Evidence for right-handed neutrinos at a neutrino factory

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\section*{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 \(10\%\) 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.

\section{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 do 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 seen from a neutrino factory \[3\]. In the following we show that present experimental constraints leave room, corresponding to a $10\%$ de$^2$icit in the expected number of events in appropriate processes, for observing NI involving light right-handed (RH) neutrinos. Such a scenario can be easily realized with a mild extension of the SM, through the addition of additional scalar weak isodoublet (to be denoted by $\Phi$), and a scalar iso-triplet of hypercharge $1$, (denoted by $\eta^\pm$), besides three RH neutrinos, $\nu^\ell_R$. But it requires an adequate choice of Yukawa couplings to suppress lepton flavour (LF) violation; moreover, for this particular model, the de$^2$icit allowed by current electroweak precision data (EWPD) is reduced by a factor of $3$ compared to the general case above where arbitrary NI are parametrized by gauge-invariant dimension six operators with unrelated coefficients.

Within the SM m uons 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 lim its 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 neutrino-less double decay, 0$^\nu$). 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 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$^\nu$. Though we could consider LN violating NI, their effects can be ignored in our analysis\[2\], as can be the

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

\[2\] An effective theory only involving the light SM fields and invariant under the SM gauge group has only one dimension-six operator violating LN [4], the famous Weinberg operator [5] giving a Majorana mass 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 role in our analysis. If the effective theory also includes RH neutrinos, as in our case, there are two additional dimension five LN violating operators [7]. One of them generates a magnetic coupling for the RH neutrinos and is very strongly constrained when the $\nu^\ell_R$ are light [8], the other generates a correction to the $\nu^\ell_R$ Majorana mass term.
LN violating effects from neutrino masses.

Hence, the NI eects we are interested in will be relatively large and LN conserving; whereas neutrino masses will be safely taken to vanish. The eective 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 eective Hamiltonian describing muon decay reads

$$H = \frac{4G_F}{2} \sum_{i=L,R} \sum_{\ell=e,\mu,\tau} g_{LL}^{\ell} (\ell \ell \ell \ell) + h.c. \tag{1}$$

where \(\ell\) label the chirality of the neutrinos, while \(\ell\) 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 \(g_{LL}^{\ell}\); \(g_{LR}^{\ell}\) 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 rst one consists in rewriting (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 eective 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 and 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 eective theory of the type mentioned above, except that the eective-operator coefficients are all expressed in terms of a small number of parameters and can be constrained more tightly.

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.
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 (I) 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

\[
\begin{align*}
\frac{4G_F}{v^2} & \left[ g^{\sigma}_{LR} \bar{u} (\bar{D}^\sigma_L \ell) \right] \frac{1}{M_R} \bar{l} R + g^{\sigma}_{RR} \bar{u} (\bar{D}^\sigma_R \ell) \frac{1}{M_R} \bar{l} R \\
g^{\nu}_{LR} \bar{u} (\bar{D}^\nu_L \ell) \frac{1}{M_R} \bar{l} R + g^{\nu}_{RR} \bar{u} (\bar{D}^\nu_R \ell) \frac{1}{M_R} \bar{l} R \\
g^{\nu}_{LL} \bar{u} (\bar{D}^\nu_L \ell) \frac{1}{M_R} \bar{l} R + g^{\nu}_{RR} \bar{u} (\bar{D}^\nu_R \ell) \frac{1}{M_R} \bar{l} R \\
\end{align*}
\]  
(2)

where \( \ell \) stands for one of the three SM lepton doublets (\( e; \mu; \tau \)) and \( \nu \) is the SM Higgs doublet, and \( v' = 246 \text{ GeV} \) is the vacuum expectation value. We do not use the basis proposed in [8] 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 \([9]\). E1 is manifestly invariant under SU(3)_c SU(2)_L U(1)_Y, and LF (and LN) conserving. Note that \( g^{\sigma}_{LR} \) and \( g^{\nu}_{RR} \) are associated with operators of dimension eight: \( \frac{1}{M_R} \bar{l} R \) and \( \bar{u} (\bar{D}^\sigma_R \ell) \) have hypercharges \( Y = 1 \) and 1, respectively, and we cannot construct a gauge invariant vector with opposite hypercharge using only \( \frac{1}{M_R} \bar{l} R \) at low energies; (2) reduces to \( \hat{\mathcal{O}} \) with the definition \( g^{\nu}_{LR} = 1 + g^{\nu}_{RL} \).

