Charged Higgs contribution to $0\nu2\beta$ decay

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The singly charged Higgs boson contribution to $0\nu2\beta$ is neglected on grounds of couplings involving small masses and small nuclear matrix elements. We reconsider such contributions, but now in the light of QCD corrections and renormalization group evolution. It is found that the charged Higgs contribution is generically as large as (and at times significantly larger than) the other contributions and there can be large cancellations between contributions coming from different sources. This observation will have an important impact on the phenomenology.

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Charge neutrality of the neutrinos opens up the possibility of them being Majorana particles. Neutrinoless double beta ($0\nu2\beta$) decay, $(A, Z) \rightarrow (A, Z + 2) + 2e^-$ is an unambiguous signature of the Majorana nature of the neutrinos. Such a process violates lepton number by two units. Experimental confirmation of the mixing of different neutrinos and the fact that neutrinos are massive particles (see [3] for best fit values of the parameters) already implies physics beyond the standard model (SM). $0\nu2\beta$ decay is a powerful probe of physics beyond SM since it has the potential to discriminate between the two hierarchies of the neutrino masses. This becomes particularly important and effective in the context of models which involve TeV scale particles, like low scale seesaw models or low energy supersymmetric models including models with R-parity violation or leptoquark models. More interestingly, $0\nu2\beta$ diagrams in such low scale models can have distinctive signatures at the large hadron collider (LHC). For an incomplete list discussing various aspects of $0\nu2\beta$ decay and impact on other phenomenological issues see e.g. [4].

Experimentally, studies have been carried out on several nuclei ([5]-[11]). Only one of the experiments (HM) has claimed observation of $0\nu2\beta$ signal in $^{76}\text{Ge}$. The half-life at 68% confidence level is: $T_{1/2}^{0\nu}(^{76}\text{Ge}) = 2.22^{+0.31}_{-0.34} \times 10^{25}$ yr. A combination of the Kumland-Zen and EXO-200 results, both using $^{136}\text{Xe}$, yields a lower limit on the half-life $T_{1/2}^{0\nu}(^{136}\text{Xe}) > 3.4 \times 10^{25}$ yr which is at variance with the HM claim. Very recently, the GERDA experiment reported the lower limit on the half-life based on the first phase of the experiment: $T_{1/2}^{0\nu}(^{76}\text{Ge}) > 2.1 \times 10^{25}$ yr. A combination of all the previous limits results in a lower limit $T_{1/2}^{0\nu}(^{76}\text{Ge}) > 3.0 \times 10^{25}$ yr at 90% confidence level. The new GERDA result and the combination both are again at odds with the positive claim of HM. Higher statistics in future will shed more light. One can think of comparing $0\nu2\beta$ predictions for different nuclei in order to study the sensitivity of of theoretical calculations on the nuclear matrix elements (NMEs) used.

It is practically useful to separate the $0\nu2\beta$ decay amplitude into the so called long-range and short-range parts (for a review of theoretical and experimental issues and the sources of uncertainties and errors, see [12] and references therein). The long range contribution is the one arising when a light neutrino is exchanged while the short range part gets its name from the fact that the intermediate particles are all very massive and therefore the effective interaction becomes point-like once the heavy degrees of freedom are integrated out. This distinction between the long range and the short range contributions to $0\nu2\beta$ amplitude is also natural and appropriate from the point of view of renormalization and evolution under renormalization group equations (RGEs). The last piece of input is the non-perturbative NMEs, which are nothing but properly normalized matrix elements of the quark level operators sandwiched between the nucleon states. At present, the biggest source of uncertainty stems from the NMEs, and the predictions can vary up to a factor of two or more depending upon the specific NMEs employed (see [13]).

