Majorana Neutrino Masses from Neutrinoless Double-Beta Decays and Lepton-Number-Violating Meson Decays (postprint)

Authors: Jun-Hao Liu, Jue Zhang, Shun Zhou

Date: 2016-09-06T00:00:00+00:00

Abstract

The Schechter-Valle theorem states that a positive observation of neutrinoless double-beta (0) decays implies a finite Majorana mass term for neutrinos when any unlikely fine-tuning or cancellation is absent. In this note, we reexamine the quantitative impact of the Schechter-Valle theorem, and find that current experimental lower limits on the half-lives of 0 -decaying nuclei have placed a restrictive upper bound on the Majorana neutrino mass | mee |<7.43×10−29 eV radiatively generated at the four-loop level. Furthermore, we generalize this quantitative analysis of 0 decays to that of the lepton-number-violating (LNV) meson decays M−→M ++ − + − (for , = e or). Given the present upper limits on these rare LNV decays, we have derived the loop-induced Majorana neutrino masses | mee |<9.7×10−18 eV, | me |<1.6×10−15 eV and | m |<1.0×10−12 eV from K−→ ++e−+e−, K−→ ++e−+ − and K−→ ++ −+ −, respectively. A partial list of radiative neutrino masses from the LNV decays of D, Ds and B mesons is also given.

Full Text

Preamble

Majorana Neutrino Masses from Neutrinoless Double-Beta Decays and Lepton-Number-Violating Meson Decays

Jun-Hao Liua,∗, Jue Zhangb,a,†, and Shun Zhoua,b,‡

aInstitute of High Energy Physics, Chinese Academy of Sciences, Beijing 100049, China
bCenter for High Energy Physics, Peking University, Beijing 100871, China

∗E-mail: liujunhao@ihep.ac.cn
†E-mail: zhangjue@ihep.ac.cn

‡E-mail: zhoush@ihep.ac.cn

Abstract

The Schechter-Valle theorem states that a positive observation of neutrinoless double-beta (0) decays implies a finite Majorana mass term for neutrinos when any unlikely fine-tuning or cancellation is absent. In this note, we reexamine the quantitative impact of the Schechter-Valle theorem and find that current experimental lower limits on the half-lives of 0 -decaying nuclei have placed a restrictive upper bound on the radiatively generated Majorana neutrino mass | mee | < 7.43 × 10−29 eV at the four-loop level. Furthermore, we generalize this quantitative analysis of 0 decays to that of the lepton-number-violating (LNV) meson decays M− → M + − − (for , = e or μ).

Given the present upper limits on these rare LNV decays, we have derived the loop-induced Majorana neutrino masses | mee | < 1.0 × 10−12 eV, | meμ | < 1.6 × 10−15 eV, and | mμμ | < 9.7 × 10−18 eV from K− → + + e− + e−, K− → + + e− + μ−, and K− → + + μ− + μ−, respectively. A partial list of radiative neutrino masses from the LNV decays of D, Ds, and B mesons is also given.

PACS number(s): 11.30.Fs, 12.15.Lk, 14.60.Lm

Introduction

It remains an open question whether massive neutrinos are Majorana particles, whose antiparticles are themselves [?]. The final answer to this fundamental question will tell us whether lepton number is conserved in nature and help us explore the origin of neutrino masses.

Currently, the most promising way to determine if massive neutrinos are their own antiparticles is to observe 0 decays N(Z, A) → N(Z + 2, A) + e− + e−, where Z and A stand respectively for the atomic and mass numbers of a nuclear isotope N(Z, A) [?], [?]. Over the last few decades, numerous dedicated experiments have been carried out to search for this type of decay [?], [?]. So far, no positive signals have been observed, and lower bounds on the half-lives of the implemented nuclear isotopes can be drawn from experimental data. The GERDA Phase-I experiment [?] has disproved the signals of 0 decays claimed by the Heidelberg-Moscow experiment [?], and the joint lower bound from all previous 76Ge-based experiments on the half-life turns out to be T1/20 > 3.0 × 1025 yr at the 90% confidence level [?], [?]. A combined analysis of the 136Xe-based experiments KamLAND-Zen [?] and EXO-200 [?] yields a lower bound T1/20 > 3.4 × 1025 yr at the 90% confidence level. If the neutrino mass ordering is inverted (i.e., m3 < m1 < m2), next-generation 0 experiments with a few tons of target mass will be able to discover a remarkable signal in the near future [?].

