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# Loop bounds on non-standard neutrino interactions

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Abstract: We reconsider the bounds on non-standard neutrino interactions with matter which can be derived by constraining the four-charged-lepton operators induced at the loop level. We find that these bounds are model dependent. Naturalness arguments can lead to much stronger constraints than those presented in previous studies, while no completely model-independent bounds can be derived. We will illustrate how large loop-contributions to four-charged-lepton operators are induced within a particular model that realizes gauge invariant non-standard interactions and discuss conditions to avoid these bounds. These considerations mainly affect the $\mathcal{O}(10^{-4})$ constraint on the non-standard coupling strength $\varepsilon_{e\mu}$, which is lost. The only model-independent constraints that can be derived are $\mathcal{O}(10^{-1})$. However, significant cancellations are required in order to saturate this bound.

Keywords: Beyond Standard Model, Neutrino Physics

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## Contents





## 1 Introduction

Neutrino non-standard interactions (NSI) were originally proposed [1] as a mechanism to produce neutrino flavour conversion in matter and considered as a possible explanation for the solar and atmospheric neutrino deficits [1–9]. Although it is now clear that such NSI cannot fully account for the observed neutrino flavour conversion, the increasing sensitivity of neutrino oscillation experiments to sub-leading effects has triggered a new interest in them and their interference with neutrino oscillations at present (e.g., K2K, MINOS, OPERA [10–17]) and future (e.g., SuperBeams, $\beta$Beams or Neutrino Factories [18–30]) facilities. In particular, the determination of the leptonic mixing angle $\theta_{13}$ could be severely affected by degeneracies with the non-standard parameters [31–33].

Neutrino NSI can be described by effective four-fermion operators of the form

$$\mathcal{L}_{\rm NSI} = -2\sqrt{2}G_F \varepsilon^{ff' L,R}_{\alpha\beta} (\overline{\nu_\alpha}\gamma^\mu \nu_\beta)(\overline{f_{L,R}}\gamma_\mu f'_{L,R})\,, \tag{1.1}$$

where $f$ and $f'$ are charged fermions with the same quantum numbers and $L, R$ represent the chirality. For the rest of the paper we will denote with $\nu$ the left-handed neutrinos, $\ell$ the left-handed charged leptons, $L$ the left-handed lepton doublets and $E$ the right-handed charged leptons. For NSI of neutrinos with normal matter, $f = f'$ can be either an electron, an up-quark or a down-quark. While strong experimental bounds are present on the corresponding four-charged-fermion interactions, $(\overline{\ell_\alpha}\gamma^\mu\ell_\beta)(\overline{f_{L,R}}\gamma_\mu f'_{L,R})$, neutrino NSI are much less constrained. Constraints on them have been derived in [34–36] making use of both tree level and one-loop processes. Here we reconsider the one-loop bounds, focusing on the necessity of implementing the neutrino NSI in a gauge invariant way in order to obtain a gauge independent result.

The simple promotion of the neutrino fields in eq. (1.1) to lepton doublets in order to construct a gauge invariant operator would imply that the strong bounds stemming

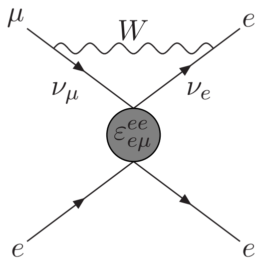


**Figure 1**. One-loop contribution to the four-charged-fermion vertex arising from the operator of eq. (1.1) via $W$ exchange.

from flavour violating four-charged-fermion processes will apply to neutrino interactions as well. In order to avoid this and allow large NSI, the simultaneous presence of tree level flavour violating four-charged-fermion interactions must be forbidden [37–39]. However, even if this requirement is satisfied, these interactions can be generated at one loop from the operator of eq. (1.1) via a $W$ exchange between the neutrino legs. In [34] this has been exploited to set bounds on some $\varepsilon$. Notably, an $\mathcal{O}(10^{-4})$ bound on $\varepsilon^{ff}_{e\mu}$ was derived through loop contributions to the decay $\mu \to 3e$ and the $\mu - e$ conversion in nuclei. Consequently, the non-standard coupling strength $\varepsilon^{ff}_{e\mu}$ was neglected in a large number of studies (see for example [12–16, 29, 40–44]). However, as we will show, when gauge invariance is imposed, only naturalness arguments can be invoked. These arguments can lead to even stronger constraints on the NSI strength, but no completely model-independent bounds can be derived. The relevant diagram contributing to $\mu \to 3e$ is depicted in figure 1. The computation of this diagram using only the operator of eq. (1.1) with $f = f' = e$ renders a gauge dependent result due to the fact that the operator itself is not gauge invariant. Thus, the necessity of considering a gauge invariant formulation of NSI is manifest. In the following, we will consider gauge invariant realisations of NSI both when the operator of eq. (1.1) is realised from a dimension-six ($d = 6$) and from a $d = 8$ operator. We will discuss here the NSI with charged leptons, i.e., $f = f' = \ell$. However, similar arguments as those presented here are applicable in the case of neutrino NSI with quarks. While gauge invariance has to be carefully taken care of in the cases mentioned above, one-loop processes like the $W$ and $Z$ decays (such as the ones depicted in figures 5-7 of [34]) do not contain gauge boson propagators in the loops and the bounds derived from them still apply.

