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Quark see-saw, Higgs mass and vacuum stability

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Abstract. The issue of vacuum stability of standard model (SM) is discussed by embedding it within the TeV scale left–right quark see-saw model. The Higgs potential in this case has only two coupling parameters (λ_1, λ_2) and two mass parameters. There are only two physical neutral Higgs bosons (h, H), the lighter one being identified with the 126 GeV Higgs boson. We explore the range of values for (λ_1, λ_2) for which the vacuum is stable for all values of the Higgs fields till 10^{16} GeV. Combining with the further requirement that the scalar self-couplings remain perturbative till 10^{16} GeV, we find (i) an upper and lower limit on the second Higgs (H) mass to be within the range: $0.4 \leq (M_H/v_R) \leq 0.7$, where v_R is the parity breaking scale and (ii) the masses of heavy vector-like top, bottom and τ partner fermions (P_3, N_3, E_3) have an upper bound $\leq v_R$. These predictions can be tested at LHC and future higher energy colliders.

Keywords. Left–right symmetry; see-saw mechanism; vacuum stability; Higgs.

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1. Introduction

The discovery of the 126 GeV Higgs boson at the LHC [1] has provided the striking final confirmation of the standard electroweak model of Glashow, Weinberg and Salam. But the observed value of the Higgs mass has raised an interesting issue that if there is no new physics below 10^{10} GeV or so, the scalar self-coupling of the Higgs boson, λ, turns negative above this scale, making the SM vacuum unstable at high temperatures [2]. This near critical value of λ can cause the Universe, at some far future epoch, to make a transition to the deeper minimum [3], a not very desirable prospect and has led to speculations that there must be new physics nearby that would stabilize this vacuum and avoid this possibility. In this paper, we discuss salient features of one such minimal possibility which was presented by us in a recent paper [4].

We consider an extension of the Standard Model where quark and charged lepton masses arise from a generalized see-saw mechanism (we call it quark see-saw or universal see-saw) [5], via the introduction of a new set of TeV or higher mass vector-like SM singlet fermions, that provide the see-saw 'counterweight'. A natural setting for the universal see-saw is not the Standard Model but one with an extended gauge sector based on the gauge group $G_{LR} \equiv SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ with parity symmetry [6]. Symmetry breaking in this model is implemented by two Higgs doublets – one, a doublet under $SU(2)_L$ and a second one which is a doublet under $SU(2)_R$. This set-up prevents direct Yukawa couplings between the left and right chiral SM quarks, making quark see-saw an essential element of the model. The left–right quark see-saw model (denoted here by SLRM) has the advantage that it has a particularly simple Higgs sector, i.e., only one extra right-handed doublet Higgs boson beyond the SM Higgs field. It is therefore different from many multi-Higgs extension of SM discussed in the literature. After symmetry breaking, the model has only two neutral Higgs fields, one of which can be identified with the SM Higgs field (the 126 GeV Higgs boson). This model has the additional advantage that it also provides a solution to the strong CP problem without an axion [7] and for a low right-handed scale ($\leq$100 TeV), protects [8] this solution from possible large Planck scale effects [9].

As the model is based on the gauge group $G_{LR} \equiv SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ with parity symmetry, the Higgs potential of the model has only one extra scalar coupling parameter compared to SM. The parity symmetry is assumed to be softly broken by the mass terms of the Higgs doublets, so that parity is a technically 'natural' symmetry [10]. As noted, this model has only two physical neutral Higgs fields and no extra charged ones. We denote the two Higgs self-scalar couplings by (λ_1, λ_2) and analyse the renormalization group evolution of these couplings to address the stability of the ground state of the theory that breaks the full gauge symmetry down to $U(1)_{\rm em}$. We find a stable vacuum and a perturbative theory all the way upto 10^{16} GeV, which therefore presents a solution to the vacuum stability problem.

As far as the masses of the heavy vector-like fermions go, in principle, the masses of all but the top partner fermion field could be large but in this paper, we consider both the right-handed scale and all the vector-like fermion masse to be in the TeV range in analysing the vacuum stability issue. This makes the model amenable to experimental tests at the Large Hadron Collider [10a].

