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NEUTRINOLESS DOUBLE BETA DECAY: ELECTRON ANGULAR CORRELATION AS A PROBE OF NEW PHYSICS

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Abstract. The angular distribution of the final electrons in the so-called long range mechanism of the neutrinoless double beta decay ($0\nu 2\beta$) is derived for the general Lorentz invariant effective Lagrangian. Possible theories beyond the SM are classified from their effects on the angular distribution, which could be used to discriminate among various particle physics models inducing $0\nu 2\beta$ decays. However, additional input on the effective couplings will be required to single out the light Majorana-neutrino mechanism. Alternatively, measurements of the effective neutrino mass and angular distribution in $0\nu 2\beta$ decays can be used to test the correlations among the parameters of the underlying physics models. This is illustrated for the left-right symmetric model, taking into account current phenomenological bounds.

1 Introduction

Neutrinoless double beta decay ($0\nu 2\beta$) is forbidden in the Standard Model (SM) by lepton number (LN) conservation, which is a consequence of the renormalizability of the SM. However, being the low energy limit of a more general theory, an extended version of the SM could contain nonrenormalizable terms (tiny to be compatible with experiments), in particular, terms that violate LN and allow the $0\nu 2\beta$ decay. Probable mechanisms of LN violation can include exchange by: Majorana neutrinos ν_Ms [1–3] (one of the main candidates after the observation of neutrino oscillations [4]), SUSY Majorana particles [5–10], scalar bilinears (SBs) [11], e.g. doubly charged dileptons (the component ξ^{--} of the $SU(2)_L$ triplet Higgs scalar etc.), leptoquarks (LQs) [12], right-handed W_R bosons [3, 13] etc. From these particles light νs are much lighter than the electron and others are much heavier than the proton, giving rise to the two possible classes of mechanisms for the $0\nu 2\beta$ decay called the long range and the short range mechanism, respectively. For both the classes, the separation of the lepton physics from the nuclear physics takes place [14] which simplifies calculations. For the first class, in contrast to the second class, the pion exchange

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mechanism is suppressed, and in this case operators (in the effective field theory) stemming from the light neutrino exchange have precisely the same form as the leading order heavy particle exchange $0\nu 2\beta$ decay operators, enabling a precise comparison among models [15]. According to the Schechter–Valle theorem [16] any mechanism of the $0\nu 2\beta$ decay produces an effective Majorana mass for the neutrino, which must therefore contribute to this decay in any case. These various contributions will have to be disentangled to extract information from the $0\nu 2\beta$ decay on the characteristics of the sources of LN violation, in particular, on the neutrino masses and mixing.

Despite a lack of confirmation for the claimed observation of the $0\nu 2\beta$ decay [17], the restrictions on the decay half-life [18] make it possible to get bounds on the parameters of the models with LN violation (see [14] for a recent discussion). Once the $0\nu 2\beta$ decay has been established with good accuracy in the forthcoming experiments, the characteristics of this decay (half-life and the angular correlation between the electrons) could be combined with all the updated information from other experiments on the neutrino mixing and masses (neutrino oscillations, tritium beta decay [19], cosmology (WMAP [20]) etc.) to perform a best fit in the multidimensional space of parameters of a general underlying particle physics model. The fit parameters also include the masses and couplings of the nonstandard particles that could be involved in various LN-violating processes mentioned earlier. This would allow to determine the dominant mechanism (or a set of competing ones) of the $0\nu 2\beta$ decay.

Our aim in this paper is to examine the possibility of the determination of the decay mechanism from the angular correlation of the final electrons in this process. We restrict ourselves to the long-range mechanism and derive the angular distribution for the general Lorentz invariant effective Lagrangian. The experimental facilities that can measure the electron angular distribution in the $0\nu 2\beta$ decay are NEMO3 [21], ELEGANT V and others [22]. We argue that the measurement of the angular correlation coefficient in these experiments would provide discrimination among the various competing scenarios of the $0\nu 2\beta$ decay. We illustrate this by parametrizing the angular distribution as $d\Gamma/d\cos\theta \sim 1 - K\cos\theta$ ($K = 1$ for the light Majorana mass scenario). Using the example of the left-right symmetric model [23], we work out the correlation among the angular coefficient K, the mass of the right-handed W_R boson, m_{W_R}, and either the effective Majorana neutrino mass $\langle m\rangle = \sum_i U_{ei}^2 m_i$ or the half-life $T_{1/2}$, taking into account the current bounds on the various parameters. It is shown that for values of $|\langle m\rangle|$ below 10 meV, the angular correlation between the electrons distinguishes the left-right symmetric models from the SM + light Majorana mass scenario.

2 Angular distribution for the long range mechanism of $0\nu 2\beta$ decay

2.1 General effective Lagrangian

For the decay mediated by light ν_Ms, the most general effective Lagrangian is the Lorentz invariant combination of the leptonic j_α and the hadronic J_α currents of definite tensor structure and helicity [24]

$$\mathcal{L} = \frac{G_F}{\sqrt{2}}[(U_{ei} + \epsilon^{V-A}_{V-A,i})j^{\mu i}_{V-A} J^{+}_{V-A,\mu} + {\sum_{\alpha,\beta}}' \epsilon^{\beta}_{\alpha i} j^{i}_{\beta} J^{+}_{\alpha} + \text{h.c.}] , \qquad (1)$$

