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LOOKING FOR SIGNALS BEYOND THE NEUTRINO STANDARD MODEL*

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Any new neutrino physics at the TeV scale must include a suppression mechanism to keep its contribution to light neutrino masses small enough. We review some seesaw model examples with weakly broken lepton number, and comment on the expected effects at large colliders and in neutrino oscillations.

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1. Introduction

Lepton flavour-changing processes have been only observed in neutrino oscillations [1]. These can be explained by introducing nonzero neutrino masses and the corresponding charged current mixing matrix (MNS) [2] which relates neutrino mass and current eigenstates. This defines the minimal neutrino Standard Model νSM [3], which can be realised with the addition of a Majorana mass term or introducing three light right-handed neutrinos with Yukawa couplings to the SM ones.

The small size of the light neutrino masses, $m_{\nu_i} \sim 1$ eV, makes the observation of neutrino mixing very difficult. In neutrino oscillations the long baseline distance L enhances the small ratio $\Delta m^2_{ij}/E_\nu$, where $\Delta m^2_{ij} =$

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$m^2_{\nu_i} - m^2_{\nu_j}$, making the relevant quantity for neutrino oscillations

$$\frac{\Delta m^2_{ij}\ [\mathrm{eV}^2]\ L\ [\mathrm{km}]}{E_\nu\ [\mathrm{GeV}]} \tag{1}$$

of order unity. However, in high energy collider experiments the available luminosities cannot sufficiently enhance the small mass ratios $m^2_{\nu_i}/E^2$, with E the relevant energy scale in the process, and then lepton flavour violating (LFV) effects are negligible. Hence, the observation of lepton flavour violation at colliders will imply new physics near the TeV scale, which is the scale to be probed at LHC. Conversely, it is also expected that if there is new physics at this scale, it violates lepton flavour because the new interactions do not need to be aligned with the neutrino current eigenstates in general. Any extended model with new neutrino physics near the electroweak scale must include a mechanism for decoupling the generation of light neutrino masses from the physics at the new scale. In Section 2 we discuss how this works in the three types of seesaw. The symmetry protecting light neutrino masses appears to be in all three cases lepton number conservation. We also discuss the limits on the coefficients of the dimension 6 operators parameterising the new physics at low energy. Limits on masses and mixings of heavy neutrinos at large colliders like ILC, CLIC and LHC are presented in Section 3. Section 4 is devoted to new possible effects in neutrino oscillations.

2. Low energy physics

At energies much smaller than the mass of any new resonance, the departures from the SM can be parameterised by an effective Lagrangian, which is determined by the light field content and the required symmetries. The precision is given by the order considered in the momentum expansion. In the case of light neutrinos the effective Lagrangian depends on their Dirac or Majorana nature. In the Dirac case we have to introduce at least three new right-handed neutrinos to pair with the SM left-handed counterparts, and lepton number is in principle conserved. The smallness of SM neutrino masses stems from the smallness of the Yukawa couplings, which requires a satisfactory explanation. In the Majorana case the field content is the same as in the SM, light neutrinos are Majorana particles and lepton number is broken. The smallness of neutrino masses is related to the large scale of this symmetry breaking.

We will concentrate on the second possibility. The most general effective Lagrangian invariant under $\mathrm{SU}(3)_\mathrm{C} \times \mathrm{SU}(2)_\mathrm{L} \times \mathrm{U}(1)_\mathrm{Y}$,

$$\mathcal{L}^{\mathrm{eff}} = \mathcal{L}_4 + \frac{1}{\Lambda}\mathcal{L}_5 + \frac{1}{\Lambda^2}\mathcal{L}_6 + \dots\ , \tag{2}$$

is explicit up to dimension 6 in Ref. [4]. $\mathcal{L}_4$ stands for the SM Lagrangian, $\mathcal{L}_5$ contains the only dimension 5 operator allowed by gauge symmetry[1],

$$\mathcal{O}_5 = \overline{L^c}\tilde{\phi}^*\tilde{\phi}^\dagger L\,, \tag{3}$$