We find that the solution to the stability problem of SM vacuum, puts a lower limit on the mass of second neutral Higgs boson of the model. The requirement that the scalar self-couplings do not 'blow up' till the GUT scale of 10^{16} GeV, imposes an upper bound on the second Higgs mass. Combining these we get, $0.4 \leq (M_H/v_R) \leq 0.7$, where v_R is the parity breaking scale. A second consequence of vacuum stability requirement is that the masses of heavy vector-like top, bottom and τ partner fermion (P_3, N_3, E_3) have an upper bound, i.e. $M^{\rm max}_{P_3,N_3,E_3} \leq v_R$. We then discuss some aspects of the heavy and light Higgs boson phenomenology in the model. We find an interesting relation between the heavy Higgs boson decay modes and hh, WW, ZZ which is characteristic of the model and may be used to test it.

This paper is organized as follows: in §2, we present the basic ingredients of the model including the scalar potential and the neutral Higgs masses in the unitary gauge; in §3, we present the renormalization group equations for different couplings of the model; in §4, we present the phenomenology of the heavy Higgs field including its production at LHC and its decay modes. In §5, we give some comments on the neutrino mass profiles in our model. We summarize our results in §6.

2. Left–right see-saw model (SLRM)

2.1 *Particle assignment*

As this model is a TeV scale embedding of SM in the left–right model with quark and charged lepton see-saw [5], the SM fermions plus the right-handed neutrinos are assigned to doublets of the left- and right-handed $SU(2)$s, according to their chirality as in standard left–right models. We add four kinds of vector-like fermions (P, N, E, $\mathcal{N}$), one set per generation, to our model to generate fermion masses

$$\begin{aligned} & Q_L\left(2,\ 1,\ \frac{1}{3}\right);\ Q_R\left(1,\ 2,\ \frac{1}{3}\right); \\ & \Psi_L\ (2,\ 1,\ -1)\,;\ \Psi_R\ (1,\ 2,\ -1)\,; \\ & P_{L,\,R}\left(1,\ 1,\ \frac{4}{3}\right);\ N_{L,\,R}\left(1,\ 1,\ -\frac{2}{3}\right); \\ & E_{L,\,R}\ (1,\ 1,\ -2)\,;\ \mathcal{N}_{L,\,R}\ (1,\ 1,\ 0)\,, \end{aligned} \tag{1}$$

where Q and Ψ are the quark and lepton doublets, respectively, and (Q, P, N) are colour $SU(3)_c$ triplets, while the remaining fields are singlets. The scalar field content of the left–right see-saw model [5] consists of only one additional Higgs doublet. They transform under the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ as follows:

$$\chi_L = \begin{pmatrix} \chi_L^+ \\ \chi_L^0 \end{pmatrix} \in (2,\ 1,\ 1), \quad \chi_R = \begin{pmatrix} \chi_R^+ \\ \chi_R^0 \end{pmatrix} \in (1,\ 2,\ 1).$$

The scalar potential in our model is given by

$$\begin{aligned} V =& -\mu_L^2 \chi_L^\dagger \chi_L - \mu_R^2 \chi_R^\dagger \chi_R \\ & +\lambda_1[(\chi_L^\dagger \chi_L)^2 + (\chi_R^\dagger \chi_R)^2] + \lambda_2(\chi_L^\dagger \chi_L)(\chi_R^\dagger \chi_R)\,. \end{aligned} \tag{2}$$

Note that parity symmetry in the above equation is softly broken so that left–right symmetry is natural [10]. When $\mu_{L,R}^2 > 0$, the full gauge symmetry breaks down to $U(1)_{\rm em}$ at the minimum of the potential:

$$\chi_L = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_L \end{pmatrix}, \quad \chi_R = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_R \end{pmatrix} \tag{3}$$

and we obtain the minimization conditions

$$\frac{v_L^2}{2} = \frac{\lambda_2 \mu_R^2 - 2\lambda_1 \mu_L^2}{\lambda_2^2 - 4\lambda_1^2}, \quad \frac{v_R^2}{2} = \frac{\lambda_2 \mu_L^2 - 2\lambda_1 \mu_R^2}{\lambda_2^2 - 4\lambda_1^2}\,. \tag{4}$$

