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New Leptoquark Mechanism of Neutrinoless Double Beta Decay

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Abstract

A new mechanism for neutrinoless double beta ($0\nu\beta\beta$) decay based on leptoquark exchange is discussed. Due to the specific helicity structure of the effective four-fermion interaction this contribution is strongly enhanced compared to the well-known mass mechanism of $0\nu\beta\beta$ decay. As a result the corresponding leptoquark parameters are severely constrained from non-observation of $0\nu\beta\beta$-decay. These constraints are more stringent than those derived from other experiments.

PACs: 11.30, 12.30, 13.15, 14.80, 23.40

Neutrinoless double beta decay ($0\nu\beta\beta$) is forbidden in the standard model (SM) of electro-weak interactions since it violates lepton number conservation. Therefore, experimental observation of this exotic process would be an unambiguous signal of physics beyond the SM (see refs. [1]-[4] for reviews).

Essential progress in the exploration of $0\nu\beta\beta$-decay both from theoretical and experimental sides has been achieved in the last few years (see, for instance [4] and references therein). The considerably improved experimental lower bounds on the half lives of various isotopes enhance the potential of $0\nu\beta\beta$ experiments in testing different concepts of physics beyond the SM such as supersymmetry (SUSY) and leptoquarks (LQ).

The SUSY mechanisms of $0\nu\beta\beta$-decay were comprehensively investigated in a series of papers [5]-[9]. It turned out that constraints on certain SUSY-parameters from non-observation of $0\nu\beta\beta$-decay [7] are stronger than those from current and near future accelerator and non-accelerator experiments. Therefore, it is useful to investigate other possible contributions of physics beyond the SM to $0\nu\beta\beta$-decay to obtain $0\nu\beta\beta$ constraints on the corresponding parameters.

In this note we present a new mechanism of $0\nu\beta\beta$-decay associated with the leptoquark contribution to the effective low-energy charged current lepton-quark interactions. The diagrams describing this contribution are presented in fig. 1.

The SM symmetries allow 5 scalar (S) and 5 vector (V^μ) LQs with the following $LQ(SU(3)_c \otimes SU(2)_L \otimes U(1)_Y)$ assignments: $S_0(3_c, 1; -2/3)$, $\tilde{S}_0(3_c, 1; -8/3)$, $S_{1/2}(\bar{3}_c, 2; -7/3)$, $\tilde{S}_{1/2}(\bar{3}_c, 2; -1/3)$, $S_1(3_c, 3; -2/3)$, $V_0(\bar{3}_c, 1; -4/3)$, $\tilde{V}_0(\bar{3}_c, 1; -10/3)$, $V_{1/2}(3_c, 2; -5/3)$, $\tilde{V}_{1/2}(3_c, 2; 1/3)$, $V_1(\bar{3}_c, 3; -4/3)$, where $Y = 2(Q_{em} - T_3)$.

The most general form of the renormalizable LQ-quark-lepton interactions consistent with $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry can be written as [10]

$$\begin{aligned}
\mathcal{L}_{LQ-l-q} &= \lambda^{(R)}_{S_0} \cdot \overline{u^c_R} e_R \cdot S_0^{R\dagger} + \lambda^{(R)}_{\tilde{S}_0} \cdot \overline{d^c_R} e_R \cdot \tilde{S}_0^{\dagger} + \lambda^{(R)}_{S_{1/2}} \cdot \overline{u_R} \ell_L \cdot S_{1/2}^{R\dagger} + \\
&+ \lambda^{(R)}_{\tilde{S}_{1/2}} \cdot \overline{d_R} \ell_L \cdot \tilde{S}^{\dagger}_{1/2} + \lambda^{(L)}_{S_0} \cdot \overline{q^c_L} i\tau_2 \ell_L \cdot S_0^{L\dagger} + \lambda^{(L)}_{S_{1/2}} \cdot \overline{q_L} i\tau_2 e_R \cdot S_{1/2}^{L\dagger} + \\
&+ \lambda^{(L)}_{S_1} \cdot \overline{q^c_L} i\tau_2 \hat{S}_1^{\dagger} \ell_L + \lambda^{(R)}_{V_0} \cdot \overline{d_R} \gamma^\mu e_R \cdot V_{0\mu}^{R\dagger} + \lambda^{(R)}_{\tilde{V}_0} \cdot \overline{u_R} \gamma^\mu e_R \cdot \tilde{V}^{\dagger}_{0\mu} + \\
&+ \lambda^{(R)}_{V_{1/2}} \cdot \overline{d^c_R} \gamma^\mu \ell_L \cdot V^{R\dagger}_{1/2\mu} + \lambda^{(R)}_{\tilde{V}_{1/2}} \cdot \overline{u^c_R} \gamma^\mu \ell_L \cdot \tilde{V}^{\dagger}_{1/2\mu} + \lambda^{(L)}_{V_0} \cdot \overline{q_L} \gamma^\mu \ell_L \cdot V_{0\mu}^{L\dagger} + \\
&+ \lambda^{(L)}_{V_{1/2}} \cdot \overline{q^c_L} \gamma^\mu e_R \cdot V^{L\dagger}_{1/2\mu} + \lambda^{(L)}_{V_1} \cdot \overline{q_L} \gamma^\mu \hat{V}^{\dagger}_{1\mu} \ell_L + h.c.
\end{aligned} \tag{1}$$

