Constraints on heavy Majorana neutrino phenomenology from the Vacuum Stability of the Higgs

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Abstract: The vacuum stability condition of the Standard Model Higgs potential with mass in the range of 124-127 GeV puts an upper bound on the Dirac mass of the neutrinos. We study this constraint with the right-handed neutrino masses upto TeV scale. The heavy neutrinos contribute to $\Delta L = 2$ processes like neutrinoless double beta decay and same-sign-dilepton production in the colliders. The vacuum stability criterion also restricts the light-heavy neutrino mixing and constrains the branching ratio of lepton flavour violating process, like $\mu \rightarrow e\gamma$ mediated by the heavy neutrinos. We show that neutrinoless double beta decay with a lifetime $\sim 10^{25}$ years can be observed if the the lightest heavy neutrino mass is $< 4.5$ TeV. We show that the vacuum stability condition and the experimental bound on $\mu \rightarrow e\gamma$ together put a constrain on heavy neutrino mass $M_R > 3.3$ TeV. Finally we show that the observation of same-sign-dileptons (SSD) associated with jets at the LHC needs much larger luminosity than available at present. We have estimated the possible maximum cross-section for this process at the LHC and show that with an integrated luminosity $100 \text{ fb}^{-1}$ it may be possible to observe the SSD signals as long as $M_R < 400$ GeV.
1. Introduction

The recent measurement of ATLAS and CMS \[1, 2, 3\] have confirmed the existence of a new boson which has mass in the range 126.5 GeV (ATLAS at 5.0\(\sigma\)) and 125.3 ± 0.6 GeV (CMS at 4.9\(\sigma\)), and it is expected to be a Standard Model Higgs. This mass range implies that quartic coupling \(\lambda_h\) of the Higgs has a value close to the vacuum stability limit \[4, 5, 6, 7, 8, 9, 10\]. The top-quark loop makes a negative contribution to the \(\beta\)-function of \(\lambda_h\) while the gauge couplings give a positive contribution. If the quartic coupling \(\lambda_h(\mu)\) becomes negative at large renormalization scale \(\mu\), it implies that in the early universe the Higgs potential would be unbounded from below and the vacuum would be unstable in that era. It has been pointed out that the Higgs mass in the 126 GeV range being close to the vacuum stability limit, one can put stringent constraints on new physics which affects the running of the Higgs quartic coupling.

One class model which can be constrained from the stability criterion of the Higgs coupling is the see-saw models of neutrino masses \[11, 12, 13, 14, 15, 16\]. In Type-I see-saw models \[17\] one introduces a number of heavy gauge singlet Majorana neutrinos which have Yukawa couplings with the Higgs and lepton doublets. The electroweak symmetry breaking gives rise to the Dirac mass matrix \(M_D\),

\[-\mathcal{L} = \bar{N}_R M_D \nu_L + \frac{1}{2} \bar{N}_R M_R N_R^c + \text{h.c.} \tag{1.1}\]

If \(M_D \ll M_R\) in the pure Type-I models \[17\] the light neutrino masses are given by \(M_\nu = M_D^T M_R^{-1} M_D\). It has been discussed earlier in many papers that light neutrino
masses which can explain the solar and atmospheric neutrino oscillations are obtained by assuming the eigenvalues $M_D \sim 100 \text{ GeV}$ and $M_R > 10^{14} \text{ GeV}$. By a suitable choice of $M_D$ and $M_R$ one can set $M_D^T M_R^{-1} M_D = 0$ and the light neutrino masses are given by higher order terms in $M_D^T M_R^{-1}$ [18, 19, 20]. In this way it is possible to generate viable light neutrino masses while reducing the scale $M_R$ to less than a TeV. In [13, 14], the constraints on various TeV-scale Type-I neutrino mass models from the vacuum stability criterion of Higgs coupling has been checked.