Diagonalization of the CP-even Higgs mass matrix (in the limit of $v_R \gg v_L$) leads to two mass eigenvalues

$$M_h^2 = 2\lambda_1\left(1 - \frac{\lambda_2^2}{4\lambda_1^2}\right) v_L^2, \quad M_H^2 = 2\lambda_1 v_R^2\,. \tag{5}$$

2.2 *Yukawa interactions and fermion masses*

The Yukawa interactions responsible for fermion masses in this model are given by

$$- \mathcal{L}_{\rm Y} = \bar{Q}_L Y_u \tilde{\chi}_L P_R + \bar{Q}_L Y_d \chi_L N_R + \bar{\Psi}_L Y_e \chi_L E_R + (L \leftrightarrow R)$$
$$+ \bar{P}_L M_P P_R + \bar{N}_L M_N N_R + \bar{E}_L M_E E_R + \text{h.c.}, \tag{6}$$

where $\tilde{\chi}_{L,R} = i\tau_2 \chi^*_{L,R}$. Note that due to the left–right gauge invariance, there is no direct coupling between the left- and right-handed chiral light quarks as would have been the case for the Standard Model gauge group with heavy vector-like quarks e.g., [12]. We do not include the $\mathcal{N}$ couplings and discuss it at the end of the paper separately. In the above equation, Y_a and M_a ($a = u, d, e$) are matrices with complex elements, so that theory has CP violation. For simplicity of discussion, we assume all Yukawa couplings to be real and note that our discussion of the Higgs sector and vacuum stability is not affected by this.

In the SLRM, all the quarks obtain their masses from the see-saw mechanism, e.g., for the top sector alone,

$$\begin{pmatrix} 0 & (1/\sqrt{2}) Y_t v_L \\ (1/\sqrt{2}) Y_t v_R & M_{P_3} \end{pmatrix}, \tag{7}$$

which leads to generic see-saw-type mass relations:

$$m_{q_a} \simeq \frac{Y_a^2 v_L v_R}{2 M_a} . \tag{8}$$

Most interesting consequence of the see-saw relation is for the top quark. First of all, the relevant Yukawa coupling Y_t for top quark can differ from that in SM, depending on v_R and the mass of the heavy P_3 fermion. For example, if $M_{P_3} \gg v_R$, then Y_t can be much larger than one. In addition to making the theory non-perturbative, large values of Y_t will also lead to gross instability of the vacuum, the very problem we are addressing. We therefore carefully analyse the dependence of Y_t for different values of v_R and M_{P_3}. As we are exploring TeV scale physics, we shall keep v_R also in the few TeV range. As shown in figure 1, for v_R and M_P in the range of few TeV, Y_t is generally larger than its corresponding SM value at v_R scale. In combination with RGE analysis, this helps us to put an upper bound on Y_t and hence an upper bound on the top partner mass M_{P_3}. We find that in the entire allowed parameter space of our model, $M_{P_3} \leq v_R$.

3. Renormalization group evolutions (RGE) of couplings and vacuum stability

In this section, first we present the RGE equations, below and above the heavy fermion mass M_F and $SU(2)_R$ symmetry scale (v_R) and then we study their implications for vacuum stability. For simplicity, both M_F and v_R are chosen to be very near to each other and in the TeV range. We use only one matching scale v_R as by virtue of our assumption, all new particles beyond SM start contributing at this scale to the RGEs.

