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NEUTRINOLESS DOUBLE BETA DECAY AND THE LIMITS TO THE RIGHT-HANDED CURRENTS ☆

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We consider the neutrinoless double beta decay within two extended electroweak gauge models with right-handed weak interactions, the left–right symmetric model and the mirror fermion model. We show that if the two electron-like neutrinos with opposite chiralities appearing in these models are lighter than a few MeV, the double beta decay data do not restrict the magnitude of right-handed currents.

If the electron neutrino ν_e is a Majorana particle, it can propagate as a virtual state between two neutrons in a nucleus and cause the reaction $(A, Z) \to (A, Z+2) + 2e^-$, called the neutrinoless nuclear double beta decay $[(\beta\beta)_{0\nu}]$ [‡1]. To proceed in the standard framework of left-handed (LH) charged weak interactions, this reaction requires a non-zero neutrino mass m_{ν_e}. The non-observation of $(\beta\beta)_{0\nu}$ has been used to put an upper bound $m_{\nu_e} \leqslant 5.6$ eV on this mass [2].

On the other hand, if the chirality of the charged weak interactions is not standard, the double beta decay can also proceed through the right-handed (RH) component of the electron-neutrino current [3]. In that case the experimental upper bound for $(\beta\beta)_{0\nu}$ gives a limit on the deviation from the pure V–A structure. If the weak current is parametrized as

$$J_\lambda = \bar{e}\gamma_\lambda[(1-\gamma_5) + \eta(1+\gamma_5)]\nu_e\,, \qquad (1)$$

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‡1 For a review and general references, see ref. [1].

the data implies, assuming $m_{\nu_e} = 0$,

$$\eta \leqslant 2.4 \times 10^{-5}\,. \qquad (2)$$

It is remarkable that this bound is almost four orders of magnitude more restrictive than the constraints obtained from the analysis of all the other weak processes [4]. If $m_{\nu_e} \neq 0$, the limit (2) can become even more severe. Therefore, it would seem that in the case of Majorana neutrino, the RH currents could not play any role in charged weak processes.

Because of the severity of the bound (2), it is important to carefully examine the general validity of the arguments leading to it. It should be emphasized that the bound (2) is derived by presuming that only one Majorana neutrino contributes to the $(\beta\beta)_{0\nu}$ process. In the present paper, we will show that if the theory contains two moderately light $[m_\nu \leqslant O(10)$ MeV] neutrinos coupling to the electron *with opposite chiralities*, $(\beta\beta)_{0\nu}$ will not give a bound on the strength of RH currents. In comparison, one should keep in mind the well-known modifications [5] to the limit on the neutrino masses when two or more (instead of one) light *left-handed* neutrinos are supposed to couple to the electron.

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Two light, opposite chirality neutrinos appear naturally in many models. We will consider two such extensively discussed models, both having RH currents. The first one is the $SU(2)_L \times SU(2)_R \times U(1)$ left–right symmetric (LR) model, in which the two Majorana neutrinos are the two orthogonal combinations of ν_L and ν_R $(\sim\nu_L^c)$ [‡2]. The second model is based on the standard $SU(2)_L \times U(1)$ symmetry, where RH currents could arise e.g. from the mixing of fermions with mirror fermions [7]. We will call this the mirror mixing (MM) model. We remind the reader that all $N \geqslant 2$ supersymmetries [8] inevitably contain mirror fermions. This fact follows from the reality of the $N \geqslant 2$ SUSY fermion representation. Since the mirror fermion mass terms break $SU(2)_L$, their mass scale cannot be heavier than O(100) GeV. In mirror models the left-handed neutrino ν_L has a right-handed mirror partner N_R. In both of these models, currents of the type (1) are produced if the LH state (ν_L) and the RH state (ν_R or N_R) will mix with each other. In the MM model one also needs some mixing between the electron and its mirror partner, since the pure electron does not have RH couplings in $SU(2)_L \times U(1)$.

Let us mention that there are no phenomenological objections against having two light electron-like neutrinos (we denote the mass eigenstates by ν_1 and ν_2). This is reflected by the well-known fact that the neutrino may well be a light Dirac particle, i.e. ν_1 and ν_2 are degenerate in mass. Moreover, Gronau and Nussinov [9] have pointed out that even within the "minimal" LR model, where the $SU(2)_R$ breaking and the mass of ν_R are intimately related, the RH neutrino may be as light as a few tens of MeV. In the LR models with a different Higgs sector, no lower bounds for neutrino masses exist [‡3].

Let us consider charged weak interactions described by the currents

$$J_L^\mu = \bar{e}_L\gamma^\mu\nu_L, \quad J_R^\mu = \bar{e}_R\gamma^\mu N_R, \tag{3}$$

where ν_L and N_R are two-component neutrino states. From here on, if not otherwise stated, N_R will denote the RH state for both (LR and MM) models. In the LR model J_R^μ couples to W_R, in the MM model to W_L. We will treat the two models within the same formalism.