## 2 $d = 6$ realisations

### 2.1 Flavour antisymmetric operator

There is only one gauge invariant $d = 6$ effective operator that can induce the effective interaction of eq. (1.1) in a direct way while avoiding the generation of four-charged-fermion operators. This operator is of the form:

$$\mathcal{O}^a_6 = (\bar{L}_\gamma i\tau_2 L^c_\alpha)(\overline{L^c_\beta} i\tau_2 L_\delta)\,, \tag{2.1}$$



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where $L^c = C\bar{L}^T$ and $C$ is the charge conjugation operator. It can be generated upon integrating out a heavy charged scalar singlet with lepton number violating couplings to the lepton doublets, as in [45–48]. This operator induces NSI only for leptons and with a very characteristic flavour structure:[1]

$$2\mathcal{O}_6^a = (\bar{\ell}_\alpha\gamma^\mu\ell_\beta)(\bar{\nu}_\gamma\gamma_\mu\nu_\delta)+(\bar{\ell}_\gamma\gamma^\mu\ell_\delta)(\bar{\nu}_\alpha\gamma_\mu\nu_\beta)-(\bar{\ell}_\alpha\gamma^\mu\ell_\delta)(\bar{\nu}_\gamma\gamma_\mu\nu_\beta)-(\bar{\ell}_\gamma\gamma^\mu\ell_\beta)(\bar{\nu}_\alpha\gamma_\mu\nu_\delta)\,. \quad (2.2)$$

This structure implies the antisymmetry relations

$$\varepsilon^{\alpha\beta}_{\gamma\delta} = -\varepsilon^{\gamma\beta}_{\alpha\delta} = -\varepsilon^{\alpha\delta}_{\gamma\beta} = \varepsilon^{\gamma\delta}_{\alpha\beta}\,. \quad (2.3)$$

Thus, for each process of the type $\ell \to 3\ell'$ (i.e., a decay of a heavy lepton into three lighter leptons), there will be four contributing diagrams at one loop. For example, the process $\tau^- \to \mu^-\mu^+e^-$ will receive contributions from $\varepsilon^{\mu\tau}_{e\mu}$, $\varepsilon^{e\mu}_{\mu\tau}$, $\varepsilon^{\mu\mu}_{e\tau}$ and $\varepsilon^{e\tau}_{\mu\mu}$. In the approximation of massless fermions, using the antisymmetry relations of eq. (2.3), the four contributions exactly cancel. Thus, any loop contribution from $d = 6$ operators to the four-charged-fermion interactions will be suppressed at least by a factor $\mathcal{O}(m_\ell^2/M_W^2)$, where $m_\ell$ is the mass of the heaviest lepton involved in the process. As a consequence, the bound on the corresponding $\varepsilon$ will be increased by the inverse of this factor. However, the antisymmetry relations of eq. (2.3) also imply that NSI with electrons that change the neutrino flavour are related to charged lepton flavour changing interactions, and thus stringent constraints can be derived in this case (see [38, 49]). Note that $\mu \to 3e$ is forbidden for $d = 6$ operators, since there are no $\varepsilon$ with the given symmetries which can contribute to this process.

### 2.2 Non-standard interactions from non-unitarity

The second possibility to generate neutrino NSI avoiding four-charged-fermion interactions is in an indirect way via the dimension six operator

$$\mathcal{O}_6^{\rm kin} = -(\bar{L}_\alpha\tilde{H})\, i\partial\!\!\!/ (\tilde{H}^\dagger L_\beta)\,, \quad (2.4)$$

where $H$ is the Higgs doublet (we choose the hypercharge of $H$ to be 1/2) and $\tilde{H} = i\tau_2 H^*$. This operator induces non-canonical neutrino kinetic terms. After diagonalising and normalising them, a non-unitary leptonic mixing matrix is produced and, upon integrating out the $W$ and $Z$ bosons, neutrino NSI are induced. The tree level generation of this operator, avoiding similar contributions to charged leptons (see, e.g., [50]), involves the addition of Standard Model-singlet fermions (right-handed neutrinos) which couple to the Higgs and lepton doublets via Yukawa couplings like in the standard seesaw model [51–54].

This second realisation of NSI at $d = 6$ is also quite constrained due to the effects of a non-unitary mixing matrix (see [38, 55, 56]). The bounds on the $\varepsilon_{e\mu}$ element are particularly strong due to the enhancement of the $\mu \to e\gamma$ process if the GIM mechanism is not realised given the non-unitarity of the mixing matrix. In [38], a bound of $|\varepsilon_{e\mu}| < 5.9\times10^{-5}|n_n/n_e - 1|$ for $M_{N_R} > M_W$ or $|\varepsilon_{e\mu}| < 9.1\times10^{-4}|n_n/n_e - 1|$ for $M_{N_R} < M_W$ was computed, where $n_n$ ($n_e$) is the neutron (electron) density in matter. Notice that, since $n_n \simeq n_e$, this means an additional suppression of $\mathcal{O}(10^{-2})$. These bounds are already much stronger than the loop bounds discussed here and thus we will not consider this possibility further.

[1] NSI with right-handed fields avoiding four-charged-fermion interactions can be realised through higher-dimensional operators as we will discuss in the next section.