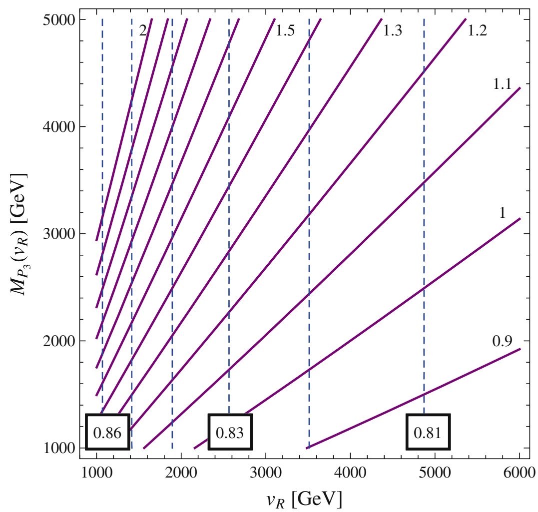

Figure 1. The purple solid lines indicate the values of Y_t as a function of v_R and M_{P_3} and the vertical blue dashed lines are the top quark Yukawa coupling in the SM as a function of v_R. For this plot, we have used the top quark mass $m_t(m_t) = 163.3$ GeV.

3.1 *RGEs below and above the heavy fermion and right-handed scale*

Below the heavy fermion and right-handed scale, the SM can be viewed as the effective theory of the SLRM. We therefore use the SM β functions till v_R as follows [13,14]. Note that our $U(1)_Y$ gauge coupling is not normalized as in GUT theories.

Case 1. $\mu \leq v_R, M_F$

$$
\begin{aligned}
\beta(g') &= \frac{1}{16\pi^2}\left[\left(\frac{10}{9}n_f + \frac{1}{6}\right) g'^3\right],\\
\beta(g) &= \frac{1}{16\pi^2}\left[-\left(\frac{43}{6} - \frac{2}{3}n_f\right) g^3\right],\\
\beta(g_s) &= \frac{1}{16\pi^2}\left[-\left(11 - \frac{2}{3}n_f\right) g_s^3\right],\\
\beta(\lambda) &= \frac{1}{16\pi^2}\left[\frac{9}{8}\left(\frac{1}{3}g'^4 + \frac{2}{3}g'^2 g^2 + g^4\right) + 24\lambda^2 - 2Y_4\right.\\
&\qquad \left. - \lambda(3g'^2 + 9g^2) + 4\lambda Y_2\right],\\
\beta(h_t) &= \frac{1}{16\pi^2}\left[-h_t\left(\frac{17}{12}g'^2 + \frac{9}{4}g^2 + 8g_s^2\right) + \frac{3}{2}h_t(h_t^2 - h_b^2) + h_t Y_2\right],\\
\beta(h_b) &= \frac{1}{16\pi^2}\left[-h_b\left(\frac{5}{12}g'^2 + \frac{9}{4}g^2 + 8g_s^2\right) + \frac{3}{2}h_b(h_b^2 - h_t^2) + h_b Y_2\right],\\
\beta(h_\tau) &= \frac{1}{16\pi^2}\left[-\frac{9}{4}h_\tau\left(\frac{5}{3}g'^2 + g^2\right) + \frac{3}{2}h_\tau^3 + h_\tau Y_2\right],
\end{aligned} \tag{9}
$$

with n_f, the number of flavours, and

$$Y_2 = 3h_t^2 + 3h_b^2 + h_\tau^2 \,,$$
$$Y_4 = 3h_t^4 + 3h_b^4 + h_\tau^4 \,. \tag{10}$$

Above the (v_R, M_F) scales (which we assume to be nearly the same), due to the extended gauge interaction and the heavy vector-like fermions, the β functions are substantially different (note that we have a different set of Yukawa couplings from the effective SM theory, though they are closely correlated, and see the matching conditions below for the normalization of g_{BL}).