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

\[
\begin{align*}
g^{\nu}_{LR} (v = \text{few h})^4 & \text{ for } g^{\nu}_{LR}; \\
g^{\nu}_{RR} (v = \text{few h})^2 & \text{ otherwise;}
\end{align*}
\]  
(3)

where \( \hat{\mathcal{O}} \) denotes the heavy scale of the physics responsible for generating the corresponding operator.

A slight in that reference the RH neutrinos are assumed to have masses of few hundreds of GeV. (Note that the operator \( \hat{\mathcal{O}} \) in Eq. (7) of that reference is redundant.)
A characteristic feature of this particular extension is that, except for the eect on the muon decay constant and inverse muon decay, EWPD are blind to the operators in Eq. (2). For instance, although $g_{LR}^S$ and $g_{L}^2$ contribute to $e^+e^-$ and a eect the $Z^0$ invisible width, the eects is negligible compared to the $Z^0$ pole contribution. E1 does not contribute to LF violating processes either, because it does not include LF violating operators. Similarly, universality is preserved since the gauge couplings to leptons remain the same as in the SM.

It must be emphasized, however, that in specific models (such as E2, described in section 3 below) the various coe cients in Eq. (1) are written in terms of the fundamental parameters of the theory, these parameters appear in all other interactions, and in general will contribute to other observables, which in turn can be used to stringently constrain $g$. In fact, it is non-trivial to nd models where these limits are not so strict they exclude further observable eects from (1).

As we argue below the largest departure from the SM predictions is parameterised by $g_{LL}^V$ and $g_{R}^S$. Constraints on these two operators $\bar{\Gamma}_{LR}^{s}$ and $\bar{\Gamma}_{LR}^{l}$ in Eq. (2) can be obtained using current data, as described in the Appendix. At 90% C.L. we nd

$$\begin{align*}
\text{case} & \quad g_{LL}^V \quad 1 + g_{L}^2 \quad 1 + g_{LR}^S \\
(a) & \quad > 0.960 \quad < 0.550 \\
(b) & \quad > 0.957 \quad < 0.579 \\
(c) & \quad > 0.998 \quad < 0.954
\end{align*}$$

where the limits in case (a) are obtained directly from (4) using muon decay and $e^+e^-$ data [1,11]; in case (b) from a global t using precision data to the operator coe cients with the SM parameters xed at their best- t values [1]; in case (c) as in case (b) but taking one e ective-operator coe cient to be non-zero at a time. This last possibility, though often adopted for simplicity is seldom realistic: as noted previously one must expect the heavy physics to generate several operators with related coe cients, so that allowing for several non-zero interactions become compulsory. Even more, if the new physics is to have sizable indirect eects, then it must also conspire to preserve the excellent t of the SM to the EW PD at the 0.1% level [1]. In Fig. 1 we plot the bounds for case (b). The results in (4) indicate that the two parameters in the t are highly correlated. This is expected since the t is dominated by the constraint on the strength of the muon decay constant $G_m$ which is proportional to the SM Fermi

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See [12] and reference therein for a t to operators involving only left-handed neutrinos.

These limits stay mainly unchanged if the SM parameters are also left free in the global t.
Figure 1: 90% C.L. bounds for E1 case (b) in Eq. (4); and the same for the SM extension E2 in next section. The narrow bands between the origin and the crosses define the 90% confidence region for the global t to the two parameters for E1 (left cross) and E2 (right cross) respectively.

