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Neutrinoless double beta decay

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Abstract

We review the potential to probe new physics with neutrinoless double beta decay $(A, Z) \rightarrow (A, Z + 2) + 2e^-$. Both the standard long-range light neutrino mechanism as well as non-standard long-range and short-range mechanisms mediated by heavy particles are discussed. We also stress aspects of the connection to lepton number violation at colliders and the implications for baryogenesis.

1. Introduction

Neutrinoless double beta decay ($0\nu\beta\beta$) experiments are not simply neutrino mass experiments, but have a much more fundamental goal, namely the quest for lepton number violation (LNV). The basic decay mode is

$$(A, Z) \rightarrow (A, Z + 2) + 2e^-, \tag{1}$$

i.e. the transition of a nucleus with mass and atomic numbers A and Z to a nucleus with A and $Z + 2$ under emission of two electrons only. This process obviously violates electron lepton number L_e by two units. At present this endeavor is entering a particular exciting stage, with numerous experiments operating or being under development, using different isotopes and experimental techniques (see table 1). The previous best limit on the decay, set by the Heidelberg–Moscow experiment in 2001 [1], has finally been improved from 2012 on [2–5], and the limits will be further and further increased, with the potential of discovery always present. A large number of reviews has been written in the last few years [6–16], adding to the important earlier ones [17–21], and emphasizing the importance of the decay and the strong interest of various communities.

In this review we discuss the main physics potential and the conceptual implications that neutrinoless double beta decay brings along. We consider not only the standard three neutrino paradigm, but also different frameworks, including situations associated with heavy particle exchange, so-called short range mechanisms. Tests of such mechanisms are possible for instance in collider experiments. In turn, observation of LNV, either in $0\nu\beta\beta$ decay or at colliders, has important ramifications for baryogenesis, which we will outline as well.

Why is it important to look for LNV? One could give several reasons, for instance:

- lepton number (as well as baryon number) is only an accidentally conserved global symmetry in the standard model (SM)[4], and its conservation in extended theories seems very unlikely. Indeed, the lowest higher dimensional operator one can write down, $\mathcal{L} = 1/\Lambda\,(\Phi L)\,(\Phi L)$, immediately violates lepton number and generates neutrino mass. In this language, neutrino mass and LNV are the leading order new physics effects that one might expect to appear, as all other operators are suppressed by additional powers of the cut-off scale Λ. As neutrino mass has been observed in the form of neutrino oscillations, hopes are high that LNV is present as well;

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[4] Though not really connected to double beta decay or Majorana neutrinos that require LNV by two units, one should note that even within the SM lepton number is actually not conserved: chiral anomalies related to instanton tunneling break global lepton and baryon number by three units each.

Table 1. Overview of present and future $0\nu\beta\beta$ decay experiments, their energy resolution and sensitivity to event topology (i.e. the individual energy of the electrons and/or their angular correlation, useful to distinguish mechanisms). Timescales, references and details can be found in [7, 14, 16].

Name	Isotope	Source = detector			Source ≠ detector
		ΔE high	ΔE low	Topology	Topology
AMoRE	^{100}Mo	✓	—	—	—
CANDLES	^{48}Ca	—	✓	—	—
COBRA	^{116}Cd (and ^{130}Te)	—	—	✓	—
CUORE	^{130}Te	✓	—	—	—
CUPID	^{82}Se / ^{100}Mo / ^{116}Cd / ^{130}Te	✓	—	—	—
DCBA/MTD	^{82}Se / ^{150}Nd	—	—	—	✓
EXO	^{136}Xe	—	—	✓	—
GERDA	^{76}Ge	✓	—	—	—
KamLAND-Zen	^{136}Xe	—	✓	—	—
LUCIFER	^{82}Se / ^{100}Mo / ^{130}Te	✓	—	—	—
LUMINEU	^{100}Mo	✓	—	—	—
MAJORANA	^{76}Ge	✓	—	—	—
MOON	^{82}Se / ^{100}Mo / ^{150}Nd	—	—	—	✓
NEXT	^{136}Xe	—	—	✓	—
SNO+	^{130}Te	—	✓	—	—
SuperNEMO	^{82}Se / ^{150}Nd	—	—	—	✓
XMASS	^{136}Xe	—	✓	—	—

- the Universe contains more matter than antimatter. In order to generate this baryon asymmetry of the Universe, baryon number conservation has to be violated. Unless nature treats baryon and lepton number in a completely different manner, also LNV can be expected;

- in grand unified theories lepton and baryon number are often connected, based on the fact that their difference can be gauged in an anomaly free way when right-handed neutrinos are introduced. Thus baryon number violation typically implies LNV. Moreover, GUTs usually implement a seesaw mechanism and thus Majorana neutrinos, leading eventually to $0\nu\beta\beta$ decay;

- almost all mechanisms that generate and suppress neutrino masses result in Majorana neutrinos and thus eventually induce $0\nu\beta\beta$ decay;

- all theories beyond the SM that violate lepton number by one or two units lead to neutrinoless double beta decay. Those include supersymmetric theories with *R*-parity violation, left–right symmetry theories, models with spontaneously broken lepton number, etc[5];

- in general, global symmetries are not expected to be conserved in quantum gravity theories. One could thus gauge lepton number, and in order to avoid long range forces one would need to break the gauge symmetry, leading again typically to LNV.

All in all, lepton number is not expected to be conserved, and the observation of LNV would be as important as baryon number violation, e.g. proton decay. The decay width of double beta decay for a single operator inducing the decay can always be written as

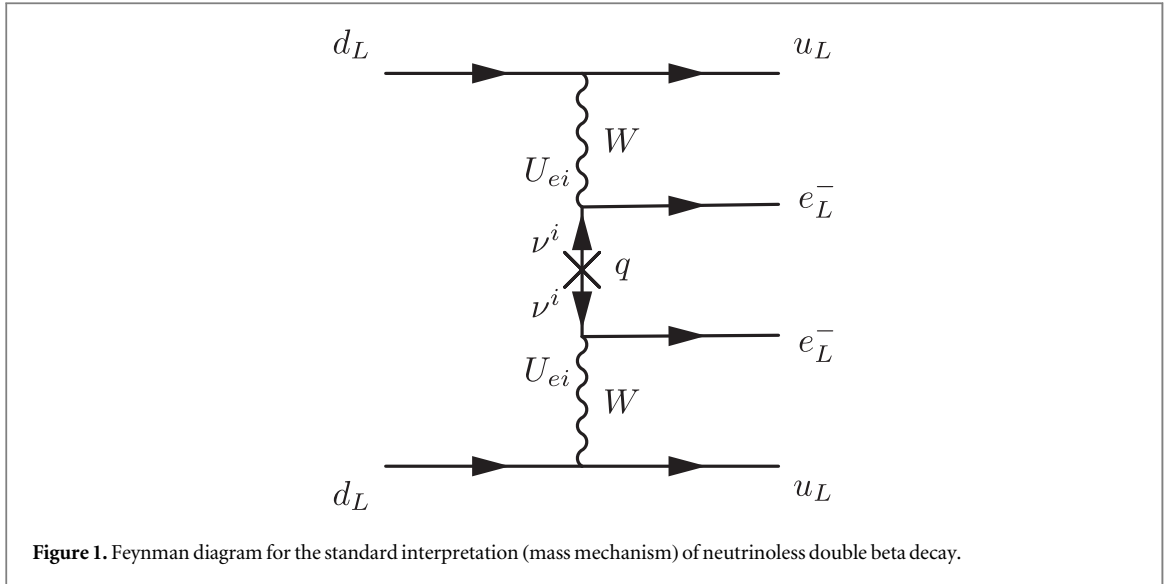$$\Gamma = G(Q, Z)|\mathcal{M}\,\varepsilon|^2, \tag{2}$$

where $G(Q, Z)$ is a calculable phase space factor typically scaling with the endpoint energy as Q^5 and $\mathcal{M}$ is the nuclear matrix element, which is notoriously difficult to calculate. The particle physics parameter ε, which depends on particle masses, mixing parameters etc, is most important from the point of view of this review. Note that more than one mechanism can contribute, hence the amplitude of the decay can actually be

$$\mathcal{A} = \sum_x \mathcal{M}_x\,\varepsilon_x, \tag{3}$$

i.e. a sum over different mechanisms, which can potentially interfere with each other.