Recently, for the very first time, it has been shown [14] that perturbative QCD corrections to the short range part can have an important effect on the $0\nu2\beta$ rate. The main effect is related to the fact that QCD corrections generate operators with colour mis-matched structure. These operators have effective couplings, called the Wilson coefficients encoding the relevant information about the heavy degrees of freedom, which very roughly speaking are $1/N_c$ of the colour matched operators, $N_c$ being the number of colours. Though accompanied by smaller coefficients, such operators when Fierz transformed can lead to different Dirac structures whose nuclear matrix elements are way large compared to others usually considered. This observation is expected to have a huge impact on the phenomenological studies in a given model. The operators generated due to the mediation of a scalar fall under this category. Another important outcome of the QCD corrections and RG evolution to the low scale is that there is a large cancellation between some of the colour matched and colour mis-matched operators. Operators of the form $V - A \otimes V - A$ or $V + A \otimes V + A$ exhibit this feature. But these are the operators that appear naturally in most of the theories of interest, thereby making the impact of QCD corrections an important feature that
should be included in the calculation of the $0\nu2\beta$ rate. An issue of concern is the possibility of large cancellations among the various short range contributions, thereby significantly altering the limits on the masses and couplings in a given underlying model. Such cancellations (or large enhancements) will also change the phenomenological aspects while studying the same models (applicable to low scale models) at LHC. What is important here is the fact that such cancellations or enhancements do not depend on specific NMEs chosen.

A large class of models have an extended Higgs sector, popular examples being two Higgs doublet models, supersymmetry, left-right symmetric models. In these models, apart from other particles, one has at least one physical singly charged scalar (denoted by $H^+$) that mediates charged-current interactions. We shall assume for the present that in each of the models considered, there are heavy right handed neutrinos ($N$) present. The contribution of $H^+$ to $0\nu2\beta$ amplitude can be obtained by replacing the $W$’s by the $H^\pm$ and appropriately changing the couplings, which typically depend on the masses of the fermions at the relevant vertex. This feature holds in all the models mentioned above. These contributions to $0\nu2\beta$ are simply ignored since the vertices are dependent on masses of light quarks and/or suppression due to charged Higgs mass in the propagators. Further, the NMEs relevant for a contribution arising due to $H^+$ are smaller than the ones for $W$’s. All of these have prompted one to totally discard the $H^+$ contributions. However, as argued above, QCD corrections can change the picture completely. In the present note we consider the minimal left-right symmetric model (see for example [14] for the details and features of the model) for concreteness but we emphasize again that the features studied here remain true in all the models mentioned above.

We begin by recapitulating the essentials of the left-right model. For consistency of notation, we follow [16]. The smallest gauge group implementing the left-right symmetry is $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. The left handed and right handed fields transform as doublets under $SU(2)_L$ and $SU(2)_R$ respectively and therefore the two types of fields are treated at the same footing. The model naturally contains right handed neutrinos, appearing as a component of the doublet along with the right handed leptons. The fermionic content (and charge assignment, with $Y = B - L$) of the model is thus:

\begin{align*}
L_{iL} & = \begin{pmatrix} \nu_i^L \\ \ell_i^L \end{pmatrix}_L, \quad L_{iR} = \begin{pmatrix} \nu_i^L \\ \ell_i^L \end{pmatrix}_R, \quad Q_{iL} = \begin{pmatrix} \nu_i^L \\ \ell_i^L \end{pmatrix}_L, \quad Q_{iR} = \begin{pmatrix} \nu_i^L \\ \ell_i^L \end{pmatrix}_R
\end{align*}

The gauge couplings and the gauge fields are denoted as $g_L, g_R$ ($g_L = g_R = g$), $g'$, $W_L, W_R$, $B$. The scalar sector of the model contains a bi-doublet and two triplets:

\begin{align*}
\phi & = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^0 & \phi_2^+ \end{pmatrix} (2, 2, 0) \\
\Delta_{L,R} & = \begin{pmatrix} \delta_{L,R}^+ & \delta_{L,R}^{++} \\ \delta_{L,R}^0 & \delta_{L,R}^{0+} \end{pmatrix} (3, 1, 1(3), 2)
\end{align*}