## 3 $d = 8$ realisations

Dimension eight realisations of the NSI offer more freedom. Gauge invariant $d = 8$ operators can be generated by adding two Higgs doublets to the four-lepton operators. The vev of the Higgs field can then be exploited to break the SU(2) symmetry between the charged leptons and the neutrinos and can thus induce neutrino NSI avoiding their charged-lepton counterparts. A basis for operators involving two Higgs doublets, two left-handed lepton doublets and two fermions, either left or right-handed (for the matter components), is given in [37, 39]:

$$\mathcal{O}^{\mathbf{111}}_{LLH} = (\bar{L}_\beta \gamma^\rho L_\alpha)(\bar{L}_\delta \gamma_\rho L_\gamma)\left(H^\dagger H\right), \tag{3.1}$$

$$\mathcal{O}^{\mathbf{331}}_{LLH} = (\bar{L}_\beta \gamma^\rho \vec{\tau} L_\alpha)(\bar{L}_\delta \gamma_\rho \vec{\tau} L_\gamma)\left(H^\dagger H\right), \tag{3.2}$$

$$\mathcal{O}^{\mathbf{133}}_{LLH} = (\bar{L}_\beta \gamma^\rho L_\alpha)(\bar{L}_\delta \gamma_\rho \vec{\tau} L_\gamma)\left(H^\dagger \vec{\tau} H\right), \tag{3.3}$$

$$\mathcal{O}^{\mathbf{313}}_{LLH} = (\bar{L}_\beta \gamma^\rho \vec{\tau} L_\alpha)(\bar{L}_\delta \gamma_\rho L_\gamma)\left(H^\dagger \vec{\tau} H\right), \tag{3.4}$$

$$\mathcal{O}^{\mathbf{333}}_{LLH} = (-\mathrm{i}\epsilon^{abc})(\bar{L}_\beta \gamma^\rho \tau^a L_\alpha)(\bar{L}_\delta \gamma_\rho \tau^b L_\gamma)\left(H^\dagger \tau^c H\right), \tag{3.5}$$

$$\mathcal{O}^{\mathbf{111}}_{LEH} = (\bar{L}_\beta \gamma^\rho L_\alpha)(\bar{E}_\delta \gamma_\rho E_\gamma)\left(H^\dagger H\right), \tag{3.6}$$

$$\mathcal{O}^{\mathbf{313}}_{LEH} = (\bar{L}_\beta \gamma^\rho \vec{\tau} L_\alpha)(\bar{E}_\delta \gamma_\rho E_\gamma)\left(H^\dagger \vec{\tau} H\right). \tag{3.7}$$

In addition to the two left-handed lepton doublets, the two last operators contain two right-handed charged leptons and the first five two additional left-handed lepton doublets. The generalisation to operators involving interactions with quarks is straightforward replacing these fields by their quark counterparts. Generically, after electro-weak symmetry breaking (EWSB), these operators generate both neutrino NSI and non-standard four-charged-fermion interactions at tree level. In order to avoid the latter, the following conditions have to be met [39]:[2]

$$\mathcal{C}^{\mathbf{111}}_{LEH} = -\mathcal{C}^{\mathbf{313}}_{LEH}\,, \quad \mathcal{C}^{\mathbf{111}}_{LLH} + \mathcal{C}^{\mathbf{331}}_{LLH} + \mathcal{C}^{\mathbf{133}}_{LLH} + \mathcal{C}^{\mathbf{313}}_{LLH} = 0\,, \quad \mathcal{C}^{\mathbf{333}}_{LLH} \text{ arbitr.}. \tag{3.8}$$

The second condition involves four operators of the basis and we can thus choose three independent combinations satisfying it in order to form a basis for the $d = 8$ operators that induce neutrino NSI avoiding four-charged-lepton interactions:

$$\mathcal{A} = \frac{1}{4}(\mathcal{O}^{\mathbf{111}}_{LLH} - \mathcal{O}^{\mathbf{331}}_{LLH}) = (\overline{L^c_\alpha} i\tau_2 L_\gamma)(\bar{L}_\delta i\tau_2 L^c_\beta)(H^\dagger H)\,, \tag{3.9}$$

$$\mathcal{B} = (\mathcal{O}^{\mathbf{111}}_{LLH} - \mathcal{O}^{\mathbf{133}}_{LLH}) = 2(\bar{L}_\beta \gamma_\rho L_\alpha)(\bar{L}_\delta \tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\gamma)\,, \tag{3.10}$$

$$\mathcal{C} = (\mathcal{O}^{\mathbf{111}}_{LLH} - \mathcal{O}^{\mathbf{313}}_{LLH}) = 2(\bar{L}_\beta \tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\alpha)(\bar{L}_\delta \gamma_\rho L_\gamma)\,, \tag{3.11}$$

$$\mathcal{D} = \mathcal{O}^{\mathbf{333}}_{LLH} = 2(\bar{L}_\beta \gamma_\rho L_\gamma)(\bar{L}_\delta \tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\alpha) - 2(\bar{L}_\beta \tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\gamma)(\bar{L}_\delta \gamma_\rho L_\alpha)\,, \tag{3.12}$$

$$\mathcal{E} = (\mathcal{O}^{\mathbf{111}}_{LEH} - \mathcal{O}^{\mathbf{313}}_{LEH}) = 2(\bar{L}_\beta \tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\alpha)(\bar{E}_\delta \gamma_\rho E_\gamma)\,. \tag{3.13}$$

All of these operators have to be divided by $M^4$, where $M$ is the new physics scale from which the NSI originate, and multiplied by a flavour-dependent coefficient: $a^{\alpha\beta\gamma\delta}$ multiplies

[2] Notice that our relations differ from those in [39] due to different conventions for the Higgs hyper-charge.

