Neutrinoless Double Beta Decay: Searching for New Physics with Comparison of Different Nuclei

A. Ali[1], A. V. Borisov[2], D. V. Zhuridov[3]

[1]DESY, Hamburg, Germany

[2]Faculty of Physics, Moscow State University, Russia

[3]Department of Physics, National Tsing Hua University, Hsinchu, Taiwan

14th Lomonosov Conference on Elementary Particle Physics

Moscow, 19–25 August 2009

Abstract

The neutrinoless double beta decay is analyzed using a general Lorentz invariant effective Lagrangian for various decaying nuclei of current experimental interest: ^{76}Ge, ^{82}Se, ^{100}Mo, ^{130}Te, and ^{136}Xe.

We work out the half-lives and the electron angular correlations in several scenarios for new physics, in particular, in the left-right symmetric models, the R-parity-violating SUSY and models with leptoquarks.

The theoretical uncertainty in the nuclear matrix elements is discussed.

Contents

I. Introduction: Lepton number L and mechanisms of its violation in theories beyond the SM

In the Standard Model (SM), the lepton L and baryon B numbers are conserved due to the accidental $U(1)_L \times U(1)_B$ symmetry.

But the L and B nonconservation is a generic feature of various extensions of the SM.

That is why lepton-number violating processes are sensitive tools for testing theories beyond the SM.

The following LNV processes have been extensively studied:

• neutrinoless double beta decay $(A, Z) \to (A, Z+2) + e^- + e^-$ [W. H. Furry, Phys. Rev. **56**, 1184 (1939)];

• rare meson decays $M^+ \to M'^- \ell^+ \ell'^+$ $(\ell, \ell' = e, \mu)$ $(K^+ \to \pi^- \mu^+ \mu^+$ etc.);

• same-sign dilepton production in high-energy hadron-hadron, lepton-hadron, and lepton-lepton collisions:

$pp \to \ell^\pm \ell'^\pm X$, $e^\pm p \to \overset{(-)}{\nu_e} \ell^\pm \ell'^\pm X$, $e^- e^- \to W^- W^- \to \bar{\nu}_\ell \ell^- \bar{\nu}_{\ell'} \ell'^-$;

• (μ^-, e^+) conversion in nuclei $(A, Z) + \mu_b^- \to e^+ + (A, Z-2)^*$.

Introduction

There are three lepton families (generations) in the SM:

$$\begin{pmatrix} \nu_{eL} \\ e_L \end{pmatrix} \quad \begin{pmatrix} \nu_{\mu L} \\ \mu_L \end{pmatrix} \quad \begin{pmatrix} \nu_{\tau L} \\ \tau_L \end{pmatrix}$$

By definition, the lepton family number (LFN) $L_\ell = +1(-1)$ for particles $\ell = e^-, \nu_e, \ldots$ (for antiparticles $\bar{\ell} = e^+, \bar{\nu}_e, \ldots$), and $L_\ell = 0$ for leptons $\ell' \neq \ell$, $\bar{\ell}' \neq \bar{\ell}$. The total lepton number (LN)

$$L = L_e + L_\mu + L_\tau,$$

so that $L = +1(-1)$ for each ℓ ($\bar{\ell}$) and $L = 0$ for other particles (nonleptons).

In the minimal SM (with massless neutrinos), each LFN is conserved separately. For example, in the muon decay:

$$\mu^- \ \rightarrow \ e^- \ + \ \bar{\nu}_e \ + \ \nu_\mu \quad (Br \simeq 100\ \%)$$

$$\mu^- \rightarrow e^- \bar{\nu}_e \nu_\mu \gamma \ (Br = (1.4 \pm 0.4)\ \%)$$

$$\mu^- \rightarrow e^- \bar{\nu}_e \nu_\mu e^+ e^- \ (Br = (3.4 \pm 0.4) \times 10^{-5})$$

• LFN violating modes (upper bounds, $CL = 90\ \%$):

$$Br(\mu^- \to e^- \nu_e \bar{\nu}_\mu) < 1.2\ \%$$

$$Br(\mu^- \to e^- \gamma) < 1.2 \times 10^{-11}$$

$$Br(\mu^- \to e^- 2\gamma) < 7.2 \times 10^{-11}$$

$$Br(\mu^- \to e^- e^+ e^-) < 1.0 \times 10^{-12}$$

Introduction

The SM has three active neutrinos $\nu_{\ell L} (\ell = e, \mu, \tau)$ taking part in charged current (CC) and neutral current (NC) weak interactions:

$$\mathcal{L}_{\rm CC} = -\frac{g}{\sqrt{2}} \sum_{\ell} \left(\bar{\ell}_L \gamma^\mu \nu_{\ell L} W_\mu^- + \bar{\nu}_{\ell L} \gamma^\mu \ell_L W_\mu^+ \right),$$

$$\mathcal{L}_{\rm NC} = -\frac{g}{2 \cos \theta_W} \sum_{\ell} \bar{\nu}_{\ell L} \gamma^\mu \nu_{\ell L} Z_\mu .$$

The SM contains no sterile neutrinos $\nu_{\ell R}$.

In the SM, the lepton L_ℓ and baryon B numbers are conserved to all orders of perturbation theory due to the accidental global symmetry:

$$G_{\rm SM}^{\rm global} = U(1)_{L_e} \times U(1)_{L_\mu} \times U(1)_{L_\tau} \times U(1)_B ,$$

existing at the level of renormalizable operators.

The symmetry $G_{\rm SM}^{\rm global}$ is called accidental because we do not impose it intentionally. It is a direct consequence of the gauge symmetry and the choice of the representations of the physical fields.

Introduction

The SM is a chiral gauge theory, since there are L-doublets and R-singlets of the gauge group $SU(2)_L$ (they have different electroweak interactions):

$$\begin{pmatrix} \nu_{e_L} \\ e_L \end{pmatrix}, \; e_R; \quad \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \; u_R, \; d_R \text{ etc.}$$

The left-handed and right-handed chiral components of a Dirac field ψ are defined as:

$$\psi_L = P_L\psi, \quad \psi_R = P_R\psi, \quad \psi = \psi_L + \psi_R,$$

the chirality projection operators

$$P_{L,R} = \frac{1 \mp \gamma^5}{2} = P^2_{L,R}, \quad P_L P_R = 0, \quad \gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3.$$

Chirality is eigenvalue of the operator γ^5:

$$\gamma^5\psi_L = -\psi_L, \quad \gamma^5\psi_R = +\psi_R.$$

Useful relations:

$$\bar\psi\gamma^\mu\partial_\mu\psi = \bar\psi_L\gamma^\mu\partial_\mu\psi_L + \bar\psi_R\gamma^\mu\partial_\mu\psi_R,$$
$$\bar\psi\psi = \bar\psi_R\psi_L + \bar\psi_L\psi_R.$$

Introduction

We see:

- Chiral components interact with gauge fields independently.
- The Dirac mass term ($\mathcal{L}_D = -m_D\bar{\psi}\psi$) in the Lagrangian relates different chiral components and violates chirality conservation that takes place for massless (Weyl) fermions.

In the SM, neutrinos are massless due to absence of $\nu_{\ell R}$.

The only possible neutrino mass

$$\mathcal{L}_{ML} = -\frac{1}{2}m_L(\bar{\nu}_L\nu_L^c + \bar{\nu}_L^c\nu_L)$$

violates the LN: $\Delta L = \pm 2$. The global symmetry $G_{\rm SM}^{\rm global}$ prevents generation of the Majorana mass term $\mathcal{L}_{ML}$ by loop corrections.

- The $B - L$-violating terms cannot be induced even nonperturbatively because the $U(1)_{B-L}$ subgroup of $G_{\rm SM}^{\rm global}$ is non-anomalous.

Introduction

Discovery of neutrino oscillations (1998-2002) [predicted by B. Pontecorvo in 1957]:

$$\nu_\ell \to \nu_{\ell'}\,(\ell \neq \ell').$$

Here ν_ℓ is the neutrino of flavor $\ell = e, \mu, \tau$. It is created in association with the charged lepton ℓ^+ in the decay

$$W^+ \to \ell^+ + \nu_\ell.$$

Neutrino oscillations clearly demonstrate the LFN violation:

$$\Delta L_{\ell'} = -\Delta L_\ell = 1.$$

Up to now the oscillations have been observed unambiguously for

- solar (Homestake, SAGE, GALLEX-GNO, SNO: $\nu_e \to \nu_\mu(\nu_\tau)$),
- atmospheric (Super-Kamiokande: $\nu_e \to \nu_\mu(\nu_\tau)$) ,
- reactor (KamLAND: $\bar\nu_e \to \bar\nu_\mu$) [$\Leftrightarrow$ sol],
- accelerator (K2K: $\nu_\mu \to \nu_\tau$)[$\Leftrightarrow$ atm]

neutrinos [for a review, see PDG-2008].

Introduction

The neutrino oscillations imply that neutrinos are massive and mixed particles, i.e. the neutrino flavor state is a coherent superposition of neutrino mass eigenstates:

$$|\nu_\ell\rangle = \sum_i U^*_{\ell i} |\nu_i\rangle \quad (\ell = e, \mu, \tau),$$

$U = (U_{\ell i}) \equiv U_{PMNS}$ is the Pontecorvo–Maki–Nakagawa–Sakata lepton mixing matrix, ν_is are neutrinos with definite masses m_i.

The neutrino mass spectrum is nontrivial: $\Delta m^2_{jk} \equiv m^2_j - m^2_k \neq 0$.

The oscillation probability

$$P(\nu_\ell \to \nu_{\ell'}) = |\langle \nu_{\ell'} | \nu_\ell(L) \rangle|^2 \simeq \left| \sum_j U_{\ell' j} U^*_{\ell j} \exp\left(-i m^2_j \frac{L}{2E}\right) \right|^2$$

$$= \sum_j |U_{\ell j}|^2 |U_{\ell' j}|^2 + 2\mathrm{Re} \sum_{j>k} U^*_{\ell j} U_{\ell' j} U_{\ell k} U^*_{\ell' k} \exp\left(-i \Delta m^2_{jk} \frac{L}{2E}\right),$$

where E is the neutrino energy, L is the distance between a source and a detector, and the expansion of the neutrino momentum $p_j = (E^2 - m^2_j)^{1/2}$ in $(m_j/E)^2$ has been used:

$$\exp(i p_j L) \simeq e^{iEL} \exp\left(-i \frac{m^2_j}{2E} L\right).$$

Introduction

The minimal scheme of three neutrino mixing,

$$\nu_\ell = \sum_{j=1}^{3} U_{\ell j} \nu_j; \quad m_1 < m_2 < m_3,$$

provides two independent Δm^2_{jk} and allow to describe all the neutrino oscillation data (except the LSND anomaly).

Due to (PDG-2008)

$$\Delta m^2_{\rm sol} \simeq 7.6 \times 10^{-5}\ {\rm eV}^2 \ll \Delta m^2_{\rm atm} \simeq 2.7 \times 10^{-3}\ {\rm eV}^2,$$

in the leading approximation, neutrino oscillations in atmospheric and solar ranges of Δm^2 are described by two-neutrino formulas with the mixing matrix

$$U = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix}$$

For this case, the oscillation probability

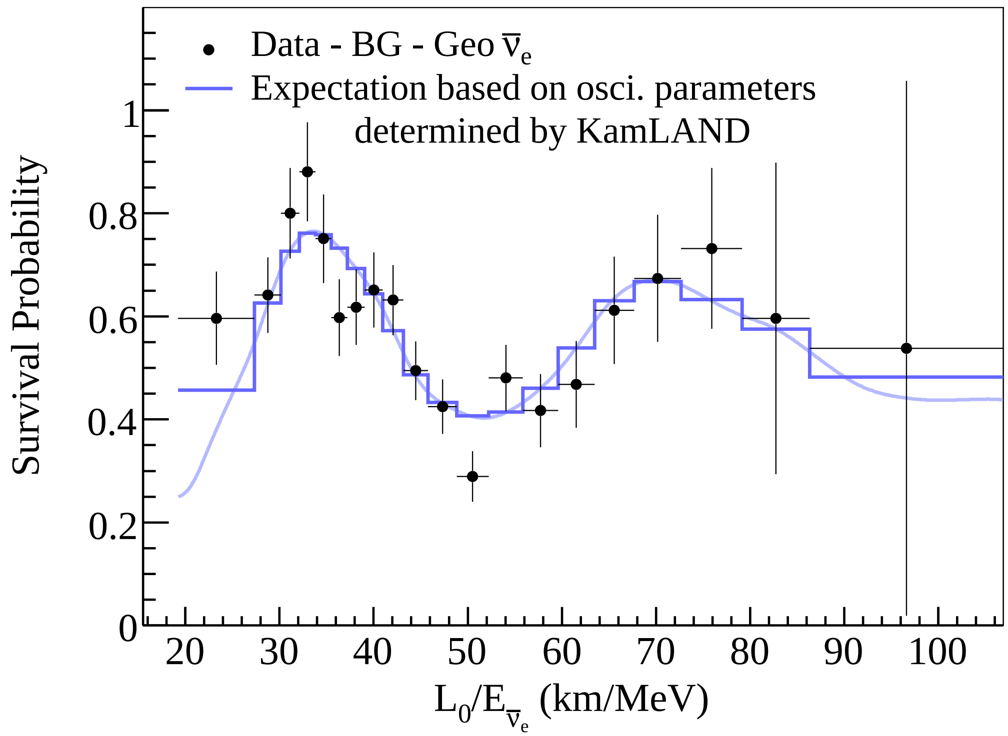
$$P(\nu_\ell \to \nu_{\ell'}) = \sin^2 2\theta \sin^2\left(1.27 \Delta m^2({\rm eV}^2) \frac{L({\rm km})}{E({\rm GeV})}\right) \quad (\ell \neq \ell'),$$

Introduction

the survival probability

$$P(\nu_\ell \to \nu_\ell) = 1 - P(\nu_\ell \to \nu_{\ell'}) = 1 - \sin^2 2\theta \sin^2 \left(1.27\Delta m^2(\mathrm{eV}^2)\frac{L(\mathrm{km})}{E(\mathrm{GeV})}\right).$$

The sinusoidal L/E dependence of the survival probability is plainly visible in the experiment: The KamLAND Collaboration (S. Abe et al.), PRL **100**, 221803 (2008), arXiv:0801.4589.