The neutral components of the Higgs fields acquire vacuum expectation values (VEVs), assumed to be all real here:

\begin{align*}
\langle \phi_{1,2}^0 \rangle = \frac{\kappa_{L,R}}{\sqrt{2}}, \quad \langle \phi_{1,2}^0 \rangle = \frac{v_{L,R}}{\sqrt{2}}
\end{align*}

For what is relevant below, there are two charged gauge bosons $W_{1,2}$ with masses

\begin{align*}
M_{W_{1,2}}^2 = \frac{g'^2}{4} \left( \kappa_+ + v_R^2 \pm \sqrt{v_R^4 + 4\kappa_+^2 \kappa_-^2} \right)
\end{align*}

where $\kappa_\pm = \sqrt{\kappa_1^2 \pm \kappa_2}$ and the angle $\xi$ parametrizes the mixing between the left and right $W$ fields: $\tan 2\xi = -2\kappa_1 \kappa_2 / v_R^2$. In what follows, we shall always assume $v_R >> \kappa_+$. In this limit, various expressions simplify a lot. Further, again for simplicity we assume that $\xi$ is small and to bring out the main points relevant for the present study, we shall set it to zero while writing the relevant interactions. In the scalar sector there are 14 physical Higgs bosons: four neutral scalars, two neutral pseudo-scalars, two singly charged scalars ($H^\pm_{1,2}$) and two doubly charged scalars. Among the singly charged scalars, one of them, $H^+_{1,2}$, is lighter and does not couple to quarks. It therefore does not participate in $0\nu2\beta$ process. The other singly charged scalar, $H^+_{2,3}$, is somewhat heavier but has the desired interactions. Let us assume that we have a TeV scale left-right model, implying that the heavy particles, including the heavy right handed neutrinos, are all at TeV range (the heavier among them like $H^+_{2,3}$ would be at a few TeV scale). The exact values of the masses of the particles will depend on the details of the parameters of the model. We choose to stay somewhat generic at this point. This has the advantage that the analysis below can be easily carried over to other models of interest where the interactions have the same form. For complete details about the particle spectrum, masses and interactions, the reader is referred to [16]. One has the following relevant interactions:

\begin{align*}
\mathcal{L}_{ffW} = \frac{g}{\sqrt{2}} [U^C_{LM}(\bar{U}D)_{V-A} + K_L(\bar{N}\ell)_{V-A}] W_{1\mu}^+ + \frac{g}{\sqrt{2}} [U^C_{RM}(\bar{U}D)_{V+A} + K_R(\bar{N}\ell)_{V+A}] W_{3\mu}^0
\end{align*}

where the fields are now written in the mass basis and $U^C_{LM}, U^C_{RM}, K_L$ and $K_R$ are the various mixing matrices. The fermion-charged Higgs interactions are:

\begin{align*}
\mathcal{L}_{ffH^+_{1,2}} = -\bar{U}[P_L(m_u U^C_{LM} B^+ - m_d U^C_{LM} A^+)] D + \bar{N}_a[P_L(m_N(K_L)_{at} B^+ - m_L(K_R)_{at} A^+)] D + \bar{N}_a[\Omega_L a m_N_a (K_R)_{at} B^+ - (K_L)_{at} m_L B^+)] \bar{\ell} \ell
\end{align*}
where summation over the indices is implicit and so is the hermitian conjugate part, and

\[ A^+ \sim \frac{\sqrt{2} \kappa_+}{\kappa_-} H_2^+, \quad B^+ \sim \frac{2 \sqrt{2} \kappa_1 \kappa_2}{\kappa_+ \kappa_-} H_2^+ \] (7)