Case 2. $\mu \geq v_R, M_F$

$$\begin{aligned}
\beta(g_{BL}) &= \frac{1}{16\pi^2}\left[\frac{41}{2}g_{BL}^3\right],\\
\beta(g) &= \frac{1}{16\pi^2}\left[-\frac{19}{6}g^3\right],\\
\beta(g_s) &= \frac{1}{16\pi^2}\left[-3g_s^3\right],\\
\beta(\lambda_1) &= \frac{1}{16\pi^2}\left[\frac{9}{8}\left(\frac{3}{4}g_{BL}^4 + g_{BL}^2 g^2 + g^4\right) + \left(24\lambda_1^2 + 2\lambda_2^2\right) - 2\tilde{Y}_4\right.\\
&\qquad \left. - \lambda_1\left(\frac{9}{2}g_{BL}^2 + 9g^2\right) + 4\lambda_1 \tilde{Y}_2\right],\\
\beta(\lambda_2) &= \frac{1}{16\pi^2}\left[\frac{27}{16}g_{BL}^4 + (24\lambda_1\lambda_2 + 4\lambda_2^2)\right.\\
&\qquad \left. - \lambda_2\left(\frac{9}{2}g_{BL}^2 + 9g^2\right) + 4\lambda_2 \tilde{Y}_2\right],\\
\beta(Y_t) &= \frac{1}{16\pi^2}\left[\frac{3}{2}Y_t\left(Y_t^2 - Y_b^2\right) - Y_t\left(\frac{17}{8}g_{BL}^2 + \frac{9}{4}g^2 + 8g_s^2\right) + Y_t \tilde{Y}_2\right],\\
\beta(Y_b) &= \frac{1}{16\pi^2}\left[\frac{3}{2}Y_b\left(Y_b^2 - Y_t^2\right) - Y_b\left(\frac{5}{8}g_{BL}^2 + \frac{9}{4}g^2 + 8g_s^2\right) + Y_b \tilde{Y}_2\right],\\
\beta(Y_\tau) &= \frac{1}{16\pi^2}\left[\frac{3}{2}Y_\tau^3 - \frac{9}{4}Y_\tau\left(\frac{5}{2}g_{BL}^2 + g^2\right) + Y_\tau \tilde{Y}_2\right],
\end{aligned} \tag{11}$$

with

$$\tilde{Y}_2 = 3Y_t^2 + 3Y_b^2 + Y_\tau^2 \,,$$
$$\tilde{Y}_4 = 3Y_t^4 + 3Y_b^4 + Y_\tau^4 \,. \tag{12}$$

In order to run the couplings to ultrahigh energy scales, we have to match all the seemingly effective SM couplings to that in the full scenario of SLRM. For simplicity and concreteness, the following matchings are considered at the right-handed scale v_R:

(1) Let us start with the gauge couplings. The matching conditions for strong and weak couplings are trivial, while the matching of $U(1)$ gauges is as follows:

$$\frac{1}{\alpha_Y(v_R)} = \frac{3}{5}\frac{1}{\alpha_{I_{3R}}(v_R)} + \frac{2}{5}\frac{1}{\alpha_{BL}(v_R)} \tag{13}$$

with

$$\alpha_Y = \frac{\tilde{g}'^2}{4\pi}, \quad \alpha_{I_{3R}} = \frac{g^2}{4\pi}, \quad \alpha_{BL} = \frac{\tilde{g}_{BL}^2}{4\pi}, \tag{14}$$

where $\tilde{g}'$ and $\tilde{g}_{BL}$ are the normalized couplings in the context of GUT,

$$\tilde{g}' = \sqrt{\frac{5}{3}}g', \quad \tilde{g}_{BL} = \sqrt{\frac{2}{3}}g_{BL}\,. \tag{15}$$

Below, the normalized $\tilde{g}_{BL}$ is denoted in eq. (11) simply as g_{BL}.

(2) To obtain matching conditions for the quartic scalar couplings λ and $\lambda_{1,2}$, we integrate out the heavy scalar at the scale of its mass (approximately the right-handed scale) from the potential [11]. To the linear order of v_L/v_R, the mass term, triple coupling term and quartic coupling term point have the same matching relationship as implied in eq. (5),

$$\lambda(v_R) = \lambda_1(v_R)\left(1 - \frac{\lambda_2^2(v_R)}{4\lambda_1^2(v_R)}\right). \tag{16}$$

This simple relation has deeper phenomenological implications than just being superficially the matching condition: it means evidently that, at the right-handed scale, λ_1 is always larger than the SM quartic coupling λ (or we can roughly say that λ is increased by the SM scalar interacting with its 'right-handed' partner), which potentially help to solve the stability problem of the SM vacuum.