Introduction

So neutrino oscillations require extension of the SM (New Physics) to include nonzero neutrino masses and violation of LFNs.

• Including neutrino masses

The nature of neutrino masses: Dirac or Majorana?

To be Dirac or Majorana? That is one of the main unsolved questions of particle physics.

The neutrino oscillations do not probe the nature of the mass.

The Dirac neutrino carries the lepton number that distinguishes it from the antineutrino:

$$\nu^D \neq \bar{\nu}^D.$$

The Dirac neutrino mass term $\mathcal{L}_D$ is generated just like the quark and charged lepton masses via the standard Higgs mechanism with addition of right-handed neutrinos $\nu_{\ell R}$:

$$-\mathcal{L}_{\rm Yuk} = y_{\ell\ell'}\bar{L}_\ell\tilde{\varphi}\nu_{\ell' R} + \text{H.c.},$$

$$\bar{L}_\ell = (\bar{\ell}_{\ell L}, \bar{\nu}_{\ell L}), \quad \tilde{\varphi} = i\tau_2\varphi \Rightarrow \tilde{\varphi}_0 = \begin{pmatrix} v/\sqrt{2} \\ 0 \end{pmatrix};$$

Introduction

$$-\mathcal{L}_{Yuk} \Rightarrow -\mathcal{L}_D = (M_D)_{\ell\ell'}\, \bar{\nu}_{\ell L}\nu_{\ell' R} + \text{H.c.}$$

The Dirac mass matrix is complex and nondiagonal:

$$(M_D)_{\ell\ell'} = \frac{v}{\sqrt{2}} y_{\ell\ell'}.$$

Therefore $\mathcal{L}_D$ violates LFNs L_e, L_μ, L_τ, but it conserves the total LN $L = L_e + L_\mu + L_\tau$.

After the standard diagonalization:

$$\mathcal{L}_D = -\sum_{i=1}^{3} m_i \bar{\nu}_i \nu_i,$$

ν_i is the 4 component field of Dirac neutrinos with mass m_i.

Flavor fields

$$\nu_{\ell L}(x) = \sum_{i=1}^{3} U_{\ell i}\nu_{iL}(x),$$

U is the PMNS mixing matrix.

Introduction

- The Majorana mass terms (for a simple case of one flavor):

$$\mathcal{L}_{ML} = -\frac{1}{2} m_L (\bar{\nu}_L^c \nu_L + \bar{\nu}_L \nu_L^c), \quad \mathcal{L}_{MR} = -\frac{1}{2} m_R (\bar{\nu}_R^c \nu_R + \bar{\nu}_R \nu_R^c).$$

Here the charge conjugated field is defined as follows:

$$\psi^c = C\bar{\psi}^T = C\gamma^{0T}\psi^* (\psi^* = (\psi^+)^T), \quad \bar{\psi}^c \equiv \overline{\psi^c} = \psi^T C = -\psi^T C^{-1};$$

$$C = i\gamma^2\gamma^0, \quad C^+ = C^{-1}, \quad C^T = -C.$$

$$C\gamma^{\mu T} C^{-1} = -\gamma^\mu, \quad C\gamma^{5T} C^{-1} = \gamma^5.$$

Useful relations ($C : L \to R, R \to L$):

$$\psi_L{}^c \equiv (\psi_L)^c = \frac{1}{2}\left(1 + \gamma^5\right)\psi^c = (\psi^c)_R \, ,$$

$$\psi_R{}^c \equiv (\psi_R)^c = \frac{1}{2}\left(1 - \gamma^5\right)\psi^c = (\psi^c)_L \, .$$

The Majorana mass term violates lepton number by two units, $\Delta L = \pm 2$.

The Majorana neutrino is a true neutral particle identical to its antiparticle (E. Majorana, 1937): $\nu^M \equiv \bar{\nu}^M$, for the field of Majorana neutrinos:

$$\nu(x) = \nu^c(x).$$

Introduction

The Majorana field, $\psi = \psi^c$, depends only on the two independent components of ψ_L (or ψ_R):

$$\psi = \psi_L + \psi_R = \psi_L^c + \psi_R^c = \psi_L + \psi_L^c = \psi^c,$$

because $\psi_L^c = (\psi^c)_R = \psi_R$.

The total Dirac-Majorana mass term:

$$\mathcal{L}_{D+M} = \mathcal{L}_D + \mathcal{L}_{ML} + \mathcal{L}_{MR} = -\frac{1}{2}\bar{N}_L^c M N_L + \mathrm{H.c.},$$

$$N_L = \begin{pmatrix} \nu_L \\ \nu_R^c \end{pmatrix}, \quad M = \begin{pmatrix} m_L & m_D \\ m_D & m_R \end{pmatrix}.$$

where ν_L and ν_R are independent.

The diagonalized Dirac-Majorana mass term:

$$\mathcal{L}_{D+M} = -\frac{1}{2}\sum_k m_k(\bar{\nu}_{kL}^c \nu_{kL} + \bar{\nu}_{kL}\nu_{kL}^c) = -\frac{1}{2}\sum_k m_k \bar{\nu}_k \nu_k,$$

$$\nu_k = \nu_{kL} + \nu_{kL}^c = \nu_k^c.$$

The two massive neutrinos are Majorana particles.

Introduction

- The seesaw mechanism for neutrino masses

From experimental data ($0.04\ \mathrm{eV} < \mathrm{Mass}\ [\mathrm{Heaviest}\ \nu_i] < (0.07 \div 0.7)\ \mathrm{eV}$ [PDG-2008]):

$$m_\nu \ll m_\ell,\ m_q.$$

The dominant paradigm for the origin of finite but tiny neutrino mass is the seesaw mechanism [P. Minkowski (1977); T. Yanagida (1979); M. Gell-Mann, P. Ramond, R. Slansky (1979); S. Glashow (1979)]: beyond the SM (at ultra-high energies) there exists a mechanism generating the right-handed Majorana mass term, and the Dirac mass term is generated with the standard Higgs mechanism:

$$M = M_{\text{seesaw}} = \begin{pmatrix} 0 & m_D \\ m_D & m_R \end{pmatrix}, \quad m_R \gg m_D, \quad m_D \sim m_\ell \text{ or } m_q,$$

where $m_L = 0$ (no Higgs triplets!).

The neutrino ν_R is completely neutral under the SM gauge group $SU(2)_L \times U(1)_Y$, and m_R is not connected with the SM symmetry breaking scale

$$v = \left(\sqrt{2} G_F\right)^{-1/2} \simeq 246\ \mathrm{GeV},$$

but is associated to a different higher mass scale, e.g., the GUT-scale:

$$m_R \sim \Lambda_{\mathrm{GUT}} \gg m_D.$$

Introduction

Diagonalization of the mass matrix M_{seesaw} gives

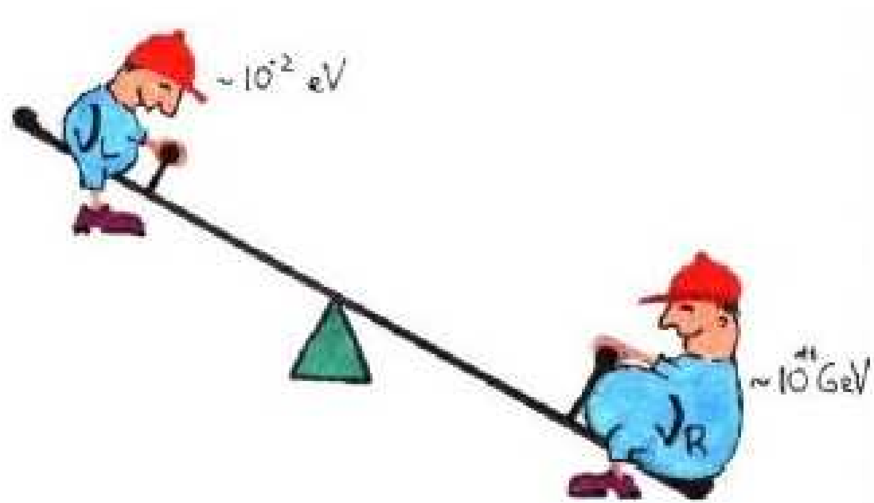
$$U^T M U = \begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix}, \quad m_1 \simeq \frac{m_D^2}{m_R} \ll m_D, \quad m_2 \simeq m_R \gg m_D;$$

$$\nu_L = i\nu_{1L}\cos\theta + \nu_{2L}\sin\theta,$$

$$\nu_R^c = -i\nu_{1L}\sin\theta + \nu_{2L}\cos\theta,$$

$$\tan 2\theta = 2m_D/m_R \ll 1.$$

The heavier ν_2 the lighter ν_1 $\Rightarrow$ Seesaw Mechanism:



Introduction

- Tiny neutrino mass from non-renormalizable operator

It is natural to assume that the SM is an effective low energy theory for $E \ll \Lambda \equiv \Lambda_{\rm NewPhysics}$:

$$\mathcal{L}_{\rm eff} = \mathcal{L}_{\rm SM} + \sum_{n>4} \frac{1}{\Lambda^{n-4}} \mathcal{O}_n \quad (\mathcal{L}_{\rm SM} \equiv \mathcal{O}_4),$$

where $\mathcal{O}_n$ (with mass dimension $n \geq 5$) includes all possible non-renormalizable operators constructed from the SM fields (the fields of heavy particles are integrated out). [The classification of $\mathcal{O}_n$'s was done by S. Weinberg.]

There is a single set of dimension-five terms that is consistent with the gauge symmetry:

$$\mathcal{O}_5 = z_{\ell\ell'} \left(\bar{L}_{\ell L} \tilde{\varphi}\right) \left(\tilde{\varphi}^T L^c_{\ell' L}\right) + \text{H.c.},$$

which violates the accidental global symmetry $G^{\rm global}_{\rm SM}$: $\Delta L = \pm 2$. After spontaneous symmetry breaking $\mathcal{O}_5$ leads to the Majorana mass term:

$$-\mathcal{L}_{ML} = \frac{1}{2} \left(M_L\right)_{\ell\ell'} \bar{\nu}_{\ell L} \nu^c_{\ell' L} + \text{H.c.},$$

$$\left(M_L\right)_{\ell\ell'} = \frac{v^2}{\Lambda} z_{\ell\ell'}.$$

Introduction

Therefore a typical neutrino mass scale

$$m_\nu \sim \frac{v^2}{\Lambda} = 6 \times 10^{-3} \left(\frac{10^{16}\ \mathrm{GeV}}{\Lambda} \right)\ \mathrm{eV}.$$

The neutrino mass is the first evidence of New Physics.

•The seesaw mechanism is a particular example of this scheme.