Note that \( \kappa_- \to 0 \) would give rise to singular behaviour of observables and thus this strict limit needs to be avoided. On the other hand, for choice of parameters, there can be enhancement due smaller values of \( \kappa_- \). Compare the above interaction with that in 2HDMII or supersymmetry. There the charged Higgs couples to the up and down type members of the doublets as (as an example the quarks, but the same structure will follow for the leptons with appropriate changes):

\[-\frac{\sqrt{2}}{v} V_{UD} U(m_u \cot \beta P_L + m_d \tan \beta P_R) D \, H^+ + H.C. \] (8)

The above form is simpler than the explicit one given above for the left-right model. To gain a clear and quick understanding of the situation, let us momentarily work with this form. Further, recalling that \( m_u \sim m_d/2 \), let us choose to take both of them to be equal for simplicity, and denote it by \( m_q \). Further, in the minimal left-right model, \( U_C^{KM} = U_R^{KM} = V^{KM} \). The quark part of the 0\( \nu \beta \) amplitude will have the following structures: \((S \pm P) \otimes (S \pm P)\) and \((S \pm P) \otimes (S \mp P)\). The structures will be weighted by (in the left-right model, for some choices of the parameters, there can be a large enhancement as discussed above - this is not explicitly displayed for the time being though for detailed numerical analysis this will play a crucial role)

\[ C_H = V_{ud}^2 T_{ea}^2 (m_q m_{N_a})^2 / (m_{N_a} m_N^2) \] (9)

where the factors \((m_q m_{N_a})^2\) and \(m_{N_a} m_N^2\) in numerator and denominator respectively arise from the vertices and propagators and we have denoted the electron-heavy neutrino mixing by \( T_{ea} \). It is to be noted that the smallness of the quark mass at the vertex is compensated by the large heavy neutrino mass. QCD corrections will lead to quark operators with colour mis-matched structure. The weight (magnitude) of these operators after the RG evolution to the relevant low scale is typically 0.1-0.5 of the colour matched operator.

The following scalar-pseudoscalar operators are of interest for the present study:

\[ O_1^{SP \pm \pm} = \bar{u}_i (1 \pm \gamma_5) d_i \bar{u}_j (1 \pm \gamma_5) d_j \bar{c} (1 + \gamma_5) e^c \]
\[ O_2^{SP \pm \pm} = \bar{u}_i (1 + \gamma_5) d_i \bar{u}_j (1 + \gamma_5) d_j \bar{c} (1 + \gamma_5) e^c \] (10)

In addition, the following tensor operators are required for RG purposes [17]:

\[ O_1^{T \pm \pm} = \bar{u}_i \sigma_{\alpha \beta} (1 \pm \gamma_5) d_i \bar{u}_j \sigma_{\alpha \beta} (1 \pm \gamma_5) d_j \bar{c} (1 + \gamma_5) e^c \]
\[ O_2^{T \pm \pm} = \bar{u}_i \sigma_{\alpha \beta} (1 + \gamma_5) d_i \bar{u}_j \sigma_{\alpha \beta} (1 + \gamma_5) d_j \bar{c} (1 + \gamma_5) e^c \] (11)

It is sufficient to only consider operators with \( LL \) structure since the \( RR \) operators will have the same properties under QCD renormalization. Following [17], and adapting to the present case, the Wilson coefficients at the low scale approximately read (in units of \( C_H \), the only non-zero high scale coefficient):

\[ C_1^{SP \pm \pm} \sim 3, \quad C_2^{SP \pm \pm} \sim 0.16 \] (12)
\[ C_1^{T \pm \pm} \sim 0.06, \quad C_2^{T \pm \pm} \sim -0.17 \]