Here q and ℓ are the quark and the lepton doublets. Following [10, 11] we distinguish $S(V)^{L,R}$ being LQs coupled to the left-handed and right-handed quarks respectively (see, however, the discussion on chiral couplings in [12]). For LQ triplets $\Phi_1 = S_1, V_1^\mu$ the notation $\hat{\Phi}_1 = \vec{\tau} \cdot \vec{\Phi}_1$ is used.

On the same footing the LQ fields couple to the SM Higgs doublet field H. A complete list of the renormalizable LQ-Higgs interactions is given in ref. [12]. These new interactions are especially important for $0\nu\beta\beta$-decay, since after electro-weak symmetry breaking they lead to mixing between different LQ multiplets. In turn this mixing generates the effective 4-fermion interactions involving right-handed leptonic currents. In combination with the ordinary SM left-handed charged current interactions the latter produce the contribution to $0\nu\beta\beta$-decay shown in the diagrams of fig. 1 with large enhancement factors. This type of contribution is absent in the case of decoupled LQ and Higgs sectors [12].

Under electro-weak symmetry breaking the neutral component of the SM Higgs field acquires a non-zero vacuum expectation value, $\langle H^0 \rangle$, which creates via LQ-Higgs interaction terms non-diagonal mass matrices for LQ fields with the same electric charge but from different $SU(2)_L$ multiplets. To obtain observable predictions from the LQ-lepton-quark interaction Lagrangian in eq. (1), the LQ fields ($I = S, V$) with non-diagonal mass matrices have to be rotated to the mass eigenstate basis I'. This can be done in the standard way: $I(Q) = \mathcal{N}^{(I)}(Q) \cdot I'(Q)$, where $\mathcal{N}^{(I)}(Q)$ are orthogonal matrices such that

$\mathcal{N}^{(I)T}(Q_I)\cdot\mathcal{M}_I^2(Q)\cdot\mathcal{N}^{(I)}(Q) = Diag\{M_{I_n}^2\}$, with the M_{I_n} being the mass of the relevant mass eigenstate field I'.

Now it is straightforward to derive the effective 4-fermion $\nu - u - d - e$ interaction terms generated by the LQ exchange in the upper parts of the diagrams in fig. 1. After Fierz rearrangement they take the form [12]

$$\begin{aligned}\mathcal{L}_{LQ}^{eff} &= (\bar{\nu}P_R e^c)\left[\frac{\epsilon_S}{M_S^2}(\bar{u}P_R d) + \frac{\epsilon_V}{M_V^2}(\bar{u}P_L d)\right] - \\ &- (\bar{\nu}\gamma^\mu P_L e^c) \\ &\times \left[\left(\frac{\alpha_S^{(R)}}{M_S^2} + \frac{\alpha_V^{(R)}}{M_V^2}\right)(\bar{u}\gamma_\mu P_R d) - \sqrt{2}\left(\frac{\alpha_S^{(L)}}{M_S^2} + \frac{\alpha_V^{(L)}}{M_V^2}\right)(\bar{u}\gamma_\mu P_L d)\right],\end{aligned} \tag{2}$$

where

$$\epsilon_I = 2^{-\eta_I}\left[\lambda_{I_1}^{(L)}\lambda_{\tilde{I}_{1/2}}^{(R)}\left(\theta_{43}^I(Q_I^{(1)}) + \eta_I\sqrt{2}\theta_{41}^I(Q_I^{(2)})\right) - \lambda_{I_0}^{(L)}\lambda_{\tilde{I}_{1/2}}^{(R)}\theta_{13}^I(Q_I^{(1)})\right], \tag{3}$$

$$\alpha_I^{(L)} = \frac{2}{3+\eta_I}\lambda_{I_{1/2}}^{(L)}\lambda_{I_1}^{(L)}\theta_{24}^I(Q_I^{(2)}), \quad \alpha_I^{(R)} = \frac{2}{3+\eta_I}\lambda_{I_0}^{(R)}\lambda_{\tilde{I}_{1/2}}^{(R)}\theta_{23}^I(Q_I^{(1)}). \tag{4}$$

$\eta_{S,V} = 1, -1$ for scalar and vector LQs. $\theta_{kn}^I(Q)$ is a mixing parameter defined by

$$\theta_{kn}^I(Q) = \sum_l \mathcal{N}_{kl}^{(I)}(Q)\mathcal{N}_{nl}^{(I)}(Q)\left(\frac{M_I}{M_{I_l}(Q)}\right)^2, \tag{5}$$

where $\mathcal{N}^{(I)}(Q)$ are mixing matrix elements for the scalar $I = S$ and vector $I = V$ LQ fields with electric charges $Q = -1/3, -2/3$. Common mass scales M_S of scalar and M_V of vector LQs are introduced for convenience.

