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Evidence for right-handed neutrinos at a neutrino factory

F. del Aguila [a], J. de Blas [a], R. Szafron [c], J. Wudka [b,*], M. Zrałek [c]

[a] Departamento de Física Teórica y del Cosmos and CAFPE, Universidad de Granada, E-18071 Granada, Spain
[b] Department of Physics and Astronomy, University of California, Riverside, CA 92521-0413, USA
[c] Department of Field Theory and Particle Physics, Institute of Physics, University of Silesia, Uniwersytecka 4, 40-007 Katowice, Poland

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ABSTRACT

We emphasize that a muon based neutrino factory could show the existence of light right-handed neutrinos, if a deficit in the number of detected events is observed at a near detector. This could be as large as $\sim 10\%$ if the size of the new interactions saturates the present limits from electroweak precision data, what is not excluded by the oscillation experiments performed up to now. A simple model realizing such a scenario can be obtained adding right-handed neutrinos to the minimal Standard Model, together with an extra scalar doublet and a triplet of hypercharge 1. In this case, however, the possible deficit is reduced by a factor of ~ 3, and the Yukawa couplings must be adequately chosen. This is also generically required if lepton flavour violation must be below present bounds.

1. Introduction

Neutrino oscillations provide the only observational evidence of new physics beyond the minimal Standard Model (SM). At present these oscillations can be fully explained by introducing neutrino masses and the corresponding charged current mixing [1]. These observations, however, cannot distinguish between Dirac or Majorana neutrinos, nor they require new interactions (NI).[1] The obvious question is then, where do we have to look in order to determine the neutrino character and/or to observe possible NI involving light neutrinos? This will become especially relevant when we face the need to interpret new data with higher statistics and precision, as foreseen from a neutrino factory [3]. In the following we show that present experimental constraints leave room, corresponding to a $\sim 10\%$ deficit in the expected number of events in appropriate processes, for observing NI involving light right-handed (RH) neutrinos. Such a scenario can be easily realised with a mild extension of the SM, through the addition of additional scalar weak isodoublet (to be denoted by η), and a scalar isotriplet of hypercharge 1, (denoted by Δ), besides three RH neutrinos, ν_R^i. But it requires an adequate choice of Yukawa couplings to suppress lepton flavour (LF) violation; moreover, for this particular model, the deficit allowed by current electroweak precision data (EWPD) is reduced by a factor of ~ 3 compared to the general case above where arbitrary NI are parameterized by gauge-invariant dimension six operators with unrelated coefficients.

Within the SM muons only decay into left-handed (LH) neutrinos. Even if the spectrum is enlarged to include their RH counterparts, these are not produced in such decays because they have no gauge interactions, and neutrino masses are negligible. On the other hand, if other interactions are present in nature, a muon based neutrino factory could inject an admixture of neutrinos with both chiralities. We will show below that the limits on NI involving RH neutrinos are to a large extent those derived from (inverse) muon decay, and therefore relatively weak. This then is a promising reaction where to look for new physics effects in the neutrino system.

Let us first, however, state our setup. We assume that the three light neutrinos are Dirac-type neutrinos, i.e. that there are three light neutrino singlets beyond the minimal SM, and that lepton number (LN) is conserved. In practice this is not a restriction on the light neutrino character, but on the type of NI. Notice that neutrino masses are negligible in all experiments performed up to now, except in neutrino oscillations (and eventually in neutrinoless double β decay, $0\nu\beta\beta$). Thus, we can assume that all interactions conserve LN because neutrino masses are much smaller than the energy relevant in the processes considered and/or the

* Corresponding author.
E-mail addresses: faguila@ugr.es (F. del Aguila), deblasm@ugr.es (J. de Blas), szafron@us.edu.pl (R. Szafron), jose.wudka@ucr.edu (J. Wudka), zralek@us.edu.pl (M. Zrałek).

[1] In what follows we will ignore the LSND data [2].

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experimental precision is much lower than the size of the effects proportional to them, as it is the case for all foreseen experiments not involving neutrino oscillations and excluding $0\nu\beta\beta$. Though we could consider LN violating NI, their effects can be ignored in our analysis,[2] as can be the LN violating effects from neutrino masses.