(3) The matching relation for the Yukawa couplings is somewhat straightforward due to the see-saw mechanism eq. (8),

$$\frac{h_f(v_R)}{\sqrt{2}} \simeq \frac{Y_f^2(v_R)v_R}{2M_F}, \tag{17}$$

with $f = t,\ b,\ \tau$ and F their corresponding heavy partners. In the numerical running of the RGEs, we shall resort to the exact relations, as large Yukawa couplings, especially for the top quark, would invalidate such a simple approximation.

3.2 *Vacuum stability and universal see-saw*

The Standard Model has only one Higgs field and the stability vacuum requires that the scalar coupling λ must satisfy the positivity condition $\lambda(\mu) > 0$ for all values of the mass μ. However, when λ is extrapolated to large μ using renormalization group equations, the negative contribution of the top quark coupling turns it negative around 10^{10} GeV for $M_h = 126$ GeV for which $\lambda(m_W) \simeq 0.131$. This is the vacuum stability problem. In contrast, in the SLRM, the presence of the extra 'right-handed' Higgs doublet χ_R implies a new scalar coupling λ_2 and the vacuum stability condition requires that not only $\lambda_1 > 0$ but also $2\lambda_1 + \lambda_2 > 0$ and both conditions must be maintained for all values of μ (or Higgs field). As mentioned above, by choosing λ_2 appropriately, we can increase the value of λ_1 at the v_R scale without conflicting with the observed Higgs mass. However, it cannot be made arbitrarily large because it would then hit the Landau pole when extrapolated to the GUT scale. This means that λ_1 must have an upper bound.

We assume that the left–right symmetric theory at the TeV scale that we consider here, is a 'low-energy' effective phenomenological manifestation of some GUT theory at ultra-high energy scales. We therefore assume that the couplings remain perturbative only up to generic GUT scale (10^{16} GeV) but not to the higher Planck scale. Note that we do not mean that our model necessarily unifies to a single GUT group at 10^{16} GeV. Unification of this model is a highly model-dependent issue.

To be specific, in the numerical running, we set the heavy mass parameters for the third-generation to be the same, i.e.,

$$M_F = M_{P_3} = M_{N_3} = M_{E_3} \,. \tag{18}$$

Note that this does not necessarily mean that the three third-generation partners have the same mass eigenvalues (especially the mass eigenvalue of the top quark partner is significantly different from the other two), as they also get contribution from mixing with the SM fermions. At v_R scale, with v_R fixed, the Yukawa couplings are solely determined by the value of M_F (figure 2).

Given a value of v_R, we have only two free parameters in the SLRM: the quartic coupling λ_1 (λ_2 is fixed by the SM Higgs mass) and the universal heavy fermion mass parameter M_F. We also assume the masses of the other generation vector-like fermion masses to be the same as the third-generation one but their Yukawa couplings are small