In this paper we assume a Yukawa couplings of heavy Majorana neutrinos with the lepton doublets with the Standard Model Higgs and that the heavy neutrinos masses are in the 100 GeV-10 TeV range. We obtain the constraints on the Higgs-neutrino Yukawa couplings by calculating the renormalisation group evolutions (RGEs) of $\lambda_h(\mu)$ (which is fixed at the electroweak scale by the Higgs mass). The vacuum stability condition is the requirement that $\lambda_h(M_W \leq \mu \leq M_P) \geq 0$. We find that this leads to the constraint $Y_\nu \leq 0.14$ on the elements of the Yukawa coupling matrix. We then apply this condition (which implies that the Dirac neutrino masses $M_D \leq 24.36 \text{ GeV}$) on the phenomenology of TeV scale heavy neutrinos [21, 22].

We study three aspects of the heavy neutrino phenomenology in the light of the vacuum stability condition on $Y_\nu$: (1) Neutrino-less double beta decay ($0\nu \beta\beta$), (2) Lepton flavor violating decays like $\mu \to e\gamma$, and (3) Same-sign-dilepton signals at the LHC. All these process depend upon the mixing of the light neutrino gauge eigenstates with the heavy neutrino mass eigenstates which is given by the mixing matrix $V \simeq M_D^T M_R^{-1}$. The Dirac mass $M_D$ gets an upper bound from the vacuum stability criterion which in turn puts constraints on the processes listed above. Our analysis is independent of the specific light neutrino mass model. We shall put upper bounds on various processes by assuming that the elements of $M_D$ which contribute to that particular signal are as large as can be allowed by the vacuum stability condition. Our choices for $M_D$ and $M_R$ may not give realistic neutrino mass through the usual Type-I seesaw which means that these signals can be further restricted by specific choice of the flavour structures.

In Section (2) we discuss the running of the Higgs quartic coupling in the Standard Model for the 125 GeV Higgs. In Section (3) we introduce the Yukawa couplings $Y_\nu$ between the Higgs and heavy neutrinos in the context of SM extended by heavy singlet fermion and study the effect of these neutrino Yukawa’s on the running of the Higgs quartic coupling. In this section we establish the bound on $Y_\nu$ from the stability criterion. In Section (4) we study $0\nu \beta\beta$, in Section (5) the lepton flavor violations, and in Section (6) we estimate the same-sign-dilepton signals at the LHC. Our results are summarized in the concluding section.

2. Vacuum stability of the Standard Model Higgs potential

The Higgs mass measured by ATLAS and CMS collaborations [1, 2, 3] is in the mass range 124.7 GeV-126.5 GeV is close to the bound on Higgs mass from of electro-weak vacuum stability condition [4]. In [5], it has been shown that in Standard Model, the Higgs boson quartic coupling $\lambda_h$ can remain positive upto Planck scale with appropriate choice of top
quark mass $m_t$, strong coupling constant $\alpha_s$ etc. The coupling $\lambda_h > 0$ ensures the stable vacuum and the bounded potential from below. Details of this study is performed recently in [10].

The RG-improved Higgs quartic term can be written as,

$$V_{\text{eff}} = \frac{\lambda_h(t)}{4!} [\xi(t)\phi]^4,$$

where $\xi$ signifies the wave function renormalisation and $t \sim \log(\mu/M_Z)$, $\mu$ is the scale of renormalisation. Here $\lambda_h(t)$ is the effective Higgs quartic coupling with loop corrections. Loop corrections can cause an instability of the potential if $\lambda(t)$ becomes negative at any scale $\mu < M_P$. The gauge boson loop makes a positive contribution whereas the top quark makes a negative contribution to the $\beta$ function of $\lambda_h$. Hence the instability of the potential mainly comes from loop-correction of top quark.

We compute the RG running of $\lambda_h(\mu)$ using the two loop RG equations available for the Standard Model [5, 9, 23]. We have also included the proper matching conditions at top pole mass [24]. The Fig. 1 shows the variation of $\lambda_h$ with different values top quark mass $m_t$. With increase of $m_t$, $\lambda_h$ becomes negative even before the Planck scale. For subsequent calculations, we have chosen different sets of top mass keeping Higgs mass constant, and vice-versa.

Now we move beyond the Standard Model by adding a heavy neutrino Yukawa coupling.