The Schechter-Valle theorem [?] states that a clear signal of 0 decays will unambiguously indicate a finite Majorana mass of neutrinos, if neither a fine-tuning among parameters nor a cancellation among different contributions is assumed. This theorem signifies the physical importance of searching for 0 decays experimentally. The quantitative impact of the Schechter-Valle theorem has already been studied by Duerr, Lindner, and Merle in Ref. [?], where it was found that the Majorana neutrino masses implied by the Schechter-Valle theorem are too small to explain neutrino oscillations. Explicitly, assuming one short-range operator to be responsible for 0 decays, they found that current experimental lower bounds on the half-lives of 0 -decaying isotopes indicate an upper bound on the Majorana neutrino mass | m | < O(10−24) eV, where m denotes the effective neutrino mass term associated with L Lc for , = e, μ, . In this paper, we reexamine this problem and obtain that the upper bound is actually | mee | < 7.43 × 10−29 eV, which is about five orders of magnitude smaller. Furthermore, we generalize the analysis of 0 decays to that of LNV rare decays of B, D, and K mesons. For instance, we obtain | mee | < 1.0 × 10−12 eV from current upper bounds on the LNV rare decays of K mesons. The radiative Majorana neutrino masses related to other LNV decays are also tabulated. Therefore, we confirm the conclusion from Ref. [?] that although the Schechter-Valle theorem generally implies a tiny Majorana neutrino mass, we must explore other mechanisms to generate the observed neutrino masses at the sub-eV level.

The remainder of this work is organized as follows. In Sec. 2, we recall the calculation of Majorana neutrino masses from the four-loop diagram mediated by the effective operator that is also responsible for 0 decays. The generalization to LNV meson decays is performed in Sec. 3, where the corresponding Majorana masses are computed. Finally, we summarize our main conclusions in Sec. 4.

2 Majorana Masses from 0 Decays

In this section, we present a brief review of the calculation of Majorana neutrino masses radiatively generated from the operator that leads to 0 decays, following Ref. [?] closely. Such a calculation can be readily generalized to the case of Majorana neutrino masses induced by LNV meson decays, as shown in the next section.

At the elementary-particle level, 0 decays can be expressed as d + d → u + u + e− + e−, where the up quark u, the down quark d, and the electron e− are all massive fermions. If 0 decays take place, they can be effectively described by the LNV operator O0 = đđuuee, in which the chiralities of charged fermions have been omitted and will be specified later. As already pointed out by Schechter and Valle [?], this operator will unambiguously result in a Majorana neutrino mass term mee eL eLc. The relevant Feynman diagrams are given in Fig. 1. It is worthwhile to notice that quark and charged-lepton masses are indispensable for the Schechter-Valle theorem to be valid, as emphasized in Ref. [?]. In the Standard Model (SM), only left-handed neutrino fields participate in weak interactions, so electron masses can be implemented to convert right-handed electron fields into left-handed ones, which are then coupled to left-handed neutrino fields via the charged weak gauge boson W+. This makes sense, since the chirality of electrons in the operator O0 can in general be either left-handed or right-handed. For the same reason, quark masses are also required to realize the hadronic charged-current interactions in the SM. In this case, the operator O0 in Fig. 1(a) can be attached to left-handed neutrinos through two propagators of W+, leading to the neutrino self-energy diagram in Fig. 1(b).

Assuming that 0 decays are mediated by short-range interactions, one can write down the most general Lorentz-invariant Lagrangian containing various point-like operators as follows [?]:

$$\mathcal{L}_{0\nu\beta\beta} = \frac{G_F^2}{2m_p}\left(\epsilon_1 JJj + \epsilon_2 J^{\mu\nu}J_{\mu\nu}j + \epsilon_3 J^\mu J_\mu j + \epsilon_4 J^\mu J_{\mu\nu}j^\nu + \epsilon_5 J^\mu J j_\mu\right)$$

where G_F = 1.166 × 10−5 GeV−2 and m_p = 938.27 MeV denote respectively the Fermi constant and the proton mass, and _i (for i = 1, 2, ⋯, 5) are effective coupling constants. In Eq. (1), the hadronic currents are defined as [?]:

$$J \equiv \bar{u}(1\pm\gamma_5)d, \quad J^\mu \equiv \bar{u}\gamma^\mu(1\pm\gamma_5)d, \quad J^{\mu\nu} \equiv \bar{u}[\gamma^\mu,\gamma^\nu](1\pm\gamma_5)d,$$

while the leptonic currents are given by:

$$j = \bar{e}(1\pm\gamma_5)e^c, \quad j^\mu = \bar{e}\gamma^\mu(1\pm\gamma_5)e^c, \quad j^{\mu\nu} = \bar{e}[\gamma^\mu,\gamma^\nu](1\pm\gamma_5)e^c,$$

where ec CēT with C i 2 0 is the charge-conjugated electron field. According to the fact that fermion fields are Grassmann numbers, one can immediately verify that the tensor leptonic current j automatically vanishes. Different chiralities of hadronic and leptonic currents in Eqs. (2) and (3) should be distinguished by the left- and right-handed projection operators PL,R = (1 5)/2. For instance, we have defined JL,R = 2ūPL,Rd, and similarly for the other types of currents in Eqs. (2) and (3), in which the corresponding subscripts "L" or "R" are omitted without causing any confusion. In this connection, the effective coupling constants _i should also be regarded as xyz (for x, y, z = L, R), which carry superscripts for different chiralities of hadronic and leptonic currents.

Given one of the five operators in Eq. (1), one can set an upper limit on its coupling _i by assuming that it is responsible for the 0 decay and saturates the experimental lower bound on the half-life, as done in Ref. [?]. Recently, some of those limits have been recalculated in Ref. [?], using more recent results for the nuclear matrix elements. The effective coupling constants for the operators JLJLjL and J RJ RjL have been found to be 1 < 2 × 10−7 and 3 < 1.5

× 10−8, respectively. Having obtained these couplings, we are then ready to evaluate the induced neutrino mass by inserting the dimension-nine effective operators into the butterfly diagram, as depicted in Fig. 1. The authors of Ref. [?] have demonstrated that the operator JLJLjL leads to a vanishing neutrino mass term via the butterfly diagram, while the other one J RJ RjL does lead to a tiny Majorana neutrino mass, which we revisit below.

Now that the operator 3J RJ RjL is responsible for the 0 decays, the radiatively induced Majorana mass term for electron neutrinos can be extracted from the self-energy in Fig. 1(b) by setting the external momentum to zero as [?]:

$$\delta m^{\nu}_{ee} = \frac{128 g^4 G_F^2 \epsilon_3 m_u m_d m_e^2 I_{0\nu\beta\beta}}{(4\pi)^4 m_W^2},$$

where g is the weak gauge coupling, and m_u, m_d, and m_e are the up-quark, down-quark, and electron masses, respectively. In addition, the loop integral is given by:

$$I_{0\nu\beta\beta} = \frac{2}{(2\pi)^4(2\pi)^4}\int d^4k_1 d^4q_1 \frac{1}{(k_1^2-m_e^2)[(k_1+q_1)^2-m_u^2](q_1^2-m_d^2)(k_2^2-m_W^2)},$$

where m_W = 80.4 GeV is the W-boson mass, and k1 and q1 stand for the four-momenta of internal particles running in the loop and can be easily identified via the integrand on the right-hand side of Eq. (5) and from Fig. 1(b). To evaluate this integral, we employ the technique for massive two-loop diagrams from Ref. [?] and arrive at:

$$I_{0\nu\beta\beta} = \frac{1}{(4\pi^2)^{4+\epsilon}}\mu^{2\epsilon}\int_0^1 dz \frac{(1-z)M_W^2}{[M_W^2+zm_e^2]^2} I(m_e^2, m_u^2, m_d^2; M_W^2+zm_e^2),$$

with 4 − n as usually introduced in dimensional regularization and being the renormalization scale. The relevant function reads as [?]:

$$I(m_i^2, m_j^2, m_k^2; k^2) \equiv \int \frac{d^npd^nq}{(2\pi)^{2n}} \frac{1}{(p^2+m_i^2)[(q+k)^2+m_j^2][(p+q)^2+m_k^2]} = \pi^4(\pi m_i^2)^{n-4}\Gamma(2-\frac{n}{2})\Gamma(3-\frac{n}{2})\int_0^1 dx$$