There exists a large number of seesaw models in which both m_D and m_R vary over many orders of magnitude, with m_R ranging somewhere between the TeV scale and the GUT scale ($\sim 10^{15} \div 10^{16}$ GeV). For example, from the atmospheric neutrino oscillations,

$$m_3 > \sqrt{\Delta m^2_{\mathrm{atm}}} \equiv m_a \simeq 5 \times 10^{-2}\ \mathrm{eV}.$$

Assuming

$$m_a = m_D^2 / m_R,$$

we obtain

$$m_R = m_D^2 / m_a \simeq 0.6 \times 10^{15}\ \mathrm{GeV} \quad \text{for } m_D = m_t \simeq 174\ \mathrm{GeV},$$

and

$$m_R \simeq 5\ \mathrm{TeV} \quad \text{for } m_D = m_e \simeq 0.511\ \mathrm{MeV}.$$

Introduction

- General case of an arbitrary number $n_s(\geq 3)$ of electroweak-singlet (sterile) neutrinos:

$$-\mathcal{L}_{D+MR} = \bar{\nu}_L M_D \nu_R + \frac{1}{2}\bar{\nu}_R^c M_R \nu_R + \mathrm{H.c.},$$

where M_D is a $3\times n_s$ Dirac mass matrix and M_R is a $n_s \times n_s$ Majorana mass matrix. Diagonalization of the neutrino mass matrix by means of a unitary $(3+n_s)\times(3+n_s)$ matrix V gives 3 light and n_s heavy Majorana neutrinos:

$$\nu_{\ell L} = \sum_{k=1}^{3} V_{\ell k}\nu_{kL}^{\mathrm{light}} + \sum_{k=4}^{n_s+3} V_{\ell k}\nu_{kL}^{\mathrm{heavy}}.$$

- A possible scenario of the generation of the Dirac-Majorana mass term $\mathcal{L}_{D+MR}$ suitable for the seesaw mechanism: the grand unified group $G_{\mathrm{GUT}} = SO(10)$ can be broken to the SM group $G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$ through the chain

$$SO(10) \xrightarrow{\Lambda_{\mathrm{GUT}}} G_{\mathrm{SM}} \times U(1)_{B-L} \xrightarrow{V} G_{\mathrm{SM}} \xrightarrow{v} SU(3)_c \times U(1)_{em},$$

where the breaking scales:

$$\Lambda_{\mathrm{GUT}} \sim 10^{15} \div 10^{16}\ \mathrm{GeV}, \quad V \sim 1 \div 10\ \mathrm{TeV}, \quad v = \left(\sqrt{2}G_F\right)^{-1/2} \simeq 246\ \mathrm{GeV}.$$

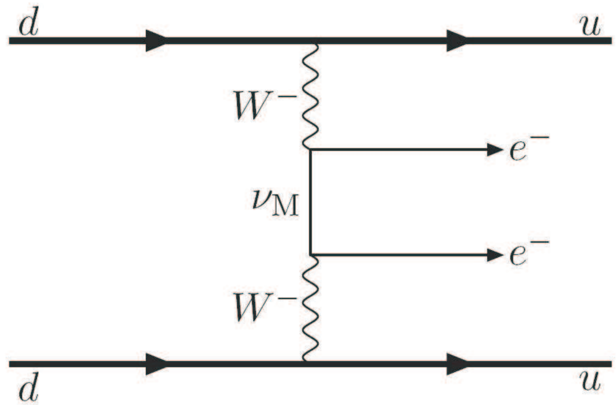
The generated mass matrices: $M_R = YV/\sqrt{2}$, $M_D = yv/\sqrt{2}$; Y and y are matrices of Yukawa couplings.

Introduction

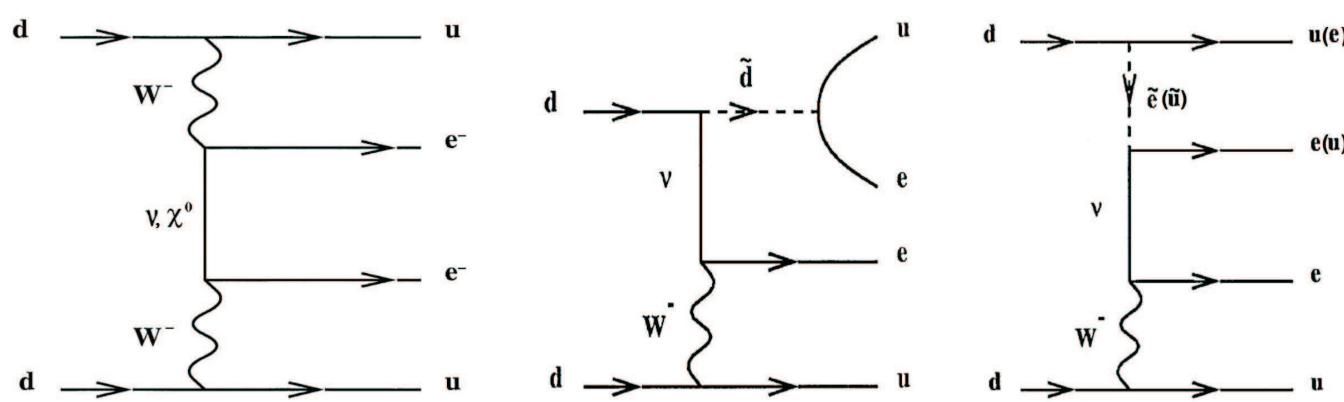
• Mechanisms of LN violation

Probable mechanisms of LN violation (for illustration, in the generic quark-level process for the $0\nu2\beta$ decay: $d + d \to u + u + e^- + e^-$) may include exchange by:

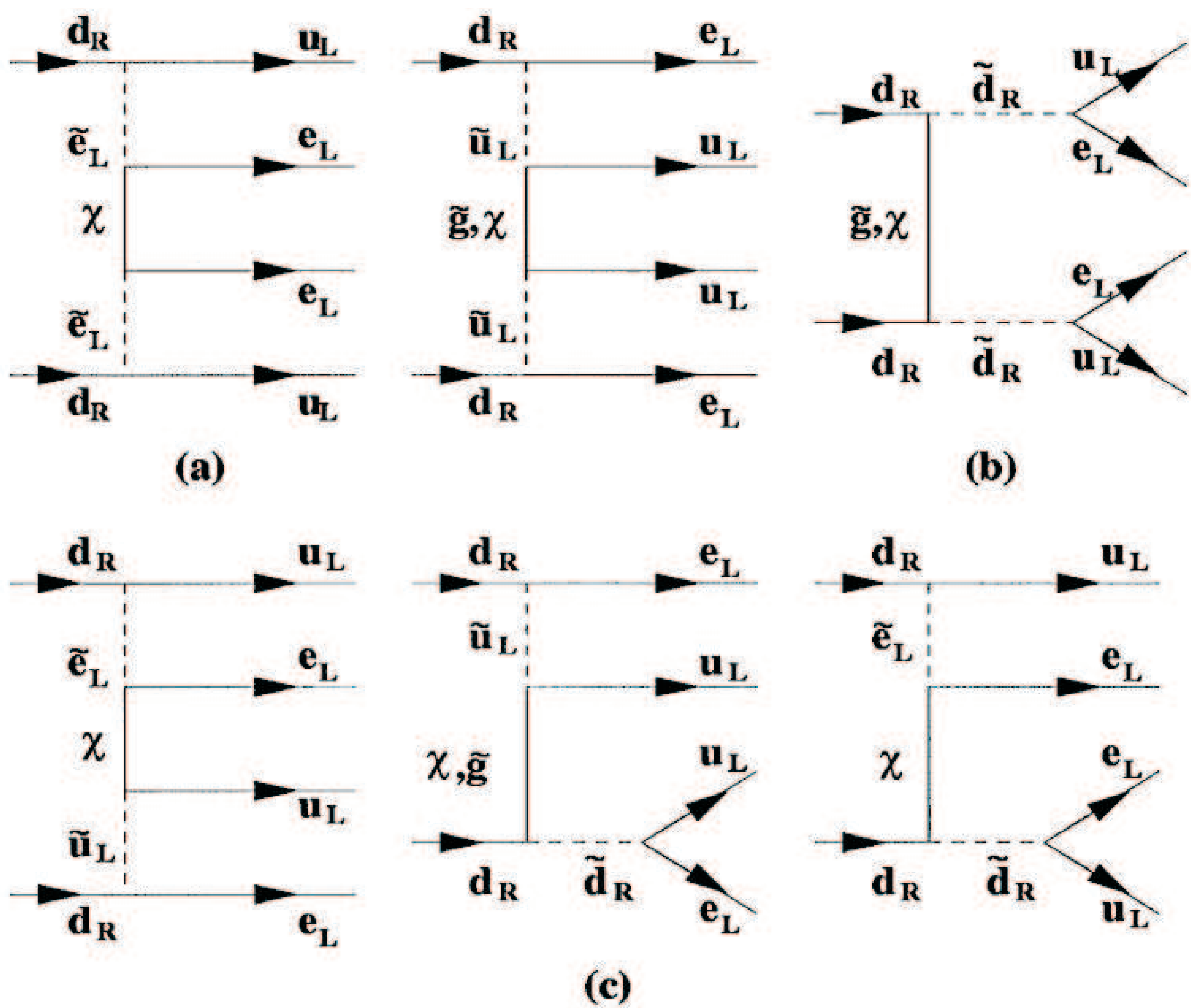
• Majorana neutrinos (the preferred mechanism after the discovery of neutrino oscillations)[SM + ν_M]



• SUSY particles [RPV MSSM: neutralinos χ^0, sleptons $\tilde{\ell}$, squarks $\tilde{q}$, gluinos $\tilde{g}$]

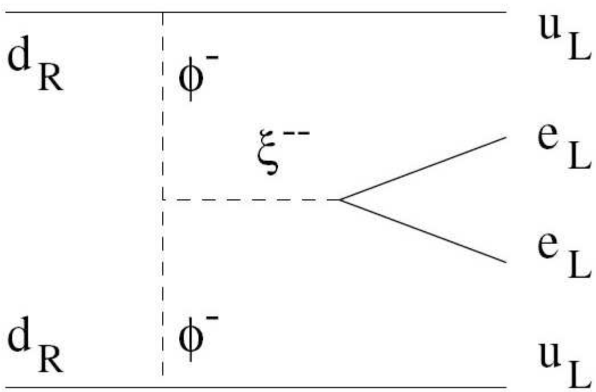


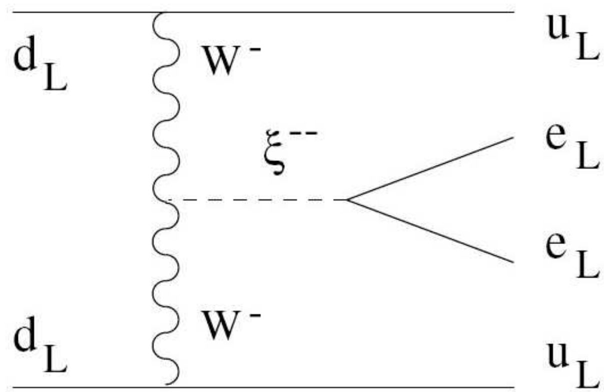
Introduction



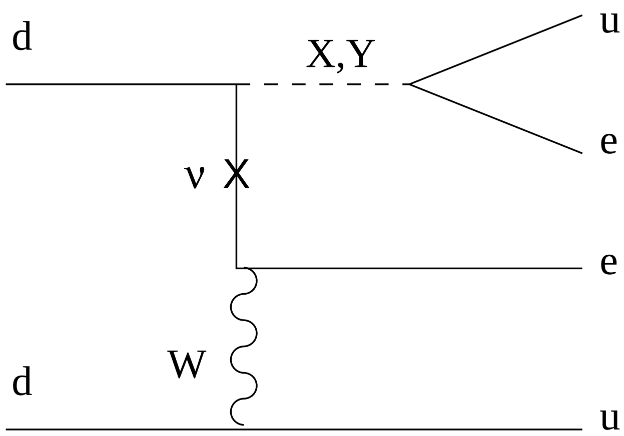
Introduction

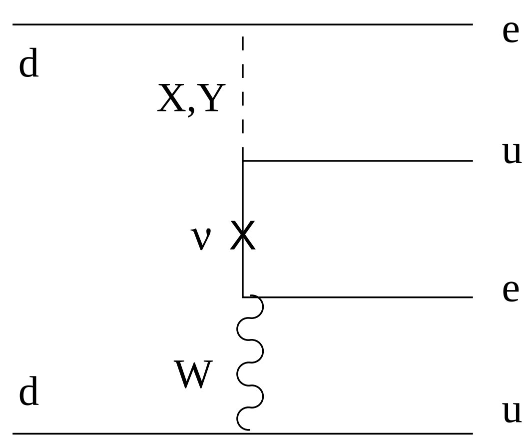
• Scalar bilinears(the component ξ^{--} of the $SU(2)_L$ triplet Higgs scalar, doubly charged dileptons etc.)





