Neutrino masses and electroweak symmetry breaking

F. Bazzocchi and J. W. F. Valle

Instituto de Física Corpuscular - C.S.I.C./Universitat de València
Campus de Paterna, Apt 22085, E-46071 València, Spain

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Neutrino mass generation may affect the basic structure of the electroweak symmetry breaking sector. We consider a broad class of elementary particle theories where neutrinos get mass at a low mass scale. We show how these can be made natural up to few TeV or so, in the absence of supersymmetry or other possible stabilizing mechanisms. Although the standard signatures for which LHC has been optimized are absent, others are expected. A generic one among these is the possibility of an invisibly decaying Higgs boson which is characteristic of models with spontaneous breaking of lepton number symmetry below TeV or so.

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Two important theoretical challenges of particle physics are elucidating the nature of electroweak symmetry breaking and the origin of neutrino masses. While the neutrino oscillation data confirmed the existence of neutrino masses and mixing [1], the Higgs boson has not yet given an irrefutable proof of its existence.

It has been suggested by Barbieri and Hall [2] that the physics that tames the quadratic Higgs boson mass divergence may be much less ambitious than supersymmetry. Here we follow up on this idea and propose that the physics responsible for “postponing” naturalness is the same one providing neutrinos their mass.

We point out that neutrino masses, however small, may drastically affect the electroweak sector. While some of our considerations hold for generic models with explicitly broken lepton number, we will focus on models with spontaneous violation of lepton number. In such models a lepton number symmetry is broken by an $SU(3)_c \otimes SU(2)_L \otimes U(1)$ singlet vacuum expectation value $\langle \sigma \rangle$ below a few TeV [3]. It has long been noted that in all these models the Higgs has the “invisible” mode

$$h \rightarrow JJ$$

as a sizeable decay channel [3], which may dominate over the SM modes, such as $b\bar{b}$ or $\tau\bar{\tau}$. Here $J$ denotes the associated pseudoscalar Goldstone boson, called majoron.

Since it is weakly interacting with all other particles, this leads to events with large missing energy that could be observable at collider experiments [3, 4, 5].

Here we note that, apart from changing the low energy theory by the existence of the new Higgs decay in Eq. (1), such models generically improve the naturalness of the electroweak symmetry breaking sector even in the absence of supersymmetry or some other kind of physics capable of canceling the quadratically divergent top loop contribution to the Higgs boson mass. Therefore they provide a “worse-case” scenario for LHC [2] in which naturalness in the electroweak symmetry breaking sector is effectively “lifted” to a scale higher than testable at the LHC.

However, the novel properties of the scalar sector itself could give visible “signs of life” in collider experiments, for example, the mode in Eq. (1) and many other signatures which depend on how neutrinos get their masses by coupling to the scalars.

Consider the simplest tree level scalar potential that can simultaneously account for electroweak symmetry breaking and the generation of naturally small neutrino masses,

$$V[\Phi, \sigma] = \mu_{\Phi\Phi}^2 (\Phi^\dagger \Phi) + \mu_{\sigma\sigma}^2 (\sigma^\dagger \sigma) + \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 (\sigma^\dagger \sigma)^2 + \lambda_{\Phi\sigma} (\Phi^\dagger \Phi)(\sigma^\dagger \sigma)$$

(2)
The potential in Eq. (2) has a U(1) global symmetry under which $\sigma$ gets rephased, but not $\Phi$. This will become lepton number once leptons are coupled (see example below).

Once we include the one-loop quadratic divergent corrections to the scalar potential the previous $\mu^2_\Phi$, are replaced by

$$
\mu^2_\Phi = \mu^2_\Phi + \frac{\Lambda^2}{16\pi^2} \gamma_\Phi
$$

$$
\mu^2_\sigma = \mu_\sigma^2 + \frac{\Lambda^2}{16\pi^2} \gamma_\sigma ,
$$

where $\gamma_\Phi$ takes into account the gauge, the top and scalar contributions and $\gamma_\sigma$ only the scalar one. The parametrization given by Eq. (3) allows us to discuss the stabilization of the vevs in terms of their dependence on the cut off scale $\Lambda^2$, i. e., in terms of the fine tuning.

The vacuum configuration that breaks the electroweak gauge symmetry and the $U(1)$ global symmetry through

$$
(\Phi) = v_\Phi/\sqrt{2} \quad (\sigma) = v_\sigma/\sqrt{2}.
$$

By solving the extremization conditions we can write the expressions of the vevs in terms of the parameters of the potential

$$
v_\Phi^2 = 2\frac{\lambda_\Phi \mu_\Phi^2 - 2\lambda_\sigma \mu_\Phi^2}{4\lambda_\Phi \lambda_\sigma - \lambda_\Phi^2} 
$$

$$
v_\sigma^2 = 2\frac{\lambda_\Phi \mu_\sigma^2 - 2\lambda_\sigma \mu_\sigma^2}{4\lambda_\Phi \lambda_\sigma - \lambda_\Phi^2}.
$$