The review is organized as follows: in section 2 we summarize double beta decay mediated by light massive Majorana neutrinos while section 3 deals with alternative and short-range mechanisms, including potential

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[5] We note here that if lepton number is violated not by two units but by three or more, there will be no neutrinoless double beta decay, but rather processes with $\Delta L = 3, 4, \ldots$ One explicit example is 'neutrinoless quadruple beta decay' presented in [22].

Figure 1. Feynman diagram for the standard interpretation (mass mechanism) of neutrinoless double beta decay.

tests. The connection between $0\nu\beta\beta$ decay, LNV at colliders and baryogenesis is discussed in section 4, before we conclude[6] in section 5.

2. Neutrinoless double beta decay and neutrino masses

We begin with the arguably best motivated possibility for the decay, the 'standard interpretation' or 'mass mechanism', namely that the light massive neutrinos that we observe to oscillate in terrestrial experiments mediate double beta decay. In this case, searches for the process are searches for neutrino mass, complementing the other approaches to determine neutrino masses. Those approaches include direct searches in classical Kurie-plot experiments like the upcoming KATRIN [23], Project 8 [24], ECHo [25] or MARE [26] experiments, and cosmological observations, see [27] for a review in this Focus Issue. Cosmology probes the sum of neutrino masses

$$\Sigma = \sum m_i, \tag{4}$$

Kurie-plot experiments test the incoherent sum

$$m_\beta = \sqrt{\sum \left| U_{ei} \right|^2 m_i^2} \, , \tag{5}$$

whereas neutrinoless double beta decay in the standard interpretation tests the quantity (see figure 1)

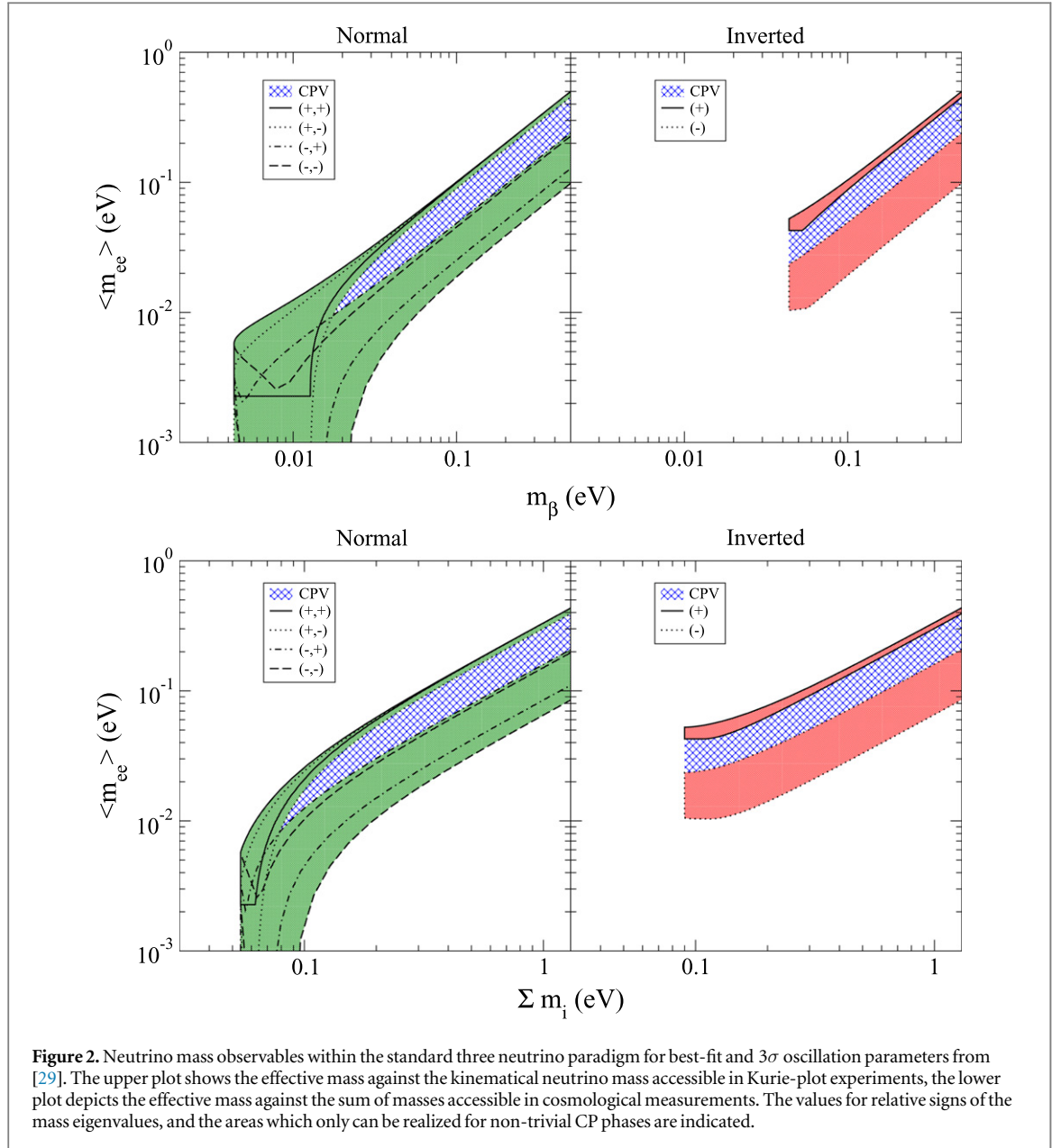$$\left| m_{ee} \right| = \left| \sum U_{ei}^2 m_i \right|, \tag{6}$$

which is usually called the effective mass and coincides with the *ee* element of the neutrino mass matrix in flavor space.

Here m_i are the neutrino masses, and U_{ei} are elements of the leptonic mixing, or PMNS, matrix that is usually parametrized as

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}\,s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}\,s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} P, \tag{7}$$

where $s_{ij} = \sin\theta_{ij}$, $c_{ij} = \cos\theta_{ij}$ and δ is the 'Dirac phase' responsible for CP violation in neutrino oscillation experiments. The diagonal phase matrix $P = \mathrm{diag}(1, e^{i\alpha}, e^{i(\beta+\delta)})$ contains the two Majorana phases α and β, which are associated with the Majorana nature of neutrinos and thus only show up in lepton number violating processes (a review on properties of Majorana particles can be found in [28]). For three neutrinos we have therefore 9 physical parameters, three masses $m_{1,2,3}$, three mixing angles θ_{12}, θ_{13}, θ_{23} and three phases δ, α, β. The effective mass depends thus on 7 out of those 9 physical neutrino parameters:

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[6] Topics that are not covered in this review are the experimental and nuclear physics aspects, where the interested reader should consult e.g. the review articles [7, 14, 16] and [13], respectively.
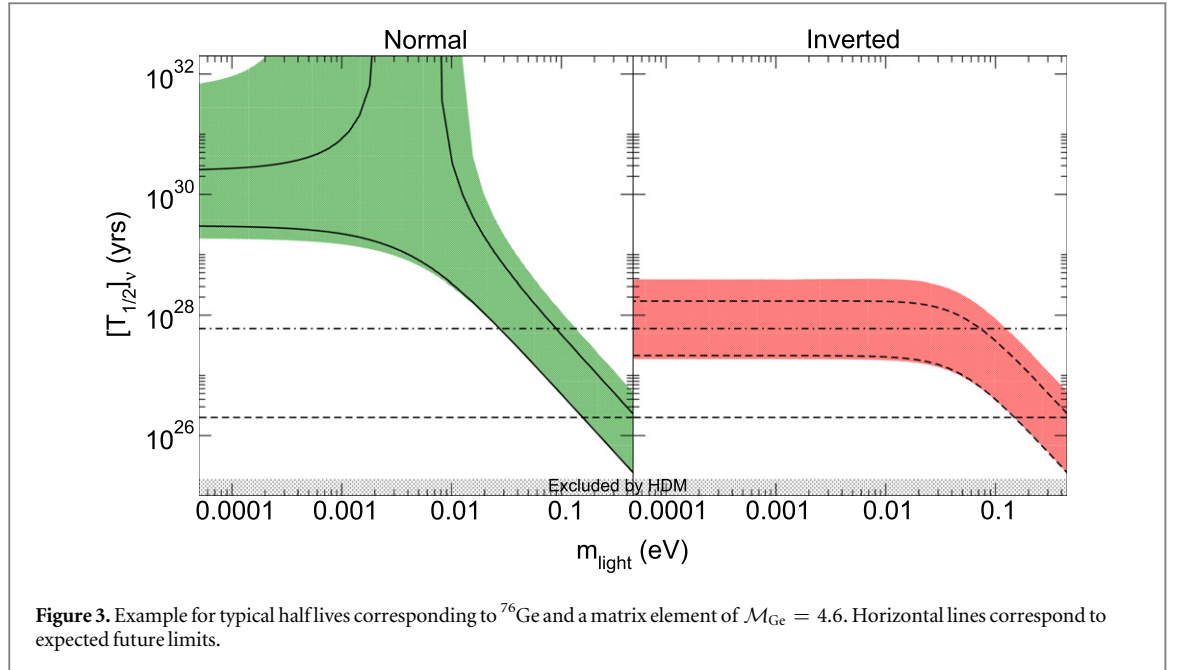
Figure 2. Neutrino mass observables within the standard three neutrino paradigm for best-fit and 3σ oscillation parameters from [29]. The upper plot shows the effective mass against the kinematical neutrino mass accessible in Kurie-plot experiments, the lower plot depicts the effective mass against the sum of masses accessible in cosmological measurements. The values for relative signs of the mass eigenvalues, and the areas which only can be realized for non-trivial CP phases are indicated.

$$\left| m_{ee} \right| = f\left(\theta_{12}, \theta_{13}, \alpha, \beta, m_1, m_2, m_3 \right). \tag{8}$$

Of these seven parameters, we currently do not know the phases and the lightest mass, where in addition the mass ordering is unknown, i.e. it could be either $m_3 > m_2 > m_1$ (normal ordering) or $m_2 > m_1 > m_3$ (inverted ordering). Global fits of all available neutrino data can be found in [29–31]. One can then use equations (4)–(6) to plot the three neutrino mass observables against each other [32], see figure 2, and interpret potential current and future experimental results [33].

For instance, in case one finds positive results for m_β and $|m_{ee}|$ in any of the green or red areas in the upper plot of figure 2, then this would be a convincing confirmation of the paradigm that there are three massive Majorana neutrinos mixing among each other. Even more spectacular would be if inconsistencies arise, e.g. a measurement of the effective mass that is incompatible with limits from KATRIN or cosmology. This would imply that something in our interpretation of double beta decay goes amiss, i.e. that another mechanism causes the decay. Therefore, the complementarity of the various approaches to determine the neutrino mass offers exciting possibilities, since different assumptions enter their interpretation. KATRIN-like experiments are essentially model-independent, as only bizarre things like tachyonic neutrinos could spoil the results, and moreover the interpretation is 'clean' as beta decay is theoretically well under control. However, in terms of numbers the limits are and will be the weakest, and further improvement beyond 0.1 eV seems impossible. Cosmology yields the best limits in terms of numbers, and can even contribute to the question of mass ordering. However, it suffers from difficult systematics and relies on model input, e.g. departures from simple ΛCDM

Table 2. Summary of the main approaches to neutrino mass.

Method	Observable	Now (eV)	Near (eV)	Far (eV)	Pro	Con
Kurie	$\sqrt{\sum |U_{ei}|^2\, m_i^2}$	2.3	0.2	0.1	Model-indep.; theo.clean	Final?; worst
Cosmo.	$\sum m_i$	0.7	0.3	0.05	Best; NH/IH	Systemat.; model-dep.
$0\nu\beta\beta$	$|\sum U_{ei}^2 m_i|$	0.3	0.1	0.05	Fundament.; NH/IH	Model-dep.; theo.dirty

Figure 3. Example for typical half lives corresponding to ^{76}Ge and a matrix element of $\mathcal{M}_{\mathrm{Ge}} = 4.6$. Horizontal lines correspond to expected future limits.

models can weaken limits considerably. Double beta decay is the most fundamental approach as it is connected to LNV, and can even say something about the mass ordering (see below). However, it is very model-dependent as many mechanisms apart from the standard neutrino mass mechanism can mediate the decay. Furthermore, the process is theoretically 'dirty', as nuclear matrix element introduce a sizable uncertainty. The pros and cons of the different approaches and their current as well as near and far future limits are summarized in table 2.

It is important to note that for the normal mass ordering the effective mass can vanish, whereas for the inverted ordering the effective mass cannot vanish [34]. Hence, the lifetime in this latter case is necessarily finite, though of course an experimental challenge. The lower limit is given by

$$\left| m_{ee} \right|_{\min}^{\mathrm{IH}} = \sqrt{\Delta m_{31}^2}\ c_{13}^2\ (1 - 2\sin^2\theta_{12}) \simeq (0.01\ldots0.02)\ \mathrm{eV}, \tag{9}$$

corresponding to half-lives around 10^{27} yrs, see figure 3. This minimal value depends rather strongly on the solar neutrino mixing angle θ_{12}. Hence, a more precise determination of θ_{12} in future oscillation experiments would be rather welcome [35]. Within the well-motivated three Majorana neutrino paradigm the upper and lower value of the effective mass in the inverted ordering are the natural medium-term goal for neutrinoless double beta decay searches. In case the mass ordering turns out to be normal, this motivation is lost. However, the value of neutrino mass remains unknown, and consistency checks with cosmological or Kurie-plot limits are necessary. Moreover, as argued in the introduction, the highly important search for LNV needs to be pursued further.

What is the current limit on the effective mass? To answer this question, a comparison of different isotopes and matrix elements is necessary. One of the most competitive lifetime limits is set by GERDA [4], $T_{1/2}^{\mathrm{Ge}} > 2.1 \times 10^{25}$ yrs, or, combined with earlier Germanium experiments [1, 36], $T_{1/2}^{\mathrm{Ge}} > 3.0 \times 10^{25}$ yrs. A similarly strong limit is obtained by the KamLAND-Zen experiment [3], namely $T_{1/2}^{\mathrm{Xe}} > 2.6 \times 10^{25}$ yrs. Using equation (2), one finds that experiments using ^{136}Xe give a better limit than experiments with ^{76}Ge if their lifetime limit fulfills the condition:

Table 3. Limits on the effective mass m_{ee} (in eV) from Germanium and Xenon experiments and different matrix element calculations. The calculations listed in bold face yield a better limit for ^{136}Xe, the ones in italic give a better limit for ^{76}Ge. Adapted from [39].

NME	^{76}Ge		^{136}Xe
	GERDA	Combined	KamLAND-Zen
EDF(U) [41]	**0.32**	**0.27**	**0.13**
ISM(U) [42]	**0.52**	**0.44**	**0.24**
IBM-2 [43]	**0.27**	**0.23**	**0.16**
pnQRPA(U) [44]	**0.28**	**0.24**	**0.17**
SRQRPA-B [45]	**0.25**	**0.21**	**0.15**
SRQRPA-A [45]	**0.31**	**0.26**	**0.23**
QRPA-A [46]	*0.28*	*0.24*	*0.25*
SkM-HFB-QRPA [47]	*0.29*	*0.24*	*0.28*

$$T_{1/2}^{\text{Xe}} > T_{1/2}^{\text{Ge}} \, \frac{G_{\text{Ge}}}{G_{\text{Xe}}} \left| \frac{\mathcal{M}_{\text{Ge}}}{\mathcal{M}_{\text{Xe}}} \right|^2. \tag{10}$$

Using the phase space factors of [37, 38], and the matrix elements of various groups, the limits on the effective mass in table 3 are obtained, adapted from [39]. Some matrix element approaches have a better limit from Germanium, others from Xenon. Taking correctly the conservative values, both isotopes give essentially the same limit of[7]

$$\left| m_{ee} \right| \lesssim 0.3 \text{ eV.} \tag{11}$$

Future improvement of this limit goes with the square root of lifetime limits.

