Neutrinoless double beta decay in the context of seesaw models

Jacobo López-Pavón

IPPP, Department of Physics, Durham University, South Road, Durham, DH1 3LE, UK

E-mail: `jacobo.lopez-pavon@durham.ac.uk`

Abstract. The phenomenology of neutrinoless double beta decay in the context of seesaw models is studied. We focus on the interplay between the contribution of the light mostly active neutrinos and the extra degrees of freedom of the models, clarifying why the light neutrino contribution should always be considered. We show how considering the relation between the mostly active neutrino and the New Physics contribution leads to strong constraints on the New Physics parameters or to a cancellation, which would result into unobservable process even if neutrinos are Majorana particles, depending on the mass regime of the extra states. Finally, we also discuss how non canonical versions of the seesaw could reconcile large rates of neutrinoless double beta decay with more stringent cosmological bounds on neutrino masses. In order to perform this analysis we compute the nuclear matrix elements as a function of the mass of the mediating fermion using the Interacting Shell Model (ISM) and estimating the associated uncertainty.

Neutrino oscillations, implying the massive nature of neutrinos, constitute an evidence for physics beyond the Standard Model (SM). Thus, models accommodating this neutrino masses become an important component in the search for New Physics. In this context, one of the most promising processes is the neutrinoless double beta decay ($0\nu\beta\beta$ decay), whose detection is the purpose of several ongoing and upcoming experiments. Since this process is lepton number violating, its observation would imply that neutrinos are Majorana fermions.

Among the more popular models for neutrino masses, we find the different types of seesaw mechanisms: type-I, where the SM field content is extended with fermion singlets, type-II, where the extension is done with heavy scalar triplets, and type-III, where the extra heavy particles included are fermion triplets. They are the extensions of the SM particle content that lead to the Weinberg $d = 5$ effective operator after the heavy fields have been integrated out.

All these extra degrees of freedom, required to induce the Majorana nature of the SM neutrinos, can also contribute to the $0\nu\beta\beta$ process. In the literature, it is common to study the effects of the SM neutrinos or the extra states introduced to account for their masses independently. However, these contributions are indeed related through the generation of neutrino masses. Taking into account this relation, a precious information about the origin of neutrino masses can be obtained.

In order to study the interplay between the contributions to the $0\nu\beta\beta$ decay rate of the SM neutrinos and the extra degrees of freedom, we have computed the nuclear matrix element
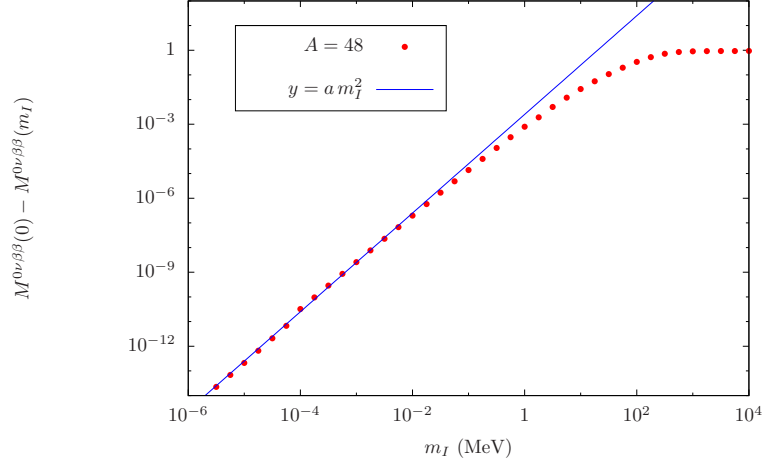
Figure 1. NME cancellation in the light mass regime. The example is taken for the ^{48}Ca decay, but the same quadratic dependence is seen for the remaining nuclei studied. The value of the parameter is $a = 2.51 \cdot 10^{-3}$ MeV^{-2}.

(NME) without any assumption on the mass of the neutrinos mediating the process:

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G_{01} \left| \sum_j U_{ej}^2 \frac{m_j}{m_e} M^{0\nu\beta\beta}(m_j) \right|^2 , \tag{1}$$

where G_{01} is a well-known kinematic factor, m_e the electron mass, m_j the mass of the j neutrino, $M^{0\nu\beta\beta}(m_j)$ the NME and U_{ej} are elements of the neutrino mixing matrix.

