Left-right symmetry and leading contributions to neutrinoless double beta decay

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We study the impact of the mixing (LR mixing) between the standard model W boson and its hypothetical, heavier right-handed parter WR on the neutrinoless double beta decay (0νββ-decay) rate. Our study is done in the minimal left-right symmetric model assuming type-II dominance scenario with charge conjugation as the left-right symmetry. We then show that the 0νββ-decay rate may be dominated by the contribution proportional to this LR mixing, which at the hadronic level induces the leading-order contribution to the interaction between two pions and two charged leptons. The resulting long-range pion exchange contribution can significantly enhance the decay rate compared to previously considered short-range contributions. Finally, we find that even if future cosmological experiments rule out the inverted hierarchy for neutrino masses, there are still good prospects for a positive signal in the next generation of 0νββ-decay experiments.

Determining the properties of the light neutrinos under charge conjugation is a key challenge for particle and nuclear physics. As the only electrically neutral fermions in the Standard Model (SM) of particle physics, neutrinos are the sole SM candidates for possessing a Majorana mass. The corresponding term in the Lagrangian breaks the conservation of total lepton number (L) by two units: $\mathcal{L}_M = -y_\nu \overline{C}H^TH\ell/\Lambda$, where $\ell$ and $H$ are the SM left handed lepton doublet and Higgs doublet, respectively, and $\Lambda$ is a mass scale whose presence is needed to maintain dimensionality. After the neutral component of the Higgs doublet obtains a vacuum expectation value (vev) $v/\sqrt{2}$, the resulting Majorana mass operator is $\mathcal{L}_M = -(m_\nu/2)\overline{\nu}\nu/\Lambda$. For $y_\nu \sim \mathcal{O}(1)$, the observed scale of light neutrino masses consistent with oscillation experiments [1] and cosmological bounds [2, 3] would imply $\Lambda \gtrsim 10^{15}$ GeV.

An experimental determination that neutrinos are Majorana fermions could, thus, provide circumstantial evidence for L-violating processes at ultra-high energy scales involving new particles not directly accessible in the laboratory. In the widely-considered see-saw mechanism, the L-violating, out-of-equilibrium decays of these particles (fermions) could generate the cosmic matter-antimatter asymmetry [4]. Neutrino oscillation experiments are agnostic regarding the existence of a Majorana neutrino mass term. However, the observation of 0νββ-decay in the nuclear transition $(A,Z) \rightarrow (A,Z+2) + e^- + e^-$ [5] – a process that also violates $L$ by two units – would provide conclusive evidence that light neutrinos are Majorana fermions [6].

The recent 0νββ-decay search in the KamLAND-Zen experiment [7] provides the most stringent upper limit on the effective Majorana mass $|m_{\beta\beta}|$, which is $0.061 - 0.165$ eV at 90% confidence level (C.L.), where the range reflects the uncertainty in nuclear matrix element (NME) computations. In the three-neutrino framework [8], $|m_{\beta\beta}|$ depends on the neutrino mass spectrum. In the inverted hierarchy (IH) it is bounded below $|m_{\beta\beta}| \gtrsim 0.01$ eV, while in the normal hierarchy (NH) it can be vanishingly small. The next generation of 0νββ-decay searches with ton-scale detectors [9–14] aim for sensitivities for $|m_{\beta\beta}|$ as low as $0.01$ eV. If neutrinos are Majorana fermions, and if the IH is realized in nature, one would thus expect a non-zero result in the ton-scale experiments.

Cosmological observations provide complementary information on neutrino masses, currently constraining the sum of neutrino masses (dubbed $\Sigma m_\nu$) to be smaller than $0.12$ eV at the $2\sigma$ level [15]. Global fits [2, 3] of neutrino oscillation data, 0νββ-decay search results, and cosmological surveys show that the NH is favored over the IH at about 2$\sigma$ level. For future cosmological surveys [16–20], it is possible to exclude the IH, while the favored $|m_{\beta\beta}|$ may be out the reach of ton-scale 0νββ-decay experiments [9–14] in the three-neutrino framework. Then, it is natural to ask how one could interpret a 0νββ-decay signal if cosmological measurements and/or future oscillation experiments demonstrate conclusively that the light neutrino mass ordering is in the NH.