So far the effective mass has simply been used as a phenomenological parameter. Of course, in case one has a model at hand, one can predict $|m_{ee}|$ to some extent. One example are popular flavor symmetry models to explain the peculiar features of lepton mixing [48, 49]. While the neutrino mass itself cannot be predicted in this framework, relations between neutrino masses are possible to predict, so-called neutrino mass sum-rules such as $\tilde{m}_1 + \tilde{m}_2 = \tilde{m}_3$. Here the masses are understood to be complex, i.e. including the Majorana phases. These relations exclude some possible combinations of masses and phases, and thus only certain areas in parameter space are possible, which allows to rule out certain models. Many sum-rule examples have been discussed in the literature [50–53]. Even more predictive are some grand unified theories, where the Yukawa matrices of all fermions are related and fitting the constrained matrices to the observed mass and mixing parameters allows to predict unknown parameters such as $|m_{ee}|$, see [54].

While the three neutrino paradigm is very attractive and robust, there are longstanding hints that light sterile neutrinos with mass around an eV and mixing around 10% exist, see [55] for a review of the various hints and ongoing as well as future tests. Such a fourth neutrino would modify all neutrino mass observables, in particular the effective mass:

$$\left| m_{ee} \right| = \left| \underbrace{\left| U_{e1} \right|^2 m_1 + \left| U_{e2} \right|^2 m_2 \, e^{2i\alpha} + \left| U_{e3}^2 \right| m_3 \, e^{2i\beta}}_{m_{ee}^{\text{act}}} + \underbrace{\left| U_{e4} \right|^2 m_4 \, e^{2i\gamma}}_{m_{ee}^{\text{st}}} \right|, \tag{12}$$
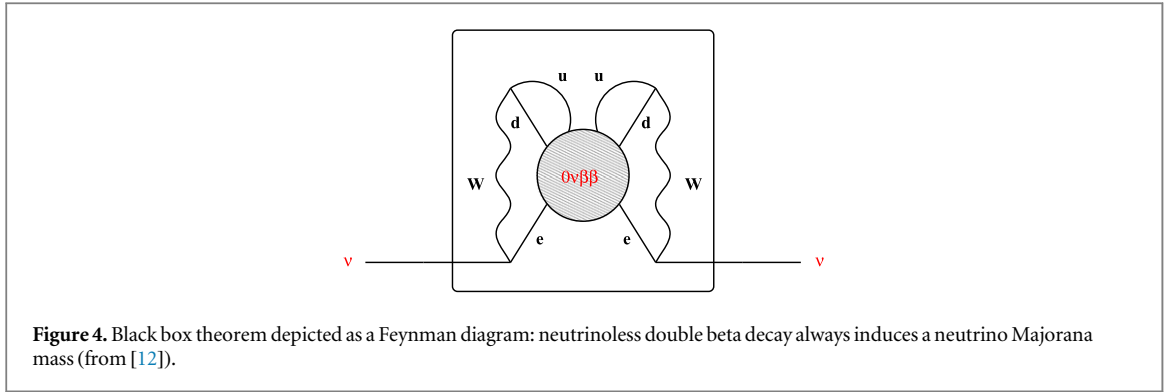
where γ is an additional Majorana phase and m_{ee}^{act} the three neutrino contribution discussed so far. The sterile contribution $|m_{ee}|^{\text{st}}$ to $0\nu\beta\beta$ (assuming a $1+3$ scenario) generates typical values of the same order as m_{ee}^{act} for the inverted ordering:

$$\left| m_{ee} \right|^{\text{st}} \simeq \sqrt{\Delta m_{\text{st}}^2} \, \left| U_{e4} \right|^2 \begin{cases} \gg \left. \left| m_{ee} \right|^{\text{act}} \right|_{\text{NH}} \\ \simeq \left. \left| m_{ee} \right|^{\text{act}} \right|_{\text{IH}} \end{cases}. \tag{13}$$

Thus, in contrast to the three-generation case, for a normal mass ordering of the active neutrinos the effective mass cannot vanish anymore, whereas for an inverted ordering of the active neutrinos the effective mass can vanish now [56–59]. The phenomenology has completely turned around! This demonstrates that any physics output of neutrinoless double beta decay depends dramatically on the assumptions.

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[7] See also [40] for an approach to combine different experiments in a statistical manner.

Figure 4. Black box theorem depicted as a Feynman diagram: neutrinoless double beta decay always induces a neutrino Majorana mass (from [12]).

We are thus naturally lead to discuss alternative mechanisms of double beta decay, to be addressed in the following section.

3. Neutrinoless double beta decay and short-range mechanisms

Apart from the standard interpretation where a massive Majorana neutrino is being exchanged between SM $V - A$ vertices, in principle any operator converting two d-quarks into two u quarks, two electrons and nothing else and thereby violating lepton number by two units will trigger the decay. This does not mean, however, that neutrinoless double beta decay and the question whether the neutrino possesses a Majorana mass are totally decoupled: the observation of neutrinoless double beta decay demonstrates that lepton number is violated by two units. Such LNV implies that neutrinos have to be Majorana particles. This has been proven by the so-called black box theorem [60–64] which states that the $0\nu\beta\beta$ diagram can always be inserted in an SM loop diagram giving rise to radiatively generated neutrino masses (see figure 4): thus if neutrinoless double beta decay is observed, a Majorana neutrino mass term is generated at four-loop order, even if the underlying particle physics model does not contain a tree-level neutrino mass. Of course this contribution to the neutrino mass is rather small [65], namely of order $G_F^4/(16\pi^2)^4 m_{u,d,e}^5 \sim 10^{-25}$ eV, and thus clearly neither the dominant contribution to neutrinoless double beta decay nor to neutrino mass itself. Note that this four-loop contribution is only the minimal, guaranteed connection between neutrino mass and double beta decay arising in any scenario with $\Delta L_e = 2$ LNV. Explicit models leading to $0\nu\beta\beta$ can generate neutrino mass at tree, 1-, 2- or 3-loop level. Depending on the model, the neutrino masses generated in this way can lead to a comparable, sub-dominant or dominant neutrino contribution to the decay, and/or to a main, sub-leading or negligible contribution to neutrino mass. For a comparative analysis of all scalar-mediated models based on the SM gauge group see [66].

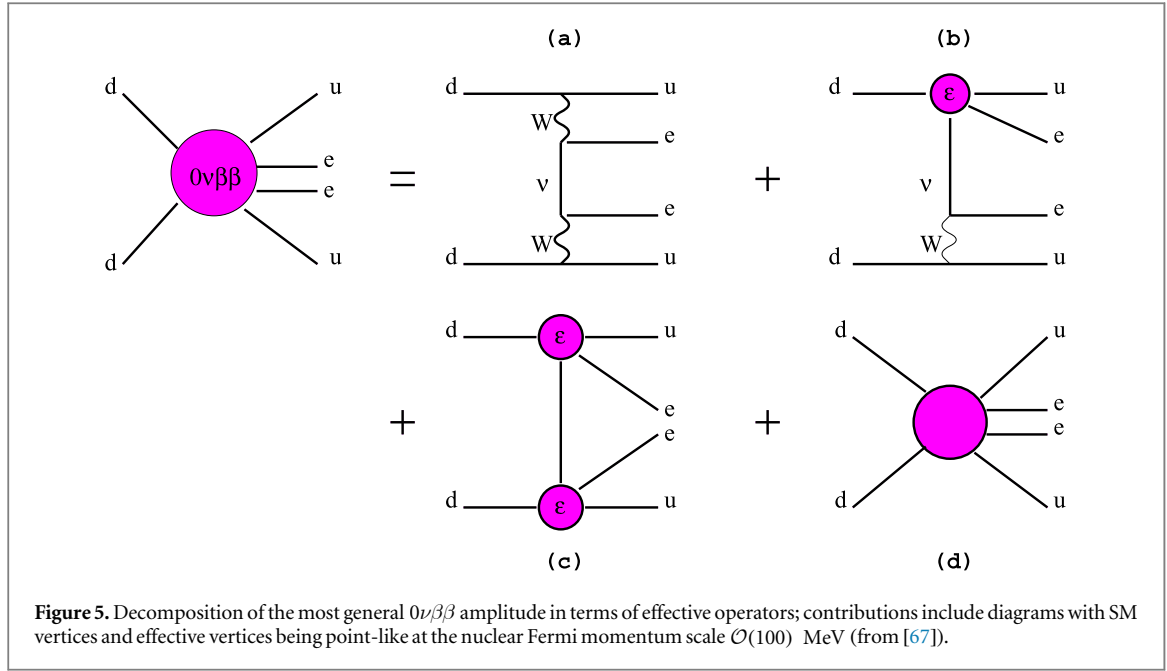The most general decay rate contains all combinations of leptonic and hadronic currents induced by the operators