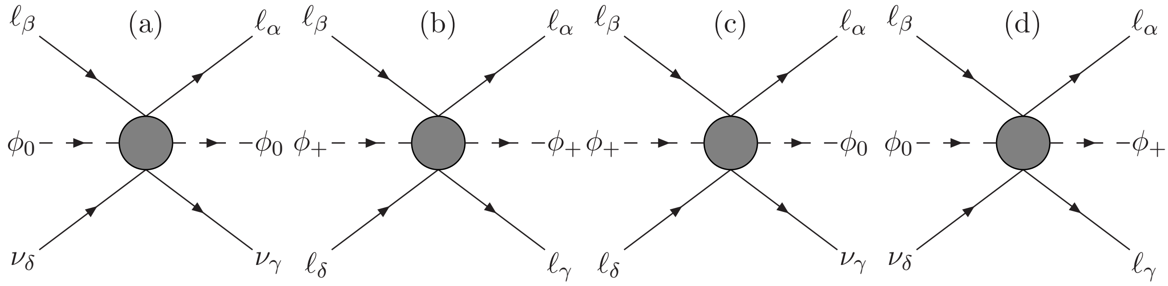


**Figure 2**. The effective interactions induced by $\mathcal{B}$.

$\mathcal{A}$, $b^{\alpha\beta\gamma\delta}$ multiplies $\mathcal{B}$ and so on. Note that, apart from $\mathcal{E}$, which is independent since it is the only one containing right-handed fields, the other operators are all related. Indeed $\mathcal{C}$ is obtained from $\mathcal{B}$ simply by rearranging the flavour indexes, $\mathcal{D}$ is equivalent to $\mathcal{B}-\mathcal{C}$ after a Fierz transformation and again a rearrangement of two indexes and also $\mathcal{A}$ can be written in terms of combination of $\mathcal{B}$ operators with different flavour indexes. Thus, the only new structure that $d = 8$ operators offer to select neutrino NSI avoiding four-charged-fermion operators is the combination $(\bar{L}\tilde{H})\gamma^\rho(\tilde{H}^\dagger L)$ present in eqs. (3.10)–(3.13).

We will therefore compute the loop of figure 1 for this new structure. From the combination $(\bar{L}\gamma_\rho L)(\bar{L}\tilde{H})\gamma^\rho(\tilde{H}^\dagger L)$, "opening" the operators in components, it is easy to check which operators are generated together with the one in eq. (1.1). Taking all of them into account, a gauge invariant computation can be performed. In this case, the generated $\varepsilon$ are completely independent and the effective interactions shown in figure 2 are all generated with the same strength. Out of these diagrams, diagram (a) will result in the effective neutrino NSI when the Higgs acquires a vev. There are now several one-loop contributions to the four-charged-lepton vertex. In diagram (a), we can connect the neutrino lines with either a $W$ or a $\phi_+$, resulting in a conversion of the neutrinos into charged leptons. For diagrams (c) and (d), we can connect the charged Goldstones to the neutrino lines, and finally, for diagram (b) we can close the $\phi_+$-loop. It is now possible to check the gauge independence of the result explicitly. We have done this by performing the computation in the $R_\xi$ gauge and splitting the $W$ propagator in the unitary gauge part and the $\xi$-dependent part and checked that the gauge dependence introduced by the $W$ and all the diagrams with the Goldstones cancels, as it should. The remaining gauge independent contribution is then the one from the $W$ exchange with the propagator in the unitary gauge.

The most widespread way of estimating bounds from loop processes of an unknown high energy theory through its effective description is exploiting the logarithmic divergence, as in [34]. Indeed the coefficient of this term and that of the logarithmic running of the full theory are the same and the mild scale dependence is just an $\mathcal{O}(1)$ correction. On the other hand, the finite and quadratic contributions of the effective theory are less reliable, since they depend on the matching with the unknown full theory. Moreover, these contributions can be fine-tuned away, while for the logarithmic running this can only be true at a given scale.[3]

[3] Unless the cancellation is implemented through an operator with exactly the same running.



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We find that, neglecting lepton masses, the logarithmic divergences coming from the two parts of the $W$-propagator exactly cancel at one loop.[4] Then, also in the case of this $d = 8$ operator, only very weak bounds on $\varepsilon$ can be derived through the logarithmic divergence using the decay of a heavy lepton into three lighter leptons. It should be noted that, even if the logarithmic divergence is not present, a quadratic divergence is. We will devote the next subsection to the physical interpretation of this quadratic divergence.

We argue that the discussion presented here can be also applied to higher-dimensional operators where only Higgs doublets are added. In order to preserve gauge invariance in the NSI, it is necessary to include the effects of the internal Goldstone loops. In order to avoid specifying the underlying theory, we must therefore compute the loop diagrams of the effective theory in the unitary gauge, where the Goldstone propagators vanish.

### 3.1 The quadratic divergence

The computation of the loop with a $W$ exchange between the neutrino legs of diagram (a) of figure 2 in the unitary gauge turns out to give a vanishing logarithmic contribution, while a quadratic divergence is present. It is easy to check that a similar diagram with a $Z$ exchange plus a diagram where the physical Higgses of diagram (a) of figure 2 are closed in a loop give exactly the same quadratic divergence to the neutrino NSI operator.