The neutrino lagrangian is given by

$$\mathcal{L} = i\bar{\mathcal{N}}_L \not\partial \mathcal{N}_L - \tfrac{1}{2}(\mathcal{N}_L^T C M \mathcal{N}_L + \text{h.c.}), \tag{4}$$

where

$$\mathcal{N}_L = \begin{pmatrix} \nu \\ N^c \end{pmatrix}_L, \tag{5}$$

and $N^c = C\gamma_0^T N^*$. The mass matrix M is

$$M = \begin{pmatrix} m_1 & m_{12} \\ m_{12} & m_2 \end{pmatrix}, \tag{6}$$

which we take to be real. (For simplicity, and clarity of the argument, we will neglect here the possible intergenerational neutrino mixings as well as mixing between the W_L and W_R bosons in the LR model.) The neutrino mass terms are then

$$-2\mathcal{L}_M = m_1\nu_L^T C\nu_L + m_2 N_R^T C N_R + m_{12}(\bar{\nu}_R^c N_L^c + \bar{N}_R\nu_L) + \text{h.c.}. \tag{7}$$

Here m_1 and m_2 are lepton number violating Majorana masses, whereas m_{12}-terms are of Dirac type. If $m_{12} = 0$, the two states decouple. In accordance with our assumption of two light neutrinos, we take $|M_{ij}| \ll \langle p\rangle$, where $\langle p\rangle \simeq O(10)$ MeV is the average momentum of the virtual neutrino.

Let us record for completeness also the expressions for the parameters m_1, m_2 and m_{12} in terms of the physical (positive) masses m_{ν_1} and m_{ν_2} and the mixing angle ϕ:

$$m_1 = \tfrac{1}{2}[m_{\nu_1}(1+\cos 2\phi) + \eta m_{\nu_2}(1-\cos 2\phi)],$$
$$m_2 = \tfrac{1}{2}[m_{\nu_1}(1-\cos 2\phi) + \eta m_{\nu_2}(1+\cos 2\phi)],$$
$$m_{12} = \tfrac{1}{2}(m_{\nu_1} - m_{\nu_2}\eta)\sin 2\phi. \tag{8}$$

Here η is the product of the *CP* quantum numbers of the two mass eigenstates ν_1 and ν_2, and its value $\eta = \pm 1$ depends on the relative magnitudes of m_1, m_2 and m_{12}.

The leptonic part of the effective lagrangian for $(\beta\beta)_{0\nu}$ will contain two types of tensors, which we will call "current tensors" and "mass tensors". The

‡2 The double beta decay in the LR models was discussed in ref. [6]. There, the right-handed neutrino was assumed to be heavy (≥1 GeV).

‡3 It has recently been argued that a second (mirror) neutrino with a mass of about 100 eV, coupled subdominantly to the electron, would improve the fit to the tritium decay spectrum [10].

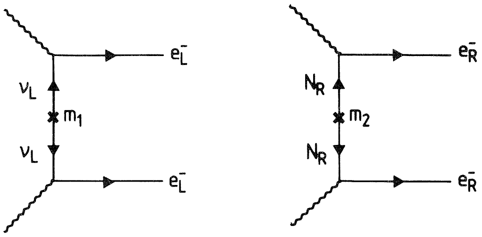

Fig. 1. Lowest order diagrams giving rise to a mass tensor contribution to the $(\beta\beta)_{0\nu}$ amplitude.

latter are given by

$$J_L^\lambda J_L^{T\rho} = \bar{e}_L \gamma^\lambda \nu_L \nu_L^T C \gamma^\rho e_R^c ,$$

$$J_R^\lambda J_R^{T\rho} = \bar{e}_R \gamma^\lambda N_R N_R^T C \gamma^\rho e_L^c . \quad (9)$$

These correspond to the situation when the two electron–neutrino vertices have the same (LH and RH, respectively) chiral structure and, accordingly, the virtual neutrino has to flip its helicity. The lowest order contributions to (9) are depicted in fig. 1. Higher order contributions are obtained by adding an even number of mass insertions to the propagators, but these will be suppressed by powers of $M_{ij}^2/\langle p\rangle^2 \ll 1$.

Current contributions, on the other hand, are given by the tensors $J_L^\lambda J_R^{T\rho}$ and $J_R^\lambda J_L^{T\rho}$. Here the two interaction vertices have opposite chiral structure. However, to first order in $M_{ij}/\langle p\rangle$ these contributions vanish, simply because the kinetic terms in (4) are diagonal in the (ν_L, N_R) basis [‡4]. The leading (second order) contributions are depicted in fig. 2. Two mass insertions are needed, one of the Majorana type to break the lepton number and the other of Dirac type to restore the helicity. Again, higher orders are suppressed by powers of $M_{ij}^2/\langle p\rangle^2$.