$$\mathcal{O}_{V\mp A} = \gamma^\mu(1 \mp \gamma_5), \quad \mathcal{O}_{S\mp P} = (1 \mp \gamma_5), \quad \mathcal{O}_{T_{L/R}} = \frac{i}{2}[\gamma_\mu, \gamma_\nu](1 \mp \gamma_5), \tag{14}$$

allowed by Lorentz invariance. The most general operator inducing the decay can be parametrized in terms of effective couplings ε parametrizing interactions which appear point-like at the nuclear Fermi momentum scale (the inverse size of the nucleon) $\mathcal{O}(100)$ MeV. Figure 5 displays the decomposition of the general decay amplitude into the standard interpretation (contribution (a) with a light Majorana neutrino being exchanged between two SM weak interaction vertices), contribution (b) with a light Majorana neutrino being exchanged between one effective operator vertex and an SM weak interaction vertex, contribution (c), which contains two non-SM vertices and can be neglected when compared to contribution (b), and contribution (d) with a single point-like dimension nine operator [67, 68].

We can estimate the energy scale of short-range diagrams which can lead to comparable double beta decay lifetimes compared with the standard interpretation. The standard diagram discussed in section 2 has an amplitude of order $G_F^2 |m_{ee}|/q^2$. If the decay is mediated by particles heavier than the characteristic momentum scale of $q \simeq 100$ MeV, then the corresponding amplitude is c/M^5, where M is the mass of those particles and c a combination of flavor and possible gauge coupling parameters. Hence, for c of order one and M of order TeV this amplitude equals the current limit on the standard amplitude (ignoring here a small suppression of the nuclear matrix elements for short-range diagrams):

$$T_{1/2}^{0\nu\beta\beta}(m_\nu = 1 \text{ eV}) \simeq T_{1/2}^{0\nu\beta\beta}(M = 1 \text{ TeV}). \tag{15}$$

In the case where various mass scales enter the decay amplitude the sensitivity of $0\nu\beta\beta$ decay can be significantly enhanced, for example the bound of the standard interpretation extrapolated into the heavy mass regime

Figure 5. Decomposition of the most general $0\nu\beta\beta$ amplitude in terms of effective operators; contributions include diagrams with SM vertices and effective vertices being point-like at the nuclear Fermi momentum scale $\mathcal{O}(100)$ MeV (from [67]).

translates into a lower bound of 10^7 GeV on the effective mass of a super-massive neutrino eigenstate mixing with the ordinary SU(2) doublet neutrinos. We thus can test short-range diagrams for double beta decay with the LHC or lepton flavor violation experiments, which are also sensitive to the TeV scale. For the possibility to study the inverse neutrinoless double beta decay at linear colliders, see e.g. [69].

There exist various UV complete realizations for the different contributions, for example leptoquark and *R*-parity violating SUSY accompanied decay modes (contribution (b)) or short-range decay modes where only SUSY particles or heavy neutrinos and gauge bosons in left–right-symmetric models are exchanged between the decaying nucleons (contribution (d)). Present experiments have a sensitivity to the effective couplings of

$$\varepsilon < \text{few} \cdot (10^{-7}–10^{-10}). \tag{16}$$

For a more detailed, recent overview on this approach to double beta decay see [12].

As has been pointed out above any observation of the short-range contribution d) will typically involve TeV scale particles and thus may be probed in the search for LNV interactions at the LHC. To discuss such a phenomenon of '$0\nu\beta\beta$-LHC complementarity' though it is necessary to distinguish between the various mechanisms which may be responsible for the decay. While this is a difficult task there exist several ideas which at least in principle and for some of the mechanisms may allow for an experimental discrimination of $0\nu\beta\beta$ decay mechanisms. These include the observation of neutrinoless double beta decay in multiple isotopes [45, 70–72], or measuring the decay distribution, for example in the SuperNEMO experiment [72]. In principle one could distinguish mechanisms also by comparing double beta decay rates to processes such as double positron decay or double electron capture, whose observation is however unlikely as their rates are heavily suppressed [74]. Another possibility exsits at at the LHC itself is to identify the invariant mass peaks of particles produced resonantly in the intermediate state or to analyze the charge asymmetry between final states involving particles and/or anti-particles [75, 76].
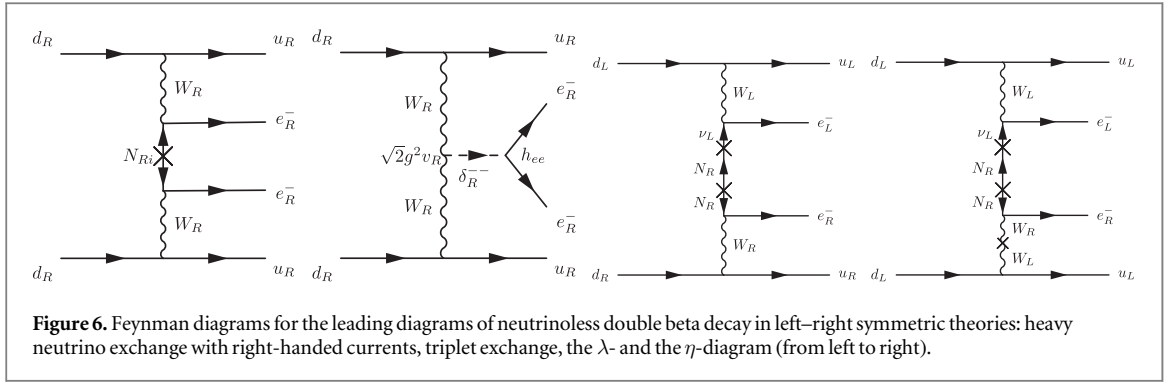
3.1. Left–right symmetry

In this section we will discuss various contributions to double beta decay in left–right symmetric models which embed the SM gauge group into an $\mathrm{SU(3)}_c \otimes \mathrm{SU(2)}_L \otimes \mathrm{SU(2)}_R \otimes \mathrm{U(1)}_{B-L}$. Such theories necessarily predict right-handed neutrinos which are included in an $\mathrm{SU(2)}_R$ doublet as a consequence of the left–right symmetry. The extended Higgs sector contains a bidoublet ϕ and two triplets Δ_L and Δ_R. The VEV v_R of the neutral component of Δ_R breaks $\mathrm{SU(2)}_R \otimes \mathrm{U(1)}_{B-L}$ to $\mathrm{U(1)}_Y$ and generates masses for the right-handed W_R and Z_R gauge bosons, and the heavy neutrinos. To make these new gauge bosons and interactions compatible with experimental constraints, they have to be sufficiently heavy and weak, respectively, resulting from a sufficiently large vev v_R.

Neutrino masses are then generated within a type-I+II seesaw,

$$m_\nu = m_L - m_D M_R^{-1} m_D^T, \tag{17}$$

Figure 6. Feynman diagrams for the leading diagrams of neutrinoless double beta decay in left–right symmetric theories: heavy neutrino exchange with right-handed currents, triplet exchange, the λ- and the η-diagram (from left to right).

where $m_{\mathrm{L}} = f v_{\mathrm{L}}$ and $M_{\mathrm{R}} = f v_{\mathrm{R}}$ are the VEVs of the triplets. Within left–right symmetric models several diagrams mediating double beta decay exist, see figure 6.

