

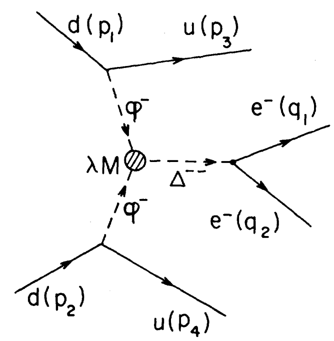

FIG. 1. New Higgs exchange contribution to neutrinoless double β decay.

Higgs multiplets

$$\phi:\ (\tfrac{1}{2}, \tfrac{1}{2}, 0);\quad \Delta_L:\ (1, 0, +2);\quad \Delta_R:\ (0, 1, +2).$$

It is then straightforward to observe that the existence of the scalar coupling $\text{Tr}(\vec{\tau}\cdot \Delta_L \phi \vec{\Delta}_R{}^\dagger \cdot \vec{\tau}\phi^\dagger)$ upon substituting $\langle \Delta_R{}^0 \rangle = v_R \neq 0$ leads to the Feynman diagram in Fig. 1, with M replaced by v_R. Since v_R is bigger than M_w, the value G_2 is expected to be bigger in this model.

The amplitude associated with the Feynman diagram of Fig. 1 can be written down. Assuming that all Higgs particles involved are quite heavy, we obtain

$$\mathfrak{M} = G_2 \bar{u}(p_3)(1-\gamma_5) d(p_1) \bar{u}(p_4)(1-\gamma_5) d(p_2)(1-P_{12}) \bar{e}(q_2)(1+\gamma_5) e^c(q_1), \tag{6}$$

where P_{12} is the permutation operator (interchanging spins and momenta of the two electrons). At the nucleon level the above amplitude can be written as

$$\mathfrak{M} = G_2 \bar{u}(p_3)[F_S{}^{(3)} - F_P{}^{(3)}\gamma_5]\tau_+ u(p_1) \bar{u}(p_4)[F_S{}^{(3)} - F_P{}^{(3)}\gamma_5]\tau_+ u(p_2)(1-P_{12}) \bar{e}(q_2)(1+\gamma_5) e^c(q_1), \tag{7}$$

where $F_S{}^{(3)}$ and $F_P{}^{(3)}$ are the isovector scalar and pseudoscalar form factors. We use the parametrization for them suggested by Adler *et al.*[12]

$$F_P{}^{(3)}(k^2) = \frac{F_P{}^{(3)}(0)}{(1-k^2/M_A{}^2)(1-k^2/M_\pi{}^2)}, \tag{8}$$

$$F_S{}^{(3)}(k^2) = \frac{F_S{}^{(3)}(0)}{(1-k^2/M_A{}^2)^2}, \tag{9}$$

where $M_A = 0.85\ \text{GeV}/c^2$.

For the reasons discussed in Ref. 9, the momentum dependence of the form factors cannot be ignored; namely, if we neglect such dependence, the transition operator in coordinate space would behave like a δ function in the separation of two interaction nucleons. Because of the Pauli principle, it has the following consequences: First, the transition rate is greatly suppressed; and second, the results obtained depend crucially on not so well understood physics like short-range correlations. It has been argued in Ref. 9 that introduction of the form factor cures both these problems.

It is now tedious but straightforward to make a nonrelativistic reduction of Eq. (8) and obtain the following result in coordinate space[9]:

$$\mathfrak{M} = \tfrac{1}{2} G_2 \frac{\hbar c}{m_p c^2 R_0} m_e m_p{}^2 f_A{}^2 [(1-P_{12})\bar{e}(q_2)(1+\gamma_5)e^c(q_1)] \langle f | \Omega_\Delta | i \rangle, \tag{10}$$

$$\Omega_\Delta = \frac{m_p}{m_e} \sum_{i\neq j} \tau_+(i)\tau_+(j) \frac{R_0}{r_{ij}} \{\alpha_S F_1(r_{ij}) + \alpha_P [\vec{\sigma}_i \cdot \vec{\sigma}_j F_2(r_{ij}) + (\tfrac{24}{5}\pi)^{1/2} Y^2(\hat{r})(\vec{\sigma}_i \times \vec{\sigma}_j)^2 F_3(r_{ij})]\}, \tag{11}$$