where the hadronic and leptonic currents are defined as: $J^{+}_{\alpha} = \bar{u}O_\alpha d$ and $j^{i}_{\beta} = \bar{e}O_\beta \nu_i$; the leptonic currents contain neutrino mass eigenstates and the index i runs over the light eigenstates. Here and thereafter, a summation over the repeated indices is assumed; $\alpha, \beta = V \mp A, S \mp P, T_{L,R}$ ($O_{T_\rho} = 2\sigma^{\mu\nu} P_\rho$, $\sigma^{\mu\nu} = \frac{i}{2}[\gamma^\mu, \gamma^\nu]$, P_ρ is the projector, $\rho = L, R$); the prime indicates the summation over all Lorentz invariant contributions, except for $\alpha = \beta = V - A$, and U_{ei} is the PMNS mixing matrix [25]. Note that in Eq. (1) the currents have been scaled relative to the strength of the usual $V - A$ interaction with G_F being the Fermi coupling constant. The coefficients $\epsilon^{\beta}_{\alpha i}$ encode new physics, parametrizing deviations of the Lagrangian from the standard $V - A$ current-current form and mixing of the non-SM neutrinos.

In discussing the extension of the SM for the $0\nu 2\beta$ decay, Ref. [3] considered explicitly only nonstandard terms with

$$\epsilon^{V-A}_{V+A,i} = \kappa \frac{g'_V}{g_V} U'_{ei}, \quad \epsilon^{V+A}_{V-A,i} = \eta V'_{ei}, \quad \epsilon^{V+A}_{V+A,i} = \lambda \frac{g'_V}{g_V} V_{ei} \ . \qquad (2)$$

Implicitly, also the contributions encoded by the coefficients $\epsilon^{V-A}_{V-A,i}$ are discussed arising from the non-SM contribution to U_{ei} in $SU(2)_L \times SU(2)_R \times U(1)$ models with mirror leptons (see Ref. [3], Eq. (A.2.17)). Here V, U' and V' are the 3×3 blocks of mixing matrices for non-SM neutrinos, e.g., for the usual $SU(2)_L \times SU(2)_R \times U(1)$ model V describes the lepton mixing for neutrinos from right-handed lepton doublets; for $SU(2)_L \times SU(2)_R \times U(1)$ model with mirror leptons [26] U' (V') describes the lepton mixing for mirror left(right)-handed neutrinos [3] etc. The form factors g_V and g'_V are expressed through the mixing angles for left- and right-handed quarks. Thus, $g_V = \cos\theta_C$ and $g'_V = e^{i\delta}\cos\theta'_C$, with θ_C being the Cabibbo angle, θ'_C is its right-handed mixing analogoue, and the CP violating phase δ arises in these models due to both the mixing of right-handed quarks and the mixing of left- and right-handed gauge bosons (see Ref. [3], Eq. (3.1.11)). The parameters κ, η, and λ characterize the strength of nonstandard effects. Below, we give some illustrative examples relating the couplings $\epsilon^{V-A}_{V-A,i}$, $\epsilon^{V+A}_{V\pm A,i}$ and the particle masses, couplings and the mixing parameters in the underlying theoretical models.

In the R-parity-violating (RPV) SUSY accompanying the neutrino exchange mechanism [5–10], SUSY particles (sleptons, squarks) are present in one of the two effective 4-fermion vertices. (The other vertex contains the usual W_L boson.) The nonzero parameters are

$$\epsilon^{V-A}_{V-A,i} = \frac{1}{2}\eta^{n1}_{(q)RR}U_{ni}, \quad \epsilon^{S-P}_{S+P,i} = 2\eta^{n1}_{(l)LL}U_{ni},$$
$$\epsilon^{S+P}_{S+P,i} = -\frac{1}{4}\left(\eta^{n1}_{(q)LR} - 4\eta^{n1}_{(l)LR}\right)U^*_{ni}, \ \epsilon^{T_R}_{T_R,i} = \frac{1}{8}\eta^{n1}_{(q)LR}U^*_{ni}, \tag{3}$$

where the index n runs over e, μ, τ (1, 2, 3), and the RPV Minimal Supersymmetric Model (MSSM) parameters ηs depend on the couplings of the RPV MSSM superpotential, the masses of the squarks and the sleptons, the mixings among the squarks and among the sleptons. Concentrating on the dominant contributions $\epsilon^{S+P}_{S+P,i}$ and $\epsilon^{T_R}_{T_R,i}$ (as the others are helicity-suppressed), one can express $\eta^{n1}_{(q)LR}$ and $\eta^{n1}_{(l)LR}$ as follows [9]

$$\eta^{n1}_{(q)LR} = \sum_k \frac{\lambda'_{11k}\lambda'_{nk1}}{2\sqrt{2}G_F}\sin 2\theta^d_{(k)}\left(\frac{1}{m^2_{\tilde{d}_1(k)}} - \frac{1}{m^2_{\tilde{d}_2(k)}}\right),$$
$$\eta^{n1}_{(l)LR} = \sum_k \frac{\lambda'_{k11}\lambda_{n1k}}{2\sqrt{2}G_F}\sin 2\theta^e_{(k)}\left(\frac{1}{m^2_{\tilde{e}_1(k)}} - \frac{1}{m^2_{\tilde{e}_2(k)}}\right), \tag{4}$$

where k is the generation index, $\theta^d_{(k)}$ and $\theta^e_{(k)}$ are the squark and slepton mixing angles, respectively, $m_{\tilde{f}_1}$ and $m_{\tilde{f}_2}$ are the sfermion mass eigenvalues, and λ_{ijk} and λ'_{ijk} are the RPV-couplings in the superpotential.