and $\mathcal{L}_6$ includes all dimension 6 operators (81 without taking into account flavour indices) which preserve lepton and baryon number. This Lagrangian is valid for energies below Λ, the cut-off scale. After spontaneous symmetry breaking $\mathcal{O}_5$ generates light neutrino Majorana masses $m_\nu = -x_5 v^2/\Lambda$, being x_5/Λ the coefficient of this dimension 5 operator and $v = 246$ GeV the Higgs vacuum expectation value. For $m_\nu \sim 1$ eV, as required by experimental data, $\Lambda \sim 10^{14}$ GeV if $x_5 \sim 1$, or $x_5 \sim 10^{-11}$ if $\Lambda \sim 1$ TeV. In the first case new physics cannot manifest itself in any high energy experiment considered up to now. In the latter, new effects can show up in the new generation of accelerator experiments if the coefficients of the dimension 6 operators are relatively large. However, in this scenario one has to explain why the coefficient of $\mathcal{O}_5$ is so small. The simplest models including such a decoupling mechanism distinguish between the cut-off scale Λ and the effective lepton number violating (LNV) parameter entering in the definition of x_5 (see for examples Refs. [5,6]). In the rest of this section we partly review the results in Refs. [5, 7], following approximately the notation in Ref. [5]. The different values of the coefficients reflects the different normalisation and the different operator basis used.

The minimal SM extension exhibiting this decoupling only requires the addition of heavy Dirac neutrino singlets N. In this case lepton number can be assigned so that left-handed fields $N_{\rm L}$ have quantum number $\boldsymbol{q}$ and $-\boldsymbol{q}$ the right-handed counterparts $N_{\rm R}^{\rm c}$. Then, the generic mass matrix reduces to

$$\begin{array}{cc} & \begin{array}{cc} \nu_{\rm L} & N \end{array} \\ \begin{array}{c} \nu_{\rm L} \\ N \end{array} & \begin{pmatrix} 0 & Y_N^T \frac{v}{\sqrt{2}} \\ Y_N \frac{v}{\sqrt{2}} & M_N \end{pmatrix} \end{array} \longrightarrow \begin{array}{cc} & \begin{array}{ccc} \nu_{\rm L} & N_{\rm L} & N_{\rm R}^{\rm c} \end{array} \\ \begin{array}{c} \nu_{\rm L} \\ N_{\rm L} \\ N_{\rm R}^{\rm c} \end{array} & \begin{pmatrix} 0 & 0 & \frac{y_N v}{\sqrt{2}} \\ 0 & 0 & m_N \\ \frac{y_N v}{\sqrt{2}} & m_N & 0 \end{pmatrix} \end{array} \tag{4}$$

for one family, where y_N is the Yukawa coupling between the SM neutrino and the right-handed one. If $y_N \neq 0$, $N_{\rm L}$ and $\nu_{\rm L}$ have the same lepton number, $\boldsymbol{q} = 1$, and they mix. When lepton number is broken by a small entry μ instead of some of the zeroes in the above matrix, the light neutrino gets a Majorana mass proportional to it, even if the nonzero entry is in the

[1] We use the operator basis of Ref. [4]. L stands for the lepton doublet and $\tilde{\phi} = i\sigma_2\phi^*$ is the Higgs doublet with hypercharge $Y = -1/2$. Family indices are not shown, unless otherwise stated.

$(3,3)$ position because one-loop radiative corrections also generate a nonzero mass for $\nu_{\rm L}$ proportional to μ. (A similar behaviour is found in Little Higgs models [8].)

More generally, all three types of seesaw mechanisms generating $\mathcal{O}_5$ at the tree level upon integration of heavy fields [9] can incorporate a similar decoupling. In Tables I–III we collect the operators up to dimension 6 obtained from the integration of heavy fermion singlets N (type I seesaw), scalar triplets Δ (type II) and fermion triplets Σ (type III), respectively, and the corresponding coefficients [5], where now Λ is the mass of the heavy resonance. In type II seesaw the coefficient of $\mathcal{O}_5$ is explicitly proportional to the LNV product $\mu_\Delta Y_\Delta$, while none of the other coefficients contains both

TABLE I

Operators arising from the integration of heavy Majorana fermion singlets N. Y_N is the coupling matrix in the Yukawa term $-\overline{L}\tilde{\phi}Y_N^\dagger N_{\rm R}$.