Following the well known procedure [2] one can find the LQ contribution to the $0\nu\beta\beta$-decay matrix element for the diagrams in fig. 1. The LQ exchange sectors of these diagrams are described by the point-like 4-fermion interactions specified by the effective Lagrangian in eq. (2). Their bottom parts are the SM charged current interactions. The final formula for the inverse half-life of $0\nu\beta\beta$-decay reads

$$T_{1/2}^{-1}(0\nu\beta\beta) = |\mathcal{M}_{GT}|^2\frac{2}{G_F^2}\left[\tilde{C}_1 a^2 + C_4 b_R^2 + 2C_5 b_L^2\right] \tag{6}$$

with

$$a = \frac{\epsilon_S}{M_S^2} + \frac{\epsilon_V}{M_V^2}, \quad b_{L,R} = \left(\frac{\alpha_S^{(L,R)}}{M_S^2} + \frac{\alpha_V^{(L,R)}}{M_V^2}\right), \quad \tilde{C}_1 = C_1\left(\frac{\mathcal{M}_1^{(\nu)}/(m_e R)}{M_{GT} - \alpha_2 M_F}\right)^2 \tag{7}$$

In eq. (6) the coefficients C_n are defined following [2]; m_e and R are the electron mass and nuclear radius. We kept only the dominant terms in eq. (6), neglecting, particularly, terms proportional to the neutrino mass m_ν which we assume to be very small and put $m_\nu = 0$ in eq. (6). Also mixed terms, such as $a \cdot b_{L/R}$, are not accounted for, since these are expected to only slightly affect our numerical limits. The new matrix element $\mathcal{M}_1^{(\nu)}$ was introduced and calculated in ref. [9] within the pn-QRPA framework. Calculating C_i within the same approach [13] for the particular case of ^{76}Ge we have a complete set of nuclear structure coefficients in eq. (6) (all in units of inverse years): $|\mathcal{M}_{GT}|^2\tilde{C}_1 = 1.63 \times 10^{-10}$, $|\mathcal{M}_{GT}|^2 C_4 = 1.36 \times 10^{-13}$, $|\mathcal{M}_{GT}|^2 C_5 = 4.44 \times 10^{-9}$.

Now we are ready to derive constraints on the LQ parameters $a, b_{L,R}$ in eq. (6). We use the result from the Heidelberg-Moscow ^{76}Ge experiment [14] $T_{1/2}^{0\nu\beta\beta}(^{76}Ge, 0^+ \to 0^+) > 7.4 \times 10^{24}$ *years* 90% *c.l.*

Assuming no spurious cancellations between the different terms in eq. (6) we derive the following constraints on the effective LQ parameters:

$$\epsilon_I \leq 2.4 \times 10^{-9} \left(\frac{M_I}{100\text{GeV}}\right)^2, \tag{8}$$

$$\alpha_I^{(L)} \leq 2.3 \times 10^{-10} \left(\frac{M_I}{100\text{GeV}}\right)^2, \tag{9}$$

$$\alpha_I^{(R)} \leq 8.3 \times 10^{-8} \left(\frac{M_I}{100\text{GeV}}\right)^2. \tag{10}$$

Recall $I = S, V$.

It is interesting to compare these constraints with the corresponding constraints from other processes [11]. Consider the helicity-suppressed decay $\pi \to e\nu$ which is extremely sensitive to the first two scalar-pseudoscalar terms in eq. (2), leading to a helicity-unsuppressed amplitude [11]. The following constraint from $\pi \to e\nu$-decay data was obtained in ref. [12]: $\epsilon_I \leq 5 \times 10^{-7} \left(M_I/100\text{GeV}\right)^2$. Apparently, the corresponding constraints from $0\nu\beta\beta$-decay in eq. (8) are more stringent by about two orders of magnitude. This confirms that $0\nu\beta\beta$-decay is a powerful probe of physics beyond the standard model.

In summary, non-observation of $0\nu\beta\beta$ decay can provide stringent bounds on parameters of extensions of the standard model. Moreover, the $0\nu\beta\beta$ decay bounds on some of these fundamental parameters can be much more stringent than those from other experiments. Previously such a conclusion was obtained for the case of the R-parity violating supersymmetric contribution to $0\nu\beta\beta$-

decay [7]-[9]. In this letter we have shown that the leptoquark mechanism allows similar conclusions.

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Figure Captions

Fig.1 Feynman graphs for the leptoquark-induced mechanism of $0\nu\beta\beta$ decay. S and V^μ stand symbolically for a) $Q = -1/3$ (upper part) and b) $Q = 2/3$ (lower part) scalar and vector LQs.

Figure 1