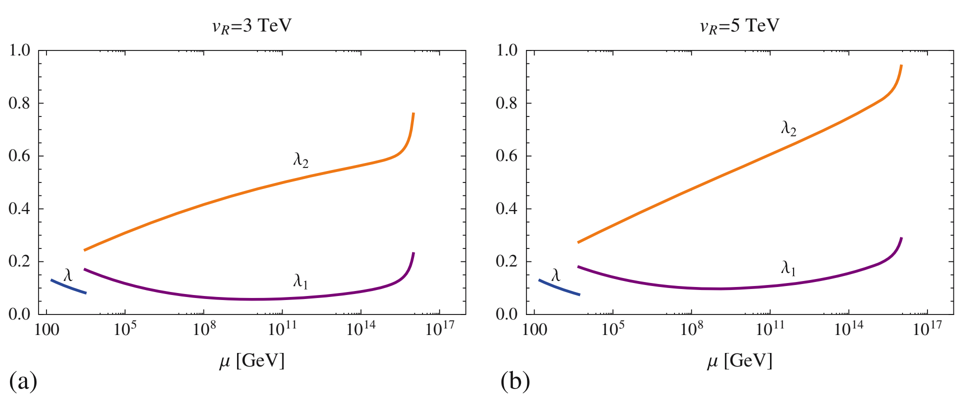

Figure 2. Examples of running of the quartic couplings λ and $\lambda_{1,\,2}$, which are allowed by both the stability and perturbativity constraints. (**a**) We set $v_R = 3$ TeV, $\lambda_1(v_R) = 0.17$ and $M_F = 1.2$ TeV and (**b**) we set $v_R = 5$ TeV, $\lambda_1(v_R) = 0.18$ and $M_F = 2$ TeV. For simplicity, we assume M_F/v_R to be nearly the same in both figures. We have of course chosen the Yukawa coupling parameters in accordance with this choice.

and therefore they do not affect our results. We scanned the full parameter space, varying v_R (near the TeV scale), $\lambda_1(v_R)$ and M_F.

Scanning of the full parameter space reveals first that at the v_R scale, the quartic coupling λ_1 is severely constrained: $\lambda(v_R) < \lambda_1(v_R) \lesssim 0.25$. As pointed out above, the value of λ_1 has to be large enough to compensate the negative contributions of Y_t to $\beta(\lambda_1)$, and yet small enough to keep out of the non-perturbative region. This constraint implies that the heavy Higgs mass is predicted to be in the range of about $[\sqrt{2\times 0.1},\ \sqrt{2\times 0.25}]\, v_R \simeq [0.4,\ 0.7]\, v_R$. We also find that the upper limit on this ratio is nearly independent of v_R, while the lower limit has a weak dependence on v_R and M_F (for smaller v_R the lower limit is increased somewhat). All these facts point to the phenomenological implication that there exists a heavy Higgs in the SLRM at the TeV scale, as explicitly depicted in figure 3. In the plot, we considered only the constraints from vacuum stability and perturbativity, but not that from the heavy fermion masses. It is interesting that the heavy Higgs boson in the SLRM is potentially detectable at the LHC (and in future high-energy colliders); in the next section, we shall study the LHC phenomenology of this predicted new particle.

We stress here that the constraints given above are obtained with a positive λ_2 from eq. (16). We also examined the case with a negative λ_2 as Higgs mass does not depend on the sign of λ_2. As expected, negative λ_2 tends to push the vacuum towards instability, worsening the SM stability problem. Thus, the allowed parameter space shrinks greatly. To keep the stability conditions up to the GUT scale, M_F is required to be small. As the aforementioned examples show, if $v_R = 3$ TeV, we require $M_F \lesssim 650$ GeV, while for $v_R = 5$ TeV, we get $M_F \lesssim 1100$ GeV. The ATLAS and CMS Collaborations have searched for vector-like quarks both with charges $2/3$ and $-1/3$ [15–18], and the most stringent bound at the moment on our model is $M_B \gtrsim 590$ GeV (B is the vector-like quark with charge $-1/3$) [17], which sets a lower limit on the negative λ_2 case: $v_R \gtrsim 2.8$ TeV. With the future search for vector-like quarks at 14 TeV LHC [19], the limit could get much stronger. Comparatively, the positive case is much less constrained and thus phenomenologically preferred. Thus, we consider mainly the positive case in this work.

It has important phenomenological significance as it predicts $M_F < v_R$, or the existence of heavy fermions, the heavy partners of b and τ fermion, below the right-handed scale. This is presented in figure 3b. This coincides with the findings of ref. [11] although these