where 2 k2/mi2 and 2 mk2/mi2. As usual, the integral is expanded with respect to = 4 − n in the limit n → 4, and the ultraviolet divergences can be separated as inverse powers of . Since the loop integral involves divergent terms proportional to −2 and −1, we need to keep terms up to 2 in I(me2, mu2, md2), namely:

$$I(m_e^2, m_u^2, m_d^2) = \frac{1}{(4\pi)^4}\left[8.02 \times 10^{-5} - 7.96 \times 10^{-4}\epsilon + 0.0041\epsilon^2 - 0.0146\epsilon^3 + 0.040\epsilon^4 + \mathcal{O}(\epsilon^5)\right],$$

so as to obtain all finite parts of I0 . In our numerical calculations, we have adopted the renormalization scale = 100 MeV, which is a characteristic scale of typical energy transfer in nuclear processes. The other implemented parameters can be found in Table 1. In the scheme of minimal subtraction, we finally obtain the induced neutrino mass from Eq. (4) as:

$$|\delta m_{ee}^{\nu}| < 7.43 \times 10^{-29}\ \text{eV},$$

which should be compared to the result | mee | < O(10−24) eV from Ref. [**?**]2. Since this mass is extremely small, one must implement other mechanisms to account for neutrino masses. In this sense, the main conclusion in Ref. [**?**] remains valid: the Schechter-Valle theorem is qualitatively correct but quantitatively irrelevant for the neutrino mass-squared differences required for neutrino oscillation experiments.

3 Majorana Masses from LNV Meson Decays

We now consider the LNV meson decays M−i → M +f − −, where M−i and M +f are the initial and final charged mesons, while the emitted same-sign charged leptons with flavors and are denoted by − and −, respectively. These processes have been extensively discussed in the presence of heavy Majorana neutrinos or a Higgs triplet [**?**]. If LNV meson decays are observed experimentally, we assume these processes are caused by short-range interactions and can be described by Lorentz-invariant operators of dimension-nine. However, to carry out an order-of-magnitude estimate of the induced Majorana neutrino masses, one can simply consider just one operator so long as it contributes dominantly. The main idea is to generalize the analysis for 0 decays to LNV meson decays. For instance, we take the operator:

$$\mathcal{L}_{MD} = \epsilon_{\alpha\beta} J_R^{\mu} j_{L\mu},$$

where (for , = e, µ) are real dimensionless couplings, and the hadronic and leptonic currents are defined similarly as before, i.e., J R = 2U PRD and jL = 2 PL c. Here U( ) and D( ) are generic up- and down-type quark fields, and we have distinguished two possibly different hadronic currents by a prime symbol. For instance, in the decay B− → D+e−µ−, the two hadronic currents are c̄ PRd and ū PRb, with b being the bottom-quark field, while the leptonic current is ēPLµc. It should be noted that the results will also be valid for the CP-conjugated channel M+i → M −f + + if CP violation in these LNV decays is negligible.

Given the operator in Eq. (12), it is straightforward to write down the Feynman amplitude for the LNV meson decay M−i(pi) → M +f(pf) +   −(p) +   −(p) as:

$$i\mathcal{M} = i\epsilon_{\alpha\beta}\frac{G_F^2}{2m_p}\langle M_f^+(p_f)\ell_\alpha^-(p_\alpha)\ell_\beta^-(p_\beta)|\bar{U}\gamma^\mu P_R D\bar{U}'\gamma_\mu P_R D'\bar{\ell}_\alpha P_L \ell_\beta^c|M_i^-(p_i)\rangle,$$

where the four-momenta of initial and final states have been specified explicitly. Since the initial and final mesons are bound states, the hadronic processes involving them generally cannot be calculated perturbatively. However, in our case, we assume that the hadronic interactions can be factorized out and related to the leptonic meson decay constants, which are defined as follows [?]:

$$\langle 0|\bar{q}\gamma^\mu\gamma_5 q'|P(p)\rangle = -ip^\mu f_P, \quad \langle 0|\bar{q}\gamma^\mu q'|V(p)\rangle = \epsilon^\mu m_V f_V,$$

where P and V denote pseudoscalar and vector mesons, respectively,   is the polarization vector for V, and mV is the vector-meson mass. For the relevant decay constants fP and fV, we adopt their numerical values from Ref. [?] and list them in Table 1.