Following we can define the fine tuning parameter

$$
D = \sqrt{D_\phi^2 + D_\sigma^2} ,
$$

where $D_\phi = \frac{\partial \ln v_\phi^2}{\partial \ln \Lambda^2}$ and $D_\sigma = \frac{\partial \ln v_\sigma^2}{\partial \ln \Lambda^2}$. Inserting these in Eq. (5) and inverting the expression we obtain an upper bound for the cut-off

$$
\Lambda^2 \leq 8\pi^2 v_\phi^2 v_\sigma^2 DF
$$

where $F$ is given by

$$
\frac{(4\lambda_\Phi \lambda_\sigma - \lambda_\Phi^2)}{\sqrt{(v_\Phi^4(-2\gamma_\Phi \lambda_\Phi + \gamma_\Phi \lambda_\Phi \sigma)^2 + v_\sigma^4(-2\gamma_\sigma \lambda_\sigma + \gamma_\sigma \lambda_\Phi \sigma)^2)}
$$

Minimizing the potential one obtains the mass matrix

$$
M = \begin{pmatrix}
2\lambda_\Phi v_\Phi^2 & \lambda_\Phi \phi v_\phi v_\sigma \\
\lambda_\Phi \phi v_\phi v_\sigma & 2\lambda_\sigma v_\sigma^2
\end{pmatrix}
$$

for the two neutral CP-even scalars, $h$ and $H$. Its diagonalization leads to a mixing angle between doublet and singlet states in the CP-even sector, given by

$$
\tan 2\alpha = \frac{\lambda_\Phi \phi v_\phi v_\sigma}{(\lambda_\sigma v_\sigma^2 - \lambda_\Phi v_\Phi^2)}
$$

and the masses $m_{H,h}^2$ of the two scalars,

$$
\lambda_\Phi v_\Phi^2 + \lambda_\sigma v_\sigma^2 \pm \sqrt{\lambda_\Phi v_\Phi^4 + \lambda_\sigma v_\sigma^4 + \lambda_\Phi v_\Phi^2 (\lambda_\Phi^2 - 2\lambda_\Phi \lambda_\sigma)}.
$$

where $m_{h_H}^2$ is associated to the plus sign, $m_h^2$ to the minus. In this scheme the vacuum breaks both electroweak symmetry and lepton number. Since the latter is a $U(1)$ global symmetry there is, in addition, a pseudoscalar Goldstone boson, the majoron, to complete the set of physical spin-less bosons.

We can use Eqs. (8)–(9) to rewrite Eq. (6), to map the restriction on the cutoff in terms of the five parameters $m_h, m_H, v_\phi, v_\sigma$ and $\alpha$. Before doing this let us consider two limiting cases. When $\cos \alpha = 1$ Eq. (6) reduces essentially to the SM case,

$$
\Lambda^2 \leq D \pi^2 m_h^2,
$$

For the case $\cos \alpha = 0$ Eq. (6) becomes

$$
\Lambda^2 \leq D 4\pi^2 m_H^2 f(x),
$$

where

$$
f(x) = \begin{cases}
\frac{x^2}{\sqrt{5 - 24x^2 + 36x^4}}, & v_\Phi \simeq v_\sigma \ll m_H, \quad x = v_\phi/v_\sigma \\
\frac{x^2}{\sqrt{4 - 24x^2 + 36x^4}}, & v_\Phi \ll v_\sigma \simeq m_H, \quad x = v_\phi/m_H \\
\frac{x^2}{\sqrt{1 + 16x^4}}, & v_\Phi \simeq m_H \ll v_\sigma, \quad x = m_H/v_\sigma
\end{cases}
$$

and the cutoff $\Lambda$, being proportional to the heaviest CP-even scalar, can be raised up to few TeV for natural choices of the parameters that give $f(x) \sim O(1)$. In Fig. I we show the regions of “extended naturalness” is this simplest model. We display the contour regions in the $m_H$-$c_\alpha$ plane leading to an increase in the effective cutoff $\Lambda$, which can reach a few TeV or so for reasonable values of parameters in this model.

The parameters of the electroweak symmetry breaking sector are constrained by electroweak precision tests.
FIG. 1: Regions in the $m_H$-cos$\alpha$ plane leading to an increase in the cutoff $\Lambda$ for 10% fine-tuning and $m_h = 115$ GeV. In the darkest contour $\Lambda \geq 6$ TeV and in the lightest $\Lambda \geq 2$ TeV, decreasing in TeV steps in between. Note that in our model we can have the lightest scalar boson with a mass of 115 GeV and a cutoff around 6 TeV, while in the SM for $m_h = 115$ GeV the cutoff is 1.2 TeV.

(EWPT) which indicate a value for the Higgs boson mass $120^{+74}_{-49}$ GeV which in our model corresponds to

$$\cos^2 \alpha \log m_h + \sin^2 \alpha \log m_H \leq \log m_h^{+2(3)\sigma},$$  

(12)

where $m_h^{+2(3)\sigma} = 277(350)$ GeV.