![Figure 1: RG of $\lambda_h$ for different values of $m_t$ in the Standard Model (with $m_h = 125$ GeV, $\alpha_s = 0.1184$).](image-url)
3. Higgs coupling with heavy neutrino

The Standard Model Higgs can couple to a singlet neutrino $N_R$ via the gauge invariant interaction term

$$-\mathcal{L}_Y = \mathcal{Y}_\nu LH N_R + \frac{1}{2} \bar{\mathcal{N}}_R M_R N_R^c + \text{h.c.,}$$

(3.1)

where $L = (\nu, l)^T$ is the lepton doublet, $H = (h^0, h^-)$ is the Higgs doublet and $N_R$ is right-handed singlet neutrino. $M_R$ are the Majorana masses for $N_R$. This interaction generates Dirac mass term, $M_D$, after electroweak symmetry breaking which reads as $M_D = Y_{\nu}v$ ($v = 174$ GeV).

In our further analysis we will not consider the flavour structures of both $Y_{\nu}$ and $M_R$, i.e., we will assume that $Y_{\nu} = Y_{\nu} \text{diag}(1, 1, 1)$ and right-handed neutrinos are degenerate, i.e. $M_R = M_R \text{diag}(1, 1, 1)$.

We will see that this new Yukawa coupling affects the RG evolutions of $\lambda_h$ and thus gets constrained from vacuum stability. This $Y_{\nu}$ also plays important role in the production and decays of $N_R$ leading to same-sign-dilepton associated with jets at the LHC.

The running of neutrino Yukawa coupling is as follows, [11, 13, 14]

$$\mu^\frac{d}{d\mu} \left( Y_{\nu}^\dagger Y_{\nu} \right) = \frac{1}{(4\pi)^2} Y_{\nu}^\dagger Y_{\nu} \left[ 6\lambda_t^2 + 2 \text{Tr} \left( Y_{\nu}^\dagger Y_{\nu} \right) - \frac{9}{10} g_1^2 + \frac{9}{10} g_2^2 \right] + 3 Y_{\nu}^\dagger Y_{\nu}. \quad (3.2)$$

\footnote{Once this neutral field $h^0$ acquires vacuum expectation value (vev) $v=174$ GeV the electroweak symmetry breaking occurs.}

\hspace{1cm}

**Figure 2:** RGE of $\lambda_h$ for different values of $m_h$ in Standard Model (with $m_t = 172.5$ GeV, $\alpha_s = 0.1184$).
The introduction of neutrino sector to Standard Model also modify the RG evolution of the Higgs quartic coupling $\lambda_h$ and Yukawa coupling of top quark $\lambda_t$ as follows.

The extra contribution for the singlet fermionic field to Higgs quartic coupling ($\lambda_h$) is

$$\hat{\beta}_{\lambda_h} = \frac{1}{(4\pi)^2} \left[ -4 \text{Tr}(Y_\nu Y_\nu Y_\nu Y_\nu^\dagger) + 4\lambda_h \text{Tr}(Y_\nu Y_\nu^\dagger) \right],$$

(3.3)

and to the top quark Yukawa coupling ($\lambda_t$) is

$$\hat{\beta}_{\lambda_t} = \frac{1}{(4\pi)^2} \left[ \text{Tr}(Y_\nu^\dagger Y_\nu) \right].$$

(3.4)

![Figure 3: Running of $\lambda_h$ for different values of neutrino Yukawa coupling $Y_\nu$ with $M_R = 0.1-1$ TeV, ($m_t = 172.5$ GeV, $\alpha_s = 0.1184$).](image)

Using these RG equations, the running of $\lambda_h$ for this model has been shown in Fig. 3. The impact of neutrino Yukawa coupling $Y_\nu$ on $\lambda_h$ is in similar fashion as the $\lambda_t$ and $\lambda_h$ becomes negative before Planck scale ($M_{\text{Planck}}$) with comparatively larger values of $Y_\nu$. We know there is an uncertainty in top mass measure measurement $173.2 \pm 0.9$ GeV [25] and $173.3 \pm 2.8$ GeV [26], and that feature has been grabbed in Fig. 1. In Fig. 2 we perform the evolutions of $\lambda_h$ for different Higgs masses choosing suitable top mass.