The calculation was performed using Interacting Shell Model (ISM) nuclear wavefunctions, one of the most popular methods employed to obtain $0\nu\beta\beta$ decay NMEs. The results [1],[2] show two distinct regions for the behavior of the NME as a function of the virtual neutrino mass: almost constant up to $m_i \simeq 100$ MeV and then decreasing quadratically as the neutrino mass increases beyond 100 MeV. The value for the mass for which the transition takes place, 100 MeV, corresponds to the typical momentum exchange of the decay, which is the momentum of the virtual neutrino $|\mathbf{p}|$. This scale is fixed by the typical distance between the two decaying nucleons, $r \simeq 1$ fm, and is usually referred as the nuclear scale. Notice that this behavior of the NME can easily be understood expanding the neutrino propagator $\frac{1}{p^2 - m_j^2}$.

Different approximations are required in order to perform the calculation, both obtaining the wavefunctions and in the treatment of the two-body transition operator, and consequently some uncertainties are induced into the NMEs. These can be estimated[1] in $\sim^{+25}_{-35}$ % for light neutrino exchange ($m_j \leq 100$ MeV) and $\sim^{+35}_{-40}$ % for heavy neutrinos ($m_j \geq 100$ MeV), more sensitive to the short range part of the transition operator.

With the NMEs we can now analyze the contributions to $0\nu\beta\beta$ decay of the different mechanisms that lead to Majorana neutrino masses. Within the type-I seesaw, depending on whether the extra mass eigenstates fall in the light or heavy neutrino mass regimes we can split their respective contributions to the amplitude:

$$A \propto \sum_i^{\text{light}} m_i U_{ei}^2 M^{0\nu\beta\beta}(0) + \sum_I^{\text{light}} m_I U_{eI}^2 M^{0\nu\beta\beta}(m_I) + \sum_I^{\text{heavy}} m_I U_{eI}^2 M^{0\nu\beta\beta}(m_I), \tag{2}$$
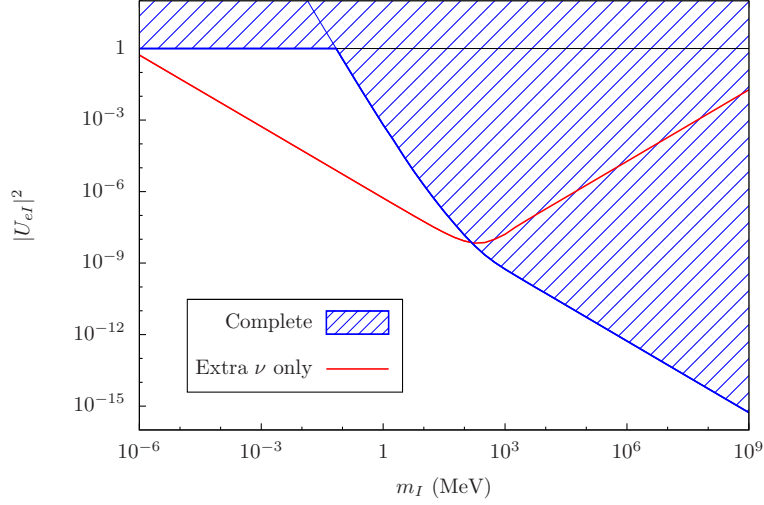
---

[1] See [1] for details.

Figure 2. Bounds from CUORICINO on the extra neutrino mixing from $0\nu\beta\beta$ decay in ^{130}Te, with a 90 % CL half-life [7]. We have assumed the extra neutrinos are non hierarchical and show the bounds as a function of their common mass. We compare the case in which the contribution from the SM neutrinos is properly taken into account (striped area) to that in which only the extra contribution is considered (above the red line).

where capital letters denote the mass index of the mostly sterile states and lowercase letters the mostly active SM states. Notice that we can safely assumed $M^{0\nu\beta\beta}(m_i) = M^{0\nu\beta\beta}(0)$ for the NMEs associated to the SM neutrinos. Moreover, the diagonalization of the neutrino mass matrix provides the following useful relation:

$$U^*\text{diag}\left\{m_1, m_2, ..., m_n\right\}U^\dagger = \begin{pmatrix} 0 & Y_N v/\sqrt{2} \\ Y_N^T v/\sqrt{2} & M_N \end{pmatrix}. \tag{3}$$

which relates the mostly active and extra degrees of freedom parameters. We can now distinguish three cases exhibiting very different phenomenologies depending on the mass regime of the extra mass eigenstates:

All extra mass states are light. In this case Eq. (3) implies

$$A \approx -\sum_I^{\text{light}} m_I U_{eI}^2 \left(M^{0\nu\beta\beta}(0) - M^{0\nu\beta\beta}(m_I) \right). \tag{4}$$

Since in this regime the NMEs are basically independent of the neutrino mass, $M^{0\nu\beta\beta}(m_I) \simeq M^{0\nu\beta\beta}(m_i) \simeq M^{0\nu\beta\beta}(0)$, the rate of $0\nu\beta\beta$ decay is very suppressed. Indeed, only the different neutrino masses in the NME prevent a full cancellation leading to a suppression of $\Delta m^2/p^2$ with $|p^2| \simeq (100 \text{ MeV})^2$. The $\Delta m^2/p^2$ dependence of $M^{0\nu\beta\beta}(0) - M^{0\nu\beta\beta}(m_I)$ that drives the suppression to the $0\nu\beta\beta$ rate is depicted in Fig. 1. As it can be observed, the $0\nu\beta\beta$ decay becomes experimentally inaccessible in this regime, even if neutrinos are Majorana particles.