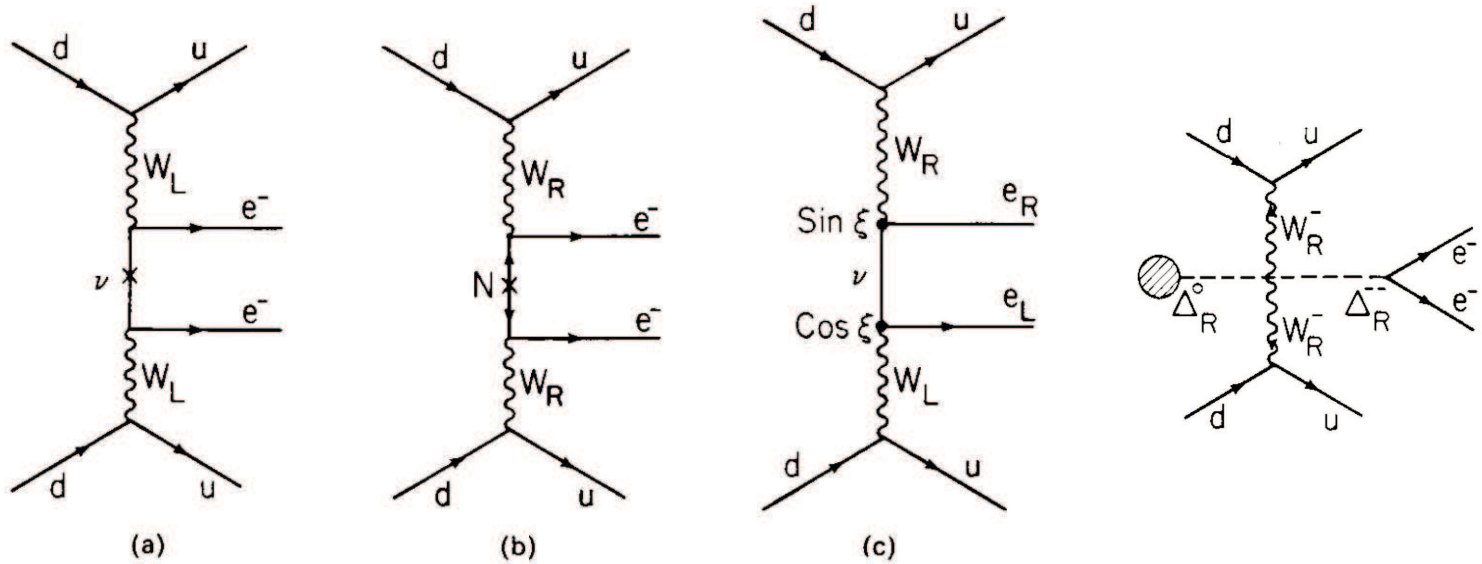
• Leptoquarks (in various extensions of the SM: scalar or vector particles carrying both L and B)





Introduction

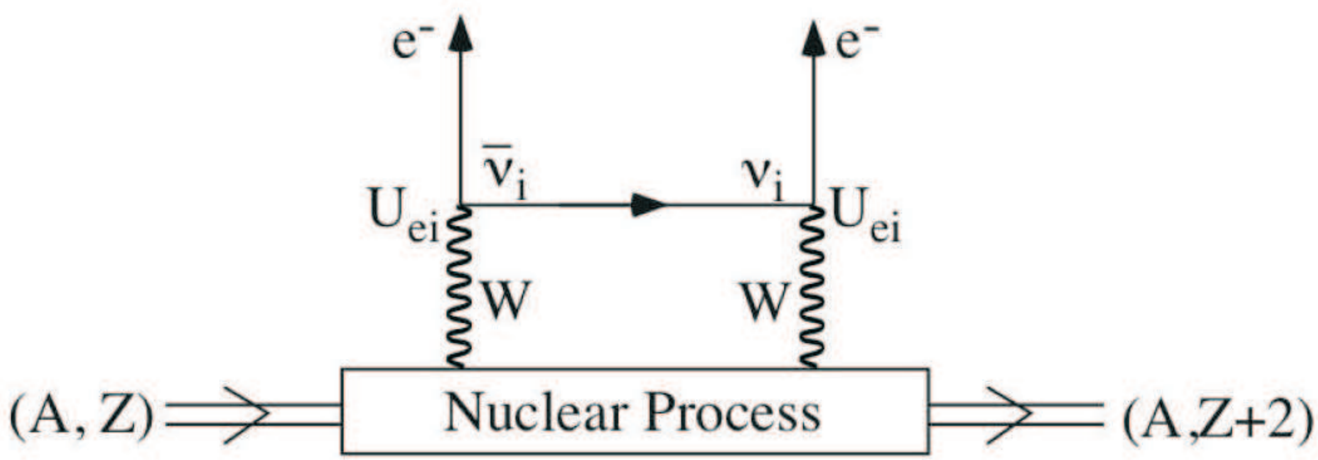
• Right-handed W_R bosons in the left-right symmetric models based on the gauge group $G_{LR} = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$
($G_{LR} \to G_{\rm SM} = SU(3)_c \times SU(2)_L \times U(1)_Y \to SU(3)_c \times (1)_{em}$), ν_R's are needed $\Rightarrow$ the seesaw mechanism:)



• Other (Kaluza–Klein sterile singlet neutrinos in theories with large extra dimensions: an infinite tower of KK neutrino mass eigenstates, ...)

II. Neutrinoless Double Beta Decay: $(A, Z) \to (A, Z+2) + e^- + e^-$

- The present bound on half-life: $T_{1/2}^{-1} > 1.6 \times 10^{25}$ yr (for ^{76}Ge).
- The conservative assumption about the dominant (?) mechanism [PDG-2008]



$$T_{1/2}^{-1} = G\,(\Delta E, Z)\,|M|^2\,|\langle m_{ee}\rangle|^2\,,$$

G is the phase space integral, M is the nuclear matrix element, the effective Majorana mass

$$\langle m_{ee}\rangle = \sum_k U_{ek}^2 m_k.$$

Neutrinoless Double Beta Decay

- The present upper bound

$$|\langle m_{ee}\rangle| < 0.3 \div 1.0 \text{ eV}.$$

A number of experiments have been proposed to reach $|\langle m_{ee}\rangle| \sim 10^{-2}$ eV.

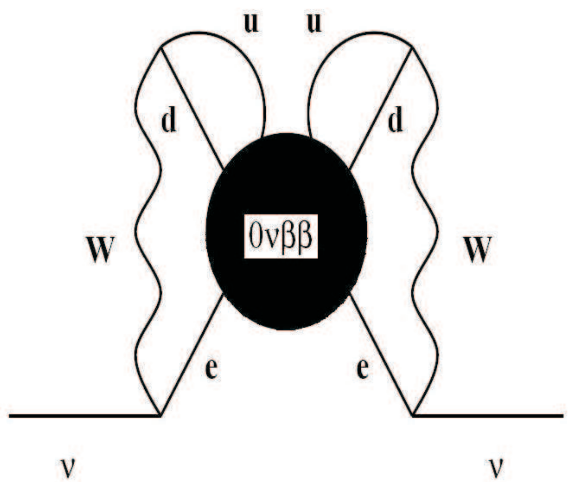
Table 1: Prospective half-life sensitivities at 90% C.L. (T_X^{90}) for different nuclei X in future projects [A. S. Barabash, 13th LCEPP (2007)]. All projects plan a second phase with lower backgrounds and higher sensitivities.

X	T_X^{90}/y	Project
^{76}Ge	2.0×10^{26}	GERDA, MAJORANA
^{82}Se	2.0×10^{26}	SuperNEMO
^{130}Te	2.1×10^{26}	CUORE
^{136}Xe	6.4×10^{25}	EXO

Neutrinoless Double Beta Decay

• Numerous LNV interactions can lead to the $0\nu 2\beta$ decay: long range and short range mechanisms $\Rightarrow$the separation of the lepton physics from the hadron physics.

• The Schechter–Valle theorem (1982): in gauge theories, any mechanism inducing the $0\nu 2\beta$ decay generates a Majorana mass for the neutrino.



• Determining the underlying physics mechanism of the $0\nu 2\beta$ decay: comparison of measurements in different nuclei [F. Deppisch and H. Päs (2007), V.M. Gehman and S.R. Elliott (2007), G.L. Fogli, E. Lisi, A.M. Rotunno (2009)], measuring the angular correlation of the final electrons [A. Ali, AVB, D. V. Zhuridov (2007)] (the proposed experimental facilities to do it: SuperNEMO, MOON, EXO).

Neutrinoless Double Beta Decay

We use the most general effective Lagrangian for the $0\nu 2\beta$ decay mediated by light Majorana neutrinos (ν_Ms)

$$\mathcal{L} = \frac{G_{\rm F} V_{ud}}{\sqrt{2}} [(U_{ei} + \epsilon^{V-A}_{V-A,i}) j^{\mu i}_{V-A} J^{+}_{V-A,\mu} + \sum_{\alpha,\beta}{}' \epsilon^{\beta}_{\alpha i} j^{i}_{\beta} J^{+}_{\alpha} + \mathrm{H.c.}] \ ,$$

the hadronic and leptonic currents: $J^{+}_{\alpha} = \bar{u} O_{\alpha} d$ and $j^{i}_{\beta} = \bar{e} O_{\beta} \nu_i$; the index i runs over the light neutrino mass eigenstates; $\alpha, \beta {=} V \mp A, S \mp P, T_{L,R}$ ($O_{T_\rho} = 2\sigma^{\mu\nu} P_\rho$, $\sigma^{\mu\nu} = \frac{i}{2}[\gamma^\mu, \gamma^\nu]$, $P_\rho = (1 \mp \gamma_5)/2$ is the projector, $\rho = L,\ R$); the prime indicates the summation over all the Lorentz invariant contributions, except for $\alpha = \beta = V - A$, U_{ei} is the PMNS mixing matrix and V_{ud} is the CKM matrix element.

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.

The nonzero coefficients

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

are given in Table 1 for several particular SM extensions: ν_Ms plus RPV SUSY, ν_Ms plus right-handed currents connected with right-handed W_R bosons or LQs.

Neutrinoless Double Beta Decay

Table 2: Nonzero coefficients ϵ_α^β for various models.

Model	Nonzero ϵs
with W_Rs	$\epsilon_{V\mp A}^{V\mp A}$
RPV SUSY	$\epsilon_{S+P}^{S\mp P}$, ϵ_{V-A}^{V-A}, $\epsilon_{T_R}^{T_R}$
with LQs	$\epsilon_{S\mp P}^{S+P}$, $\epsilon_{V\mp A}^{V+A}$

The above Lagrangian describes the effective 4-fermion vertices in the diagram representing the so-called long range mechanism of the $0\nu2\beta$ decay mediated by light neutrinos:

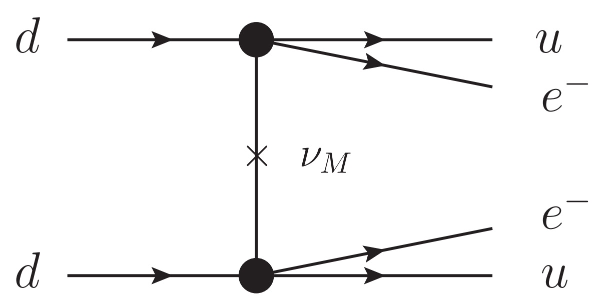


Figure 1: The long range contribution to the $0\nu2\beta$ decay.

Neutrinoless Double Beta Decay

The differential width for the $0^+(A,Z) \to 0^+(A,Z+2)e^-e^-$ transitions is

$$\frac{d\Gamma}{d\cos\theta} = \frac{\ln 2}{2}|M_{\rm GT}|^2\mathcal{A}(1-K\cos\theta),$$

where θ is the angle between the electron momenta in the rest frame of the parent nucleus, $M_{\rm GT}$ is the Gamow–Teller nuclear matrix element, and the angular correlation coefficient is

$$K = \frac{\mathcal{B}}{\mathcal{A}}\,, \quad -1 < K < 1.$$

The decay width is derived taking into account the leading contribution of the parameters ϵ_α^β, i.e. either the both 4-fermion vertices in Fig. 1 are standard $(V-A)\otimes(V-A)$ due to interaction via W_L or one of the vertices is standard and the other one is nonstandard $\alpha\otimes\beta$ (except for $\alpha=\beta=V-A$) due to interaction via W_R, LQ, a sparticle etc.