It is simpler to understand the origin of this quadratic contribution to both operators from the decomposition into singlet minus triplet of eqs. (3.10)–(3.13):

$$(\bar{L}\tilde{H})\gamma^\rho(\tilde{H}^\dagger L) = \frac{1}{2}[(\bar{L}\gamma^\rho L)(H^\dagger H) - (\bar{L}\gamma^\rho \vec{\tau} L)(H^\dagger \vec{\tau} H)]\,. \tag{3.14}$$

If the Higgs legs are closed in a loop, the resulting tadpole will contribute with a quadratic divergence only through the singlet term. The corresponding divergence is not present for the triplet term since it involves a trace of the $\tau$ matrices, which vanishes. This means that the quadratic divergence will be proportional to the singlet combination $(\bar{L}\gamma^\rho L)$ and thus will be equal for the neutrino NSI and the four-charged-fermion interaction. Since the logarithmic contributions to the process vanish, we can now try to estimate the constraints that this quadratic divergence implies. We will consider these quadratic contributions for the operator of eq. (3.13) involving the right-handed leptons $E$, but similar discussions also apply to eqs. (3.10)–(3.12), as well as to interactions with quarks, while the contributions cancel for eq. (3.9). However, as we will show later, for the left-handed fields antisymmetries similar to the ones discussed at $d = 6$ can be considered in order to avoid the bounds.

At tree level, after EWSB, the operator $\mathcal{E}$ generates neutrino NSI with strength

$$\varepsilon^{\delta\gamma}_{\beta\alpha} = -\frac{e^{\beta\alpha\delta\gamma}}{2}\frac{v^4}{M^4}\,. \tag{3.15}$$

On the other hand, if we parametrise the four-charged-lepton interactions as

$$\mathcal{L}_{\rm CF} = -2\sqrt{2}G_F\varepsilon^{\beta\alpha,\gamma\delta}_{CF,R}(\overline{\ell_\beta}\gamma_\mu\ell_\alpha)(\overline{E_\delta}\gamma^\mu E_\gamma)\,, \tag{3.16}$$

[4]Nevertheless, they might be present at the two loop level.

then $\varepsilon_{CF,R}^{\beta\alpha,\delta\gamma} = 0$. However, the contribution from the quadratic divergence arising at one loop is equal for the neutrino and the four-charged-lepton NSI. Adding this, we have:

$$\varepsilon_{\beta\alpha}^{\delta\gamma} = -\frac{e^{\beta\alpha\delta\gamma}}{2}\frac{v^2}{M^2}\left(\frac{v^2}{M^2} + \frac{\Lambda^2}{M^2}\frac{k}{8\pi^2}\right) \tag{3.17}$$

$$\varepsilon_{CF,R}^{\beta\alpha,\delta\gamma} = -\frac{e^{\beta\alpha\delta\gamma}}{2}\frac{v^2}{M^2}\frac{\Lambda^2}{M^2}\frac{k}{8\pi^2}\,, \tag{3.18}$$

where $k = \mathcal{O}(1)$ depends on the UV completion of the full theory, and $\Lambda \leq M$ is the scale at which new physics appears and the effective theory is no longer valid. Thus, it is clear that no model independent bounds can be derived from this quadratic contribution, since assumptions on the sizes of $\Lambda$ and $k$ have to be made. For example, the high energy completion of the theory could be such that $k = 0$ due to some significant fine-tunings and then no bounds would stem from this process. On the other hand, naturalness arguments can be invoked to argue that, in absence of significant fine-tunings, $k = \mathcal{O}(1)$ and $\Lambda$ should be at least as high as the electroweak scale. With these assumptions, bounds of the same order as the ones derived in [34] will be recovered. On the other hand, if no new physics is present between $v$ and $M$, $\Lambda$ could be identified with $M$ and very stringent constraints could be derived.

Below we will consider, as an example, the computation of the quadratic contribution in a complete theory whose low energy effects are described precisely by the operator of eq. (3.10), without the need of fine-tuning in order to cancel similar operators contributing to four-charged-fermion processes at tree level. In this example, $\Lambda = M$ and $k = 1/2$. Thus, even if only neutrino NSI are induced at tree level, the loop contribution only has a suppression of $8\pi^2 \sim \mathcal{O}(100)$, which could dominate the $\mathcal{O}(v^2/M^2)$ tree level contribution (unless $M \lesssim 1\,\text{TeV}$) and the four-charged-fermion operator would be induced with a strength similar to that of the neutrino NSI. This would imply strong bounds on the neutrino NSI and the main motivation for considering $d = 8$ operators would be lost.