Dimension	Operator	Coefficient
5	$\mathcal{O}_5 = \overline{L^c}\tilde{\phi}^*\tilde{\phi}^\dagger L$	$\frac{1}{2}Y_N^T M_N^{-1} Y_N$
6	$\mathcal{O}_{\phi L}^{(1)} = \left(\phi^\dagger i D_\mu \phi\right)\left(\overline{L}\gamma^\mu L\right)$	$\frac{1}{4}Y_N^\dagger (M_N^\dagger)^{-1} M_N^{-1} Y_N$
	$\mathcal{O}_{\phi L}^{(3)} = \left(\phi^\dagger i \sigma_a D_\mu \phi\right)\left(\overline{L}\sigma_a\gamma^\mu L\right)$	$-\frac{1}{4}Y_N^\dagger (M_N^\dagger)^{-1} M_N^{-1} Y_N$

TABLE II

Operators arising from the integration of heavy scalar triplets Δ. Y_Δ is the coupling matrix in the Yukawa term $\tilde{\overline{L}} Y_\Delta \left(\vec{\sigma}\cdot\vec{\Delta}\right) L$, with $\tilde{\overline{L}} = -L^T C i\sigma_2$ and C the matrix entering the spinor charge conjugation definition; and μ_Δ, λ_3 and λ_5 are the coefficients of the scalar potential terms $\tilde{\phi}^\dagger\left(\vec{\sigma}\cdot\vec{\Delta}\right)^\dagger\phi$, $-\left(\phi^\dagger\phi\right)\left(\vec{\Delta}^\dagger\vec{\Delta}\right)$ and $-\left(\vec{\Delta}^\dagger T_i \vec{\Delta}\right)\phi^\dagger\sigma_i\phi$, respectively.

Dimension	Operator	Coefficient
4	$\mathcal{O}_4 = \left(\phi^\dagger\phi\right)^2$	$2\left\|\mu_\Delta\right\|^2/M_\Delta^2$
5	$\mathcal{O}_5 = \overline{L^c}\tilde{\phi}^*\tilde{\phi}^\dagger L$	$-2\,Y_\Delta\mu_\Delta/M_\Delta^2$
6	$\mathcal{O}_{LL}^{(1)} = \frac{1}{2}\left(\overline{L^i}\gamma^\mu L^j\right)\left(\overline{L^k}\gamma_\mu L^l\right)$	$2/M_\Delta^2 (Y_\Delta)_{jl}(Y_\Delta^\dagger)_{ki}$
	$\mathcal{O}_\phi = \frac{1}{3}\left(\phi^\dagger\phi\right)^3$	$-6\left(\lambda_3+\lambda_5\right)\left\|\mu_\Delta\right\|^2/M_\Delta^4$
	$\mathcal{O}_\phi^{(1)} = \left(\phi^\dagger\phi\right)\left(D_\mu\phi\right)^\dagger D^\mu\phi$	$4\left\|\mu_\Delta\right\|^2/M_\Delta^4$
	$\mathcal{O}_\phi^{(3)} = \left(\phi^\dagger D_\mu\phi\right)\left(D^\mu\phi^\dagger\phi\right)$	$4\left\|\mu_\Delta\right\|^2/M_\Delta^4$

TABLE III

Operators arising from the integration of heavy Majorana fermion triplets Σ. Y_Σ is the coupling matrix in the Yukawa term $-\overline{\vec{\Sigma}_{\rm R}}Y_\Sigma(\tilde{\phi}^\dagger\vec{\sigma}L)$ and Y_l in $-\overline{L}Y_l\phi l_{\rm R}$.

Dimension		Operator	Coefficient
5	$\mathcal{O}_5$	$= \overline{L^c}\tilde{\phi}^*\tilde{\phi}^\dagger L$	$\frac{1}{2}Y_\Sigma^T M_\Sigma^{-1} Y_\Sigma$
6	$\mathcal{O}_{\phi L}^{(1)}$	$= \left(\phi^\dagger i D_\mu \phi\right)\left(\overline{L}\gamma^\mu L\right)$	$\frac{3}{4}Y_\Sigma^\dagger (M_\Sigma^\dagger)^{-1} M_\Sigma^{-1} Y_\Sigma$
	$\mathcal{O}_{\phi L}^{(3)}$	$= \left(\phi^\dagger i\sigma_a D_\mu \phi\right)\left(\overline{L}\sigma_a\gamma^\mu L\right)$	$\frac{1}{4}Y_\Sigma^\dagger (M_\Sigma^\dagger)^{-1} M_\Sigma^{-1} Y_\Sigma$
	$\mathcal{O}_{l\phi}$	$= \left(\phi^\dagger \phi\right)\overline{L}\phi l_{\rm R}$	$Y_\Sigma^\dagger (M_\Sigma^\dagger)^{-1} M_\Sigma^{-1} Y_\Sigma Y_l$