with $R_0 = r_0 A^{1/3}$ ($r_0 = 1.1$ fm), the nuclear radius; $\vec{r}_{ij} = \vec{r}_i - \vec{r}_j$, $\hat{r} = \vec{r}_{ij}/|\vec{r}_{ij}|$,

$$\alpha_S = (1/3f_A^2)[F_S^{(3)}(0)]^2, \quad \alpha_P = \tfrac{1}{12}(M_\pi/m_p)^4[F_P^{(3)}(0)]^2, \tag{12a}$$

$$F_1(r_{ij}) = \tfrac{1}{48}(M_A/m_p)^2 x_A(x_A^2 + 3x_A + 3)e^{-x_A}, \tag{12b}$$

$$F_2(r_{ij}) = 2[(M_\pi/M_A)^2 e^{-x_\pi} - e^{-x_A}] + \tfrac{1}{2}[(x_\pi - 2)e^{-x_\pi} + (x_A - 2)e^{-x_A}], \tag{12c}$$

$$F_3(r_{ij}) = 2[(M_\pi/M_A)^2 Z(x_A)] + \tfrac{1}{2}[(x_\pi + 1)e^{-x_\pi} + (x_A + 1)e^{-x_A}], \tag{12d}$$

$$x_A = (M_A c^2/\hbar c)|\vec{r}_{ij}|; \quad x_\pi = (M_\pi c^2/\hbar c)|\vec{r}_{ij}|; \quad Z(x) = (x^2 + 3x + 3)x^{-2}e^{-x}. \tag{12e}$$

We notice that the structure of the operator Ω_Δ is not very different from the one involved in the heavy-Majorana-neutrino–mediated process[9] except that it contains a tensor component.

If one now defines a lepton-number–nonservation parameter $\eta_\Delta = (G_2 m_p/G_F^2)$, the transition probability associated with Fig. 1 takes a familiar form[5,7,9]:

$$W_{0\nu}(i \to f e^- e^-) = K_2 |\chi(\alpha(z+2))|^2 A^{-2/3} f_2(\epsilon_0) |\eta_\Delta \langle f|\Omega_\Delta|i\rangle|^2, \tag{13}$$

where

$$K_2 = c\,\frac{m_e c^2}{\hbar c}\,\frac{(G_F m_p^2)^4}{(2\pi)^3}\left(\frac{\hbar c}{m_p c^2 r_0}\right)^2\left(\frac{m_e}{m_p}\right)^6 f_A^4 = 2.6\times10^{-15}\ \text{yr}^{-1}, \tag{14}$$

$$f_2(\epsilon_0) = \tfrac{2}{15}\epsilon_0[\epsilon_0^4 + 10\epsilon_0^3 + 40\epsilon_0^2 + 60\epsilon_0 + 30] \tag{15a}$$

(ϵ_0 is the available energy of double β decay in units of $m_e c^2$),

$$\chi(\alpha z) = \frac{2\pi\alpha Z}{1 - \exp(-2\pi\alpha Z)} \tag{15b}$$

represents the distortion of the electron wave function. In order to estimate the rate given by Eq. (13), we taken the values of α_s and $\alpha_p = 0.40$; $\eta_\Delta \approx 10^{-8}$. We then compute[9] the nuclear matrix element for the simplest and experimentally most interesting process: ^{48}Ca $\to$ ^{48}Ti, in the shell model where we find $\langle f|\Omega_\Delta|i\rangle = 530$. This value is approximately reduced by 20% when a simple step-function short-range correlation is used. From Eqs. (14), (16), and (17), we find

$$(T_{1/2})_{0\nu} \simeq 3\times10^2\ \text{yr} \tag{16}$$

which almost coincides with the present experimental lower limit. Of course, it must be pointed out that although we believe $M_\phi \approx M_\Delta \approx 100$ GeV is quite a reasonable choice for Higgs masses, $T_{1/2}$ depends on these masses rather sensitively (i.e., $T_{1/2} \sim M_\phi^8 M_\Delta^4$), and a slight increase in the values of these masses could increase the half-lives significantly.