Hence, the NI effects we are interested in will be relatively large and LN conserving; whereas neutrino masses will be safely taken to vanish. The effective Lagrangian describing such a scenario will not distinguish between (i.e. approximates equally well) the case of exact LN conservation with very small Dirac masses for the three light neutrinos, and the case of negligible Majorana masses for the six light neutral fermions, the only vestige of the very slightly broken LN in this case. This was explicitly proved in [9] for (inverse) muon decay assuming no additional constraints on NI. Note that by similar arguments we can also neglect LF violation induced by light neutrino mixing (to a very good approximation).

The most general four-fermion effective Hamiltonian describing muon decay reads

$$\mathcal{H}_{\mu\to\nu\bar{\nu}e} = \frac{4G_F}{\sqrt{2}} \sum_{\substack{a,b=L,R\\ \gamma=S,V,T}} g^{\gamma}_{ab}\big(\bar{e}\Gamma^{\gamma}\nu^{e}_{a}\big)\big(\overline{\nu^{\mu}_{b}}\Gamma^{\gamma}\ \mu\big) + \text{h.c.}, \tag{1}$$

where a, b label the chirality of the neutrinos, while γ refers to the Lorentz character of the interaction (scalar, vector and tensor). The present limits on the size of the various coefficients will be discussed in the following section, here we merely remark that the two couplings g^V_{LL}, g^S_{RR} are also associated with the largest departure from the SM predictions for the number of events to be detected by a neutrino factory.

There are many other available electroweak precision data that can be used to indirectly constrain the Hamiltonian (1), and in particular those two couplings. To derive such restrictions one can take two routes. The first one consists in re-writing (1) as a linear combination of higher-dimensional operators invariant under $SU(3)_c \times SU(2)_L \times U(1)_Y$, and adding these operators to the SM Lagrangian [10]; the coefficients of the resulting effective theory can then be bound using experimental data at all available energies. In this approach the physics responsible for generating these operators is left unspecified, except for the requirement that its characteristic scale lie beyond the electroweak scale. The advantage of this approach is its generality, since it is not tied to any specific assumption about the physics beyond the SM, its disadvantage is the proliferation of coefficients, and the fact that more than one gauge-invariant operator contributes to each term in (1).

The second route is to extend the SM by adding a specific set of new fields (such as the η and Δ mentioned previously) and interactions. This has the advantage of providing a specific scenario for the physics beyond the SM, but the results obtained are often specific to the assumptions made in constructing the model. At scales below those of the heavy particles this model will reduce to an effective theory of the type mentioned above, except that the effective-operator coefficients are all expressed in terms of a small number of parameters and can be constrained more tightly.

In the following we will examine both of these possibilities. In the next section we consider effective Lagrangian approach, which we denote by E1, where we choose a set of 4-fermion effective operators that generate (1) at low energies and have little impact on other electroweak observables. In Section 3 we will also consider a specific extension of the SM, which we denote by E2, based on an extended scalar sector. Finally, in the last two sections we discuss the implications of these SM extensions for neutrino oscillations and other experiments, respectively.

2. Electroweak precision data constraints

One effective-Lagrangian extension of the SM, which denote by E1, consists in adding to the SM the following set of effective operators