We outline the RGEs of $\lambda_h$ for different sets of $Y_\nu$ for $m_h = 124.7 - 126.5$ GeV and $m_t = 172.5$ GeV. We check the stability condition, defined as $\lambda(\mu \leq M_{\text{Planck}}) > 0$ and reveal that to avoid the instability of potential, the maximum value of the Yukawa coupling $Y_\nu$ at $\mu \sim \text{TeV}$ must be:

$$Y_\nu \leq 0.14.$$  

(3.5)
This upper limit of $Y_\nu$ sets the tolerance of the vacuum in this model. It has been noted that the light-heavy mixing parameter can be encapsulated in terms of the Dirac mass, $M_D \sim Y_\nu \nu$, see \cite{22,27}. In other words this mixing which in turn also affects the production and decay of the heavy Majorana neutrino gets constrained. Thus eventually this bound can be useful to adjudge the possibility of being probed or ruled out this TeV scale model at the LHC.

4. Gauge interactions of heavy neutrinos

We consider three generations of Standard Model $SU(2)_L$ lepton doublets $L_{iL} = (\nu_i, \ell_i)^T_L$, $(\ell = e, \mu, \tau)$ and three singlets $N_R$. The relation between the neutrino flavour and the mass eigenstates can be written as

$$\nu_{iL} = \sum_{i=1}^{3} U_{iL} \nu_i + \sum_{k=1}^{3} V_{ik} N_{kL}$$

(4.1)

$$U^\dagger U + V^\dagger V = I,$$

(4.2)

where the mixing between the light and heavy neutrinos is $V^\dagger V \simeq (M_D M_R^{-1})^2 = (v Y_\nu M_R^{-1})^2$. In terms of the mass eigenstates the charged current interaction vertices can be written as

$$-\mathcal{L}^c_{\text{int}} = \frac{g}{\sqrt{2}} W_\mu \left( \sum_{i=1}^{3} U_{iL}^* \bar{\nu}_i \gamma^\mu P_L \ell + \sum_{k=1}^{3} V_{ik}^* N_{kL} \gamma^\mu P_L \ell + \text{h.c.} \right)$$

(4.3)

Our phenomenological studies will involve $\Delta L = 2$ processes, like same-sign-dilepton (including $(0\nu \beta \beta)$) production at colliders where the source of the lepton number violation is the exchange of heavy Majorana neutrino. The coupling of the heavy neutrino to the charged leptons is parametrised by the mixing angles of $V_{ik}$. We use the upper bound of $Y_\nu$ from Eq. (3.5) to predict the parameter space where these processes may be observable. We also study lepton flavour violations like $\mu \rightarrow e \gamma$ whose upper limits are again restricted by Eq. (3.5).

4.1 Neutrinoless double beta decay

Neutrinoless double beta decay is one of the important phenomena to probe the lepton number violation. In this process, the lepton number violation occurs by two units. The half-life time of this process is also ascribed by this mixing $V_{iL}$ as following:

$$T_{1/2} = \frac{\kappa_{0\nu}}{\Gamma_{ee}} \left| \frac{(M_N)_{ee}}{(p^2)} - \frac{|V_{ei}|^2}{M_R} \right|^2,$$

(4.4)

where $\kappa_{0\nu} = G_{0\nu} (M_N m_p)^2$, nuclear matrix element (NME) for heavy neutrino, $M_N = 363 \pm 44$, $m_p$ is the proton mass, and $G_{0\nu} = 7.93 \times 10^{-15}$ yr$^{-1}$. We assume that the second term arising from the heavy neutrino mixing dominates and the mixing parameter $V_{ei}$ is explicitly related to the neutrino Yukawa coupling $Y_\nu$ via Dirac mass as:

$$|V_{ei}|^2 = \left| \left( M_D M_R^{-1} \right) \right|_{ee}^2,$$

(4.5)
and the relation Eq. (4.4) for half-life time for neutrinoless double beta decay as,

$$ T_{1/2}^{-1} \approx \frac{\kappa_{0\nu} |V_{ei}|^4}{M_R^2} = \frac{K_{0\nu}}{M_R^2} \left( M_D M_R^{-1} \right)_{ee}^4. $$

(4.6)

The experimental bound on half-life time is $T_{1/2} = 2.23^{+0.44}_{-0.31} \times 10^{25}$ yr in [28]. The study of vacuum stability gives $M_D \lesssim 24.36$ GeV. Using the values for $T_{1/2}$ and $M_D$, we can put the limit on the mass of the heavy neutrino,

$$ M_R \leq 4.5 \text{ TeV} $$

(4.7)