parameters. This allows for a relatively light scalar triplet with $M_\Delta \sim 1$ TeV and possibly observable effects at forthcoming experiments, while keeping SM neutrino masses very small (in definite models [6] there can be also extra loop suppression factors). In the other two types of seesaw the decoupling is not so explicit. In both cases the coefficient of the dimension 5 operator is proportional to $Y^T M^{-1} Y$, thus it only depends (quadratically) on Y, while the coefficients of the dimension 6 operators involve $Y^\dagger (M^\dagger)^{-1} M^{-1} Y$, with Y and $Y^\dagger$. In this way it is possible that there are cancellations in the former product which do not hold in the latter one. This is indeed what happens for quasi-Dirac neutrinos. For our one-family example in Eq. (4), if the LNV parameter μ is in the (2,2) position, the SM neutrino acquires a Majorana mass m_ν [2]

$$-Y_N^T M_N^{-1} Y_N \frac{v^2}{2} \simeq -\frac{y_N^2}{2}\left[\frac{(1-\frac{\mu}{4m_N})^2}{m_N+\frac{\mu}{2}} - \frac{(1+\frac{\mu}{4m_N})^2}{m_N-\frac{\mu}{2}}\right]\frac{v^2}{2} \simeq \frac{\mu y_N^2}{m_N^2}\frac{v^2}{2}, \quad (5)$$

where we only keep the dominant terms in μ/m_N. (Less natural cancellations are also possible in more involved models [7,10].) While m_ν is proportional to μ, the coefficients of the dimension 6 operators are not,

$$Y_N^\dagger (M_N^\dagger)^{-1} M_N^{-1} Y_N \simeq \frac{|y_N|^2}{2}\left[\frac{(1-\frac{\mu}{4m_N})^2}{(m_N+\frac{\mu}{2})^2} + \frac{(1+\frac{\mu}{4m_N})^2}{(m_N-\frac{\mu}{2})^2}\right] \simeq \frac{|y_N|^2}{m_N^2}\,. \quad (6)$$

Hence, new fermions can exist near the TeV scale with observable effects beyond the SM in future experiments, while maintaining the SM neutrinos light enough.

[2] The 2×2 bottom-right submatrix must be diagonalised before applying the seesaw formula in order to make the cancellation apparent. The masses of the two Majorana eigenstates are taken to be positive, $m_{N_1} \simeq m_N + \mu/2$, $m_{N_2} \simeq m_N - \mu/2$.

Present experimental limits on the different combinations of quadratic products of Yukawa couplings y^*y entering the dimension 6 operators range from 0.3 to 0.002 for a heavy neutrino singlet N with a mass of 1 TeV; from 1 to 10^{-5} for a heavy scalar triplet Δ of the same mass, and from 0.01 to 3×10^{-5} for a heavy fermion triplet Σ equally heavy. A detailed analysis can be found in Ref. [5].

3. Lepton signals at large colliders

The next generation of large colliders will be able to further constrain the masses and mixings of the seesaw messengers (see Ref. [11] for a review in the case of heavy neutrino singlets). Here we restrict ourselves to e^+e^- and hadron colliders.

3.1. e^+e^- colliders

The process $e^+e^- \to N\nu \to \ell^\pm W^\mp(\to q\bar{q}')\nu$ sets the most stringent limits on the mass and the mixing of a heavy neutrino singlet (seesaw type I) for a large enough center of mass energy so that N is produced [12] (see also Ref. [13]). Lepton colliders are a rather clean environment, being the irreducible background for this process the SM four-fermion $\ell\nu q\bar{q}'$ production (which includes W^+W^- plus non-resonant diagrams). The non-observation of an excess in the ℓjj invariant mass distribution will set limits on the heavy neutrino mass m_N and its mixing with the charged leptons $V_{\ell N} = Y^*_{N\ell}v/(\sqrt{2}m_N)$, $\ell = e, \mu, \tau$. Limits are rather independent of m_N up to nearly the kinematical limit, and independent of the Dirac or Majorana character of the heavy neutrino. In Fig. 1 we show the combined limits on the mixing of a new heavy neutrino singlet *(i)* at ILC, with a centre of mass energy $\sqrt{s} = 500$ GeV and an integrated luminosity $L = 345$ fb^{-1}, taking $m_N = 300$ GeV; *(ii)* at CLIC, with $\sqrt{s} = 3$ TeV, $L = 1000$ fb^{-1}, and taking $m_N = 1.5$ TeV.