We remark that in the presence of the proposed mechanism, the full double β-decay experiments impose a constraint of the form

$$|\eta_\nu\langle f|\Omega_\nu|i\rangle + \eta_N\langle f|\Omega_N|i\rangle + \eta_\Delta\langle f|\Omega_\Delta|i\rangle|^2 \leq b, \tag{17}$$

where η_ν and η_N correspond to the lepton-number–nonconservation parameters associated with the conventional light and heavy Majorana neutrinos, which in the notation of Ref. 9 can be written as $\eta_\nu = m_\nu/m_e$ and $\eta_N = \beta^2 m_p/m_N$ where β is the effective coupling of N to electrons. It has been shown in Ref. 9 that $\langle f|\Omega_\nu|i\rangle \approx 1$ and $\langle f|\Omega_N|i\rangle \approx 72$. From the data on ^{48}Ca $\to$ ^{48}Ti decay, we get $b = 0.4\times10^{-10}$, which from Eq. (16) leads to useful constraints on gauge models. For instance, if $\eta_N = \eta_\Delta = 0$, this implies $\eta_\nu \leq 0.6\times10^{-5}$ or $m_\nu \leq 3$ eV, etc.

In conclusion, we have pointed out a new contribution to double β decay, which does not involve Majorana neutrinos and makes a measurable contribution to it, provided that η_Δ is not smaller than 10^{-8}. Our results are based on the ^{48}Ca $\to$ ^{48}Ti decay but we see no reason why it should not hold in general.

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Proximity to the Rho-Meson Condensation Threshold in Nuclei

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Critical conditions for instability in the isovector unnatural-parity channel in nuclei are investigated in terms of the random-phase approximation. The particle-hole residual interaction consists of the π and ρ (2π) meson-exchange attractive components in addition to the short-range repulsive Landau term. It is found that the nonresonant part of the iterated two-pion exchange moves the characteristics of the nuclear phase transition closer to those for rho-meson condensation. This mode may dominate at high density.

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Although the possible existence of a pion-condensed mode of hadronic matter has been widely discussed,[1,2] careful studies of nucleon-nucleon correlations have led to the conclusion that the critical density is greater than 1.5 times the normal density of nuclear matter.[3-6] Recent discussions have thus focused on the possibility of observing precritical phenomena in ordinary nuclei.[7-9] The proximity of the pion condensation threshold should amplify the pionic field in finite nuclei and thus increase the cross sections of (e,e'),[8] (p,p'),[9,10] and (π,γ) or (γ,π) reactions[11,12] for isovector unnatural-parity transitions at large momentum transfer.

Quantitative discussions of precritical behavior have centered on the 1^+, $T=1$ state of ^{12}C for which the $M1$ form factor $|F_M|^2$ predicted by the Cohen-Kurath wave functions[13] is substantially smaller than the data from (e,e') reactions.[8,14] Analyses in the context of precritical behavior with π- and ρ-meson exchange interactions have led to values of the Landau-Migdal parameter $g'\lesssim 0.4$.[8] However, by inclusion also of the nonresonant part of the two-pion exchange (TPE) in the isovector channel, good agreement with the data was obtained for $g'\sim 0.55$.[14] Analyses of recent (p,p') data for the same transition have shown little evidence of precritical enhancements and are consistent with a value $g'\gtrsim 0.60$.[15,16]

In examining the previous analyses, we realized that the amount of enhancement that could be produced, and the subsequent estimate of g', were very sensitive to the strength of the $\vec{\sigma}\times\vec{q}$ correlations provided by the ρ-meson exchange (RME) and the iterated one-pion exchange (OPE). The TPE strength is commonly related to the $N\bar{N}\to 2\pi$