For the mechanism with LQs in one of the effective vertices [12], the nonzero coefficients are

$$\epsilon^{S+P}_{S-P} = -\frac{\sqrt{2}}{4G_F}\frac{\epsilon_V}{M_V^2}, \quad \epsilon^{S+P}_{S+P} = -\frac{\sqrt{2}}{4G_F}\frac{\epsilon_S}{M_S^2},$$
$$\epsilon^{V+A}_{V-A} = -\frac{1}{2G_F}\left(\frac{\alpha^{(L)}_S}{M_S^2} + \frac{\alpha^{(L)}_V}{M_V^2}\right), \quad \epsilon^{V+A}_{V+A} = -\frac{\sqrt{2}}{4G_F}\left(\frac{\alpha^{(R)}_S}{M_S^2} + \frac{\alpha^{(R)}_V}{M_V^2}\right), \tag{5}$$

where

$$\epsilon^{\beta}_{\alpha} = U_{ei}\epsilon^{\beta}_{\alpha i}, \tag{6}$$

the parameters $\epsilon_{S(V)}$, $\alpha^{(L)}_{S(V)}$, $\alpha^{(R)}_{S(V)}$ depend on the couplings of the renormalizable LQ-quark-lepton interactions consistent with the SM gauge symmetry, the mixing parameters and the common mass scale $M_{S(V)}$ of the scalar (vector) LQs [27].

The upper bounds on some of the ϵ_α^β parameters (6) from the Heidelberg–Moscow experiment were derived in Ref. [28] using the s-wave approximation for the electrons, considering nucleon recoil terms and only one nonzero parameter $\epsilon_{\alpha i}^\beta$ in the Lagrangian at a time, see Table 1.

Table 1: Upper bounds on some of ϵ_α^β parameters (6) (CL = 90%).

ϵ_{V+A}^{V+A}	ϵ_{V-A}^{V+A}	ϵ_{S+P}^{S+P}	ϵ_{S-P}^{S+P}	$\epsilon_{T_R}^{T_R}$	$\epsilon_{T_L}^{T_R}$
6×10^{-7}	4×10^{-9}	9×10^{-9}	9×10^{-9}	1×10^{-9}	6×10^{-10}

The coefficients $\epsilon_{\alpha i}^\beta$ entering the Lagrangian (1) can be expressed as

$$\epsilon_{\alpha i}^\beta = \hat{\epsilon}_\alpha^\beta \hat{U}_{ei}, \tag{7}$$

where $\hat{U}_{ei}$ are mixing parameters for non-SM neutrinos (see, e.g., Eq. (2)). As this Lagrangian describes also ordinary β-decays (without LN violation), the coefficients $\hat{\epsilon}_\alpha^\beta$ are constrained by the existing data on precision measurements in allowed nuclear beta decays, including neutron decay [29]. For example, from these data we obtain the conservative bound

$$\left|\hat{\epsilon}_{V+A}^{V+A}\right| < 7\times10^{-2}. \tag{8}$$

From Eqs. (6), (7), (8) and Table 1 we can assume that the nonstandard mixing is small:

$$\left|U_{ei}\hat{U}_{ei}\right| \lesssim 10^{-5}. \tag{9}$$

2.2 Approximations and electron angular distribution

We have calculated only the leading order in the Fermi constant and the leading contribution of the parameters ϵ_α^β to the decay width using the approximation of relativistic electrons and non-relativistic nucleons. Following Ref. [2] we describe the outgoing electrons by plane waves approximately taking into account the effect of the nuclear Coulomb field by the Fermi factors $F(\varepsilon_s)$ [2,3] with the s-th electron energy ε_s. The non-relativistic structure of the nucleon currents in the impulse approximation is taken from Ref. [30]. We neglect the nucleon recoil terms, because the accurate calculation of the corrections should also include simultaneously a precise calculation of the effect of the nuclear Coulomb field, requiring the technique of the spherical waves for the electrons. Note that in Ref. [30] the recoil terms due to the weak magnetism were calculated in the model with only $V \mp A$ currents (the recoil terms due to the pseudoscalar form factor were not taken into account) and it was shown that in the interaction proportional to ϵ_{V-A}^{V+A} the recoil effect dominates over other contributions in the $0^+ \to 0^+$ transition. However, our numerical calculations carried out in section 3 are restricted to the coefficients U_{ei} (i.e., the SM + light Majorana νs scenario) and $\epsilon_{V+A,i}^{V+A}$ (light Majorana νs in the left-right symmetric models) entering the Lagrangian (1), and in both cases the nucleon recoil effects are not dominant.