Neutrinoless Double Beta Decay

Table 3: The expressions for $\mathcal{A}$ and $\mathcal{B}$
for different choices of only one nonzero coefficient ϵ_α^β

ϵ	$\mathcal{A}$	$\mathcal{B}$
ϵ_{V-A}^{V-A}	$\mathcal{A}_0+4C_1\lvert\mu\rvert\lvert\mu_{V-A}^{V-A}\rvert c_{02}+4C_1\lvert\mu_{V-A}^{V-A}\rvert^2$	$\mathcal{B}_0+4D_1\lvert\mu\rvert\lvert\mu_{V-A}^{V-A}\rvert c_{02}+4D_1\lvert\mu_{V-A}^{V-A}\rvert^2$
ϵ_{V+A}^{V-A}	$\mathcal{A}_0+4C_0\lvert\mu\rvert\lvert\mu_{V+A}^{V-A}\rvert c_{01}+4C_{1+}\lvert\mu_{V+A}^{V-A}\rvert^2$	$\mathcal{B}_0+4D_0\lvert\mu\rvert\lvert\mu_{V+A}^{V-A}\rvert c_{01}+4D_{1+}\lvert\mu_{V+A}^{V-A}\rvert^2$
ϵ_{V-A}^{V+A}	$\mathcal{A}_0+C_3\lvert\mu\rvert\lvert\epsilon_{V-A}^{V+A}\rvert c_2+C_5\lvert\epsilon_{V-A}^{V+A}\rvert^2$	$\mathcal{B}_0+\lvert\mu\rvert\lvert\epsilon_{V-A}^{V+A}\rvert(D_3c_2+D_{3-}s_2)+D_5\lvert\epsilon_{V-A}^{V+A}\rvert^2$
ϵ_{V+A}^{V+A}	$\mathcal{A}_0+C_2\lvert\mu\rvert\lvert\epsilon_{V+A}^{V+A}\rvert c_1+C_4\lvert\epsilon_{V+A}^{V+A}\rvert^2$	$\mathcal{B}_0+\lvert\mu\rvert\lvert\epsilon_{V+A}^{V+A}\rvert(D_2c_1+D_{2-}s_1)+D_4\lvert\epsilon_{V+A}^{V+A}\rvert^2$
ϵ_{S-P}^{S-P}	$\mathcal{A}_0+4C_0^{SP}\lvert\mu\rvert\lvert\mu_{S-P}^{S-P}\rvert c_{04}+4C_1^{SP}\lvert\mu_{S-P}^{S-P}\rvert^2$	$\mathcal{B}_0+4D_{0-}^{SP}\lvert\mu\rvert\lvert\mu_{S-P}^{S-P}\rvert s_{04}+4D_1^{SP}\lvert\mu_{S-P}^{S-P}\rvert^2$
ϵ_{S+P}^{S-P}	$\mathcal{A}_0+4C_0^{SP}\lvert\mu\rvert\lvert\mu_{S+P}^{S-P}\rvert c_{03}+4C_1^{SP}\lvert\mu_{S+P}^{S-P}\rvert^2$	$\mathcal{B}_0+4D_{0-}^{SP}\lvert\mu\rvert\lvert\mu_{S+P}^{S-P}\rvert s_{03}+4D_1^{SP}\lvert\mu_{S+P}^{S-P}\rvert^2$
ϵ_{S-P}^{S+P}	$\mathcal{A}_0+C_2^{SP}\lvert\mu\rvert\lvert\epsilon_{S-P}^{S+P}\rvert c_4+C_3^{SP}\lvert\epsilon_{S-P}^{S+P}\rvert^2$	$\mathcal{B}_0+\lvert\mu\rvert\lvert\epsilon_{S-P}^{S+P}\rvert(D_2^{SP}c_4+D_{2-}^{SP}s_4)+D_3^{SP}\lvert\epsilon_{S-P}^{S+P}\rvert^2$
ϵ_{S+P}^{S+P}	$\mathcal{A}_0+C_2^{SP}\lvert\mu\rvert\lvert\epsilon_{S+P}^{S+P}\rvert c_3+C_3^{SP}\lvert\epsilon_{S+P}^{S+P}\rvert^2$	$\mathcal{B}_0+\lvert\mu\rvert\lvert\epsilon_{S+P}^{S+P}\rvert(D_2^{SP}c_3+D_{2-}^{SP}s_3)+D_3^{SP}\lvert\epsilon_{S+P}^{S+P}\rvert^2$
$\epsilon_{T_L}^{T_L}$	$\mathcal{A}_0+4C_0^T\lvert\mu\rvert\lvert\mu_{T_L}^{T_L}\rvert c_{06}+4C_1^T\lvert\mu_{T_L}^{T_L}\rvert^2$	$\mathcal{B}_0+4D_{0-}^T\lvert\mu\rvert\lvert\mu_{T_L}^{T_L}\rvert s_{06}+4D_1^T\lvert\mu_{T_L}^{T_L}\rvert^2$
$\epsilon_{T_R}^{T_L}$, $\epsilon_{T_L}^{T_R}$	$\mathcal{A}_0$	$\mathcal{B}_0$
$\epsilon_{T_R}^{T_R}$	$\mathcal{A}_0+C_2^T\lvert\mu\rvert\lvert\epsilon_{T_R}^{T_R}\rvert c_5+C_3^T\lvert\epsilon_{T_R}^{T_R}\rvert^2$	$\mathcal{B}_0+D_2^T\lvert\mu\rvert\lvert\epsilon_{T_R}^{T_R}\rvert c_5+D_3^T\lvert\epsilon_{T_R}^{T_R}\rvert^2$

$$\mathcal{A}_0=C_1|\mu|^2,\quad \mathcal{B}_0=D_1|\mu|^2,\quad \mu=\langle m\rangle/m_e,\quad \langle m\rangle\equiv\langle m_{ee}\rangle=\textstyle\sum_i U_{ei}^2 m_i.$$

The nonstandard effective Majorana masses:

$$\mu_\alpha^\beta=m_\alpha^\beta/m_e,\quad m_{S\mp P}^{S-P}=\textstyle\sum_i U_{ei}\epsilon_{S\mp P,i}^{S-P}m_i,\quad m_{V\mp A}^{V-A}=\sum_i U_{ei}\epsilon_{V\mp A,i}^{V-A}m_i,$$
$$m_{T_{L,R}}^{T_L}=\textstyle\sum_i U_{ei}\epsilon_{T_{L,R},i}^{T_L}m_i.$$

Neutrinoless Double Beta Decay

The relative phases:

$$c_i = \cos\psi_i, \quad s_i = \sin\psi_i,$$
$$\psi_{01} = \arg(\mu\mu_{V+A}^{V-A*}), \quad \psi_{02} = \arg(\mu\mu_{V-A}^{V-A*}), \quad \psi_1 = \arg(\mu\epsilon_{V+A}^{V+A*}),$$
$$\psi_2 = \arg(\mu\epsilon_{V-A}^{V+A*}), \quad \psi_{03} = \arg(\mu\mu_{S+P}^{S-P*}), \quad \psi_{04} = \arg(\mu\mu_{S-P}^{S-P*}),$$
$$\psi_3 = \arg(\mu\epsilon_{S+P}^{S+P*}), \quad \psi_4 = \arg(\mu\epsilon_{S-P}^{S+P*}), \quad \psi_{06} = \arg(\mu\mu_{T_L}^{T_L*}),$$
$$\psi_5 = \arg(\mu\epsilon_{T_R}^{T_R*}), \quad \psi_6 = \arg(\mu\epsilon_{T_L}^{T_R*}).$$

The coefficients C and D are nontrivial combinations of the integrated phase space factors and the nuclear matrix elements (for details, see: A. Ali, A. V. Borisov, D. V. Zhuridov, *Phys. Rev.* **D 76**, 093009 (2007)). For example,

$$C_0 = ({\chi_{\rm F}}^2 - 1)A_{01}, \quad A_{01} = \frac{1}{\ln 2}\frac{a_{0\nu}}{(m_e R)^2}\int a_{01} d\Omega_{0\nu},$$
$$d\Omega_{0\nu} = m_e^{-5}|\mathbf{p}_1||\mathbf{p}_2|\varepsilon_1\varepsilon_2\delta(\varepsilon_1+\varepsilon_2+E_f-E_i)d\varepsilon_1 d\varepsilon_2 d(\hat{\mathbf{p}}_1\cdot\hat{\mathbf{p}}_2);$$
$$a_{0\nu} = (G_{\rm F} g_A)^4 |V_{ud}|^4 m_e^9/(64\pi^5), \quad a_{01} = F_0(Z,\varepsilon_2)F_0(Z,\varepsilon_1),$$
$$F_0(Z,\varepsilon) = \tfrac{4}{\Gamma^2(2\gamma_1+1)}(2pR)^{2(\gamma_1-1)}|\Gamma(\gamma_1+iy)|^2 e^{\pi y}, \quad \gamma_1 = \sqrt{1-(\alpha Z)^2}, \quad y = \alpha Z\varepsilon/p;$$
$$\chi_{\rm F} = (g_V/g_A)^2 M_{\rm F}/M_{\rm GT}; \quad M_{\rm F} = \langle 0_f^+ || \sum_{a\neq b} h_+(r_{ab},\langle E_N\rangle)\tau_+^a\tau_+^b||0_i^+\rangle,$$
$$M_{\rm GT} = \langle 0_f^+ || \sum_{a\neq b} h_+(r_{ab},\langle E_N\rangle)\sigma_a\times\sigma_b\tau_+^a\tau_+^b||0_i^+\rangle.$$

Neutrinoless Double Beta Decay

The neutrino potential (in the closure approximation that is good to better than 90% due to the large average energy of the virtual neutrino $\sim 100\ \mathrm{MeV}$):

$$h_+(r, \langle E_N \rangle) = \frac{R}{4\pi^2} \int \frac{d^3k}{\omega} \left(\frac{1}{\omega + A_1} + \frac{1}{\omega + A_2} \right) e^{i\mathbf{k}\times\mathbf{r}} \simeq RH(r, \bar{A}),$$

$$H(r, \bar{A}) = \frac{1}{2\pi^2} \int \frac{d^3k}{\omega} \frac{e^{i\mathbf{k}\times\mathbf{r}}}{\omega + \bar{A}} = \frac{2}{\pi r} \int_0^\infty dk \frac{k \sin kr}{\omega \left(\omega + \bar{A}\right)}, \quad \omega = \sqrt{k^2 + m_\nu^2},$$

$$A_j = \varepsilon_j + \langle E_N \rangle - E_i,\ j = 1, 2; \quad \bar{A} = \langle E_N \rangle - (E_i + E_f)/2;$$

$r \equiv r_{ab}$ is the distance between the nucleons a and b, $\langle E_N \rangle$ is the average energy of the intermediate nucleus N; $R = r_0 A^{1/3}$ is the nuclear radius, $r_0 = 1.1$ fm.

The wavefunction of an electron with the asymptotic momentum $\mathbf{p}$ and the spin projection s can be expanded in terms of spherical waves as

$$e_{\mathbf{p}s}(\mathbf{r}) = e_{\mathbf{p}s}^{S_{1/2}}(\mathbf{r}) + e_{\mathbf{p}s}^{P_{1/2}}(\mathbf{r}) + \ldots$$

We take into account the $S_{1/2}$ and the $P_{1/2}$ waves for the outgoing relativistic electrons.

Neutrinoless Double Beta Decay

Table 4: The integrated kinematical factors $A_{0i}^{(SP,T)}$ [in 10^{-15} yr^{-1}] for the $0^+ \to 0^+$ transition of the $0\nu 2\beta$ decay.