Here, we address this question in the context of one of the most extensively studied extensions of the SM that generically implies the existence of Majorana neutrinos: the minimal left-right symmetric model (mLRSM) [21–26]. This model may have TeV scale new particles and the contributions to the 0νββ-decay from the new right-handed sector can be appreciable. The light neutrino and new physics contributions are characterized by $G_F^2|\Delta m_{\beta\beta}|/p^2$ and $c/\Lambda^5$ [27–30], respectively. Here, the virtual neutrino momentum $p \simeq 100$ MeV, $G_F$ is the Fermi constant and $c$ denotes new Yukawa and/or gauge couplings. For $c \simeq \mathcal{O}(1)$ and $|m_{\beta\beta}| \simeq 0.1$ (0.01) eV, the new physics contribution can be comparable to the
light neutrino contribution if $\Lambda \simeq 3.7\ (5.9)\ \text{TeV}$. In particular, it has been shown [31] that in the mLRSM the contributions coming from heavy neutrinos from the exchange of two right-handed $W_R$ bosons (the RR amplitude), see Fig. 1(a), are sizable. Nonetheless, the bulk of the mLRSM parameter space would remain largely inaccessible to ton-scale $0\nu\beta\beta$-decay searches if cosmological data push the bound on $\Sigma m_\nu$ below $\sim 0.1\ \text{eV}$.

In what follows, we show that this conclusion changes dramatically in the presence of mixing between the left- and right-handed gauge bosons. This mixing results in long-range contributions associated with pion exchange due to the upper bounds on the top quark masses, respectively. No such constraint exists when charge conjugation (C) is the LR symmetry [38, 39]. In Ref. [40], they also consider the LR symmetry as C but with the maximum of $\tan \beta$ being $m_W/m_t$. Constraints from kaon CP violation and neutron electric dipole moments apply when $\alpha \neq 0$ [41–46]. Here we consider C as the LR symmetry and $\alpha = 0$ since our results are rather insensitive to fundamental sources of CP violation. There is no direct experiment bound on $\tan \beta$ since we choose $\tan \beta < 0.5$ to keep the bidoublet Yukawa coupling of order unity. We will assume $M_{W_R} = 7\ \text{TeV}$, which satisfies all the aforementioned constraints as well as the requirement of $M_{W_R} > 6\ \text{TeV}$ from the renormalization group evolution (RGE) analysis [47].

For purposes of illustration, we follow Ref. [31] and assume “type-II dominance” for neutrino masses 1. In this scenario, $m_{N_L} \propto m_{\nu_\tau}$, one has $V_L = V_R$ [31]. Using the light neutrino mass difference from solar and atmospheric neutrinos [48], $M_{W_R} = g v_R$, and fixing the neutrino mass $m_{N_{\text{max}}} = m_{N_3}$ for the NH and $m_{N_{\text{max}}} = m_{N_2}$ for the IH, it is possible to obtain all the neutrino masses in terms of the lightest neutrino mass $m_{\nu_\tau\text{min}}$.

The effective Lagrangian below the electroweak scale is

$$\mathcal{L}_{\text{eff}} = \frac{G_F^2}{\Lambda_{\beta \beta}} \left[ C_{3R}(O_{4+}^{+} - O_{4-}^{+})(\bar{e}e - \bar{\nu}_\gamma \nu_e) 
+ C_{3L}(O_{3+}^{+} + O_{3-}^{+})(\bar{e}e - \bar{\nu}_\gamma \nu_e) 
+ C_{1}(O_{1+}^{+}(\bar{e}e - \bar{\nu}_\gamma \nu_e) + C_{1}^{5}(\bar{e}e - \bar{\nu}_\gamma \nu_e)) \right],$$

where [34]

$$O_{3+}^{+} = (\bar{q}_L \gamma^\tau \gamma^\mu q_L)^{(\bar{q}_L \gamma^\tau \gamma^\mu q_L) \pm (L \to R)},$$
$$O_{1+}^{+} = (\bar{q}_R \gamma^\tau \gamma^\mu q_R)^{(\bar{q}_R \gamma^\tau \gamma^\mu q_R) \pm (L \to R)},$$
$$O_{1+}^{5} = (\bar{q}_R \gamma^\tau \gamma^\mu q_R)^{(\bar{q}_R \gamma^\tau \gamma^\mu q_R) \pm (L \to R)},$$

and $\alpha, \beta$ are the color indices, $\tau^\pm = (\tau^+ \pm \tau^-)/2$, $\tau^+$ and $\tau^-$ are the Pauli matrices.