We obtain the differential width in $\cos\theta$, where θ is the angle between the electron momenta in the rest frame of the parent nucleus in the $0^+(A,Z) \to 0^+(A,Z+2)e^-e^-$ transitions,

$$\frac{d\Gamma}{d\cos\theta} = C|M_{GT}|^2 I[(a+b)(1-k\cos\theta)] \ , \qquad I[x] = \int d\varepsilon_1 \varepsilon_1^2 \varepsilon_2^2 F(\varepsilon_1)F(\varepsilon_2)x \ , \tag{10}$$

where $\varepsilon_2 = \Delta - \varepsilon_1$ and Δ is the energy release in the process. The constant

$$C = \frac{G_F^4 g_A^4 m_e^2}{64\pi^5 R_0^2} \tag{11}$$

contains the electron mass m_e and the nuclear radius R_0, included in the definition of C so that the a and b functions and the neutrino potentials are dimensionless. The Gamow–Teller nuclear matrix element,

$$M_{GT} = \langle 0_f^+ || \sum_{a\neq b} h(r_{ab},\omega) \boldsymbol{\sigma}_a \cdot \boldsymbol{\sigma}_b \tau_+^a \tau_+^b || 0_i^+ \rangle, \tag{12}$$

contains the neutrino potential $h(r,\omega) = R_0\phi_0/r$ with $\phi_0 = e^{i\omega r}$, $r = r_{ab}$ is the distance between the nucleons a and b, and ω is the average energy of the neutrino. The operator $\tau_+^a = (\tau_1 + i\tau_2)^a/2$ converts the a-th neutron into the a-th proton; $|0_i^+\rangle$ ($\langle 0_f^+|$) is the initial (final) nuclear state. The angular correlation coefficient in Eq. (10) is

$$k = \frac{a-b}{a+b} \ , \quad -1 < k \leq 1. \tag{13}$$

The expressions for a and b for different choices of ϵ_α^β, considered only one at a time, are shown in Table 2. In this table $\varepsilon_{12} = \varepsilon_1 - \varepsilon_2$ and $\langle m \rangle$ is the effective Majorana mass. The form factors $g_V(q^2)$, $g_A(q^2)$, $F_S^{(3)}(q^2)$, and $T_1^{(3)}(q^2)$ describe the following nucleon matrix elements [31]

$$\langle P(k')|\bar{u}d|N(k)\rangle = F_S^{(3)}(q^2)\bar{\psi}(k')\tau_+\psi(k), \tag{14}$$

$$\langle P(k')|\bar{u}2\gamma^\mu P_{L,R}d|N(k)\rangle = \bar{\psi}(k')\gamma^\mu \left[g_V(q^2) \mp g_A(q^2)\gamma_5\right] \tau_+\psi(k), \tag{15}$$

$$\langle P(k')|\bar{u}2\sigma^{\mu\nu} P_{L,R}d|N(k)\rangle = \bar{\psi}(k')\left[T_1^{(3)}(q^2)\sigma^{\mu\nu} \mp \frac{i}{2}\epsilon^{\mu\nu\rho\sigma}T_1^{(3)}(q^2)\sigma_{\rho\sigma}\right]\tau_+\psi(k), \tag{16}$$

where

$$\psi = \begin{pmatrix} P \\ N \end{pmatrix} \tag{17}$$

is a nucleon isodoublet. We neglect the dipole dependence of the form factors on the momentum transfer $q = k' - k$ and omit the zero argument of the form factors. Notations similar to the ones in Ref. [3] are used in Table 2.

Table 2: Expressions for a and b in Eqs. (10) and (13) for the stated choice of ϵ_α^β.

ϵ	a	b
ϵ_{V-A}^{V-A}	$\lvert[\langle m\rangle/m_e + 2\sum_i U_{ei}\epsilon_{V-A,i}^{V-A}(m_i/m_e)](1-\chi_F)\rvert^2$	0
ϵ_{V+A}^{V-A}	$\lvert(1-\chi_F)\langle m\rangle/m_e - 2\sum_i U_{ei}\epsilon_{V+A,i}^{V-A}(m_i/m_e)(1+\chi_F)\rvert^2$	0
$\epsilon_{S\mp P}^{S+P}$	$a_0 + \frac{1}{9}(\Delta/m_e)^2(F_S^{(3)}/g_V)^2\lvert\epsilon_{S\mp P}^{S+P}\chi'_F\rvert^2$	0
$\epsilon_{T_L}^{T_R}$	$a_0 + \frac{16}{81}(T_1^{(3)}/g_A)^2(\Delta/m_e)^2\lvert\epsilon_{T_L}^{T_R}(\chi'_{GT}+3\chi'_T)\rvert^2$	0
$\epsilon_{T_L}^{T_L}$, $\epsilon_{T_R}^{T_R}$	a_0	0
ϵ	b	a
ϵ_{V-A}^{V+A}	$\frac{1}{2}\lvert\epsilon_{V-A}^{V+A}\rvert^2[(\varepsilon_{12}/m_e)^2\lvert\chi_{2+}\rvert^2 + \frac{4}{9}(\Delta/m_e)^2\lvert\chi'_P\rvert^2]$	a_0
ϵ_{V+A}^{V+A}	$\frac{1}{2}(\varepsilon_{12}/m_e)^2\lvert\epsilon_{V+A}^{V+A}\chi_{2-}\rvert^2$	a_0
$\epsilon_{S\mp P}^{S-P}$	$2(F_S^{(3)}/g_V)^2\lvert\sum_i U_{ei}\epsilon_{S\mp P,i}^{S-P}(m_i/m_e)\chi_F\rvert^2$	a_0
$\epsilon_{T_R}^{T_L}$	$32(T_1^{(3)}/g_A)^2\lvert\sum_i U_{ei}\epsilon_{T_R,i}^{T_L}(m_i/m_e)\rvert^2$	a_0