The right-handed neutrinos and W_{R} bosons can mediate the right-handed analogue of the standard mechanism discussed above [18, 77–79]. As the particles exchanged are much heavier than the nuclear Fermi momentum, this is a realization of the short-range operator. The now heavy neutrino mass will appear in the denominator of the amplitude instead of the numerator. The effective coupling is denoted ε_3^{RRz}. Assuming manifest left–right symmetry, i.e. identical gauge couplings, in terms of the left–right-symmetric model parameters it is given by

$$\varepsilon_3^{RRz} = \sum_{i=1}^{3} V_{ei}^2 \frac{m_p}{m_{N_i}} \frac{m_{W_{\mathrm{L}}}^4}{m_{W_{\mathrm{R}}}^4}, \tag{18}$$

where V denotes the matrix describing the mixing among the heavy right-handed neutrinos. Searches for $0\nu\beta\beta$ yield the limit $|\varepsilon_3^{RRz}| \lesssim 1 \times 10^{-8}$.

Moreover, because of the presence of right-handed currents, the exchange of light neutrinos does not necessarily require a chirality violating mass insertion. The coupling parameters of the corresponding effective long-range operators can be written as

$$\varepsilon_{V+A}^{V+A} = \sum_{i=1}^{3} U_{ei} S_{ei} \frac{m_{W_{\mathrm{L}}}^2}{m_{W_{\mathrm{R}}}^2}, \qquad \varepsilon_{V-A}^{V+A} = \sum_{i=1}^{3} U_{ei} S_{ei} \tan\zeta, \tag{19}$$

with the current experimental limits $|\varepsilon_{V+A}^{V+A}| \lesssim 5 \times 10^{-7}$ and $|\varepsilon_{V-A}^{V+A}| \lesssim 3 \times 10^{-9}$, respectively, and where S describes the mixing between left- and right-handed neutrinos. The diagram governed by ε_{V+A}^{V+A} is often called the λ-diagram, the one governed by ε_{V-A}^{V+A} the η-diagram. While the mixing S is small in the simplest seesaw scenarios, one can easily arrange for large left–right (or equivalently light-heavy) mixing. In this case both diagrams can be expected to dominate over the heavy neutrino exchange diagram with right-handed currents [80, 81]. Analyzes of the type-I seesaw mechanism with sizable light-heavy mixing can be found in [82, 83].
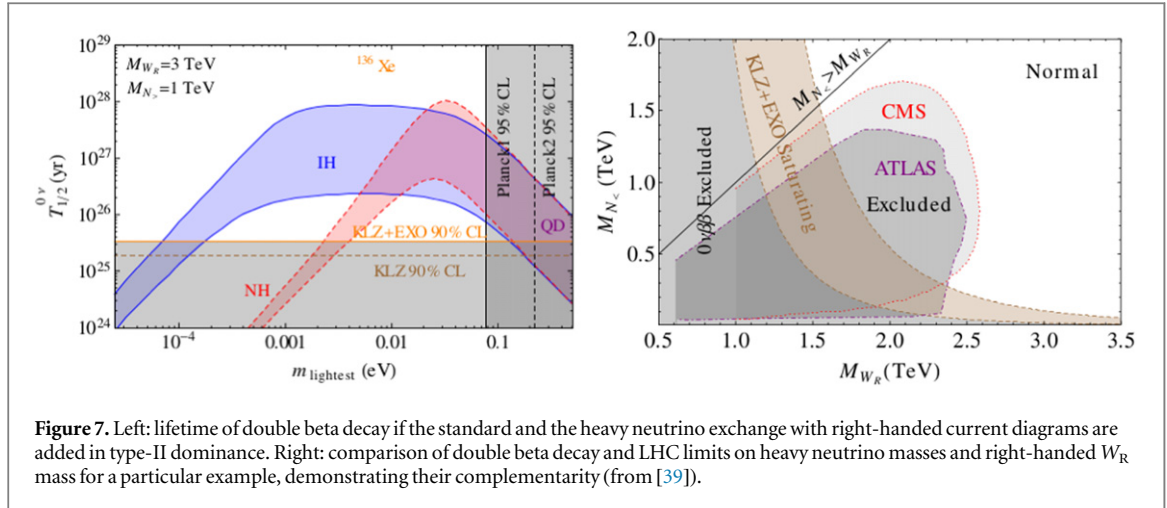
Finally, there exists a contribution from the exchange of a right-handed doubly charged Higgs triplet Δ_{R}, giving rise to the same effective operator structure as heavy neutrino exchange [84]. The corresponding effective short-range coupling is

$$\varepsilon_3^{RRz} = \sum_{i=1}^{3} V_{ei}^2 \frac{m_{N_i} m_p}{m_{\Delta_{\mathrm{R}}}^2} \frac{m_{W_{\mathrm{L}}}^4}{m_{W_{\mathrm{R}}}^4} \lesssim 1 \times 10^{-8}. \tag{20}$$

Since the Higgs triplet can mediate $\mu \to 3e$ at tree level there are strong constraints on this diagram by lepton flavor violation bounds [80].

A particularly predictive case occurs if type-II seesaw dominance holds, i.e. if the neutrino mass matrix is generated by the $SU(2)_{\mathrm{L}}$ triplet term m_{L}. Due to the discrete left–right symmetry this term is directly proportional to the heavy neutrino mass matrix, hence V in equation (20) equals the PMNS matrix U and $m_i \propto M_i$. It follows [85] that typically for a normal mass ordering the lifetime of double beta decay is finite while for an inverted mass ordering it can be infinite due to possible cancellations. Just as for the case of light sterile neutrinos (see equation (13)) the standard phenomenology has turned around.

Obviously many diagrams can contribute at the same time and interference between the different diagrams can arise. This nicely demonstrates the importance of the ideas discussed above to discriminate the various mechanisms. Another example for the consequences of several diagrams, adding for instance the heavy neutrino exchange with right-handed currents to the standard amplitude in the case of type-II dominance is illustrated in figure 7. A lower limit on the smallest neutrino mass results, in contrast to the upper limit deduced if only the

Figure 7. Left: lifetime of double beta decay if the standard and the heavy neutrino exchange with right-handed current diagrams are added in type-II dominance. Right: comparison of double beta decay and LHC limits on heavy neutrino masses and right-handed W_R mass for a particular example, demonstrating their complementarity (from [39]).

standard diagram was taken into account. We close the discussion on left–right symmetry by noting that several LHC anomalies at 2 TeV can be explained by a W_R of this mass mixing with the SM W [86].

3.2. *R*-parity violating supersymmetry

In the minimal supersymmetric extension of the standard model) (the MSSM) one typically employs a discrete Z_2 *R*-parity in order to make the lightest supersymmetric particle stable, thus providing a dark matter candidate for cosmology and to avoid too fast proton decay. Since a convincing theoretical reason for *R*-parity conservation is lacking, one can investigate the consequences of its violation. Using discrete symmetries one can avoid terms that lead to proton decay and is left with a superpotential including the LNV terms

$$W_{RPV} = \lambda'_{ijk} L_i Q_j \bar{D}_k, \tag{21}$$

where i, j, k are generation indices. Note that the LNV is by one unit, hence two vertices are required for $0\nu\beta\beta$, which occurs through long- and short-range Feynman graphs involving the exchange of superpartners [87–91]. The short-range contribution has been discussed in [87–89]. Combining the half-life limit [4] with the corrected numerical values [12] of the nuclear matrix elements first calculated in [89] leads to the limit on λ'_{111} given by

$$\lambda'_{111} \leqslant 2 \times 10^{-4} \left(\frac{m_{\tilde{q}}}{100\,\text{GeV}}\right)^2 \left(\frac{m_{\tilde{g}}}{100\,\text{GeV}}\right)^{1/2}, \tag{22}$$