Wilson coefficients $C_{3R}, C_{3L}$ and $C_1$ are obtained by integrating out the $W_{1,2}$ and $N_i$ arising respectively from the amplitudes in Fig. 1(a)(c)(d). We evolve them from the scale $\mu = M_{W_2}$ to an appropriately chosen hadronic scale $\Lambda_H = 2\ \text{GeV}$ [49]. The RGE proceeds in two steps: (a) $\mu = M_{W_2} \to M_{W_1}$; (b) $\mu = M_{W_1} \to \Lambda_H$ and it

1 The type-I seesaw scenario was studied in Ref. [40] where the new physics contribution can also dominate over standard light neutrino exchange scenario.
The hadron-lepton Lagrangian for the $\pi\pi\bar{e}e^c$, $\bar{N}N\pi\bar{e}e^c$ and $NN\bar{N}N\bar{e}e^c$ operators up to NNLO in chiral expansion is [34]

$$\mathcal{L}_{\chi\text{PT}} = \frac{G_F^2 F^2}{\Lambda_{\beta}} \left\{ \frac{\lambda}{\Lambda_{\beta}} \beta_1 \beta_2 \gamma_5 \right\} e^c + \partial^\alpha \bar{e}^\dagger \gamma^\alpha \bar{e}^c (\beta_3 + \beta_4 \gamma_5) e^c + \lambda_{\chi}/F_\pi \bar{N}i\gamma_5 \gamma^\tau \gamma^\tau N(\xi_3 + \xi_7 \gamma_5) e^c + 1/F^2 \bar{N}^7 \tau^N \bar{N}^7 \tau^N N(\xi_1 + \xi_7 \gamma_5) e^c + \text{h.c.} \right\} .$$

(7)

The first two-pion term contributes to the amplitude $A(nnn \rightarrow ppe^c e^-)$ at order of $p^{-2}$ with $p \lesssim m_\pi$ being the typical momentum transfer. When this leading-order (LO) amplitude $A^{LO}$ is present as in the mLRSM, it can give a dominant long-range contribution to the half-life of $0\nu\beta\beta$-decay [34]. The one-pion and four-nucleon and another two-pion terms, however, contribute at next-to-next to LO (NNLO) to the amplitude $A^{NNLO} \sim p^0$.

The dimensionless coefficients are expressed as [34] $\beta_1 = -\beta_2 = \frac{\xi}{2} C_{11} C_{11} + \frac{\xi}{2} C_{11} C_{13}$, $\beta_3 = -\beta_4 = \frac{\xi}{2} (C_{13} + C_{15})$, $\xi_3 = -\xi_6 = \frac{\xi}{2} (C_{15} + C_{13})$, and $\xi_1 = -\xi_4 = \frac{\xi}{2} C_{11} C_{11} + \frac{\xi}{2} C_{15} C_{13} + \frac{\xi}{2} C_{13} C_{15}$.

where $C_{11}(W_{12}) = 0$ and it appears due to the RGE of $C_{11}$. In Eq. (6), the non-vanishing Wilson coefficients at the electroweak scale are given by $C_{11}(W_{12}) = -4\lambda\xi$, $C_{33}(W_{12}) = \lambda^2$, and $C_{33}(W_{12}) = \lambda^2 (1 + 4\lambda^2/\Lambda_{\beta})$ with $1/\Lambda_{\beta} = \sum_{i=1}^3 |V_{Re}|^2/m_N$. Note that O_{33} \approx O_{11}^{1+} + O_{11}^{1-} \approx O_{11}^{1-} \approx O_{11}^{1+}$ are matched to effective operators above the electroweak scale, while operator $O_{1^+}$ does evolve under QCD running [40], so that the RGE only includes step (b).