Thus,

$$\chi_{2\pm} = \chi_{GT\omega} \pm \chi_{F\omega} - \frac{1}{9}\chi_{1\mp}, \quad \chi_{1\pm} = \chi'_{GT} - 6\chi'_T \pm 3\chi'_F, \tag{18}$$

$$\chi_F = \left(\frac{g_V}{g_A}\right)^2 \frac{M_F}{M_{GT}}, \quad \chi_P = \left(\frac{g_V}{g_A}\right)\frac{M_P}{M_{GT}}, \quad \chi_X = \frac{M_X}{M_{GT}}, \ X = T,\ GT, \tag{19}$$

with Fermi M_F, pseudoscalar M_P, and tensor M_T nuclear matrix elements:

$$M_F = \langle 0_f^+ || \sum_{a\neq b} h(r_{ab},\omega)\tau_+^a\tau_+^b || 0_i^+\rangle, \tag{20}$$

$$M_P = \langle 0_f^+ || \sum_{a\neq b} h(r_{ab},\omega)\left\{(\boldsymbol{\sigma}_a - \boldsymbol{\sigma}_b)\cdot[\mathbf{n}\times\mathbf{n}_+]\right\}\tau_+^a\tau_+^b || 0_i^+\rangle, \tag{21}$$

$$M_T = \langle 0_f^+ || \sum_{a\neq b} h(r_{ab},\omega)\left[\boldsymbol{\sigma}_a\cdot\mathbf{n}\boldsymbol{\sigma}_b\cdot\mathbf{n} - \frac{1}{3}\boldsymbol{\sigma}_a\cdot\boldsymbol{\sigma}_b\right]\tau_+^a\tau_+^b || 0_i^+\rangle, \tag{22}$$

with

$$\mathbf{r} = r\mathbf{n}, \qquad \mathbf{R} = R\mathbf{n}_+, \tag{23}$$

where $\mathbf{r} = \mathbf{r}_{ab}$ ($\mathbf{R} = \mathbf{R}_{ab}$) is the difference (half sum) of radius-vectors of the nucleons a and b; $\mathbf{n}$ and $\mathbf{n}_+$ are unit vectors. The prime and the index ω imply that the matrix element in the numerator instead of h contains the neutrino potential $h' = h + \omega R_0 h_1$ or $h_\omega = h - \omega R_0 h_1$, respectively, with $h_1 = -d\phi_0/d(\omega r)$. The quantity a for all zero ϵ_α^β is called a_0 in Table 1, and is defined as:

$$a_0 \equiv |(1-\chi_F)\langle m\rangle/m_e|^2. \tag{24}$$

For the $V \mp A$ part of the Lagrangian (1), our result agrees with Ref. [3] for the relativistic electrons ($m_e/\Delta \to 0$) that weakly interact with the nucleus ($\alpha Z \to 0$), if the recoil and P-wave effects are not taken into account.

If the effects of all the interactions beyond the SM extended by the ν_Ms, which we call the "nonstandard" effects, are zero (i.e., all $\epsilon_\alpha^\beta = 0$), then $k = 1$ and the distribution (10) is proportional to $1 - \cos\theta$. The angular coefficient deviates from 1 only for the cases $b \neq 0$ irrespective of the value of a. Therefore, the presence of the "nonstandard" *first* set of parameters in Table 2, $\epsilon_{V\mp A}^{V-A}$, $\epsilon_{S\mp P}^{S+P}$, $\epsilon_{T_L}^{T_R}$, $\epsilon_{T_L}^{T_L}$ and $\epsilon_{T_R}^{T_R}$ does not change the form of the angular distribution, but the presence of the *second* set (see lower part of Table 2), $\epsilon_{V\mp A}^{V+A}$, $\epsilon_{S\mp P}^{S-P}$, $\epsilon_{T_R}^{T_L}$, $\epsilon_{T_L}^{T_L}$ and $\epsilon_{T_R}^{T_R}$ does change this distribution. Thus, experimentally establishing $k \neq 1$ would signal the presence of beyond-the-SM contribution in the $0\nu2\beta$ decay. The converse is not true; namely establishing $k = 1$ experimentally will not single out the SM + ν_Ms as the only mechanism of the $0\nu2\beta$ decay and one would require additional input/constraints on the parameters of the underlying theory with their coefficients ϵ_α^β listed in the upper part of Table 2. The coefficient k and the set $\{\epsilon\}$ of nonzero ϵ_α^βs that change the $1 - \cos\theta$ form of the distribution for the SM plus ν_Ms are given in Table 3 (the lower two entries). They correspond to the following extensions of the SM: ν_Ms plus RPV SUSY [9], ν_Ms plus right-handed currents (RC) (connected with right-handed W bosons [3] or LQs [12]).

Table 3: The angular correlation coefficient k for various SM extensions.

SM extension	$\{\epsilon\}$	k
ν_M	—	1
ν_M+RPV SUSY	ϵ_{S+P}^{S-P}	$1 - k \ll 1$
ν_M+RC	$\epsilon_{V\mp A}^{V+A}$	$-1 < k \leq 1$

Among the models that we have listed in Table 3, the coefficient ϵ_{S+P}^{S-P} for the ν_M+ RPV SUSY case is helicity suppressed (and $\epsilon_{S-P}^{S-P} = 0$), and hence the angular coefficient $k \simeq 1$ for this model. (It is a mathematical challenge to come up with a model which lifts this chiral suppression). Among the realistic models we have discussed, only the model called ν_M+ RC can essentially change the angular coefficient k from being 1. Left-right symmetric models belong to this class and we have studied these models in detail in section 3, where the correlation among the parameters K (see Eq. (34) below), m_{W_R} and $|\langle m \rangle|$ (or $T_{1/2}$) is worked out.