In obtaining the above approximate numerical values of the Wilson coefficients, we have assumed that all the heavy particles are around TeV and have used one step integrating out of heavy degrees of freedom. Some changes are expected once the threshold effects are incorporated. The points to be noted from above are that the coefficients \( C_1^{SP \pm \pm} \) get enhanced by a factor of 3 and the corresponding colour mis-matched coefficient is 0.16\( C_H \) while the colour matched tensor operator comes with a strength 0.06\( C_H \). The operators \( O_2^{SP \pm \pm} \) are now Fierz transformed which brings the tensor structure in the picture. The sum total of all this is that there are operators \( O_i^{T \pm \pm} \) with coefficients which are \( \sim 10-15\% \) of \( C_H \). Since the NMEs associated with the tensor operators are way bigger than the scalar ones, one can not naively throw away the scalar contribution in the end. This is the big difference that is brought in by the QCD corrections and RG evolution to the low scale. For numerical purpose we shall employ \( C_1^{T \pm \pm} = 0.12 C_H \). The tensor-pseudotensor structure yields the following NME:

\[ \langle \mathcal{J}_{\mu \nu}, \mathcal{J}_{\nu \rho} \rangle \propto -\alpha_2^{SR} \mathcal{M}_{GT,N} \] (13)

with \( \alpha_2^{SR} \sim 9.6 \frac{m_q}{m_{N_a} m_{N_a}} \) which is much larger than the NME for scalar-pseudoscalar operator (note the large multiplicative factor of 9.6 appearing in \( \alpha_2^{SR} \) which will play a crucial role in eventually enhancing the contributions):

\[ \langle \mathcal{J}_{(S \pm P)}, \mathcal{J}_{(S \mp P)} \rangle \propto -\alpha_1^{SR} \mathcal{M}_{F,N} \] (14)

with \( \alpha_1^{SR} \sim 0.145 \frac{m_q}{m_{N_a} m_{N_a}} \). This, together with the fact that \( \mathcal{M}_{GT,N} > \mathcal{M}_{F,N} \) justifies the neglect of scalar contributions to 0\( \nu \beta \) in the absence of QCD corrections.

Neglecting the contribution arising due to the doubly charged Higgs bosons, the other operators of interest are:

\[ O_1^{LL} = \bar{u}_i \gamma_\mu (1 - \gamma_5) d_i \bar{u}_j \gamma_\mu (1 - \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \]
\[ O_2^{LL} = \bar{u}_i \gamma_\mu (1 - \gamma_5) d_i \bar{u}_j \gamma_\mu (1 - \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \]
\[ O_1^{RR} = \bar{u}_i \gamma_\mu (1 + \gamma_5) d_i \bar{u}_j \gamma_\mu (1 + \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \]
\[ O_2^{RR} = \bar{u}_i \gamma_\mu (1 + \gamma_5) d_i \bar{u}_j \gamma_\mu (1 + \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \]
\[ O_1^{LR} = \bar{u}_i \gamma_\mu (1 - \gamma_5) d_i \bar{u}_j \gamma_\mu (1 + \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \]
\[ O_2^{LR} = \bar{u}_i \gamma_\mu (1 - \gamma_5) d_i \bar{u}_j \gamma_\mu (1 + \gamma_5) d_j \bar{e} (1 + \gamma_5) e^c \] (15)

with the Wilson coefficients evaluated at \( \mu \sim \mathcal{O}(\text{GeV}) \) in units of the corresponding coefficients at high scale [14]:

\[ C_1^{LL,RR} \sim 1.3, \quad C_2^{LL,RR} \sim -0.6 \]
\[ C_1^{LR,RL} \sim 1.1, \quad C_2^{LR,RL} \sim 0.7 \] (16)
As noted in [14], there is substantial cancellation after Fierz rearrangement in the above set of operators: \( LL, RR \) operators effectively yield \( C_1^{LL,RR} + C_2^{LL,RR} \) as the couplings with the same NMEs involved. Explicitly:

\[
(J^{(V \pm A)}J^{(V \pm A)}) \propto \frac{m_A}{m_p m_e} (\mathcal{M}_{GT,N} + \alpha_3^{SR} \mathcal{M}_{F,N}) \quad (17)
\]