where we have assumed dominance of the gluino exchange diagram and took $m_{\tilde{d}_R} = m_{\tilde{u}_L} \equiv m_{\tilde{q}}$ for the exchanged squarks.
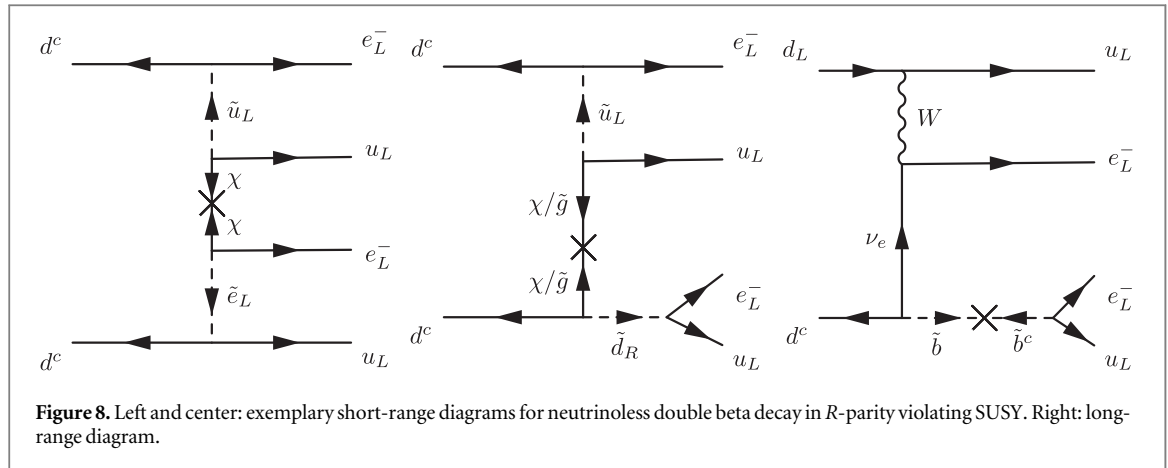
In addition $0\nu\beta\beta$ decay is also sensitive to other combinations of the couplings λ'_{ijk}. When one takes into account that the SUSY partners of the left- and right-handed quark states can mix with each other, new diagrams appear in which the neutrino-mediated double beta decay is triggered by SUSY exchange in the vertices [90–92], see figure 8 and note that this is a long-range diagram. Assuming the supersymmetric mass parameters of order 100 GeV, the present GERDA half life limit implies: $\lambda'_{113}\lambda'_{131} \leqslant 3 \times 10^{-8}$, $\lambda'_{112}\lambda'_{121} \leqslant 1 \times 10^{-6}$. Comparable bounds can be deduced from *B* and *K* physics which depend however on different superpartner masses and are thus complementary to the bounds derived here [93]. Recently, the lepton non-universality anomaly at LHCb [94] and the CMS anomaly in the search for right-handed *W* bosons have been explained within *R*-parity violating SUSY with $\lambda'_{113} = \mathcal{O}(10^{-3}–10^{-2})$ and $\lambda'_{112} = \mathcal{O}(10^{-1})$ and scalar masses in the TeV range [95].

3.3. Leptoquarks (LQs)

LQs are hypothetical bosons (scalar or vector particles) with couplings to both leptons and quarks which appear for instance in GUTs, extended technicolor or compositeness models. LQs which conserve baryon number are searched for both in accelerator experiments [96] and non-accelerator searches [97, 98]. The mixing of different LQ multiplets by a possible leptoquark-Higgs coupling [99] can lead to long-range contributions to $0\nu\beta\beta$ decay, if these couplings violate lepton number [100]. From the lower limit on the $0\nu\beta\beta$ lifetime, bounds on effective couplings can be derived [100] which are typically of order

$$Y_{LQ-Higgs} = \text{few} \cdot 10^{-6} \tag{23}$$

for LQ masses of the order of $\mathcal{O}(200)\,\text{GeV}$.

Figure 8. Left and center: exemplary short-range diagrams for neutrinoless double beta decay in *R*-parity violating SUSY. Right: long-range diagram.

3.4. Extra dimensions

Models with more than three space dimensions became popular in recent years as a way to reduce the four-dimensional Planck scale and alleviate this way the hierarchy problem. Extra dimensions have also been suggested as a way to generate small Dirac neutrino masses by utilizing the volume suppressed wave function overlap of a left-handed neutrino confined to a three-dimensional subspace called the brane and a right-handed neutrino propagating in the extra-dimensional hyperspace called the bulk [101, 102]. A minimal higher-dimensional model implementing LNV compactifies a five-dimensional theory on an S^1/Z_2 orbifold, and adds a single (bulk) sterile neutrino to the field content of the SM [103]. While the singlet neutrinos can freely propagate in the bulk, all SM particles are localized on the $(3 + 1)$-dimensional brane.

In principle the excitations of the sterile neutrino in the compactified extra dimensions, the so-called Kaluza–Klein tower of states, will contribute to the $0\nu\beta\beta$ decay rate. The masses $m(n)$ of these Kaluza–Klein states are obtained by diagonalizing the infinitely dimensional Kaluza–Klein mass matrix and result approximately as

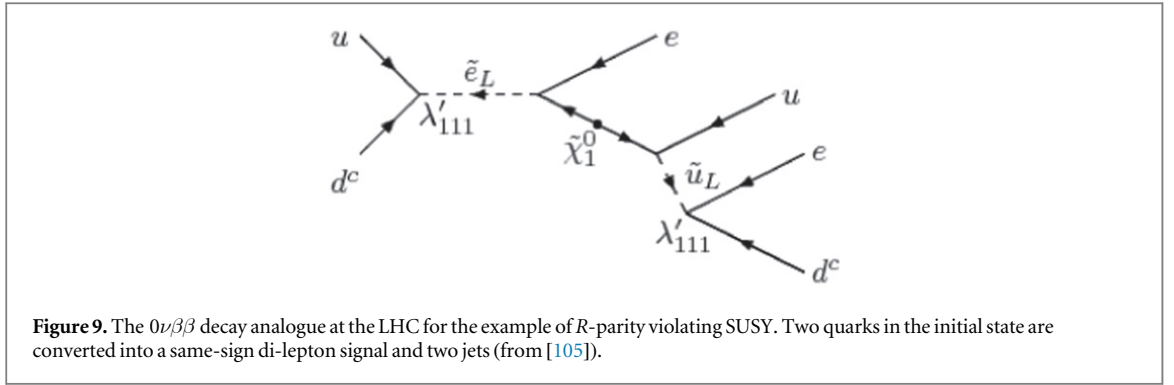$$m_{(n)} \approx \frac{n}{R} + \varepsilon. \tag{24}$$

Here n is the index denoting the Kaluza–Klein excitation, R is the radius of the extra dimension and ε is the smallest diagonal entry in the neutrino mass matrix. The tower starts with the smallest Kaluza–Klein masses being much lighter than p_F and thus giving rise to long-range contributions, continues through the 100 MeV region up to large masses with short-range contributions. Thus such extra-dimensional scenarios constitute a special case which cannot be categorized into the simple classes of contributions introduced in the effective operator parametrization described above.

Moreover, such extra-dimensional models generically predict a Kaluza–Klein neutrino spectrum with approximately degenerate masses and opposite CP parities that leads to an extremely suppressed contribution to double beta decay and only one non-vanishing Δm^2 insufficient to explain solar and atmospheric neutrino oscillations, at least if the brane is located at one of the two orbifold fixed points. In this case the lepton number violating operators thus would be absent as a consequence of the Z_2 discrete symmetry. If, however, the brane is shifted away from the orbifold fixed points, the Kaluza–Klein neutrinos can couple to the *W* bosons with unequal strength, thus avoiding CP-parity cancellations in the $0\nu\beta\beta$ amplitude. This breaking of lepton number can lead to observable effects in neutrinoless double beta decay experiments. The size of the brane-shift can then be determined from the $0\nu\beta\beta$ lifetime or its upper bound.

This leads to a nuclear matrix element depending on the Kaluza–Klein neutrino masses $m_{(n)}$, and thus to predictions for the double beta decay observable that depend on the double beta emitter isotope used in the experiment. Another interesting property of this model is that the values of the mass eigenvalues of the lightest neutrinos do not imply an upper bound on the $0\nu\beta\beta$ decay rate. The rate can be close to the experimental limit even for the case of an almost vanishing lightest neutrino mass which constitutes a rather unique property of such extra-dimensional brane-shifted scenarios.