The doubly charged scalar, depicted in Fig. 1(b), contributes solely to $C_{33}$. When the LR symmetry holds, this contribution is negligible due to collider bounds [52] and charged lepton flavor violation constraints [31]. On the contrary, when the LR symmetry is explicitly broken, these constraints are relaxed and the corresponding contribution to the $0\nu\beta\beta$-decay rate can be appreciable. For a discussion, and the possible interplay with prospective future low- and high-energy probes, see Ref. [55].

Here, we assume a LR-symmetric Lagrangian and leave the analysis of the interesting case when it is broken for a future work.

We now map the operators in Eq. (2) at GeV scale $\sim \Lambda_\chi$ onto an effective hadron-lepton Lagrangian below that scale [34, 40, 54] using chiral perturbation theory (χPT) [55, 56]. Matching entails identifying all operators at a given chiral order that transform under chiral SU(2) the same way as the four-quark factor of a given operator in Eq. (2) [34, 57]. We refer the reader to Ref. [34] for a detailed derivation, and here simply quote the results.
where

$$m_{\nu}^{ee} \simeq \sum_{i=1}^{3} |V_{Li}|^2 m_{\nu_i} (1 + \ell^N_{\nu_i} \delta_{NN}^\nu) , \tag{9}$$

and

$$|m_{\nu}^{ee}|^2 = \frac{\Lambda^4}{72 \pi^2} \frac{M_D^2}{M_0^2} \times \left[ (\beta_1 - \zeta_5 \delta_{NN} - \beta_3 \delta_{NN} + \xi_1 \delta_{NN})^2 + (\beta_2 - \zeta_6 \delta_{NN} - \beta_4 \delta_{NN} + \xi_4 \delta_{NN})^2 \right] \tag{10}$$

with $m_N = 939$ MeV and

$$\delta_{\pi \pi} = \frac{2 m_2^2 M_2}{A^2 N_0} , \quad \delta_{NN} = \frac{\sqrt{2} m_2^2 M_4}{g_A \Lambda_N m_N M_0} , \tag{11}$$

$$\delta_{NN}^\nu = \frac{2 m_2^2 M_N}{g_A^2 A^2 \Lambda_\nu} , \quad \delta_{NN}^\nu = \frac{12 m_2^2 M_N}{g_A^2 A^2 \Lambda_\nu} . \tag{12}$$

Future ton-scale experiments searching for $0\nu\beta\beta$-decay in $^{136}$Xe are considered for numerical results. The phase space factor $G_{0\nu}^{-1} = 7.11 \times 10^{24}$ eV$^2 \cdot$yr$^{-1}$ [61, 62], and the nuclear matrix elements (NMEs) $M_\nu = 2.91$, $M_0 = -2.64$, $M_1 = -5.52$ and $M_2 = -4.20$, $M_{NN} = -1.53$ are quoted [63]. We obtain that $\delta_{\pi \pi} = 0.046$, $\delta_{NN} = 0.042$, $\delta_{NN}^\nu = -0.0096$, and $\delta_{NN}^\nu = 0.063$, clearly showing the expected chiral suppression $|A^{\text{NNLO}}/A^{\text{LO}}| \sim 15 - 20$ or even larger. Again, the LEC $\ell^N_{\nu} \sim O(1)$ in NDA and is larger requiring LO $NNNN \tilde{N}\tilde{N}ee\tilde{e}$ counterterm [60].

In Fig. 2, we show the effective Majorana mass $|m_{\beta\beta}|$ as a function of $m_{\nu_\min}$ with $m_{N_{\max}} = 500$ GeV and $M_{W_R} = 7$ TeV. To illustrate the impact of the LR contribution, we give the allowed regions with $\tan \beta = 0$ (studied in Refs. [31, 64]) and $0 < \tan \beta < 0.5$ in darker and lighter colors, respectively. For most of the $\tan \beta > 0$ parameter space, the long-range pion exchange contribution dominates over other contributions. In Fig. 3, we plot the $|m_{\beta\beta}|$ as a function of $\sum m_{\nu}$ along with the current upper bound from cosmology experiments [15]. In particular, we see from Fig. 3 (upper panel) that in the NH, inclusion of the long-range contribution opens up a significant portion of parameter space accessible to ton-scale experiments. Thus, even if the future CMB and LSS data would exclude the IH [16], there are good prospects of new physics at the TeV scale giving the dominant contribution to the $0\nu\beta\beta$-decay rate in future ton-scale ex-
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