For illustration, we first deal with the decay rates of pseudoscalar mesons. The results for vector meson decays can be similarly obtained and will be given later in this section. With the help of the decay constants, the square of the amplitude in Eq. (13) can be reduced to:

$$|\mathcal{M}|^2 = \epsilon_{\alpha\beta}^2 f_i^2 f_f^2 (p_i \cdot p_f)^2 (p_\alpha \cdot p_\beta),$$

which will be inserted into the standard formula for the differential rate of three-body decays and leads to:

$$\frac{d\Gamma}{ds} = \frac{|\mathcal{M}|^2}{(2\pi)^3 32 m_i^3}\sqrt{\lambda(s, m_\alpha^2, m_\beta^2)}\sqrt{\lambda(s, m_i^2, m_f^2)}\left(1 - \frac{m_\alpha^2 + m_\beta^2}{s}\right)\left(1 + \frac{m_\alpha^2 + m_\beta^2}{2s}\right),$$

where Ef, , are the energies of final-state particles, and s   q2 is the invariant momentum square transferred to leptons such that the condition (m + m)2 s   (mi − mf)2 is satisfied. Here mi,f, , stand for the masses of the initial- and final-state particles. After direct evaluation of the integral, we obtain:

$$\Gamma = \frac{\epsilon_{\alpha\beta}^2 f_i^2 f_f^2}{16(4\pi)^3 m_i^3}\int_{(m_\alpha+m_\beta)^2}^{(m_i-m_f)^2} ds\, \lambda^{1/2}(s, m_\alpha^2, m_\beta^2)\lambda^{1/2}(s, m_i^2, m_f^2)\frac{(s - m_\alpha^2 - m_\beta^2)^2(s - m_i^2 - m_f^2)^2}{s^3},$$

where C = 1(2) for = ( ), and (a, b, c) (a − b − c)2 − 4bc is the Källén function. The LNV decay rates of vector mesons can be derived similarly, and the final results turn out to be:

$$\Gamma_V = \frac{\epsilon^2_{\alpha\beta} f_i^2 f_f^2}{16(4\pi)^3 m_i^3} \int_{(m_\alpha+m_\beta)^2}^{(m_i-m_f)^2} ds\, \lambda^{1/2}(s, m_\alpha^2, m_\beta^2)\lambda^{3/2}(s, m_i^2, m_f^2) \frac{(s-m_\alpha^2-m_\beta^2)(s-m_i^2-m_f^2)}{s^2}.$$

Finally, the partial decay width Γ can be computed by integrating the differential width dΓ/ds over the allowed range of s. By comparing current experimental bounds on LNV rare decays from Ref. [**?**] with theoretical predictions, one can extract upper limits on the corresponding coupling constants . In Table 2 , we list such upper limits for a number of LNV meson-decay processes, and these numerical values will be used to compute the neutrino masses radiatively generated at the four-loop level, as shown in Fig. 1(b).

Since we have chosen the operator in Eq. (12) for LNV meson decays, which resembles well the one 3J RJ RjL for 0 decays in the previous section, the calculation of generated Majorana neutrino mass terms from LNV meson decays follows closely that in the case of 0 decays. The only difference is the presence of two possibly different lepton flavors and different hadronic currents, which bring the CKM matrix elements into the calculation. In the case where U U and D D do not hold simultaneously, a straightforward evaluation of a similar butterfly diagram leads to an induced neutrino mass m for and lepton flavors, namely:

$$\delta m^\nu_{\alpha\beta} = \frac{64 g^4 G_F^2 \epsilon_{\alpha\beta} m_U m_D m_{U'} m_{D'} m_\alpha m_\beta C_{UU'} C_{DD'}}{m_p (4\pi)^4 m_W^4} \Big[V_{UD} V^*_{U'D'} \cdot I(m_\alpha^2, m_U^2, m_D^2) \cdot I(m_\beta^2, m_{U'}^2, m_{D'}^2) + (\alpha \leftrightarrow \beta$$