Constant \( G_F \) in Eq. (4):

\[
G = G_F \left[ 1 + g_{LL}^2 + g_{RR}^2 + g_{LR}^2 + g_{RL}^2 \right] + \frac{1}{4} \left( g_{LL}^2 + g_{RR}^2 + g_{LR}^2 + g_{RL}^2 \right) + 4 \left( g_{LR}^2 + g_{RL}^2 \right) G_F A ;
\]

(5)

with the proportionality constant A restricted by the global t to the interval

\[ 0.9997 < A < 1.0004 \text{ at 90% C.L.:} \]

(6)

This bound alone constrains the coefficients to a narrow band in the \( g_{LL}^V \quad g_{RR}^S \) plane

\[ \begin{bmatrix} g_{RR}^S \end{bmatrix}^2 \quad 8 \quad \begin{bmatrix} g_{LL}^V \end{bmatrix} ; \]

(7)

as depicted in Fig. 1.

Using the bounds above we can derive limits on the scale of new physics responsible for the two operators being considered. Using (3) and assuming the underlying physics is weakly coupled we nd the weak constraints \( > 130 \text{ GeV for } g_{RR}^S \) and \( > 500 \text{ GeV for } g_{LL}^V \).

As stressed in the introduction, we are interested in those interactions in the muon decay effective Hamiltonian that could allow for the largest deviation from the SM prediction of the number of events that may be eventually detected at a neutrino factory. In general, while it is clear that a negative \( g_{LL}^V \) is strongly favored for this to
be possible (see Eqs. (5,6)), one may wonder whether any of the other interactions involving a RH neutrino may play the role of \( g_{RR} \). This could be the case for \( g_{RR} \) or the LR and RL operators in Eq. (2). However, unlike for \( g_{RR} \), which cannot be separated from \( q_L^\gamma \) in muon decay experiments since we do not measure the polarization of the final neutrinos, all the other interactions are constrained by the absence of any significant deviation from the V-A prediction in the spectrum of the outgoing electrons. In particular, \( g_{RR}^\gamma < 0.034 \). On the other hand, the bounds for some of the LR and RL interactions are relatively weak, but they also generate radiative corrections to the neutrino masses and, assuming naturality, are constrained by the associated upper limits [13]: \( g_{LR}^\gamma ; g_{LR}^\gamma ; g_{LR}^\gamma < 10^2 \) and \( g_{RL}^\gamma ; g_{RL}^\gamma ; g_{RL}^\gamma < 10^4 \). Therefore, when we consider the data from inverse muon decay, where the incident neutrinos come from pion decays and then are LH and derive a relatively weak bound on \( q_L^\gamma \), this can be only compensated by \( g_{RR}^\gamma \) in order to satisfy (3), but in contrast with the remaining operators in Eq. (2) has no further constraints.

3 A simple SM extension

Let us discuss a simple model realising the former scenario. It extends the SM including besides three RH neutrinos \( \nu_R \) with zero LN, a second scalar doublet \( \nu_R \) with LN equal to 1, and a scalar triplet \( \nu_T \) with hypercharge 2, as in [14] and [15], respectively. They both can acquire a vacuum expectation value through very small LN violating couplings to the SM Higgs (which carries zero LN), and provide mass mixing to the light neutrinos, which must be the observed spectrum in oscillation experiments. But, these couplings, as the neutrino masses, can be safely neglected in our analysis: we can assume that light neutrinos are massless, and that LN and LF are both conserved. Thus, besides the kinetic term for the RH neutrinos, the SM fermionic Lagrangian acquires only two more terms

\[
\mathcal{L}_{\text{int}} = f_{ij} \left[ \tilde{\nu}_R^i \nu_R^j \right] + h.c.;
\]

with only \( f_{ee} \), and \( e = e^+ \) potentially large enough to produce measurable effects in current and near-future experiments. The integration of the extra scalars out results, in particular, in the following contributions to the muon decay effective Hamiltonian:

\[
\frac{4G_F}{2} \ g_{RR}^S \ e_R \ e_R \ \nu_R^L \ \bar{\nu}_L^R + \frac{g_{LL}^\gamma}{2} \ \bar{\nu}_L^R \ \nu_L^R \ \bar{\nu}_L^R + h.c.;
\]

with coefficients

\[
g_{RR}^S = \frac{1}{2G_F} \ \frac{f_{ee} F_{\text{em}}}{M^2} \quad \text{and} \quad g_{LL}^\gamma = \frac{1}{2G_F} \ \frac{F_{\text{em}}}{M^2};
\]

\[
\text{There can be also tiny RH neutrino mass term.}
\]
respectively, where $M$ stand for the scalar masses. The full set of dimension six operators arising from the integration of the scalar triplet is given in [13].