3.2. Hadron colliders

Hadron colliders produce large electroweak signals, and, in particular, they can produce new leptons with relatively large cross sections. If the usually huge SM backgrounds contribute relatively little to a specific final state, one can derive non-trivial limits on these new leptons. This is the case of heavy neutrino singlets (seesaw type I) [16–18] in like-sign dilepton final states $\ell^\pm\ell'^\pm X$. Let us summarise the analysis of Ref. [19]. At hadron colliders the heavy neutrino character plays an important role because Dirac neutrinos conserve lepton number and, in general, their signals are overwhelmed by the backgrounds. On the other hand, heavy Majorana

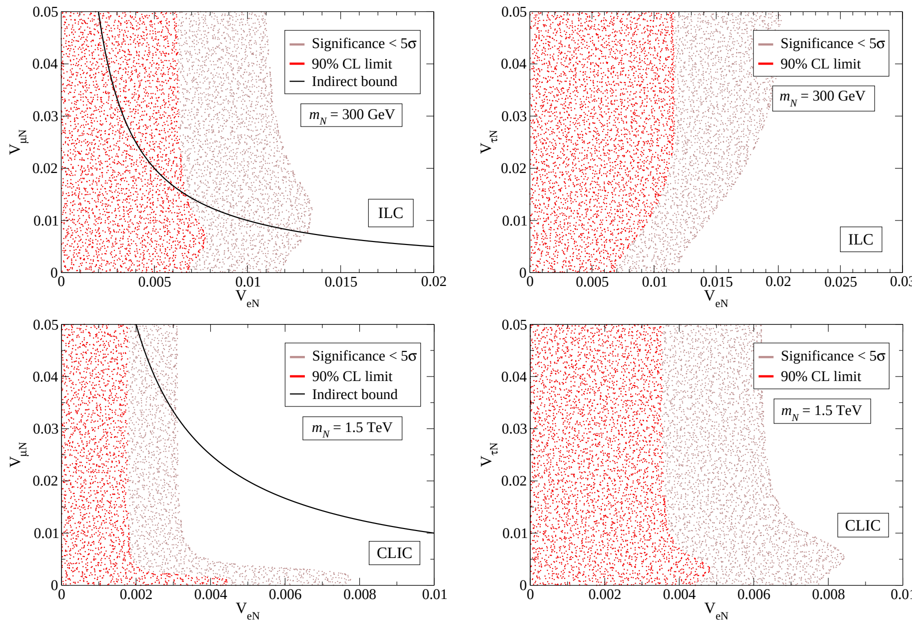

Fig. 1. Combined limits on heavy neutrino mixings at ILC (up) and CLIC (down), for the cases $V_{\tau N} = 0$ (left) and $V_{\mu N} = 0$ (right). The coloured (red) areas represent the 90% CL limits if no signal is observed. The white areas extend up to present bounds $V_{eN} \leq 0.073$, $V_{\mu N} \leq 0.098$, $V_{\tau N} \leq 0.13$ [14, 15], and correspond to the region where a combined statistical significance of 5σ or larger is achieved. The indirect limit from μ–e LFV processes is also shown.

neutrinos produce LNV signals, $p\overset{(-)}{p} \to W^\pm \to \ell^\pm N \to \ell^\pm \ell'^\pm jj$, which have smaller backgrounds, and present limits on their masses and mixings can be eventually improved. (However, realising these masses and mixings in a specific model still requires complicated cancellations to avoid generating too large SM neutrino masses, as emphasised in the former section.) At Tevatron the signal cross sections are in practice too small, but at LHC they are sizeable for heavy neutrino masses of the order of the electroweak scale (and especially for $m_N < M_W$, when the heavy neutrino is produced on its mass shell). The limits on the mixing of a heavy Majorana neutrino are plotted in Fig. 2 for the case $V_{\tau N} = 0$ and two heavy neutrino masses above and below M_W, for a luminosity $L = 30$ fb^{-1}. We point out that for $m_N = 60$ GeV the direct limit is more stringent than the indirect one from $\mu - e$ LFV processes.