where VU( )D( ) is the CKM matrix element, CUU and CDD follow the same definition of C below Eq. (17), and the loop integral I is the same as that introduced in Eq. (6). On the other hand, when U U and D D are both present, we obtain:

$$\delta m^\nu_{\alpha\beta} = \frac{16 g^4 G_F^2 \epsilon_{\alpha\beta} m_U m_D m_{U'} m_{D'} m_\alpha m_\beta}{m_p (4\pi)^4 m_W^4} \Big[V_{UD} V^*_{U'D'} \cdot I(m_\alpha^2, m_U^2, m_D^2) \cdot I(m_\beta^2, m_{U'}^2, m_{D'}^2) + V_{UD'} V^*_{U'D} \cdot I(m_\alpha^2,$$

Using numerical values of quark and lepton masses, CKM mixing angles ij (for ij = 12, 13, 23), and the Dirac phase in Table 1, we have tabulated the Majorana neutrino masses implied by various types of LNV meson decays in Table 2 .

As one can observe from Table 2 , depending on the current experimental limits, the values of Majorana neutrino masses from LNV meson decays can be quite different, spanning many orders of magnitude. The LNV meson decays may indicate Majorana neutrino mass terms meμ and mμμ , which cannot be obtained from 0 decays. For instance, if the LNV decays K− → +e−μ− and K− → +μ−μ− are observed, we arrive at | meμ | 1.0 × 10−12 eV and | mμμ | 1.6 × 10−15 eV, which are still far below the required masses from neutrino oscillation experiments.

4 Summary

Whether massive neutrinos are Majorana or Dirac particles remains an unsolved fundamental problem in particle physics. According to the Schechter-Valle theorem, if 0 decays N(Z, A) → N(Z + 2, A) + e− + e− are observed in future experiments, one can claim that neutrinos do have Majorana masses. In this short note, we have revisited the quantitative impact of the Schechter-Valle theorem and shown that the Majorana neutrino mass radiatively generated at the four-loop level is | mee | < 7.43 × 10−29 eV, which is smaller by about five orders of magnitude than that given in Ref. [**?**]. Furthermore, a similar analysis has been performed for LNV meson decays M− → M + + − + −, from which the upper bounds | mee | < 1.0 × 10−12 eV, | meμ | < 1.6 × 10−15 eV, and | mμμ | < 9.7 × 10−18 eV can be derived. A list of radiative neutrino masses from other LNV rare decays of D and B mesons is also given.

Therefore, even if 0 decays or LNV meson decays are detected and the decay rates are close to current upper bounds, we must invoke some other mechanisms to produce sub-eV neutrino masses, which can be of either Dirac or Majorana nature. In the former case, massive neutrinos should be pseudo-Dirac particles, since a small Majorana mass is implied by the LNV decays. In the latter case, compared to the sub-eV neutrino masses at leading order, the radiative Majorana masses can be neglected.

Acknowledgments

One of the authors (S.Z.) is grateful to Alexander Merle and Apostolos Pilaftsis for helpful discussions, and to the Mainz Institute for Theoretical Physics (MITP) for its hospitality and partial support during the completion of this work, which has also been supported in part by the National Recruitment Program for Young Professionals and by the CAS Center for Excellence in Particle Physics (CCEPP).

References

[1] E. Majorana, Nuovo Cim. 14, 171 (1937).
[2] M. Goeppert-Mayer, Phys. Rev. 48, 512 (1935).
[3] W. H. Furry, Phys. Rev. 56, 1184 (1939).
[4] J. J. Gomez-Cadenas, J. Martin-Albo, M. Mezzetto, F. Monrabal and M. Sorel, Riv. Nuovo Cim. 35, 29 (2012) [arXiv:1109.5515].
[5] M. Doi, T. Kotani and E. Takasugi, Prog. Theor. Phys. Suppl. 83, 1 (1985);

J. D. Vergados, Phys. Rept. 361, 1 (2002) [hep-ph/0209347]; W. Rodejohann, Int. J. Mod. Phys. E 20, 1833 (2011) [arXiv:1106.1334]; S. M. Bilenky and C. Giunti, Mod. Phys. Lett. A 27, 1230015 (2012) [arXiv:1203.5250]; J. D. Vergados, H. Ejiri and F. Simkovic, Rept. Prog. Phys. 75, 106301 (2012) [arXiv:1205.0649]; W. Rodejohann, J. Phys. G 39, 124008 (2012) [arXiv:1206.2560]; H. Päs and W. Rodejohann, New J. Phys. 17, 115010 (2015) [arXiv:1507.00170].