In this case, named E2 in the former section, the EW PD analysis presented in the Appendix implies more stringent bounds on $q_L^V$ and $g_R^{S}$ than for E1:

$$1 + q_L^V > 0.988 \text{; and } j g_R^{S} j < 0.313.$$  

(12)

The corresponding band is plotted in Fig. 1. The bounds obtained are tighter because, as emphasized previously, the integration of definite new physics also gives, in general, operators contributing to other processes, further restricting the model. In the present case the integration of the generates also the operator $F_L \; F_L \; l_L \; l_L$, which has the same coefficient $q_L^V$ as $F_L \; l_L \; l_L \; l_L$, and contributes to $e! e$, further restricting the allowed deviation from the SM predictions.$^1$

LF violation is below experimental bounds because similarly to the E1 case the only non-negligible couplings in Eq. (9) are $f_{ee}$, and $e = e$, and because the scalar doublets and triplet mix very little, as required by approximate LN conservation. The absence of new couplings and that the SM gauge couplings stay unchanged guarantee the agreement with universality constraints on the lepton sector.

4 Neutrino factory predictions

The relevant phenomenological question is where could the RH neutrinos be eventually observed if the $q_L^V$ and $g_R^{S}$ four-fermion interactions are non-zero. Obviously, they can be probed in a more precise inverse muon decay experiment: a more precise measurement of this process could give evidence for those NI (or reduce the allowed deviation from the SM in Eqs. (4,12) and Fig. 1). But a muon-based neutrino factory will also be sensitive to them. Indeed, if a substantial amount of the neutrinos produced in muon decay are RH, a near-detector sensitive to neutrino-hadron collisions will observe a deficit in the same proportion, and this deficit would be twice as large if the detector could also measure the inverse muon decay process. In Table 1 we give the maximum deficit expected in the case that the new interactions saturate the 90% C.L. bounds obtained in the previous sections. Whereas Fig. 2 we show the predicted deficit as a function of $q_L^V$ for the SM predictions considered.

If the precision in the measurement of the inverse muon decay is improved by a factor of 2 without modifying the central value, the limits on the NI would be strengthened and the allowed deficit of observed events at a neutrino factory reduced; the corresponding percentages in such a case are given in parentheses in Table 1. For the E1 extension the deficit would be reduced to 5%, being the 90% C.L. bounds in this case $1 + q_L^V > 0.975$ and $j g_R^{S} j < 0.442$. In the E2 case the improvement in the precision of inverse

$^1$In our approximation (implying negligible LN violation in the scalar sector) there are no tree-level contributions to the oblique parameters.
Table 1: Maximum deficit in the number of observed events in a near detector sensitive to neutrino-nucleon collisions (-N) and inverse muon decay (IMD) for the two SM extensions discussed in the text. In parentheses we show the deficits expected in the case that the precision on the measurement of inverse muon decay is improved by a factor of 2.

| Process | E1         | E2         |
|---------|------------|------------|
| -N      | 8.5% (5.0%) | 2.5% (2.5%) |
| IMD     | 15.4% (9.3%) | 4.8% (4.8%) |

Figure 2: Percentage of detected events in a near detector at a neutrino factory compared to the predicted number by the SM, as a function of the NI strength $q_{LL}^V$. The solid (dotted-dashed) curve corresponds to neutrino-nucleon (inverse muon decay) collisions. The vertical lines stand for the E1 and E2 limits on $q_{LL}^V$ in the text.

Muon decay would have no appreciable effect because this constraint would be weaker than the one derived from elastic scattering.