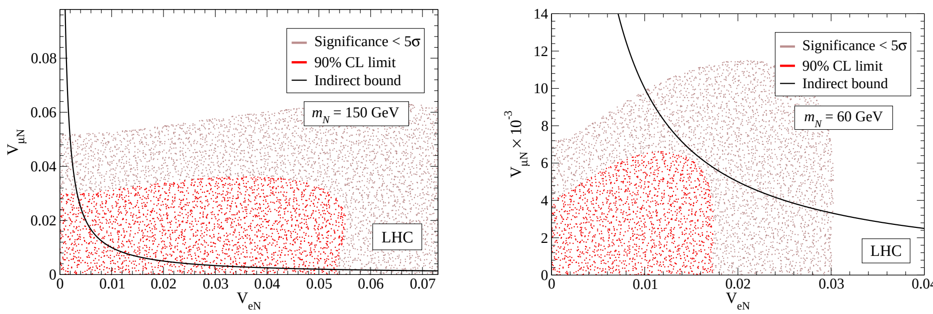

Fig. 2. Combined limits on heavy neutrino mixings at LHC for $V_{\tau N} = 0$ and two heavy Majorana neutrino masses. The meaning of the coloured areas is the same as in Fig. 1.

A Dirac neutrino does not give observable signals at LHC except if N is lighter than the W boson and couples to both electron and muon. In this situation it can produce the LFV signal $e^\pm \mu^\mp X$ with a large cross section, so that it can be observed above the large opposite sign dilepton background. In Fig. 3 we show the corresponding limits on the heavy Dirac neutrino mixings for $m_N = 60$ GeV and $L = 30$ fb^{-1}.

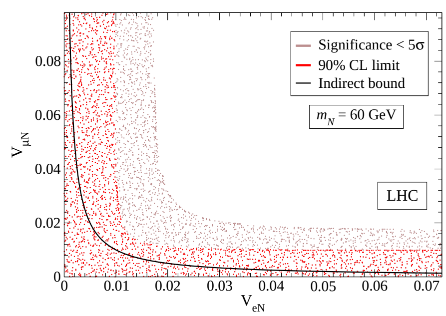

Fig. 3. The same as in Fig. 2 but for a Dirac neutrino with a mass $m_N = 60$ GeV.

Heavy neutrino limits improve significantly in the presence of new interactions, for example of a new W' [20] or a Z' [21]. In the former case, LHC is sensitive to masses up to $M_{W'} = 3$ TeV, $m_N = 2.1$ TeV [22], while in the latter it is sensitive to $M_{Z'} = 2.5$ TeV, $m_N = 800$ GeV [21] (in both cases assuming $L = 30$ fb^{-1}).

Finally, we point out that like-sign dilepton signals also arise in the other two seesaw scenarios: in the production of doubly charged scalar triplets [23], and in pair production of fermion triplets [24]. For this reason, like-sign dileptons constitute an interesting final state in which to test seesaw at LHC.

4. Neutrino oscillations beyond the νSM

Neutrino oscillation experiments will improve their precision in the future, and they may be sensitive to new physics through its effects on light neutrinos. For example, deviations from unitarity of the MNS matrix due to mixing with heavy neutrinos can manifest at the percent level in ν_μ–ν_τ transitions [15]. In the presence of new right-handed interactions, the transition probability amplitude differs if light neutrinos have Dirac or Majorana nature, as it is shown in Fig. 4 [25]. The difference (dashed line) can be at the 10% level but only for very long baseline distance L. For the examples shown, it reduces by a factor of 4 from $L = 13000$ km to $L = 6500$ km. It is also proportional to the strength of the new four-fermion interaction.

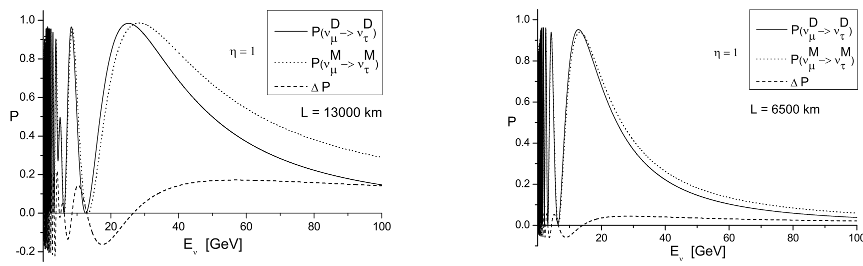

Fig. 4. Transition probabilities for Majorana (M) and Dirac (D) neutrinos and their difference ΔP as a function of the neutrino energy E_ν (in GeV) for two different baseline distances L. The new four-fermion interactions have a strength which is 1% ($\eta = 1$) of the weak interactions [26]. Note that in the Dirac case the transition amplitude with new right-handed interactions is the same as in the νSM.

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