For a quantitative understanding of the influence of the decay mechanism on the angular distribution, it is necessary to take into account the effect of the nuclear Coulomb field on the electrons in the terms with $\epsilon_{S\mp P}^{S\mp P}$ and $\epsilon_{T_{L,R}}^{T_{L,R}}$ and the nucleon recoil terms. Note that the calculation of the above-mentioned corrections for the long range mechanism of the $0\nu2\beta$ decay will change inessentially the values of a and b for ϵ_{V-A}^{V-A}, ϵ_{V+A}^{V-A} or ϵ_{V+A}^{V+A}, varied one at a time, and hence the results of the next section.

3 Electron angular correlation in left-right symmetric models

The experimental bounds on the $\epsilon_{\alpha}^{\beta}$ are connected with the masses of new particles, their mixing angles, and other parameters specific to particular extensions of the SM [2, 3, 7, 9, 11, 12]. To illustrate the kind of correlations that the measurements of $|\langle m\rangle|$ and the angular correlation coefficient k in the $0\nu 2\beta$ decay would imply, we work out the case of the left-right symmetric models [23]. In the model $SU(2)_L \times SU(2)_R \times U(1)$ the parameter λ (see Eq. (2)) is expressed through the masses m_{W_L} and m_{W_R} of the left- and right-handed W bosons [3]:

$$\lambda = \left(m_{W_L}/m_{W_R}\right)^2, \tag{25}$$

under the condition

$$m_{W_L} \ll m_{W_R}. \tag{26}$$

Eqs. (2) and (6) yield the relation

$$\epsilon_{V+A}^{V+A} = \lambda \frac{g'_V}{g_V} U_{ei} V_{ei}. \tag{27}$$

To reduce the number of free parameters, we assume the equality of the form factors of the left- and right-handed hadronic currents:

$$g_V = g'_V. \tag{28}$$

The small masses of the observable νs are likely described by the seesaw formula that in the simplest case gives

$$m_i \sim m_D^2/M_R, \quad M_R \gg m_D, \tag{29}$$

with the Dirac mass scale m_D (for the charged leptons and the light quarks $m_D \gtrsim 1$ MeV) and the mass scale M_R of right ν_Ms (in the majority of theories $M_R \gtrsim 1$ TeV). In the left-right symmetric models these scales arise usually from the two scales of the vacuum expectation values of Higgs multiplets [23]. In the seesaw mechanism, the values of the mixing parameters V_{ei} (for i numbering light mass states) have the same order of magnitude as m_D/M_R. In our discussion we use two rather conservative values (compare with Eq. (9))

$$\epsilon = 10^{-6},\ 5 \times 10^{-7} \tag{30}$$

for the mixing parameter

$$\epsilon = |U_{ei} V_{ei}|. \tag{31}$$

We recall that here the summation index i runs only over the light neutrino mass eigenstates (the summation over the total mass spectrum including also heavy states gives strictly zero due to the orthogonality condition [3]).

From Eqs. (25), (27), (28) and (31) we have

$$m_{W_R} = m_{W_L}\left(\epsilon/\left|\epsilon^{V+A}_{V+A}\right|\right)^{1/2}. \tag{32}$$

Using Eq. (26) we note the approximate equality of m_{W_L} and the mass of the observed charged gauge boson W_1 (m_{W_1}=80.4 GeV [4]).

We now turn to work out the relations among the angular correlation coefficient k, the right-handed W-boson mass m_{W_R} and the neutrino effective mass $|\langle m\rangle|$. To this end, we note that the differential width for the only nonzero nonstandard parameter ϵ^{V+A}_{V+A} can be obtained by comparing Eqs. (10), (13), (24) and the corresponding expression for b from Table 2 with the more precise result of Ref. [3], yielding:

$$\frac{d\Gamma}{d\cos\theta}(\epsilon^{V+A}_{V+A}) = |M_{GT}|^2\frac{\ln 2}{2}(A+B)(1-K\cos\theta), \tag{33}$$

which is correct for the spherical waves of electrons distorted by the Coulomb field of the nucleus in the limit of small m_e/Δ. Here,

$$K = \frac{A-B}{A+B}\ ; \quad A = a_0 G_{01}, \quad B = \left|\epsilon^{V+A}_{V+A}\chi_{2-}\right|^2 G_{02}, \tag{34}$$

with the usual phase space factors G_{0i} ($i=1,2$) defined in Ref. [3]. Note that the angular coefficient K entering in Eq. (33) differs from the coefficient k, entering Eqs. (10) and (13), as k is a function of the electron energy ε_1 and K is obtained by integrating over the energy spectrum in Eq. (10). The a_0 and χ_{2-} in Eq. (34) contain instead of ϕ_0 the neutrino potential ϕ from Ref. [3]. In the numerical calculation we have used $\chi_F = 0.274$, $\chi_{2-} = 0.551$, $M_{GT} = 1.846$, obtained in a QRPA model with p-n pairing for ^{76}Ge, with $G_{01} = 7.928\times10^{-15}$ and $G_{02} = 12.96\times10^{-15}$ (yr^{-1}) [32]. For a recent discussion on uncertainties in $0\nu2\beta$ decay nuclear matrix elements see Ref. [33].