For a neutrino factory to be sensitive to a deficit < 3% (the smallest value listed in Table 1), the neutrino flux must be known with enough precision. Besides, large fluxes are also required to have a large number of events at a near detector in order to keep the statistical error small. For a detailed study about future neutrino factories see [3]. Assuming $10^{21}$ muon decays in one year of operation, the number of expected -N events at a near detector (such as the one described in Table 1 of [3]) is of the order of $10^3$; while the number of IMD events range from $10^4$ to $10^5$, depending on the energy and polarization of the decaying muon. Then, deficits even smaller than a few per mille due to the injection of RH neutrinos could be eventually testable for such a large statistics. However, per mille deficits may be too small because the highest achievable precision in the determination of the flux is expected to be at most of $0.1\%$. But
being conservative, it could be up to a factor ten times larger, and then of the same order as the largest possible deficit for the E2 extension.

At this point one may wonder about the consistency of these large deficits with the interpretation of present neutrino oscillation experiments summarized in Table 2. As

| source                  | Experiment                          | Detection |
|-------------------------|-------------------------------------|-----------|
| Reactor (decay)         | Palo Verde [16], CHOOZ [17]         | $e$ ! $\bar{\nu}$ |
|                         | KamLAND [18]                        | $\nu$ ! $\bar{\nu}$ |
| Solar                   | SNO [19], Borexino [20]             | $e$ ! $\bar{\nu}$ |
| Atmospheric (and decays) | Super Kam iokande [21]              | ! |
| Accelerator (and K decays) | K2K [22], MINOS [23]               | ! |
|                         | CHORUS [24], NOMAD [25]             | ! |
|                         | MiniBooNE [26]                      | ! |
| Accelerator (decay at rest) | LSND [2]                           | ! $e$ |
|                         | KARMEN [27]                         | ! $\bar{\nu}$ |

Table 2: Neutrino source for the different oscillation experiments and search process.

we will argue, these seem to be largely insensitive to new four-fermion interactions involving an electron, a muon and the corresponding neutrinos. Reactor experiments are initiated by electron antineutrinos from decay, they are then LH and fully described by the SM. Similarly, solar neutrinos have electronic flavor, are also LH and produced by SM reactions. Atmospheric neutrinos are decay products of pions and muons from cosmic rays, and may include RH neutrinos. However, since the flux of cosmic rays is isotropic, and atm ospheric neutrino oscillation experiments only compare fluxes of muon neutrinos coming from different directions, they are not sensitive to a possible deficit in the total number of initial LH neutrinos from muon decays. On the other hand, in accelerator experiments looking for $\nu_e$ or $e^{-}$, neutrino beams are mainly formed by muon neutrinos originating from pion decays (with a fairly small contamination) [22], and then LH and with the SM oscillation pattern. Finally, in accelerator experiments looking for $\nu_{\mu}$ or $\mu^{-}$, the muon antineutrinos are produced in $\pi^+$ decays, and therefore they are sensitive to the NI we are interested in. However, what they measure is the number of positrons produced by inverse decay, looking for an excess of electron anti-neutrinos instead of looking for a deficit in the observed number of muon anti-neutrinos. (The excess reported by the LSND experiment [7] has no explanation in this setup, which predicts a deficit.) Hence, there appears not to be any contradiction between the significant deficit predicted by the NI considered here and the interpretation of current oscillation experiments.
5 Further phenomenological implications

In specific models that contain new fields and interactions there are in general further observable effects, as for instance the production of the new particles at large colliders. This is the case of the simple model E2 discussed in Section 3. If this type of NI saturates the EW PD limits,

$$M \lesssim 0.4 \text{ TeV}$$

implying a relatively large $\lambda$ and a light $\eta$. The LHC will be able to uncover such a scalar triplet for masses up to 900 GeV and an integrated luminosity of 30 fb$^{-1}$ [28].