[‡4] Of course, in the mass eigenstate basis the first order contributions will cancel as well due to the unitarity of the transformation matrix between the two bases.

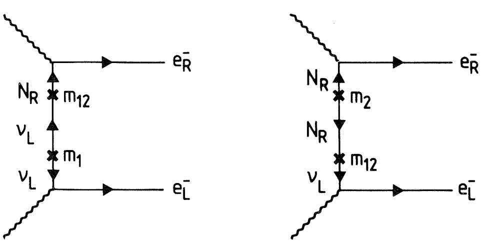

Fig. 2. Lowest order diagrams giving rise to a current tensor type contribution to the $(\beta\beta)_{0\nu}$ amplitude. Note that one of the mass insertions in both diagrams is of the Dirac type.

If we neglect the small RH modifications of the hadronic part of the $(\beta\beta)_{0\nu}$ amplitude (which would only cause small changes and hence do not have much relevance to our result), the lowest order amplitude (as given by the diagrams of figs. 1 and 2) squared is proportional to the factor

$$|A|^2 = \langle p\rangle^{-2} [m_1^2 + 2r^2 I m_1 m_2 + r^4 m_2^2 + 2\epsilon r I^{1/2}(m_1 + m_2)(m_1 + r^2 m_2) + \epsilon^2 r^2 (m_1 + m_2)^2] , \quad (10)$$

where $\epsilon = m_{12}/\langle p\rangle$. The first three terms are due to mass contributions, the fourth term is the interference between the mass and the current contributions and the last term is the pure current contribution. The factor r is in the LR model given by the mass ratio of the two gauge bosons, $r = (M_{W_L}/M_{W_R})^2$, and in the MM model by the electron–mirror electron mixing angle θ_{eE}, $r = \sin^2\theta_{eE}$. Interference effects between two contributions including electrons with different chiralities, are taken into account by the factor I, which is smaller than unity for all possible transitions. To a good approximation it is given by the factor $I = \frac{5}{4}x^2$, where $x = m_e/E_e$.

It can readily be seen from eq. (10) that if $|\epsilon| < r$, the mass contribution will dominate over the current contribution *irrespective* of the values of m_1 and m_2 (under the assumption $M_{ij}^2 < \langle p\rangle^2$). Since the present reliable upper bound on r is [4,11] $r \leqslant 0.05$. $|m_{12}|$ may be as large as about 0.5 MeV, to fulfill the above requirement. The minimum of the mass contribution is obtained for $m_1 = -m_2 I r^2$, and is

$$|A_m|^{\min} = m_1^2(1 - I^2)/I^2 . \quad (11)$$

Using the value $x \approx 0.29$ for ^{130}Te-decay and the above-mentioned experimental upper bound [2] for m_{ν_e} one finds the following limit for $|m_1|$:

$$|m_1| \leqslant 53 \text{ eV} . \quad (12)$$

Comparing this with the generous upper bound for $|m_{12}|$ one can conclude that no upper limit to the mixing angle ϕ, and thus on the RH currents, is obtained.

Note that the current contribution can vanish completely when $m_1 \simeq -m_2$. (Then, of course, $|m_{12}|$ is bounded only by $\langle p \rangle$.) The physical masses will then be equal with opposite *CP* quantum numbers and with absolute values up to $\langle p \rangle$. The mixing angle can vary between zero and $\pi/2$ depending on the ratio m_1/m_{12}.

The mass contributions will vanish only in the trivial limit $m_1 \to 0$ and $m_2 \to 0$. In this limit the $(\beta\beta)_{0\nu}$ decay cannot occur since the lepton number is not broken. It should also be noted that due to the incomplete interference between the two mass contributions with different chiralities, there cannot be a large reduction in the effective neutrino mass ‡5. This reduction could possibly be arranged with proper interfamily mixing. However, with RH currents, the situation is even more contrived since both the LH and the RH sector should be tuned.

In conclusion, we have considered the neutrinoless double beta decay in two models with right-handed charged weak interactions. We have explicitly shown that if the two electron-like neutrino states are lighter than about 10 MeV, the $(\beta\beta)_{0\nu}$ decay cannot give any constraint on the magnitude of right-handed currents in these models. This underlines the need for a more accurate measurement of the chiral structure of the weak interactions in conventional weak processes as well as of possible additional subdominantly coupled neutrino states.

‡5 This was also noted in ref. [12], where the mass eigenstate basis was used, and an arbitrary mass for ν_2 was allowed.

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