Using Eqs. (34), (24), and (32), we have

$$K = \frac{y-1}{y+1},\ y = \frac{G_{01}}{G_{02}}\left(\frac{1-\chi_F}{\chi_{2-}\epsilon}\frac{|\langle m\rangle|}{m_e}\right)^2\left(\frac{m_{W_R}}{m_{W_L}}\right)^4. \tag{35}$$

The correlation among K, m_{W_R} and $|\langle m\rangle|$ is shown in Fig. 1 for $\epsilon = 10^{-6}$ and in Fig. 2 for $\epsilon = 5\times10^{-7}$. We consider the values of $|\langle m\rangle|$, starting from the current upper bound from the Heidelberg–Moscow experiment, taken as $|\langle m\rangle| \leq 0.3$ eV, up to $|\langle m\rangle| = 0.001$ eV, covering most scenarios of neutrino mass hierarchies and mixing angles (see Ref. [34] for a recent discussion and update). Concerning the existing bounds on m_{W_R}, we note that from Eqs. (2), (7) and (28) one obtains $\hat{\epsilon}^{V+A}_{V+A} = \lambda$. With this, Eq. (25) and the constraint (8) derived from [29] yield $m_{W_R} > 300$ GeV. This bound is weaker than the one $m_{W_R} > 715$ GeV, obtained from the electroweak fits [4]. There is still a more stringent bound $m_{W_R} > 1.2$ TeV, obtained in Ref. [35] for the decay mediated by heavy Majorana neutrinos using arguments based on the vacuum stability [5], but it requires additional theory input. We assume $m_{W_R} \geq 1$ TeV in all our figures.

Using (32), (33), (34) and the relation $T_{1/2} = \ln 2/\Gamma$ we get the correlation among m_{W_R} and the measurable $0\nu 2\beta$ decay parameters, namely the half-life $T_{1/2}$ and the angular coefficient K:

$$K = 1 - 2G_{02}(|M_{GT}|\chi_{2-}\epsilon)^2 (m_{W_L}/m_{W_R})^4 T_{1/2}. \tag{36}$$

The correlation among K, m_{W_R} and $T_{1/2}$ is shown in Fig. 3 for $\epsilon = 10^{-6}$ and in Fig. 4 for $\epsilon = 5 \times 10^{-7}$.

Figs. 1 – 4 are the principal numerical results of this paper. They show that depending on the values of $|\langle m \rangle|$ (or $T_{1/2}$) and m_{W_R}, all values of the angular coefficient K are allowed. For example, Fig. 1 shows that it is possible to describe the angular distribution close to $1 + \cos\theta$ ($K \gtrsim -1$) by the (long range) mechanism with the right-handed boson W_R with the mass about 1 TeV for $|\langle m \rangle| \sim 1$ meV.

For illustration, in Fig. 5 we plot the differential width (33) vs. $\cos\theta$ for a set of values of $|\langle m \rangle|$ and m_{W_R}, assuming $\epsilon = 10^{-6}$. It is seen that the sensitivity of the electron angular distribution to the right-handed W-boson mass m_{W_R} increases with decreasing values of the effective Majorana neutrino mass $|\langle m \rangle|$, as can be seen from Fig. 5 (right), where this distribution is shown for $|\langle m \rangle| = 1$ meV, 3 meV and 5 meV.

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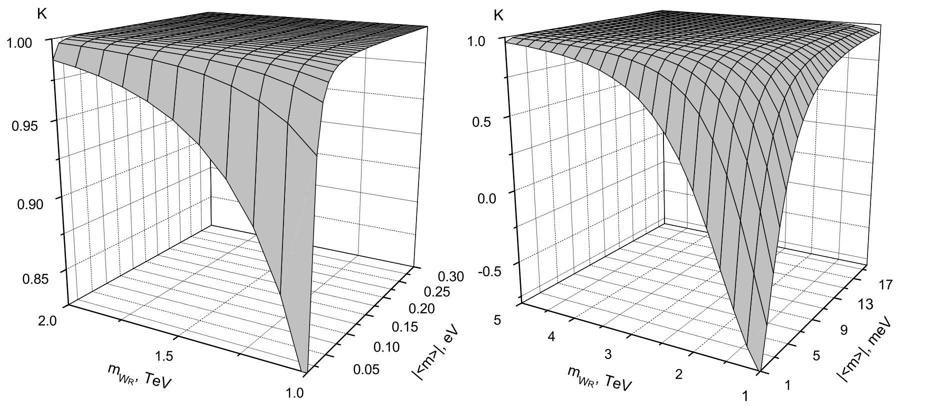
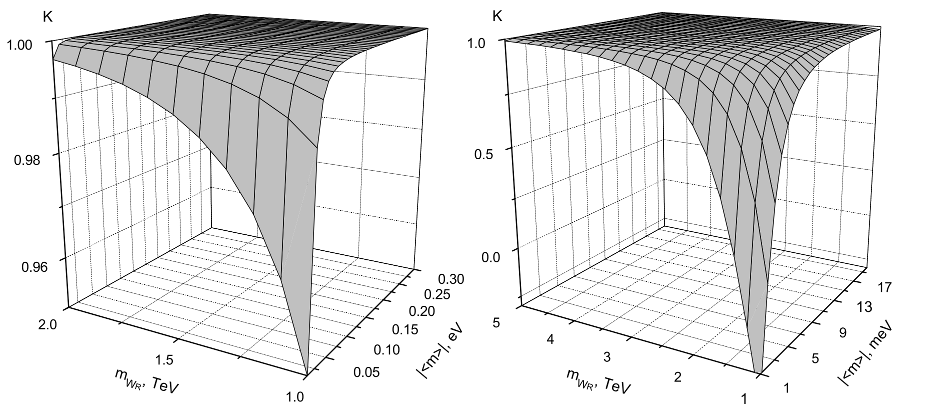


Figure 1: *Left*: Dependence of the angular correlation coefficient K on the right-handed W-boson mass m_{W_R} and the value of the neutrino effective mass $|\langle m\rangle|$ for the $0\nu 2\beta$ decay of ^{76}Ge. The mixing parameter $\epsilon = 10^{-6}$. *Right*: The same as the left figure but for smaller values of $|\langle m\rangle|$.