For this model the relation in Eq. (7) translates into a correlation between the scalar masses $M_i$ and/or the Yukawa couplings $f_{\eta e}$, namely

$$f_{\eta e} \left( \frac{M}{\Lambda} \right)^2 \gtrsim 32 \frac{p}{G_F} \frac{y_e}{M^2}$$

This implies that for these NI to have sizable effects at a neutrino factory we must have relatively large $f_{\eta e}$ and light $\eta$. The main signals in this case are missing energy plus one or two leptons $\ell = e$; because a charged (neutral) $\eta$ has only sizable decays into $\ell$. Thus, $W, WZ$ and $WW$ production can provide too large irreducible backgrounds for observing these scalars. Other SM processes like tt production can also give large backgrounds. The search for this scalar doublet is similar to left slepton searches assuming that they only decay into a LH charged lepton and the lightest supersymmetric particle [20]. (Although in general sleptons can also have cascade decays and be decay products of other supersymmetric particles [31]). In this case the LHC discovery limit is 300 GeV for an integrated luminosity of 30 fb$^{-1}$ [32], but this is assuming that the slepton doublets coupling to the rest of the charginos are degenerate. However, in our model there is only one scalar doublet, scaling the corresponding limit to 250 GeV after correcting by the factor of 2. It may also happen that the light and heavy may be too heavy to be directly observed at the LHC (e.g. $M_\eta \gtrsim 1$ TeV and $M_{\lambda} \gtrsim 250$ GeV), their effects can still be observable at a neutrino factory provided $f_{\eta e}$ and $f_{\lambda e}$, $f_{\eta e}$, $\ell = e \otimes 0.8$. Obviously, large enough lepton colliders are better suited for searching these scalars, since they couple mainly to leptons, and these colliders allow for a better kinematic reconstruction. On the other hand, the relation in Eq. (14), as the cancellation of other possible LF violating couplings, does not appear to be natural in this simple model, but they could be in more complicated frameworks.

$^8$This estimate is larger than those quoted in [28] for in this model the scalar triplet only couples to $e$ and not to taus, which are more difficult to identify and have larger backgrounds. In contrast, with type II see-saw models, this triplet is not the only source of neutrino masses, and light neutrino masses and mixings can be reproduced even when the couplings vanish.

$^9$See [28] for alternative scalar doublet models.
Finally, we note that one might consider other SM additions/extensions generating \( q^L \), as for instance, heavy neutrino singlets or triplets mixing with the electron or muon. In these cases, however, the relatively small mixings allowed by EW PD \([33]\) do not allow contributions large enough to produce a sizable deficit.

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Appendix

Our fits are performed using a $^2$ analysis for the experimental data collected in Tables 3 and 4. This data set is essentially the same used in [33], but updating the top [34].

| Quantity | Experimental Value | Standard Model |
|----------|--------------------|---------------|
| $m_\ell$ [GeV] | 173.1 $\pm$ 1.3 | 173.2 |
| $M_W$ [GeV] | 80.399 $\pm$ 0.023 | 80.373 |
| $g_L^2$ | 0.3005 $\pm$ 0.0012 | 0.3039 |
| $g_R^2$ | 0.3110 $\pm$ 0.0010 | 0.3031 |
| $L$ | 2.51 $\pm$ 0.033 | 2.46 |
| $R$ | 4.59 $\pm$ 0.41 | 5.18 |
| $g_V^e$ | 0.040 $\pm$ 0.015 | 0.0399 |
| $g_A^e$ | 0.507 $\pm$ 0.014 | 0.507 |
| $Q_W$ (Cs) | 72.74 $\pm$ 0.46 | 73.13 |
| $Q_W$ (Tl) | 116.4 $\pm$ 3.6 | 116.7 |
| $Q_W$ (e) (Miller) | 0.0403 $\pm$ 0.0053 | 0.0473 |
| $I_{Vij}^2$ | 1.0000 $\pm$ 0.0006 | 1 |

Table 3: Non-Z-pole (pseudo-) observables included in the global fits.