	^{76}Ge	^{82}Se	^{100}Mo	^{130}Te	^{136}Xe
Q [MeV]	2.039	2.9955	3.035	2.5303	2.462
A_{01}	6.69	2.95×10	4.85×10	4.63×10	4.82×10
A_{02}	1.09×10	1.01×10^2	1.74×10^2	1.20×10^2	1.19×10^2
A_{03}	3.76	2.06×10	3.48×10	3.12×10	3.22×10
A_{04}	1.30	6.07	9.97	9.28	9.63
A_{05}	2.08×10^2	7.58×10^2	1.42×10^3	1.71×10^3	1.85×10^3
A_{06}	1.69×10^3	5.70×10^3	8.72×10^3	8.62×10^3	9.01×10^3
A_{07}	1.05×10^5	4.58×10^5	8.03×10^5	7.93×10^5	8.33×10^5
A_{08}	6.59×10^3	3.16×10^4	6.76×10^4	8.07×10^4	8.77×10^4
A_{09}	4.14×10^5	1.66×10^6	2.39×10^6	1.95×10^6	1.98×10^6
A_{00}^{SP}	2.55	8.82	1.44×10	1.56×10	1.65×10
A_{01}^{SP}	3.77	1.59×10	2.61×10	2.54×10	2.65×10
A_{02}^{SP}	1.18×10^{-1}	1.01	1.84	1.45	1.47
A_{03}^{SP}	1.27×10^{-3}	2.06×10^{-2}	4.16×10^{-2}	2.71×10^{-2}	2.68×10^{-2}
A_{01}^{T}	6.03×10	2.54×10^2	4.17×10^2	4.06×10^2	4.24×10^2
A_{02}^{T}	1.50×10^3	7.42×10^3	1.39×10^4	1.43×10^4	1.52×10^4
A_{03}^{T}	7.67×10^5	4.04×10^6	8.62×10^6	9.85×10^6	1.06×10^7

Neutrinoless Double Beta Decay

Table 5: The integrated kinematical factors $B_{0i}^{(SP,T)}$ [in 10^{-15} yr^{-1}] for the $0^+ \to 0^+$ transition of the $0\nu 2\beta$ decay.

	^{76}Ge	^{82}Se	^{100}Mo	^{130}Te	^{136}Xe
Q [MeV]	2.039	2.9955	3.035	2.5303	2.462
B_{01}	5.45	2.60×10	4.28×10	3.93×10	4.07×10
B_{02}	8.95	8.80×10	1.50×10^2	1.00×100	9.91×10
B_{04}	1.21	5.79	9.51	8.74	9.04
B_{05}	7.27	3.47×10	5.70×10	5.24×10	5.43×10
B_{07}	7.72×10^4	3.82×10^5	6.69×10^5	6.24×10^5	6.48×10^5
B_{08}	4.97×10^3	2.67×10^4	5.70×10^4	6.43×10^4	6.92×10^4
B_{09}	3.00×10^5	1.36×10^6	1.96×10^6	1.51×10^6	1.52×10^6
B_{01}^{SP}	2.73	1.30×10	2.14×10	1.97×10	2.04×10
B_{02}^{SP}	7.20×10^{-2}	7.26×10^{-1}	1.33	9.66×10^{-1}	9.68×10^{-1}
B_{03}^{SP}	3.71×10^{-4}	9.67×10^{-3}	1.97×10^{-2}	1.05×10^{-2}	1.01×10^{-2}
B_{01}^{T}	4.36×10	2.08×10^2	3.42×10^2	3.15×10^2	3.26×10^2
B_{02}^{T}	1.40×10^3	7.08×10^3	1.33×10^4	1.35×10^4	1.42×10^4
B_{03}^{T}	7.16×10^5	3.85×10^6	8.22×10^6	9.27×10^6	9.98×10^6

Neutrinoless Double Beta Decay

Table 6: The numerical values of C_i and D_i multiplied by 10^{15} for various nuclei in the QRPA model without p-n pairing [G. Pantis, F. Šimkovic, J. D. Vergados and A. Faessler, *Phys. Rev.* **C 53**, 695 (1996)].

Nucleus	C_1	C_0	C_{1+}	C_2	C_3	C_4	C_5
^{76}Ge	7.208	-6.680	6.191	-4.601	1520	11.54	3.152×10^5
^{82}Se	29.97	-29.50	29.03	-15.45	2297	81.56	2.627×10^5
^{100}Mo	151.3	-20.04	2.656	-46.82	2.446×10^4	245.1	6.426×10^6
^{130}Te	46.67	-46.30	45.93	-35.89	9670	115.7	2.480×10^6
^{136}Xe	45.54	-48.16	50.94	-37.88	9896	112.0	2.581×10^6
Nucleus	D_1	D_0	D_{1+}	D_2	D_3	D_4	D_5
^{76}Ge	5.872	-5.442	5.044	-0.6179	-17.31	-9.073	-2.309×10^5
^{82}Se	26.42	-26.00	25.59	3.292	-3.428	-72.72	-2.154×10^5
^{100}Mo	133.5	-17.69	2.344	26.62	-773.7	-219.1	-5.390×10^6
^{130}Te	39.62	-39.30	38.99	-5.405	-152.7	-92.26	-1.951×10^6
^{136}Xe	38.45	-40.67	43.01	-7.952	-171.7	-86.66	-2.014×10^6

Table 7: The same as in Table 6, but in the QRPA model with p-n pairing [G. Pantis, F. Šimkovic, J. D. Vergados and A. Faessler, *Phys. Rev.* **C 53**, 695 (1996)].

Nucleus	C_1	C_0	C_{1+}	C_2	C_3	C_4	C_5
^{76}Ge	3.526	-6.188	10.86	-0.745	1065	3.037	3.153×10^5
^{82}Se	59.15	-24.39	10.06	-54.58	3206	226.1	2.627×10^5
^{100}Mo	0.1805	-5.737	182.3	-1.530	860.9	37.98	6.428×10^6
^{130}Te	30.83	-44.73	64.91	-25.58	7878	80.66	2.480×10^6
^{136}Xe	42.05	-47.99	54.77	-30.60	9509	88.69	2.581×10^6
Nucleus	D_1	D_0	D_{1+}	D_2	D_3	D_4	D_5
^{76}Ge	2.873	-5.041	8.846	0.7080	-12.21	-2.634	-2.309×10^5
^{82}Se	52.13	-21.50	8.867	-11.42	-4.815	-186.5	-2.154×10^5
^{100}Mo	0.1593	-5.062	160.9	-0.5702	-26.72	-27.92	-5.389×10^6
^{130}Te	26.17	-37.97	55.09	-4.971	-124.1	-63.12	-1.951×10^6
^{136}Xe	35.50	-40.52	46.25	-4.889	-165.0	-70.31	-2.014×10^6

Neutrinoless Double Beta Decay

Using the data on various decaying nuclei not only allow us to consider in more detail the three particular cases for the parameter space studied in [A. Ali, A. V. Borisov, and D. V. Zhuridov, *Phys. Rev.* **D 76**, 093009 (2007)]:

$$
\begin{array}{ll}
\text{A)} & \epsilon_\alpha^\beta = 0,\ |\langle m\rangle| \neq 0 \quad \text{(SM plus Majorana neutrinos)},\\
\text{B)} & \epsilon_\alpha^\beta \neq 0,\ |\langle m\rangle| = 0 \quad \text{(vanishing effective Majorana mass)},\\
\text{C)} & \epsilon_\alpha^\beta \neq 0,\ |\langle m\rangle| \neq 0,\ c_i \equiv \cos\psi_i = 0,
\end{array}
$$

but also to investigate more general situation of

$$
\text{D)} \quad \epsilon_\alpha^\beta \neq 0,\ |\langle m\rangle| \neq 0,\ c_i \equiv \cos\psi_i \neq 0,
$$

where ψ_i are the relative phases for the non-SM contributions.

We analyze only the terms with $\epsilon_{V\mp A}^{V\mp A}$ as the corresponding nuclear matrix elements have been worked out in the literature [G. Pantis, F. Šimkovic, J. D. Vergados and A. Faessler, *Phys. Rev.* **C 53**, 695 (1996)].

Neutrinoless Double Beta Decay

•A. Case $\epsilon_\alpha^\beta = 0$, $|\langle m\rangle| \neq 0$

The decay half-life and the angular correlation coefficient:

$$T_{1/2} = \left(|M_{\rm GT}|^2 \mathcal{A}\right)^{-1} = \left(|M_{\rm GT}|^2 C_1 |\mu|^2\right)^{-1} = \left(|M_{\rm GT}|^2 (\chi_{\rm F} - 1)^2 A_{01} |\mu|^2\right)^{-1},$$

$$K = D_1/C_1 = B_{01}/A_{01},$$

$\mu = \langle m\rangle/m_e$, $\chi_{\rm F} = (g_V/g_A)^2\, M_{\rm F}/M_{\rm GT}$.

For this case, K is independent on a nuclear model.

Neutrinoless Double Beta Decay

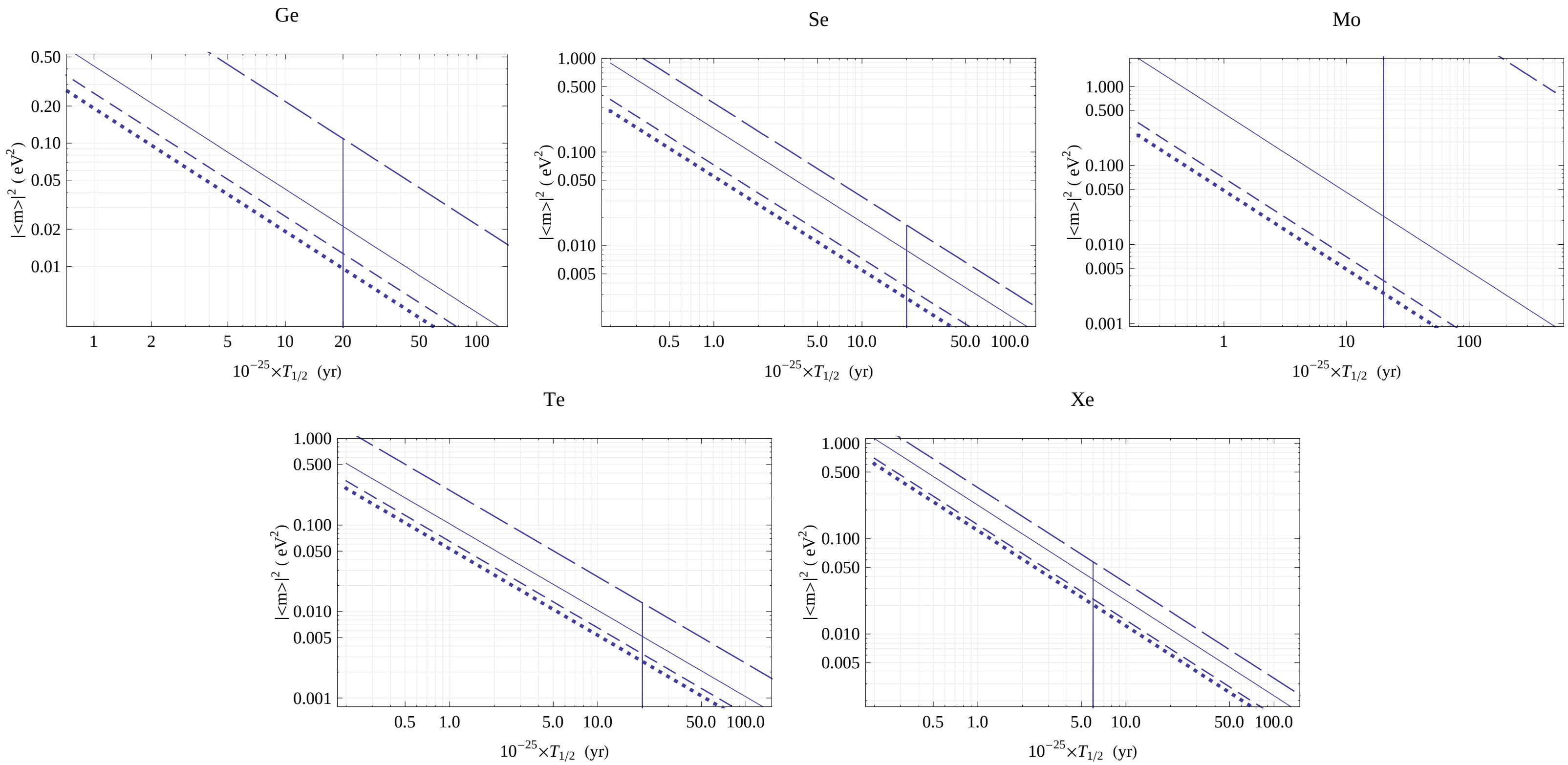


Figure 2: The squared effective Majorana mass $|\langle m\rangle|^2$ vs. $T_{1/2}$ for the QRPA model without p-n pairing (solid) and the QRPA model with p-n pairing (long-dashed) [G. Pantis at al. (1996)], $M'^{0\nu}$ in the RQRPA model (short-dashed) and $M'^{0\nu}$ in the QRPA model (dotted) [V. A. Rodin et al. (2006, 2007 (E))]. The prospective sensitivities of EXO and {GERDA-MAJORANA, SuperNEMO, MOON and COURE} to the half-lifes of ^{136}Xe and {^{76}Ge, ^{82}Se, ^{100}Mo and ^{130}Te}, respectively, are represented by the vertical lines.