### 4.2 Lepton flavor violation

The mixing of active neutrinos with heavy neutrinos can give rise to lepton flavour violations (LFV) like $\mu \rightarrow e\gamma$ as shown in Fig. 3, if we generalise $M_R$ matrix to contain off-diagonal terms. Assuming the structure of $M_R$ matrix as:

$$ M_R^{-1} = M_R^{-1} \begin{pmatrix} 1 & \epsilon_1 & \epsilon_2 \\ \epsilon_1 & 1 & \epsilon_3 \\ \epsilon_2 & \epsilon_3 & 1 \end{pmatrix}, $$

(4.8)

where $\epsilon_i$s can be chosen to satisfy the correct light neutrino mixing angles.

We know among the $\ell_i \rightarrow \ell_j \gamma$ type LFV decays the $\mu \rightarrow e\gamma$ holds the most stringent bound on its decay branching ratio (BR) which is $2.4 \times 10^{-12}$ (Present) [29], and $1.0 \times 10^{-13}$ (Future) [30].

We estimate the branching ratio of this process from vacuum stability and check its compatibility with the existing direct bounds. This branching ratio for $\mu \rightarrow e\gamma$ is accompanied by the mixing $V_{\ell_i l} (l = e, \mu)$ between light to heavy neutrino [31, 21]:

$$ \text{Br} (\mu \rightarrow e\gamma) = \frac{3\alpha}{8\pi} \left| \sum_i V_{ei} V_{\mu i}^* \hat{g}(r) \right|^2, $$

(4.9)

where $\hat{g}(r) = r \left( 1 - 6r + 3r^2 + 2r^3 - 6r^2 \ln(r) \right) / (2 (1 - r)^4)$, and $r = M_R^2 / M_W^2$. 

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**Figure 4:** Neutrinoless double beta decay diagrams involving heavy Majorana field.
Again taking the constraint from vacuum stability $M_D \simeq 24.36$ GeV

$$\text{Br} (\mu \rightarrow e\gamma) = 2.82 \times 10^{-10} \left( \frac{M_D}{24.36 \text{ GeV}} \right)^4 \left( \frac{\text{TeV}}{M_R} \right)^4.$$

(4.10)

Taking the experimental bound $[29] \text{Br} (\mu \rightarrow e\gamma) < 2.4 \times 10^{-12}$ and if $M_D \simeq 24.36$ GeV (in order to give a sizable contribution to $(0\nu\beta\beta)$ and SSD signal at LHC) then we see that $M_R \geq 3.3$ TeV. This implies that in order to observe a $(0\nu\beta\beta)$ or like-sign-dilepton signals where we need $M_R$ to be small, the texture of $Y_\nu$ and $M_R$ should be such that the $e-\mu$ flavour mixing is small.

5. Same-Sign-Dilepton signal at LHC

The processes for same-sign-dilepton (SSD) production are similar to the neutrinoless double beta decay, see Fig. 3. These processes are of phenomenological importance as it involves both $e$ and $\mu$. The signal is identified as the same-sign-dileptons + $N$ jets, $N > 2$. The interaction vertices of the heavy neutrino ($N_R$) are suppressed by the mixing parameters $\sim \mathcal{O}(Y_\nu v/M_R)$. Assuming again a flavour diagonal $Y_\nu$ and degenerate $N_R$ we estimate the cross section for SSD at the LHC.

We have implemented this SM $\oplus$ Heavy Singlet neutrino model at Calchep $[32]$, and estimated the cross-section for the process $pp \rightarrow e^\pm e^\pm + \text{jets}$ and $pp \rightarrow \mu^\pm \mu^\pm + \text{jets}$. We have considered the range of $M_R$ to be 0.1-1 TeV, and no flavour structure for the simplification of study. It has been noted that in the Fig. 4 (left) the amplitude is suppressed more $((M_D/M_R)^4)$ than the other diagram Fig. 4 (right) (here the suppression is $\mathcal{O}(M_D/M_R)^2$). The choice of our $M_R$ is such that the mixing is much smaller than 1, and that dictates us to work safely with the Fig. 4 (right).