and W [35] mass values as well as the world average for $s$ [36]. We also include light flavour data at the Z-pole [37], the Tevatron determination of the effective weak mixing angle [38], and further low energy observables, such as the weak charges for Thalium (atomic parity violation) and electron (Miller scattering) [1], the unitarity of the first row of the CKM matrix [39], and obviously, the data from inverse muon decay [40]. The SM best fit values for each observable are also included in the tables. They have been computed using ZFITTER 6.43 [41], complemented with routines for the computation of some of the low energy observables. For the minimizations of the $^2$ we use the package MINUIT [42]. The SM minimum is obtained for the following values:

$$m_H = 96^{+37}_{-28} \mathrm{ GeV} ; \quad m_\tau = 173.2 \pm 1.3 \mathrm{ GeV} ;$$
$$M_Z = 91.1875 \pm 0.0021 \mathrm{ GeV} ; \quad s(m_H^2) = 0.1184 \pm 0.0077 ;$$
$$^{(5)}_{\text{had}} (m_H^2) = (276.9 \pm 3.3) \times 10^8 .$$

The new physics effects are incorporated adding to the SM predictions the tree level contributions from the new operators. The impact of this extension on the best t
| Quantity | Experimental Value | Standard Model |
|----------|-------------------|----------------|
| $M_2^{(5)}$ | 0.02758 0.00035 | 0.02769 |
| $s$ | 0.184 0.007 | 0.184 |
| $M_Z$ [GeV] | 91.1876 0.0021 | 91.1875 |
| $Z$ [GeV] | 2.4952 0.0023 | 2.4957 |
| $\Gamma^0$ [nb] | 41.541 0.037 | 41.479 |
| $R^0_e$ | 20.804 0.050 | 20.741 |
| $R^0$ | 20.785 0.033 | 20.741 |
| $R^0_{FB}$ | 0.0145 0.0025 | 0.0164 |
| $A^0_{FB}$ | 0.0169 0.0013 | |
| $A^0_{FB}$ | 0.0188 0.0017 | |
| $A_e$ (SLD) | 0.1516 0.0021 | 0.1477 |
| $A_e$ (P ) | 0.1498 0.0049 | |
| $A_e$ (P ) | 0.1439 0.0043 | |
| $R^0_b$ | 0.21629 0.00066 | 0.2158 |
| $R^0_c$ | 0.1721 0.0030 | 0.1722 |
| $A^0_{FB}$ | 0.0992 0.0016 | 0.1035 |
| $A^0_{FB}$ | 0.0707 0.0035 | 0.074 |
| $A_b$ | 0.023 0.020 | 0.035 |
| $A_c$ | 0.070 0.027 | 0.068 |
| $A^0_{FB}$ | 0.098 0.011 | 0.1037 |
| $A_s$ | 0.095 0.091 | 0.036 |
| $R^0_{FB} = R^0_{u+d+s}$ | 0.371 0.022 | 0.359 |
| $Q_{had}^{(5)}$ | 0.0403 0.0026 | 0.0424 |
| $\sin^2 \theta_e^{\text{lep}}$ | 0.2315 0.0018 | 0.23143 |

Table 4: Z-pole (pseudo-) observables included in the global fits.
values for the SM inputs is insignificant. For the E1 completion, which allows for the largest departure from the SM,

\[
\begin{align*}
    m_{H^1} &= 92^{40}_{25} \text{ GeV}; \\
    m_{\tilde{e}_1} &= 173213 \text{ GeV}; \\
    M_{\tilde{\tau}_1} &= 9118750.0021 \text{ GeV}; \\
    s(M_{\tilde{Z}^0})_{E1} &= 0.11840.0007; \\
    (5)^{(5)}_{\text{had}}(M_{\tilde{Z}^0})_{E1} &= (276833) \times 10^7.
\end{align*}
\]

Obviously, this is because none of the new operators directly contributes to (almost) any of the considered observables; and they indirectly only through \( G \), which is in excellent agreement with the SM prediction. Thus, \( q_{L}^{\text{vec}} \) tends to be negative in order to compensate the effect of the other operators and the SM parameters do not feel the presence of the new parameters. Once we have introduced the new operators, we perform a scan over the parameter space and reconstruct the probability density function (p.d.f.) using an acceptance/rejection method. From the resulting distribution we compute the confidence regions at 90% C.L., as well as the marginal p.d.f. for each parameter, which we use to obtain the limits and confidence intervals quoted in the text.
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