Neutrinoless Double Beta Decay

Fig. 2 shows the significant nuclear model dependence of the calculated effective Majorana mass for the chosen decaying nucleus with the fixed half-life, in particular, the QRPA model without p-n pairing gives roughly 5 times more strict constraint on $|\langle m \rangle|^2$ for ^{76}Ge comparing with the QRPA model with p-n pairing.

On the other hand, Fig. 2 demonstrates the possibility to discriminate among the various nuclear models by using the results for several decaying nuclei. For 4 nuclear models, considered in Fig. 2, ^{130}Te is one of the most sensitive nuclei to the effective Majorana mass.

The differences in the half-lives and angular correlations for various nuclei in the SM extended only by Majorana neutrinos (i.e., all $\epsilon_\alpha^\beta = 0$) are described by the ratios (see the second column of Tables 8 and 9)

$$\mathcal{R}(^A\mathrm{X}) = \frac{T_{1/2}(^A\mathrm{X})}{T_{1/2}(^{76}\mathrm{Ge})}, \quad \mathcal{K}(^A\mathrm{X}) = \frac{K(^A\mathrm{X})}{K(^{76}\mathrm{Ge})}.$$

We make a comparison with ^{76}Ge as it is the best tested isotope to date.

Neutrinoless Double Beta Decay

Table 8: The ratios of the half-lives $\mathcal{R}$ and $\mathcal{R}_\alpha^\beta$ for various nuclei in QRPA without (with) p-n pairing.

Nucleus	$\mathcal{R}=\mathcal{R}_{V-A}^{V-A}$	$\mathcal{R}_{V+A}^{V-A}$	$\mathcal{R}_{V-A}^{V+A}$	$\mathcal{R}_{V+A}^{V+A}$
^{82}Se	0.4 (0.2)	0.4 (2.8)	2.1 (3.1)	0.3 (0.03)
^{100}Mo	1.1 (195.2)	52.9 (0.6)	1.1 (0.5)	1.1 (0.8)
^{130}Te	0.2 (0.1)	0.2 (0.2)	0.2 (0.1)	0.2 (0.04)
^{136}Xe	0.5 (0.2)	0.4 (0.4)	0.4 (0.2)	0.4 (0.1)

Table 9: The ratios of the angular coefficients $\mathcal{K}$ and $\mathcal{K}_\alpha^\beta$ for various nuclei in QRPA without (with) p-n pairing.

Nucleus	$\mathcal{K}=\mathcal{K}_{V\pm A}^{V-A}$	$\mathcal{K}_{V-A}^{V+A}$	$\frac{B_{09}}{A_{09}}\Big\vert_{^{A}X}\frac{A_{09}}{B_{09}}\Big\vert_{^{76}\mathrm{Ge}}$	$\mathcal{K}_{V+A}^{V+A}$	$\frac{B_{02}}{A_{02}}\Big\vert_{^{A}X}\frac{A_{02}}{B_{02}}\Big\vert_{^{76}\mathrm{Ge}}$
^{82}Se	1.08	1.12 (1.12)	1.13	1.13 (0.95)	1.06
^{100}Mo	1.08	1.15 (1.15)	1.13	1.14 (0.85)	1.05
^{130}Te	1.04	1.07 (1.07)	1.07	1.01 (0.90)	1.01
^{136}Xe	1.04	1.07 (1.07)	1.06	0.98 (0.91)	1.01

Neutrinoless Double Beta Decay

•Case B. $\epsilon_\alpha^\beta \neq 0$, $|\langle m \rangle| = 0$

The ratios for the choice of only one nonzero coefficient ϵ_α^β:

$$\mathcal{R}_\alpha^\beta({}^A\mathrm{X}) = \frac{T_{1/2}(\epsilon_\alpha^\beta, {}^A\,\mathrm{X})}{T_{1/2}(\epsilon_\alpha^\beta, {}^{76}\,\mathrm{Ge})}, \quad \mathcal{K}_\alpha^\beta({}^A\mathrm{X}) = \frac{K(\epsilon_\alpha^\beta, {}^A\,\mathrm{X})}{K(\epsilon_\alpha^\beta, {}^{76}\,\mathrm{Ge})},$$

which characterize specific alternative new physics contributions.

These ratios do not depend on ϵ_α^β directly but still depend on α and β through the different sets of C_i and D_i entering them (see Tables 8 and 9).

The entries for $\mathcal{R}_\alpha^\beta({}^A\mathrm{X})$ are dominated by the uncertainties of the nuclear models.

• In particular, the value of $\mathcal{R}_{V+A}^{V-A}$ ($\mathcal{R} = \mathcal{R}_{V-A}^{V-A}$) is accidentally enhanced by a factor of $O(100)$ for ${}^{100}\mathrm{Mo}$ comparing to other nuclei in the QRPA model without (with) p-n pairing since $\mathcal{R}_{V+A}^{V-A}({}^AX)$ ($\mathcal{R} = \mathcal{R}_{V-A}^{V-A}({}^AX)$) is proportional to the $|M_{\mathrm{GT}}|^{-2}(\chi_{\mathrm{F}}+1)^{-2}|$ ($|M_{\mathrm{GT}}|^{-2}(\chi_{\mathrm{F}}-1)^{-2}|$) combination of the matrix elements, and $|M_{\mathrm{GT}}|$ and $|\chi_{\mathrm{F}}+1|$ ($|\chi_{\mathrm{F}}-1|$) is, respectively, ~ 3 (~ 2) and ~ 5 (~ 10) times smaller for ${}^{100}\mathrm{Mo}$ comparing to other nuclei. The reason for the later is the fact that χ_F is accidentally close to -1 (1) for ^{100}Mo in contrast to the other considered nuclei. We remark that there is no similar accidental enhancement for $\mathcal{R}_{V\mp A}^{V+A}$ since the correspondent expressions are more complicate.

Neutrinoless Double Beta Decay

In contrast to $\mathcal{R}_\alpha^\beta$, the angular correlations $\mathcal{K}$ and $\mathcal{K}_{V\pm A}^{V-A}$ do not depend on the nuclear matrix elements:

$$\mathcal{K} = \mathcal{K}_{V-A}^{V-A} = \mathcal{K}_{V+A}^{V-A} = \left.\frac{B_{01}}{A_{01}}\right|_{^AX} \left.\frac{A_{01}}{B_{01}}\right|_{^{76}\mathrm{Ge}}.$$

Moreover, the coefficients $\mathcal{K}_{V\pm A}^{V+A}$ essentially do not depend on the nuclear model uncertainties:

$$\mathcal{K}_{V-A}^{V+A} \simeq \left.\frac{B_{09}}{A_{09}}\right|_{^AX} \left.\frac{A_{09}}{B_{09}}\right|_{^{76}\mathrm{Ge}}, \quad \mathcal{K}_{V+A}^{V+A} \simeq \left.\frac{B_{02}}{A_{02}}\right|_{^AX} \left.\frac{A_{02}}{B_{02}}\right|_{^{76}\mathrm{Ge}}.$$

However, the ratios of the angular correlations are not discriminating among the various underlying theories, as within the anticipated experimental uncertainty, they are all consistent with $\mathcal{K}_\alpha^\beta = 1$.

The most sensitive to the listed ratios is ^{100}Mo, except for the ratio $\mathcal{R}_{V-A}^{V+A}$ to which the most sensitive is ^{82}Se.

- From the measurements of the half-lives, the most sensitive to the effects of $\epsilon_{V+A}^{V\pm A}$ and ϵ_{V-A}^{V+A} are the pairs ^{100}Mo – ^{130}Te and ^{82}Se – ^{130}Te, correspondingly.
- From the measurements of the angular coefficients, the most sensitive to the effects of $\epsilon_{V\pm A}^{V+A}$ is the pair ^{76}Ge – ^{100}Mo.

Neutrinoless Double Beta Decay

•Case C. $\epsilon_\alpha^\beta \neq 0$, $|\langle m\rangle| \neq 0$, $c_i \equiv \cos\psi_i = 0$

For $\mu = \langle m\rangle/m_e$ and ϵ_α^β we have the equations

$$|\mu|^2 = (\lambda_1 - \lambda_2 K)/T_{1/2},$$
$$|\epsilon_\alpha^\beta|^2 = (-\lambda_3 + \lambda_4 K)/T_{1/2},$$

with the coefficients

$$\lambda_1 = \frac{D_i}{|M_{\rm GT}|^2\Delta_i}, \quad \lambda_2 = \frac{C_i}{|M_{\rm GT}|^2\Delta_i},$$
$$\lambda_3 = \frac{D_1}{|M_{\rm GT}|^2\Delta_i}, \quad \lambda_4 = \frac{C_1}{|M_{\rm GT}|^2\Delta_i},$$

$\Delta_i = C_1 D_i - D_1 C_i$. In particular, $i = 4(5)$ for ϵ_{V+A}^{V+A} (ϵ_{V-A}^{V+A}).

Neutrinoless Double Beta Decay

Table 10: The values of λ_i [in 10^{12} yr] for $\epsilon_{V+A}^{V+A} \neq 0$ for various nuclei in QRPA without p-n pairing.

Nucleus	λ_1	λ_2	λ_3	λ_4
^{76}Ge	7.942	-10.10	-5.140	-6.310
^{82}Se	3.429	-3.846	-1.246	-1.413
^{100}Mo	8.796	-9.840	-5.359	-6.072
^{130}Te	1.917	-2.403	-0.8231	-0.9698
^{136}Xe	4.112	-5.315	-1.824	-2.161

Table 11: The values of $\lambda_{1,2}$ [in 10^{12} yr] and $\lambda_{3,4}$ [in 10^{8} yr] for $\epsilon_{V-A}^{V+A} \neq 0$ for various nuclei in QRPA without p-n pairing.

Nucleus	λ_1	λ_2	λ_3	λ_4
^{76}Ge	7.657	-10.45	-1.947	-2.390
^{82}Se	3.286	-4.008	-4.030	-4.573
^{100}Mo	8.518	-10.16	-2.109	-2.390
^{130}Te	1.903	-2.419	-0.3864	-0.4553
^{136}Xe	4.130	-5.293	-0.7885	-0.9338

Large $\epsilon_{V\pm A}^{V+A}$ correspond to lower angular coefficient K, and the most sensitive to the angular correlation are ^{76}Ge, ^{82}Se and ^{100}Mo.

Neutrinoless Double Beta Decay

The left-right symmetric model $SU(2)_L \times SU(2)_R \times U(1)$

For the masses m_{W_L} and m_{W_R} of the left- and right-handed W bosons and their mixing angle ζ:

$$m_{W_R} = m_{W_L} \left(\epsilon / \left|\epsilon^{V+A}_{V+A}\right|\right)^{1/2}, \quad \zeta = -\arctan\left(\left|\epsilon^{V+A}_{V-A}\right|/\epsilon\right),$$

$m_{W_L} \simeq m_{W_1}$=80.4 GeV;

the mixing parameter

$$\epsilon = |U_{ei}V_{ei}|, \quad |V_{ei}| \sim m_D/M_R \ll 1.$$

Fig. 3 illustrates the correlation among the mass M_{W_R}, the mixing angle ζ and the half-life $T_{1/2}$ and the angular correlation K for the mixing parameter $\epsilon = |U_{ei}V_{ei}| = 10^{-6}$.