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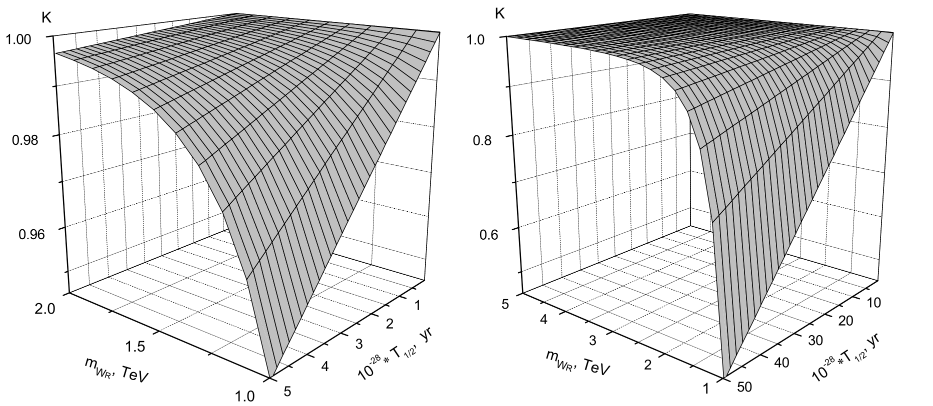

Figure 3: *Left*: Dependence of the angular correlation coefficient K on the right-handed W-boson mass m_{W_R} and the half-life $T_{1/2}$ for the $0\nu 2\beta$ decay of ^{76}Ge. The mixing parameter $\epsilon = 10^{-6}$. *Right*: The same as the left figure but for larger values of $T_{1/2}$.

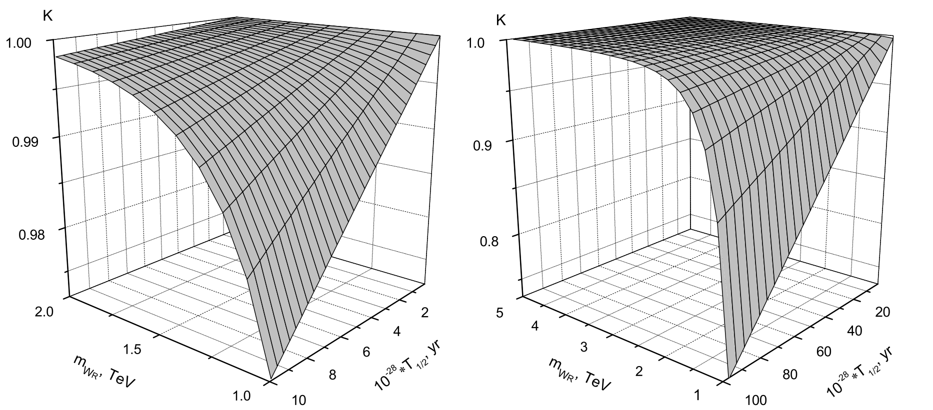

Figure 4: The same as Fig. 3 but for the smaller mixing parameter $\epsilon = 5 \times 10^{-7}$.

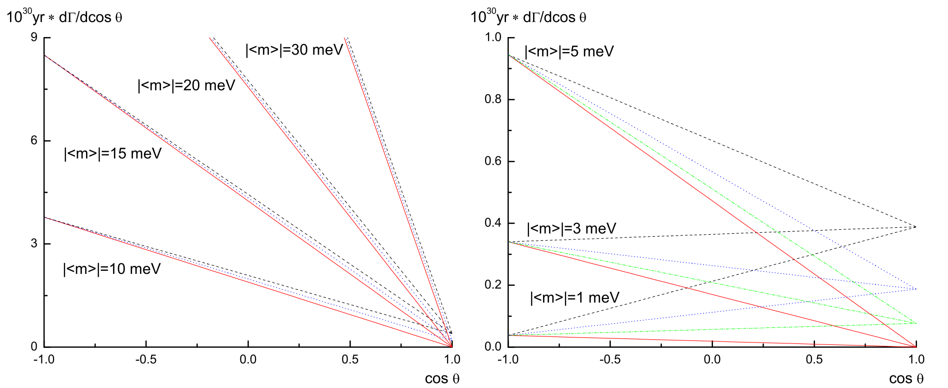

Figure 5: *Left*: Dependence of the differential width for the $0\nu2\beta$ decay of ^{76}Ge on $\cos\theta$ for a fixed value of $\epsilon = 10^{-6}$ and $|\langle m\rangle| = 10, 15, 20, 30$ meV. The dashed, dotted and straight lines correspond to $m_{W_R} = 1, 1.2, \infty$ TeV, respectively (the latter is the conventional case of the light Majorana neutrino exchange mechanism). *Right*: The same as the left figure but for smaller values of $|\langle m\rangle| = 1, 3, 5$ meV. In addition, the dash-dotted lines correspond to $m_{W_R} = 1.5$ TeV.