We will now discuss how some antisymmetries between the $\varepsilon$ parameters could avoid these corrections, in a way similar to the $d = 6$ antisymmetric realisation. As for $d = 6$, this is only possible for the case in which the matter fermions are left-handed leptons. To study this, it is more convenient to change the operator basis to a basis where these symmetries are manifest:

$$\begin{aligned}\mathcal{A} &= \frac{1}{2}(|\phi_0|^2 + |\phi_+|^2)[(\bar\nu_\beta\gamma^\rho\nu_\alpha)(\bar\ell_\delta\gamma_\rho\ell_\gamma) + (\bar\nu_\delta\gamma^\rho\nu_\gamma)(\bar\ell_\beta\gamma_\rho\ell_\alpha)\\ &\qquad - (\bar\nu_\beta\gamma^\rho\nu_\gamma)(\bar\ell_\delta\gamma_\rho\ell_\alpha) - (\bar\nu_\delta\gamma^\rho\nu_\alpha)(\bar\ell_\beta\gamma_\rho\ell_\gamma)]\,,\end{aligned} \tag{3.19}$$

$$\begin{aligned}\mathcal{S} &= \frac{\mathcal{B}+\mathcal{C}}{2} - \mathcal{A} =\\ &= \frac{1}{2}(|\phi_0|^2 + |\phi_+|^2)[(\bar\nu_\beta\gamma^\rho\nu_\alpha)(\bar\ell_\delta\gamma_\rho\ell_\gamma) + (\bar\nu_\delta\gamma^\rho\nu_\gamma)(\bar\ell_\beta\gamma_\rho\ell_\alpha)\\ &\qquad + (\bar\nu_\beta\gamma^\rho\nu_\gamma)(\bar\ell_\delta\gamma_\rho\ell_\alpha) + (\bar\nu_\delta\gamma^\rho\nu_\alpha)(\bar\ell_\beta\gamma_\rho\ell_\gamma)]\\ &\quad +2\left[|\phi_0|^2(\bar\nu_\beta\gamma^\rho\nu_\alpha)(\bar\nu_\delta\gamma_\rho\nu_\gamma) + |\phi_+|^2(\bar\ell_\beta\gamma^\rho\ell_\alpha)(\bar\ell_\delta\gamma_\rho\ell_\gamma)\right] + \ldots\,,\end{aligned} \tag{3.20}$$

$$\begin{aligned}\mathcal{X} &= \frac{-\mathcal{B}+\mathcal{C}}{2} =\\ &= (|\phi_+|^2 - |\phi_0|^2)[(\bar\ell_\beta\gamma^\rho\ell_\alpha)(\bar\nu_\delta\gamma_\rho\nu_\gamma) - (\bar\ell_\delta\gamma^\rho\ell_\gamma)(\bar\nu_\beta\gamma_\rho\nu_\alpha)] + \ldots\,,\end{aligned} \tag{3.21}$$

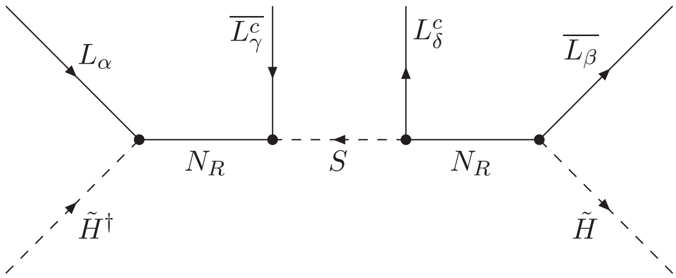


**Figure 3**. An example of a full theory inducing neutrino NSI, but not four-charged-fermion interactions, at tree level. The NSI are generated through the exchange of right-handed neutrinos $N_R$ and a scalar doublet $S$.



where ... represents terms proportional to $\phi_0^*\phi_+$ or $\phi_0\phi_+^*$, which will not contribute to our one-loop computations as long as we consider the leptons to be massless.

In the Feynman gauge, the quadratic divergences are completely determined by the loops of the Goldstones and Higgs fields. Indeed, since the quadratic divergence is independent of the mass propagating in the loop, the terms $|\phi_0|^2$ (via the Higgs and neutral Goldstone) and $|\phi_+|^2$ (via the charged Goldstones) will give the same contribution. It is thus evident, by looking at eqs. (3.19)–(3.21) that the operators $\mathcal{S}$ and $\mathcal{A}$ will contain the quadratic divergence, while $\mathcal{X}$ will not, since the two contributions cancel.

Also, a generic coupling $c^{\alpha\beta\gamma\delta}$ can be decomposed as

$$c^{\alpha\beta\gamma\delta} = s^{\alpha\beta\gamma\delta} + a^{\alpha\beta\gamma\delta} + x^{\alpha\beta\gamma\delta}\,, \tag{3.22}$$

where

$$s^{\alpha\beta\gamma\delta} = s^{\gamma\beta\alpha\delta} = s^{\alpha\delta\gamma\beta} = s^{\gamma\delta\alpha\beta} = \frac{1}{4}(c^{\alpha\beta\gamma\delta} + c^{\gamma\beta\alpha\delta} + c^{\alpha\delta\gamma\beta} + c^{\gamma\delta\alpha\beta})\,, \tag{3.23}$$

$$x^{\alpha\beta\gamma\delta} = -x^{\gamma\delta\alpha\beta} = \frac{1}{2}(c^{\alpha\beta\gamma\delta} - c^{\gamma\delta\alpha\beta})\,, \tag{3.24}$$

$$a^{\alpha\beta\gamma\delta} = -a^{\gamma\beta\alpha\delta} = -a^{\alpha\delta\gamma\beta} = a^{\gamma\delta\alpha\beta} = \frac{1}{4}(c^{\alpha\beta\gamma\delta} - c^{\gamma\beta\alpha\delta} - c^{\alpha\delta\gamma\beta} + c^{\gamma\delta\alpha\beta})\,. \tag{3.25}$$

We see that $s$ has the same symmetry as $\mathcal{S}$, $a$ has the symmetry of $\mathcal{A}$, and $x$ has the symmetry of $\mathcal{X}$, and are actually the coefficients of $\mathcal{S}$, $\mathcal{A}$, and $\mathcal{X}$.