$$\begin{aligned}
-\frac{4G_F}{\sqrt{2}}\Big[\, & g^S_{LL}\big(\overline{e_R}l^e_L\big)\big(\overline{l^\mu_L}\mu_R\big) + g^S_{RR}\big(\overline{l^e_L}\nu^e_R\big)\big(\overline{\nu^\mu_R}l^\mu_L\big) \\
&+ g^S_{LR}\big(\overline{e_R}l^e_L\big)i\sigma_2\big(\overline{\nu^\mu_R}l^\mu_L\big) - g^S_{RL}\big(\overline{l^e_L}\nu^e_R\big)i\sigma_2\big(\overline{l^\mu_L}\mu_R\big) \\
&+ \frac{\delta g^V_{LL}}{2}\big(\overline{l^e_L}\gamma^\alpha l^\mu_L\big)\big(\overline{l^\mu_L}\gamma_\alpha l^e_L\big) + g^V_{RR}\big(\overline{e_R}\gamma^\alpha\nu^e_R\big)\big(\overline{\nu^\mu_R}\gamma_\alpha\mu_R\big) \\
&- \frac{g^V_{LR}}{v^2}\big(\overline{l^e_L}\sigma_a\gamma^\alpha l^e_L\big)\big(\phi^T i\sigma_2\sigma_a\phi\big)\big(\overline{\nu^\mu_R}\gamma_\alpha\mu_R\big) \\
&+ \frac{g^V_{RL}}{v^2}\big(\overline{e_R}\gamma^\alpha\nu^e_R\big)\big(\phi^\dagger\sigma_a i\sigma_2\phi^*\big)\big(\overline{l^\mu_L}\sigma_a\gamma_\alpha l^\mu_L\big) \\
&+ g^T_{LR}\big(\overline{e_R}\sigma^{\alpha\beta}l^e_L\big)i\sigma_2\big(\overline{\nu^\mu_R}\sigma_{\alpha\beta}l^\mu_L\big) \\
&- g^T_{RL}\big(\overline{l^e_L}\sigma^{\alpha\beta}\nu^e_R\big)i\sigma_2\big(\overline{l^\mu_L}\sigma_{\alpha\beta}\mu_R\big)\Big] + \text{h.c.},
\end{aligned} \tag{2}$$

where l^f_L stands for one of the three SM lepton doublets ($f = e, \mu, \tau$) and ϕ is the SM Higgs doublet, and $v \simeq 246$ GeV its vacuum expectation value. We do not use the basis proposed in [4] for writing the beyond the SM dimension six operators; still it is important to note that this basis must be extended to include the light RH neutrinos [7].[3] E1 is manifestly invariant under $SU(3)_c \times SU(2)_L \times U(1)_Y$, and LF (and LN) conserving. Note that g^V_{LR} and g^V_{RL} are associated with operators of dimension eight: $\overline{\nu^\mu_R}\gamma^\alpha\mu_R$ and $\overline{e_R}\gamma^\alpha\nu^e_R$ have hypercharges $Y = -1$ and 1, respectively, and we cannot construct a gauge invariant vector with opposite hypercharge using only $\overline{l_L}$ and l_L; at low energies (2) reduces to (1) with the definition $g^V_{LL} \equiv 1 + \delta g^V_{LL}$.

In (2) we have chosen to write all dimensional coefficients in terms of v. This implies that the natural size for the coefficients is

$$g^\gamma_{ab} \sim \begin{cases} (v/\Lambda)^4 & \text{for } g^V_{LR,RL}, \\ (v/\Lambda)^2 & \text{otherwise}, \end{cases} \tag{3}$$

where Λ denotes the heavy scale of the physics responsible for generating the corresponding operator.

A characteristic feature of this particular extension is that, except for the effect on the muon decay constant and inverse muon decay, EWPD are blind to the operators in Eq. (2). For instance, although g^S_{LR} and g^T_{LR} contribute to $e^+e^- \to \overline{\nu^\mu}\nu^\mu$ and affect the Z^0 invisible width, the effects is negligible compared to the Z^0 pole

[2] An effective theory only involving the light SM fields and invariant under the SM gauge group has only one dimension five operator violating LN [4], the famous Weinberg operator [5] giving Majorana masses to the LH neutrinos and then negligibly small. There is no operator violating baryon minus lepton (B–L) number of dimension six [6]. Thus, operators of dimension six violating LN also violate baryon number (BN), and then involve quarks and will play no rôle in our analysis. If the effective theory also includes RH neutrinos, as in our case, there are two additional dimension 5 LN violating operators [7], one of them generates a magnetic coupling for the RH neutrinos and is very strongly constrained when the ν_R are light [8], the other generates a correction to the ν_R Majorana mass term and is therefore also negligible. There is only one dimension six operator violating B–L (and LN) but involves four RH neutrinos [7], and then is uninteresting for us. The other dimension six operators violating LN also violate BN, and can be also ignored.

[3] Although in that reference the RH neutrinos are assumed to have masses of few hundreds of GeV. (Note that the operator $\mathcal{O}_{QNdQ}$ in Eq. (7) of that reference is redundant.)