where \( |\mathcal{M}_{GT,N}| \sim (2 - 4)|\mathcal{M}_{F,N}| \) with \( \alpha_3^{SR} \sim 0.63 \). Thus, it is reasonable to say that the above matrix element is essentially governed by \( \mathcal{M}_{GT,N} \). Let us further choose to neglect \( \mathcal{M}_{F,N} \) which simplifies the discussion without having any appreciable impact on phenomenology as long as the masses of the particles are all in the TeV range. If however, the charged Higgs boson is much lighter than some of the other particles in the spectrum, considerable care needs to be taken since the scalar operator gets enhanced at the low scale by a large factor.

Next let us consider the situation in the minimal left-right model where all the above operators are present. For the constraints on the model parameters see [18] and references therein. In many of the analysis, it is quite common to assume \( k_1 \gg k_2 \). Instead, there is a large parameter space where \( k_{1,2} \) may not be this hierarchical. In such a case, many of the constraints change. In particular, even if \( m_{H^+} \) is \( \sim 10 \) TeV or so, there is a reasonable contribution to various observables due to \( \kappa_- \) appearing in the coupling. In such a case, the contributions from the \( V \pm A \otimes V \pm A \) operators and the Fierz transformed scalar operator (yielding a tensor-pseudotensor operator) can be comparable. Moreover, the relative signs between the two contributions can lead to large cancellations. In that case, the short distance contribution will be dominated by the color matched scalar contribution which is naively thrown away. Depending on the couplings, particularly if there are additional sources of CP violation, there could be significant enhancements. Either way, the phenomenological impact, i.e. effect of these on constraints on the couplings and masses, is going to be large.

Let us now briefly consider supersymmetric models (we assume as before that there are right handed neutrinos in the model), with (see for example [19]) and without R-parity (see for example [20]). With R-parity conserved, the charged scalar contribution can be the largest since the charged Higgs mass is no longer forced to be 10 TeV or so but few hundred of GeV. The NME for the colour mis-matched \( C_2^{SR-\pm} \) operator will compete with other contributions, and for a charged Higgs mass \( m_{H^+} \sim 500 \) GeV or so will provide the largest contribution. The situation is more interesting in theories with R-parity violation since in such theories, after Fierz arrangement, even without the QCD corrections, there are tensor operators. In such a case, the low scale Wilson coefficients read:

\[
C_1^{SR-\pm} \sim 3C_S + 0.75C_Y, \quad C_2^{SR-\pm} \sim 0.16C_S - 2.6C_M (18)
\]

\[
C_1^{T-\pm} \sim 0.06C_S + 0.74C_Y, \quad C_2^{T-\pm} \sim -0.17C_S + 0.1C_Y
\]

where \( C_S \) and \( C_Y \) denote the effective high scale coefficients of the scalar and tensor operators at the tree level (see [20] for analytic expressions of these). It is very likely that the charged Higgs contribution again dominates once couplings and masses of the particles involved satisfying all the experimental constraints are considered.

In this note, we have shown that the charged Higgs contribution to \( 0 \nu 2\beta \) amplitude which is usually neglected can not be ignored once QCD corrections are taken into account. In fact, the charged Higgs contribution can finally lead to large cancellations among the short range part or can completely overwhelm the other contributions. At any rate, this contribution needs to be properly accounted for in detailed numerical analysis in any model beyond SM. The impact of QCD corrections in this case is rather large and will drastically change the constraints on the model parameters. This will also change the interplay between the limits and constraints obtained from \( 0 \nu 2\beta \) and model studies at LHC and/or other observables. In view of this, it is imperative to revisit the \( 0 \nu 2\beta \) predictions in various model in the light of these corrections and obtain updated constraints, some of which will be totally new and unexpected since at least one new parameter, the charged Higgs mass, will also now need to be considered.

\[\text{References}\]

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