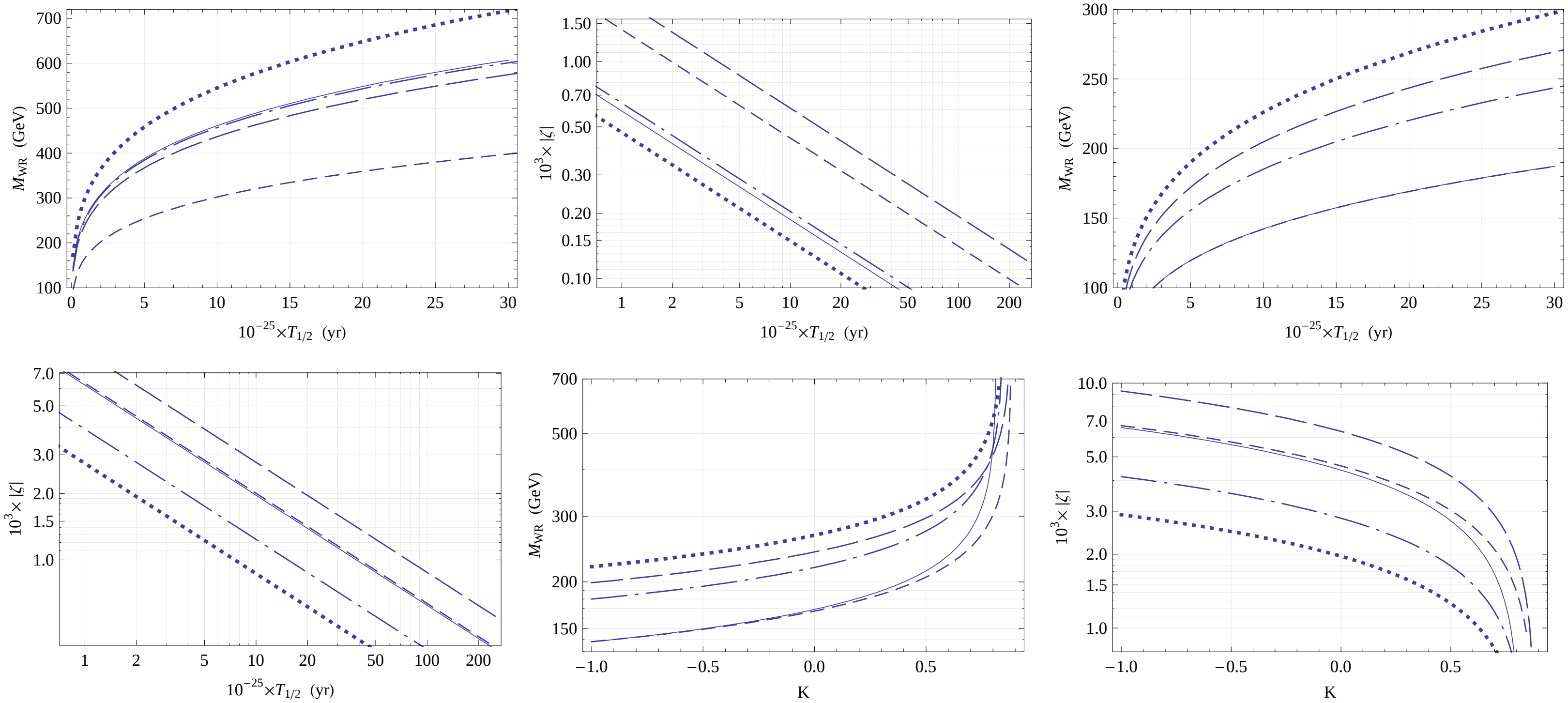


Figure 3: Correlations among: M_{W_R} (*left*), mixing angle ζ (*right*) and $T_{1/2}$ (*upper* and *middle*), K (*lower*) for the fixed $K = 0.8$ (*upper*), $K = -0.8$ (*middle*) and $T_{1/2} = 10^{26}$ yr (*lower*). Solid, long-dashed, short-dashed, dotted, and dot-dashed lines represent decays of ^{76}Ge, ^{82}Se, ^{100}Mo, ^{130}Te, and ^{136}Xe, respectively.

The higher is $T_{1/2}$ for the fixed K or the closer is K to 1 for the fixed $T_{1/2}$, the stronger is the lower bound on M_{W_R} (the upper bound on ζ).

The W_R-boson mass is stronger bounded for ^{130}Te. However the substantial effects of W_R to the $0\nu2\beta$ decays are disfavored for the half-lives below 10^{26} yr by the strong bound on M_{W_R} from the electroweak fits ($M_{W_R} > 715$ GeV [PDG-2008]).

•Case D. $\epsilon_\alpha^\beta \neq 0$, $|\langle m \rangle| \neq 0$, $c_i \neq 0$

For $\beta = S-P$, $V-A$ and T_L (except for $\alpha = T_R$, $\beta = T_L$):

$$\mathcal{A} = C_1|\mu|^2 + 4C_i|\mu||\mu_\alpha^\beta|c_{0k} + 4C_j|\mu_\alpha^\beta|^2,$$
$$\mathcal{A}K = D_1|\mu|^2 + 4D_{i'}|\mu||\mu_\alpha^\beta|c_{0k} + 4D_{j'}|\mu_\alpha^\beta|^2;$$

for $\beta = S+P$, $V+A$ and T_R (except for $\alpha = T_L$, $\beta = T_R$):

$$\mathcal{A} = C_1|\mu|^2 + C_i|\mu||\epsilon_\alpha^\beta|c_k + C_j|\epsilon_\alpha^\beta|^2,$$
$$\mathcal{A}K = D_1|\mu|^2 + D_{i'}|\mu||\epsilon_\alpha^\beta|c_k + D_{j'}|\epsilon_\alpha^\beta|^2;$$

$$\mathcal{A} = \left(|M_{\mathrm{GT}}|^2 T_{1/2}\right)^{-1}$$

.

For a nonzero ϵ_{V+A}^{V+A} (ϵ_{V-A}^{V+A}): $i = i' = 2$, $j = j' = 4$, $k = 1$ ($i = i' = 3$, $j = j' = 5$, $k = 2$).

Fig. 4 shows the correlations among $|\epsilon_{V+A}^{V+A}|$ and $|\langle m \rangle|$ for the fixed values of $c_1 = -1, 0, 1$ in QRPA model without and with p-n pairing [G. Pantis et al.].

Neutrinoless Double Beta Decay

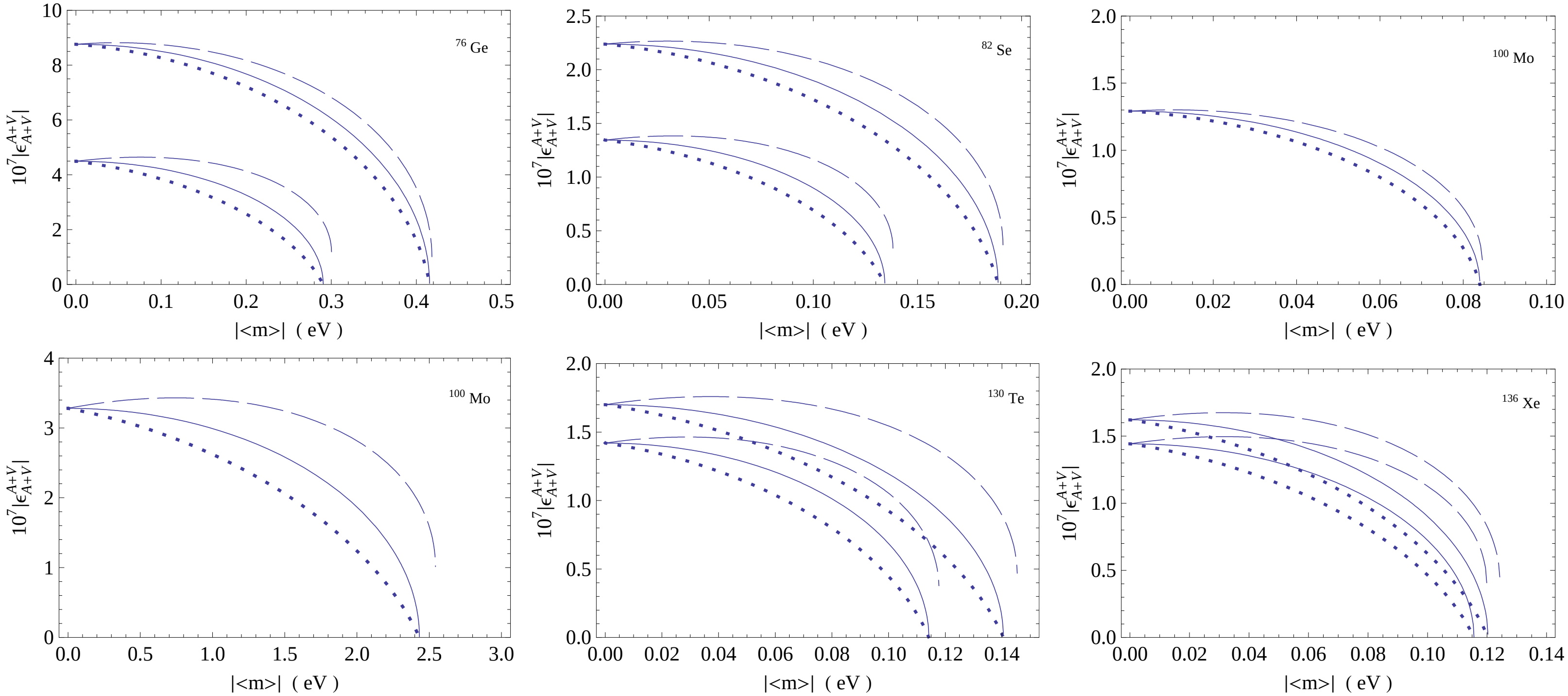


Figure 4: Correlations among $|\epsilon^{V+A}_{V+A}|$ and $|\langle m\rangle|$ for $c_1 = -1$, 0 and 1 are represented by dotted, solid, and dashed lines, respectively, for various decaying nuclei. These correlations in QRPA model without (with) p-n pairing are represented by three lower-left (upper-right) lines for each nucleus.

The variation of the QRPA model basically results in a parallel shift of the curves. The shift is maximal for ^{100}Mo and is minimal for ^{136}Xe.

Neutrinoless Double Beta Decay

It is not enough to have the system of two equations in the case of one decaying nucleus to extract the set of three variables: $\{|\mu|,\ c_{0i},\ |\mu_\alpha^\beta|\}$ or $\{|\mu|,\ c_i,\ \epsilon_\alpha^\beta\}$. To determine the needed parameters, one should have either the half-lives for three nuclei or the half lives for two nuclei with the angular coefficient K at least for one of them.

Consider the last case:

$$\begin{aligned}\mathcal{A} &= C_1|\mu|^2 + C_i|\mu||\epsilon_\alpha^\beta|c_k + C_j|\epsilon_\alpha^\beta|^2,\\ \mathcal{A}K &= D_1|\mu|^2 + D_{i'}|\mu||\epsilon_\alpha^\beta|c_k + D_{j'}|\epsilon_\alpha^\beta|^2,\\ \tilde{\mathcal{A}} &= \tilde{C}_1|\mu|^2 + \tilde{C}_i|\mu||\epsilon_\alpha^\beta|c_k + \tilde{C}_j|\epsilon_\alpha^\beta|^2.\end{aligned}$$

From this system:

$$|\mu| = \sqrt{\frac{\Delta_\mu}{\Delta}},\quad |\epsilon_\alpha^\beta| = \sqrt{\frac{\Delta_{\epsilon_\alpha^\beta}}{\Delta}},\quad c_k = \frac{\mathcal{A}\Delta - C_1\Delta_\mu - C_j\Delta_{\epsilon_\alpha^\beta}}{C_i\sqrt{\Delta_\mu\Delta_{\epsilon_\alpha^\beta}}},$$

$$\begin{aligned}&\Delta = \lambda_{11}\lambda_{22} - \lambda_{12}\lambda_{21},\quad \Delta_\mu = \eta_1\lambda_{22} - \eta_2\lambda_{12},\quad \Delta_{\epsilon_\alpha^\beta} = \eta_2\lambda_{11} - \eta_1\lambda_{21},\\ &\lambda_{11} = C_1 - C_j\tilde{C}_1/\tilde{C}_j,\quad \lambda_{12} = C_i - C_j\tilde{C}_i/\tilde{C}_j,\quad \eta_1 = \mathcal{A} - \tilde{\mathcal{A}}C_j/\tilde{C}_j,\\ &\lambda_{21} = D_1 - D_{j'}\tilde{C}_1/\tilde{C}_j,\quad \lambda_{22} = D_{i'} - D_{j'}\tilde{C}_i/\tilde{C}_j,\quad \eta_2 = \mathcal{A}K - \tilde{\mathcal{A}}D_{j'}/\tilde{C}_j.\end{aligned}$$

III. Summary

We have analyzed the half-lives and the electron angular correlations for the $0\nu2\beta$ decays of ^{76}Ge, ^{82}Se, ^{100}Mo, ^{130}Te, and ^{136}Xe in a general theoretical context.

The comparison of these characteristics for the selected decaying nuclei will help to minimize the theoretical uncertainties in the nuclear matrix elements and identify the dominant mechanism underlying these decays.

At present, no experiment is geared to measuring the angular correlations in $0\nu2\beta$ decays, as the main experimental thrust is on establishing a nonzero signal unambiguously in the first place.

We note that the running experiment NEMO3 has already measured the electron angular distributions for the two neutrino double beta decays of ^{100}Mo and ^{82}Se, and is capable of measuring these correlations in the future for the $0\nu2\beta$ decays as well, assuming that the experimental sensitivity is sufficiently good to establish these decays.

The proposed experimental facilities that can measure the electron angular correlations in the $0\nu2\beta$ decays are SuperNEMO, MOON, and EXO.