[6] M. Agostini et al. [GERDA Collaboration], Phys. Rev. Lett. 111, 122503 (2013) [arXiv:1307.4720].

[7] H. V. Klapdor-Kleingrothaus, I. V. Krivosheina, A. Dietz and O. Chkvorets, Phys. Lett. B 586, 198 (2004) [hep-ph/0404088].

[8] H. V. Klapdor-Kleingrothaus et al., Eur. Phys. J. A 12, 147 (2001) [hep-ph/0103062].

[9] M. Auger et al. [EXO-200 Collaboration], Phys. Rev. Lett. 109, 032505 (2012) [arXiv:1205.5608].

[10] A. Gando et al. [KamLAND-Zen Collaboration], Phys. Rev. Lett. 110, 062502 (2013) [arXiv:1211.3863].

[11] J. Schechter and J. W. F. Valle, Phys. Rev. D 25, 2951 (1982).

[12] M. Duerr, M. Lindner and A. Merle, JHEP 1106, 091 (2011) [arXiv:1105.0901].

[13] H. Päs, M. Hirsch, H. V. Klapdor-Kleingrothaus and S. G. Kovalenko, Phys. Lett. B 498, 35 (2001) [hep-ph/0008182].

[14] A. Ghinculov and J. J. van der Bij, Nucl. Phys. B 436, 30 (1995) [hep-ph/9405418].

[15] K. A. Olive et al. [Particle Data Group Collaboration], Chin. Phys. C 38, 090001 (2014).

[16] J. N. Ng and A. N. Kamal, Phys. Rev. D 18, 3412 (1978); J. Abad, J. G. Esteve and A. F. Pacheco, Phys. Rev. D 30, 1488 (1984); L. S. Littenberg and R. E. Shrock, Phys. Rev. Lett. 68, 443 (1992); C. Picciotto, Phys. Rev. D 56, 1612 (1997) [hep-ph/9702313]; L. S. Littenberg and R. Shrock, Phys. Lett. B 491, 285 (2000) [hep-ph/0005285]; W. Rodejohann, Phys. Rev. D 62, 013011 (2000) [hep-ph/0003149]; K. Zuber, Phys. Lett. B 479, 33 (2000) [hep-ph/0003160]; C. Dib, V. Gribanov, S. Kovalenko and I. Schmidt, Phys. Lett. B 493, 82 (2000) [hep-ph/0006277]; A. Atre, V. Barger and T. Han, Phys. Rev. D 71, 113014 (2005) [hep-ph/0502163]; Y. L. Ma, Phys. Rev. D 79, 033014 (2009) [arXiv:0901.0863]; A. Atre, T. Han, S. Pascoli and B. Zhang, JHEP 0905, 030 (2009) [arXiv:0901.3589]; J. M. Zhang and G. L. Wang, Eur. Phys. J. C 71, 1715 (2011) [arXiv:1003.5570]; J. C. Helo, S. Kovalenko and I. Schmidt, Nucl. Phys. B 853, 80 (2011) [arXiv:1005.1607]; G. Cvetic, C. Dib, S. K. Kang and C. S. Kim, Phys. Rev. D 82, 053010 (2010) [arXiv:1005.4282]; H. R. Dong, F. Feng and H. B. Li, Chin. Phys. C 39, no. 1, 013101 (2015) [arXiv:1305.3820]; F. F. Deppisch, P. S. Bhupal Dev and A. Pilaftsis, New J. Phys. 17, no. 7, 075019 (2015) [arXiv:1502.06541]; G. Cvetic, C. Dib, C. S. Kim and J. Zamora-Saa, Symmetry 7, 726 (2015) [arXiv:1503.01358]; D. Milanes, N. Quintero and C. E. Vera, Phys. Rev. D 93, no. 9, 094026 (2016) [arXiv:1604.03177].

2At this point, we are grateful to Dr. Michael Duerr for kind communications and the confirmation of a numerical error in the evaluation of the loop integral in Ref. [**?**].

Note: Figure translations are in progress. See original paper for figures.