In the plots we display the corresponding curves. Thus one sees that the SM upper limit on the Higgs boson mass is now replaced by restrictions in the $m_H$-cos$\alpha$ plane which is consistent with having $\Lambda$ of the order of few TeV, as seen in Fig. 1.

In the case in which the lightest CP-even Higgs boson is purely doublet, that is when cos$\alpha = 0$, the heavier CP-even Higgs boson mass $m_H$ is unconstrained, but we do not have an improvement with respect to the SM as suggested by Eq. (11). On the contrary, in the opposite limit when the heaviest CP-even Higgs boson is purely doublet, it is constrained as the Higgs boson in the SM, since in this case it acts as “the” effective Higgs scalar. As the value of cos$\alpha$ is varied the EWPT constraint correspondingly weakens.

Searches for invisibly decaying Higgs bosons using the LEP-II data have been performed by the LEP Collaborations. For the channel $e^+e^- \rightarrow Zh \rightarrow Zb\bar{b}$ the final state is expressed in terms of the SM hZ cross section through

$$\sigma_{hZ \rightarrow b\bar{b}} = \sigma_{hZ}^{SM} \times R_{hZ} \times BR(h \rightarrow b\bar{b})$$

$$\sigma_{hZ}^{SM} \times C_{Z(h \rightarrow b\bar{b})}^Z,$$  

(13)

where $R_{hZ}$ is the suppression factor related to the coupling of the Higgs boson to the gauge boson $Z$ (i.e. $R_{hZ}^{SM} = 1$ and for the model we have $R_{hZ} = \cos^2 \alpha$). Here $BR(h \rightarrow b\bar{b})$ is the branching ratio of the channel $h \rightarrow b\bar{b}$ which in the model is modified with respect to the SM both by the mixing angle $\alpha$ and by the presence of the invisible Higgs boson decay into the Goldstone boson $J$ associated to the breaking of the global $U(1)$ symmetry, Eq. (11). For example in Ref. [8] DELPHI gives upper bounds for the coefficients $C_{Z(h \rightarrow b\bar{b})}^Z$ corresponding to a lightest CP-even Higgs boson mass from 15 GeV up to 100 GeV. We have analysed the regions of the mixing angle $\alpha$ and of the parameters $v_\sigma$ and $m_h$ which are currently allowed by the LEP-II searches. The results are illustrated in Fig. 2. One sees that as cos$\alpha \rightarrow 1$ h must obey the LEP-II SM Higgs boson limit while, in the oppo-
site limit where it becomes fully singlet, the limit quickly deteriorates.

There are many model realizations sharing the same simplest Higgs scalar potential in Eq. (2). Models differ depending on the details of the couplings relevant for neutrino masses. We assume that we have a non-supersymmetric model of neutrino masses [13]. A simple tree-level example is the “inverse seesaw” model introduced in [9], described by

$$\mathcal{L}_Y = Y_{ij} \nu_i^c L_j \Phi + M_{ij} \nu_i^c S_j + \lambda_{ij} S_i S_j \sigma$$  \hspace{1cm} (14)$$

In addition to the $SU(3)_c \otimes SU(2)_L \otimes U(1)$ singlet “right-handed” neutrinos it contains gauge singlet leptons $S_i$ which may arise in some string models [10]. No other terms are allowed by the symmetry, e. g., a direct Majorana mass term for the singlet fields $S_i$ is forbidden by lepton number and arises only due to the singlet scalar vev $\langle \sigma \rangle$. This will then give a $9 \times 9$ neutrino mass matrix, in the basis $(\nu_i, \nu_i^c, S_i)$:

$$M_\nu = \begin{pmatrix}
0 & Y_{\nu} v_\Phi & 0 \\
Y_{\nu}^T v_\Phi & 0 & M \\
0 & M^T & \mu
\end{pmatrix}$$  \hspace{1cm} (15)$$

so that the effective left-handed light neutrino mass matrix is

$$m_\nu = m_D^T M^{-1} \mu M^{-1} m_D.$$  \hspace{1cm} (16)$$

In the limit $\mu \to 0$ lepton number is restored and neutrino masses vanish. Thus its smallness is “natural” and $\langle \sigma \rangle$ can easily be of the order of TeV or less.

We have presented only the very simplest example of a class of low-scale neutrino mass models which can be made natural up to few TeV or so, in the absence of supersymmetry or other stabilizing mechanisms. These are a viable and attractive alternative to the seesaw mechanism. As explained in [2] and briefly reviewed in [3], there are many models with spontaneous breaking of lepton number at low scale, in which case the majoron is present. All of these lead to the generic signal in Eq. (1). Thus from this point of view the possibility of invisibly decays Higgs bosons must be taken seriously from the point of view of future colliders, LHC and ILC. Additional signals associated, say, to extended Higgs sector and charged scalars may also exist. These survive even if lepton number breaking is taken as explicit. A more extended investigation of these schemes will be presented elsewhere [12].

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* Electronic address: valle@ific.uv.es
URL: http://ahep.uv.es

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