To illustrate how operators with different symmetries contribute in a different way to the quadratic divergence, we will take the theory from section 5.1 of [38] as an example (see figure 3). In this theory, the NSI are realised at tree level by two Higgs doublets selecting the neutrinos from two lepton doublets by taking a vev and the exchange of two right-handed neutrinos and a scalar doublet. For simplicity, we will here assume that the scalar has the same mass $M$ as the right-handed neutrinos. The effective tree level operator is now essentially given by

$$\frac{2c^{\beta\alpha\delta\gamma}}{M^4}(\bar{L}_\beta\tilde{H})\gamma^\rho(\tilde{H}^\dagger L_\alpha)(\bar{L}_\delta\gamma_\rho L_\gamma) = \frac{1}{M^4}(\mathcal{S}+\mathcal{A}+\mathcal{X}) \tag{3.26}$$

and thus contains all of the above mentioned operators, of which two have the quadratic divergence. At tree level, after EWSB, it generates neutrino NSI and four-neutrino interactions[5] with strength $\varepsilon^{\delta\gamma}_{\beta\alpha} = -\frac{c^{\beta\alpha\delta\gamma}}{2}\frac{v^4}{M^4}$. Notice that the introduction of right-handed neutrinos already implies that NSI will be induced at $d = 6$ through the deviations from unitary mixing. Consequently, the constraints derived from non-unitarity in ref. [38] would apply and the loop bounds would not be so relevant. However, we find this toy example useful to connect the quadratic divergence to a full theory in which it can be computed and matched in order to clarify its interpretation. Indeed, since this Standard Model extension is renormalizable and the corresponding four-charged-fermion operator does not appear at tree level, all diagrams are actually finite and we will be able to calculate them unambiguously. When computing the loops of the Higgs and Goldstones in the full theory, we will assume that they are essentially massless compared to the heavy mass scale $M$. We will also assume that the external momenta are negligible. With these approximations, the loop contribution is

$$\frac{c^{\alpha\beta\gamma\delta}}{16\pi^2 M^2}(\bar{L}_\alpha\gamma^\rho L_\beta)(\bar{L}_\gamma\gamma_\rho L_\delta) = \frac{c^{\alpha\beta\gamma\delta}}{16\pi^2 M^2}\big[(\bar{\nu}_\alpha\gamma^\rho\nu_\beta)(\bar{\nu}_\gamma\gamma_\rho\nu_\delta) + (\bar{\ell}_\alpha\gamma^\rho\ell_\beta)(\bar{\ell}_\gamma\gamma_\rho\ell_\delta) + (\bar{\nu}_\alpha\gamma^\rho\nu_\beta)(\bar{\ell}_\gamma\gamma_\rho\ell_\delta) + (\bar{\nu}_\gamma\gamma^\rho\nu_\delta)(\bar{\ell}_\alpha\gamma_\rho\ell_\beta)\big] \,. \tag{3.27}$$

It is easy to check that this is proportional to $\mathcal{S} + \mathcal{A}$, as anticipated. We note that this operator can be obtained from eq. (3.26) simply by replacing $|\phi_0|^2$ and $|\phi_+|^2$ by the factor of $M^2/(32\pi^2)$ coming from the loop integral. Thus, the complete one-loop four-fermion vertexes are given by replacing $|\phi_0|^2$ with $v^2/2$, in order to determine the tree level contribution, and replacing both $|\phi_0|^2$ and $|\phi_+|^2$ by $M^2/(32\pi^2)$ in order to determine the loop contribution. From eq. (3.27) we see that the following interactions are generated, all with similar strength: four-neutrino interactions, neutrino NSI, and four-charged-fermion interactions. However, since the four-charged-fermion and the four-neutrino interactions are completely symmetric under the exchange of flavour indexes, while neutrino NSI are not, the remaining terms for neutrino NSI and four-charged-lepton interactions are:

$$\varepsilon^{\gamma\delta}_{\alpha\beta} = -\frac{v^2}{2M^2}\left[\frac{v^2}{M^2}(s^{\alpha\beta\gamma\delta} + a^{\alpha\beta\gamma\delta} + x^{\alpha\beta\gamma\delta}) + \frac{1}{8\pi^2}(s^{\alpha\beta\gamma\delta} + a^{\alpha\beta\gamma\delta})\right], \tag{3.28}$$

$$\varepsilon^{\alpha\beta,\gamma\delta}_{CF,L} = -\frac{s^{\alpha\beta\gamma\delta}}{32\pi^2}\frac{v^2}{M^2}\,. \tag{3.29}$$

Here, $\varepsilon^{\alpha\beta,\gamma\delta}_{CF,L}$ is defined through

$$\mathcal{L}_{\rm CF} = -2\sqrt{2}G_F\varepsilon^{\alpha\beta,\gamma\delta}_{CF,L}(\bar{\ell}_\alpha\gamma^\rho\ell_\beta)(\bar{\ell}_\gamma\gamma_\rho\ell_\delta) \tag{3.30}$$

Thus, if $8\pi^2 v^2 \ll M^2$, both the neutrino NSI and the four-charged-lepton operators will be dominated by loop effects.