In earlier section we have noted the maximum $M_D = Y_\nu v$ from the vacuum stability of the Standard Model Higgs field. In this section we have use that limit and estimate the largest possible maximum cross-section for the process Fig. 6 (right) with two different sets of center of mass energy at the LHC. These two cross-sections are calculated with center of mass energy $(\sqrt{s})$ 7 TeV see Fig. 6 and 14 TeV see Fig. 7.
Figure 6: Production cross-section Fig. 4 (right) with $\sqrt{s}=7$ TeV in the LHC.

Figure 7: Production cross-section of the process Fig. 4 (right) with $\sqrt{s}=14$ TeV in the LHC.

In recent paper by ATLAS the Standard Model background has been estimated at 2.1 $fb^{-1}$ luminosity. As shown in Fig. 5 and Fig. 6 the vacuum stability puts a stringent bound on the production cross-section, through the $M_D$, the maximum allowed cross-
section is 49.02 fb at 7 TeV center of mass energy. This is the maximum cross-section that one attains using no cuts. But due to the stringent constraint from the demand of vacuum stability the allowed cross-section is quite small that yet LHC does not have enough data to see the process compared to the SM background [33]. Thus we have to wait for future data with 14 TeV center of mass energy and large integrated luminosity ($L = \int \mathcal{L} dt = \sim 100 \ f b^{-1}$). The cross section for the SSD process at the LHC with $\sqrt{s}=14$ TeV is shown in Fig. 5 and the region above the ‘thick (red)’ line is disallowed by the vacuum stability. In Fig. 5 we see that ‘shaded (cyan)’ region is the one accessible at the LHC with $\sqrt{s}=14$ TeV at $L=100 \ f b^{-1}$ considering atleast 3 events over the zero background, i.e, at 95% C.L. Taking into account the vacuum stability condition it may be possible to observe SSD signal at LHC if $M_R < 400$ GeV.

6. Conclusion

In this paper we have focused on the vacuum stability of the Higgs field in a specific scenario where the Standard Model is extended by singlet Majorana fermions. We have studied the impact of such new field that couples to the light neutrinos via the SM Higgs doublet on the RG evolution of the Higgs quartic coupling ($\lambda_h$), and we show that expectedly this new coupling ($Y_{\nu}$) lowers the scale $\mu$ at which $\lambda_h(\mu)$ becomes negative. In this study the aim is to find the maximum value of $Y_{\nu}$ which is compatible with the vacuum stability with heavy neutrino field having mass $M_R \sim$ TeV.

We showed that the vacuum stability condition constrains the Dirac mass (which we have taken to be degenerate) to be $M_D \leq 24.36$ GeV. We studied $\Delta L = 2$ processes like $(0\nu\beta\beta)$ and same-sign-dileptons at LHC and lepton flavour violating processes like $\mu \rightarrow e\gamma$ taking into account the vacuum stability bound on $M_D$ which restrict the mixing between the light and heavy neutrinos which goes as $M_D/M_R$.

We find that in order to observe $(0\nu\beta\beta)$ signal which saturates the experimental bound $T_{1/2} = 2.23^{+0.44}_{-0.31} \times 10^{25} \ \text{yr}$ [28] the heavy neutrinos must have a mass $M_R < 4.5$ TeV.

For the LFV process $\mu \rightarrow e\gamma$ if we assume $M_D$ at the largest possible value 24.36 GeV from vacuum stability (to maximise the chances for other signals), then we get the constraint $M_R > 3.3$ TeV. It may be possible to evade this bound on $M_R$ by choosing the texture of $M_D$ and $M_R$ matrices such that $e - \mu$ mixing is suppressed.

Finally estimated the maximal cross-section for the signal, like same-sign-dilepton associated with jets imposing the vacuum stability condition. We show that the data attained with $2.1 \ f b^{-1}$ integrated luminosity cannot rule out right-handed neutrinos as the vacuum stability criterion shows that the dilepton signal would be way below the SM background. It may be possible to observe the SSD at the LHC with $\sqrt{s}=14$ TeV and integrated luminosity of $100 \ f b^{-1}$ as long as $M_R < 400$ GeV. If a larger signal is seen at the LHC then it would be a sign of new physics beyond SM + sterile right-handed neutrinos.

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