All extra mass states in the heavy regime. Now the NMEs for this extra states are very suppressed compared to the SM ones. Furthermore, Eq. (3) implies that

$$A \approx -\sum_I^{\text{heavy}} m_I U_{eI}^2 \left(M^{0\nu\beta\beta}(0) - M^{0\nu\beta\beta}(m_I) \right) \approx -\sum_I^{\text{heavy}} m_I U_{eI}^2 M^{0\nu\beta\beta}(0) \approx \sum_i^{\text{light}} m_i U_{ei}^2 M^{0\nu\beta\beta}(0). \tag{5}$$

Thus, the contribution from the light active neutrinos (first term) dominates the transition rate but can be used to set a strong bound on the mixing of the heavy neutrinos [3]. This leads to a much stronger constraint than the one usually considered in the literature, as it is shown in Fig. 2.

Extra mass states in the light and heavy regimes. In this scenario the leading terms stem from the light states:

$$A \propto \sum_i^{\text{light}} m_i U_{ei}^2 M^{0\nu\beta\beta}(m_i) + \sum_I^{\text{light}} m_I U_{eI}^2 M^{0\nu\beta\beta}(m_I). \tag{6}$$

Now it is possible to satisfy Eq. (3) even in a situation where $m_i U_{ei}^2 \ll m_I U_{eI}^2$, by canceling the contribution of the extra heavy states against that of the extra light ones while keeping the light neutrino masses small. This implies a certain level of fine-tuning. However, in such a situation, the contribution of the light extra states could dominate over that of the active. As an example, if we consider the controversial Heidelberg-Moscow claim for a positive $0\nu\beta\beta$ decay signal [4], the accommodation of this signal through only SM neutrinos would require, using our ISM NMEs, 0.24 eV $< m_{\beta\beta} < 0.89$ eV at 2σ. The interpretation of this claim as light active SM neutrinos is very disfavored by the constraints from cosmology and neutrino oscillation data [5, 6]. However, this signal could be accommodated in a model with heavier neutrinos (which are not bounded by cosmology) mediating the process. Indeed, we could reinterpret the result as 0.24 eV $< \left| \sum_I^{\text{heavy}} m_I U_{eI}^2 \right| < 0.89$ eV.

As for the type-II and type-III seesaws, current bounds from accelerator experiments place the extra degrees of freedom in the heavy regime. This effectively reduces the phenomenology to that which appears for the type-I seesaw with only heavy extra states[2]. In the same manner, combined type-I and type-II or type-III seesaw models resemble the situation of type-I seesaw with both light and heavy extra states. For instance, the Heidelberg-Moscow claim can be interpreted as 0.24 eV $< \left| m_{ee}^{\Delta,\Sigma} \right| < 0.89$ eV in this context, where m_{ee}^{Δ} and m_{ee}^{Σ} are the Majorana neutrino masses induced by the type-II and type-III seesaw models, respectively.

Acknowledgments

References

[1] Blennow M, Fernandez-Martinez E, Lopez-Pavon J and Menendez J, *JHEP* **1007** (2010) 096 arXiv:1005.3240 [hep-ph].

[2] Blennow M, Fernandez-Martinez E, Lopez-Pavon J and Menendez J, http://wwwth.mppmu.mpg.de/members/blennow/nme mnu.dat, (2010).

[3] Xing Z z, *Phys. Lett. B* **679** (2009) 255 arXiv:0907.3014 [hep-ph].

[4] Klapdor-Kleingrothaus H V, Krivosheina I V, *Mod. Phys. Lett.* **A21** (2006) 1547-1566.

[5] Komatsu E *et al.* [WMAP Collaboration], *Astrophys. J. Suppl.* **192** (2011) 18. arXiv:1001.4538 [astro-ph.CO].

[6] Hannestad S, Mirizzi M, Raffelt G G, Wong Y Y Y, *JCAP* **1008** (2010) 001. arXiv:1004.0695 [astro-ph.CO].

[7] Arnaboldi C *et al.* [CUORICINO Collaboration], *Phys. Rev.* **C78** (2008) 035502. arXiv:0802.3439 [hep-ex].

---

[2] See [1] for details.