As we have just seen in the above example, a physical meaning can be attributed to the quadratic divergences obtained in the effective theory. Essentially, they can be regulated by reinserting the missing propagators of the heavy particles in the full theory, leaving a contribution to the effective NSI which is suppressed by $v^2/(8\pi^2 M^2)$ instead of the tree

[5] These are generally hard to constrain directly and their effects are usually relatively weak, see, e.g., [39, 57, 58].



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level $v^4/M^4$. In addition, the completely symmetric contribution from $s^{\beta\alpha\delta\gamma}$ will generate an additional four-charged-lepton operator at the one-loop level. Thus, the $s^{\beta\alpha\delta\gamma}$ of a model could be severely constrained by the strong bounds on decays such as $\mu \to 3e$, while the parameters $a^{\beta\alpha\delta\gamma}$ and $x^{\beta\alpha\delta\gamma}$ only contribute to the neutrino NSI and cannot be constrained from these processes. However, it is challenging to build a full theory which generates the antisymmetric couplings, but not the symmetric one, in a natural way.

An important caveat: in the above example of a full theory, we assumed that the masses of the different heavy particles were the same. In general, the tree and loop level contributions to the $\varepsilon$s may be different functions of the mass ratios and couplings, meaning that it could be possible to fine-tune these functions in such a way that the loop-contribution is zero, while still maintaining a non-zero tree level contribution.

## 4 Summary and conclusions

We have reconsidered the bounds on neutrino NSI from one-loop processes. We have shown that, in order to have non-ambiguous bounds, a gauge-invariant realisation of the NSI must be considered. We explicitly studied $d = 6$ and $d = 8$ operators and have shown that, in both cases, the logarithmic divergences of the one-loop contributions are suppressed by the factor $m_\ell^2/M_W^2$, which severely weakens the bounds.

In particular, for $d = 6$ operators, the anti-symmetry relations that arise as a consequence of the requirement of the absence of four-charged-fermion interactions at tree level force the one-loop processes to be zero in the absence of leptons masses.

In the $d = 8$ case, we have shown that the loop processes involving NSI should be calculated in the unitary gauge in order to obtain a gauge invariant result if only the $W$ exchange diagram is considered. In this way, the Goldstone loops present in gauge invariant realisations of the NSI with extra Higgs doublets are automatically taken into account. The result is that the logarithmic divergence is proportional to the factor $m_\ell^2/2M_W^2$. However, a quadratic divergence is present and can be exploited to set bounds on NSI. The use of the quadratic divergence in such a way implies model-dependent naturalness assumptions, in particular on the coefficient of the divergence and the size of the cut-off scale.

For a coefficient $k \simeq \mathcal{O}(1)$ and a cut-off $\Lambda$ of the order of the electroweak scale, the bounds presented in [34] are recovered. Pushing the cut-off scale to $M$, where the effective operators are generated, effectively assuming that no new physics appears between the electroweak scale and $M$ to cancel the quadratic contribution, implies that the loop processes can dominate over the tree level contributions inducing four-charged-fermion interactions of a strength similar to that of the neutrino NSI, unless $M \lesssim 1\,\text{TeV}$. This allows the derivation of strong bounds on the latter, but only based on naturalness arguments and not model-independently. All these considerations apply to NSI of neutrinos with both leptons and quarks. However, for neutrino NSI with left-handed leptons, we have shown that, decomposing the NSI in parts with different flavour symmetries, only the symmetric one contributes to the four-charged-fermion process. Thus, if this part is not present in the full theory from which NSI are realised, the loop constraints can be avoided and large neutrino NSI are still viable if generated via the antisymmetric couplings $a$ and $x$. It remains



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an open question whether there are natural models which can realise this. Otherwise an extra fine-tuning beyond the one required at tree level would be necessary to cancel the loop induced four-charged-fermion operators.

On the other hand, these naturalness arguments can always be evaded if one allows fine-tuning of the theory and (given that [38, 39] have shown that large NSI avoiding four-charged-fermion operators already require a significant amount of fine-tuning at tree level) invoking naturalness arguments at the loop level will not make the model more natural. Thus, we emphasise the fact that no model-independent bounds can be derived from the loop processes studied here, but that the only viable models of large NSI avoiding four-charged-fermion interactions require significant cancellations not only at tree level, but also at the one-loop level.

For practical purposes, in order to realise NSI without ad hoc cancellations, the bounds derived in [38] have to be respected or the neutrino NSI will be of the same order than the four-charged-fermion operators. On the other hand, in a model independent approach considering possible cancellations both at tree and loop levels, the strongest effect is the loosening by three orders of magnitude of the bound on $\varepsilon_{e\mu}^{ff}$, since the constraints in the other flavours are dominated by tree level considerations [34]. These tree level constraints of $\mathcal{O}(10^{-1})$ still apply to $\varepsilon_{e\mu}^{ff}$, but the stringent radiative bounds of $\mathcal{O}(10^{-4})$ from $\mu \to 3e$ ($f = e$) and $\mu \to e$ conversion in nuclei ($f = q$), is lost. Given the strength of this bound, $\varepsilon_{e\mu} = 0$ has been assumed for simplicity in many phenomenological studies. Therefore, it could be of interest to consider larger values for this parameter in order to determine its impact on future neutrino oscillation experiments. However, we would like to stress that this kind of large neutrino NSI would require significant fine-tunings at both the tree and loop levels.

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