4. Lepton number violation at colliders, double beta decay and the baryon asymmetry of the universe

In this section we deal with the links between neutrinoless double beta decay and LNV processes at colliders and in cosmology, with the latter ones having important consequences for baryogenesis. As mentioned already in the

Figure 9. The $0\nu\beta\beta$ decay analogue at the LHC for the example of *R*-parity violating SUSY. Two quarks in the initial state are converted into a same-sign di-lepton signal and two jets (from [105]).

last section, while $0\nu\beta\beta$ decay provides the best possibility to search for light massive Majorana neutrinos, LNV as featured in the short-range contributions can in general be probed also in collider processes. For example, as discussed for left–right symmetric models [85, 104] (see figure 7) and *R*-parity violating supersymmetry [93, 105], the short-range contribution can easily be crossed into a diagram with two quarks in the initial state where resonant production of a heavy particle leads to a same-sign dilepton signature plus two jets at the LHC, see figure 9. The arguments we present here rely on the possibility to resonantly produce those heavy particles. If one wants to discuss the LHC bounds in a model-independent way it is necessary to specify which particles are propagating in the inner legs, which requires a decomposition of the $d = 9$ operator in the effective mass approach discussed above. Such a decomposition has been worked out in [106] where two different possible topologies have been identified. While topology 1 contains two bosons and a fermion in the internal lines (like the right-handed analogue of the standard diagram), topology two contains an internal three-boson-vertex (like the triplet exchange diagram). This procedure was applied to diagrams for the LHC analogue of $0\nu\beta\beta$ decay and first results for a general analysis based on this decomposition for topology one were presented in [75, 76]. The result was that the LHC is typically more sensitive than $0\nu\beta\beta$ decay for short-range contributions, with the exception of leptoquark exchange. Thus—generally and with some exceptions—one can conclude that either an observation of $0\nu\beta\beta$ decay would imply an LNV signal at the LHC as well. In turn, no sign of LNV at the LHC would exclude an observation of $0\nu\beta\beta$ decay, or $0\nu\beta\beta$ decay would have to be triggered by a long-range mechanism.
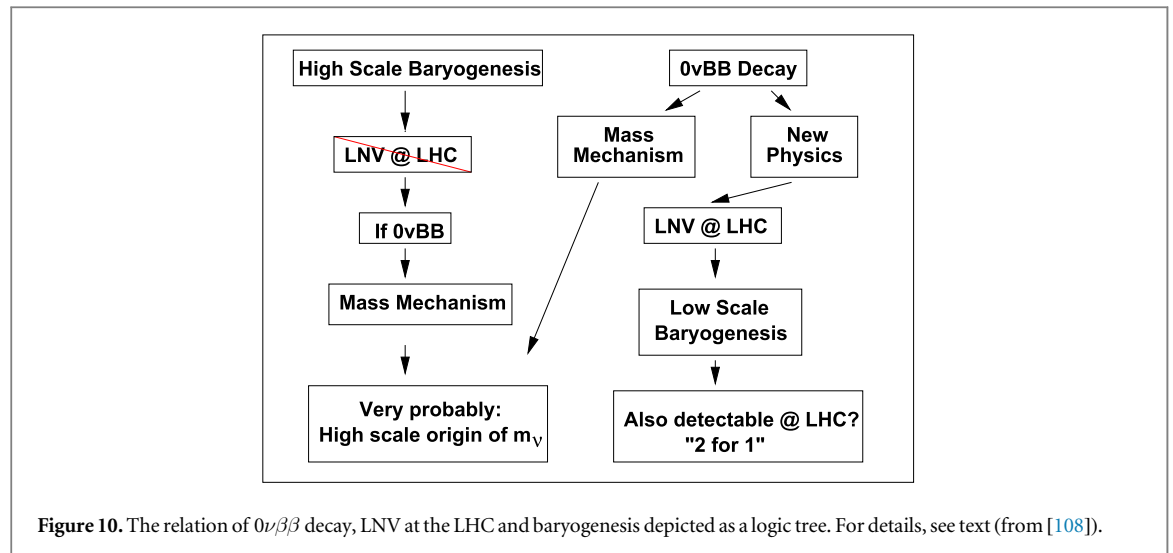
In addition, as has been mentioned before, LNV and baryon number violation are closely interrelated. More concretely, an observation of LNV at low energies has important consequences for a pre-existing lepton asymmetry in the Universe as the observation of LNV at the LHC will yield a lower bound on the washout factor for the lepton asymmetry in the early Universe. In [107] it has thus been pointed out that any observation of LNV at the LHC will falsify high-scale leptogenesis. It is easy to see that this argument can be extended even further (for further details see [108, 109]).

Similarly to what happens in leptogenesis, where the coaction of $B - L$ violating heavy neutrino decays with $B + L$ violating sphaleron processes generates a baryon asymmetry, here we stress that an observation of low energy $B - L$ violation at the LHC or elsewhere in combination with $B + L$ violating sphaleron processes will wash out any pre-existing baryon asymmetry, whatever the concrete mechanism of baryogenesis is.

By combining this argument with the results of [75, 76] mentioned above, one can conclude that an observation of short-range $0\nu\beta\beta$ decay will typically imply that LNV processes should be detected at the LHC as well. This in turn will falsify standard thermal leptogenesis and in general any high-scale scenario of baryogenesis. While the observation that low-energy LNV is dangerous for baryogenesis is not new (see e.g. [110–115]), only quite recently it has been realized in [109] that the argument applies for all short range contributions (d) and also for the long-range contribution (b) in figure 5.

It should be stressed however that this reasoning is rather general and various loopholes can exist in specific models or realizations of baryogenesis:

- these include scenarios where lepton number is broken not universally but only for specific flavors. As $0\nu\beta\beta$ decay only probes $\Delta L_e = 2$, i.e. LNV in the electron sector, it may be possible that lepton number could still be conserved for example in the τ flavor being not necessarily in equilibrium with the *e* and μ flavors in the early Universe [107]. As has been discussed in [109], however, an observation of lepton flavor violating (LFV) decays such as $\tau \rightarrow \mu\gamma$ may require LFV couplings large enough to wash out such a flavor specific lepton asymmetry when combined with LNV observed in a different flavor sector;

- models with hidden sectors, new symmetries and/or conserved charges may protect a baryon asymmetry against LNV washout as proposed for the example of hypercharge by [116];

Figure 10. The relation of $0\nu\beta\beta$ decay, LNV at the LHC and baryogenesis depicted as a logic tree. For details, see text (from [108]).

- models where lepton number is broken at a scale below the electroweak phase transition where sphalerons are no longer active.

As in general an observation of low energy LNV would invalidate any high-scale generation of the baryon asymmetry though, such protecting mechanisms should be introduced and discussed explicitly in any model combining low-scale LNV with high-scale baryogenesis.

By building up on the arguments given above, one can conclude, keeping the above mentioned loopholes in mind, that if $0\nu\beta\beta$ decay is observed, it is either triggered by a long-range mechanism, such as the standard interpretation with a light Majorana neutrino mass, or due to a short-range operator. In the latter case chances are high that lepton number is observed at the LHC as well. This further implies that baryogenesis is a low-scale phenomenon which also may be observable at the LHC or other experiments.

If, inversely, the baryon asymmetry is produced at a high scale, LNV will not be observable at the LHC. If, aditionally, $0\nu\beta\beta$ decay will be found, it will typically be triggered by a long-range operator. In combination with the assumptions that we did not see any other explicit LNV processes at experimentally accessible energies and that a high scale production mechanism is the source the baryon asymmetry, this case will probably point also towards a neutrino mass generated at a high scale, such as a 'vanilla' type-I seesaw mechanism in combination with leptogenesis.

To summarize this discussion, an observation of $0\nu\beta\beta$ decay will (see figure 10) either imply LNV at the LHC and low-scale baryogenesis and thus a possible observation of both processes in the near future, or very probably point towards a high-scale origin of both neutrino masses and baryogenesis.

5. Conclusions

The discovery of LNV would have far-reaching consequences affecting deeply our thinking about fundamental physics, including our ideas about unification and our understanding of the generation of the baryon asymmetry of the Universe. Neutrinoless double beta decay and LNV thus remain fields that enjoy large interest from both experimental and theoretical communities in nuclear and particle physics. In this review we have tried to summarize the multifaceted relations between neutrinoless double beta decay, neutrino physics and new physics beyond the SM. The continuous theoretical and experimental efforts around the world justify the hope that we may not be too far away from identifying the origin of LNV.

Acknowledgments

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