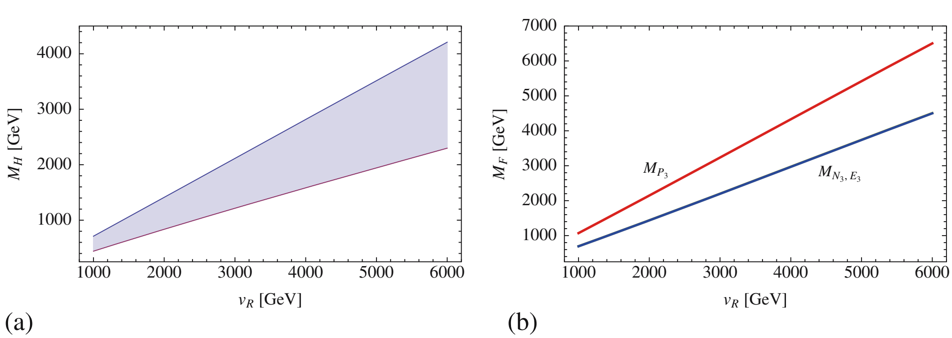

Figure 3. (**a**) Constraints on the heavy Higgs mass M_H as function of v_R (the shaded region is allowed) from vacuum stability and perturbativity. (**b**) Upper bounds on the masses M_{P_3} and M_{N_3,E_3} of heavy vector-like fermions as a function of v_R.

are strictly two different scenarios within the left–right framework. The mass of the top quark partner is significantly larger than the other two partners because of the large top quark Yukawa coupling, which contribute substantially to the top partner mass. In contrast, the contribution to the masses of N_3 and E_3 from mixing with the SM partners are much smaller and can be safely neglected.

4. Light and heavy Higgs phenomenology

In this section, we discuss the implications of the model for heavy (H) and light 126 GeV Higgs boson (h) for collider phenomenology.

4.1 *h-decay*

As mentioned in the previous section, below the right-handed scale, all the new heavy particles beyond SM (the gauge bosons, the heavy Higgs and the vector-like fermions) are integrated out, and the SM is left as the low-energy effective theory. The effects of new physics on SM Higgs decay can be generally neglected, at least to the next-to-leading order of v_L/v_R, e.g., for the bottom quark channel,

$$\begin{aligned} -\mathcal{L} &\simeq \frac{1}{\sqrt{2}}\bar{b}_L Y_b h B_R + \frac{1}{\sqrt{2}}\bar{B}_L Y_b H b_R + \text{h.c.} \\ &\Rightarrow \frac{1}{\sqrt{2}}\sin\alpha_R^b\, \bar{b}_L^m Y_b h b_R^m + \text{h.c.} \end{aligned} \tag{19}$$

Here $B = N_3$ is the heavy partner, b^m is the bottom mass eigenstate and α_R^b is the right mixing angle of the bottom quark with its heavy partner. Approximately, $\sin\alpha_R^b \simeq (1/\sqrt{2})$ $Y_b v_R/M_F$ and we recover the SM bottom quark Yukawa coupling via the see-saw relation $(1/\sqrt{2})y_b = Y_b^2 v_R/2M_F$. For the top quark coupling, although the see-saw relation might not be a good approximation (for $Y_t v_R \sim M_F$), a more exact formula reveals that we can obtain again the same Yukawa coupling as in SM. Phenomenologically, the gluon fusion production and diphoton production processes, in which the top quark loop plays an important role, are not affected in the SLRM [19a].

4.2 *Triple Higgs coupling*

Another possible effect of beyond the Standard Model physics is on the triple Higgs coupling [20]. To see if there is any such effect, let us define the unitary mixing matrix that diagonalizes the mass matrix of the two Higgs bosons as

$$\begin{pmatrix} h \\ H \end{pmatrix} = U \begin{pmatrix} h_L \\ h_R \end{pmatrix}. \tag{20}$$

The equation giving U is

$$U \cong \begin{pmatrix} 1 & -\dfrac{\lambda_2}{2\lambda_1}\dfrac{v_L}{v_R} \\ \dfrac{\lambda_2}{2\lambda_1}\dfrac{v_L}{v_R} & 1 \end{pmatrix}. \tag{21}$$

From this, we get the triple couplings from the potential

$$\lambda_1 \left[v_L h_L^3 + v_R h_R^3\right] + \frac{1}{2}\lambda_2 \left[v_L h_L h_R^2 + v_R h_L^2 h_R\right]$$
$$\Rightarrow \lambda_1 v_L h^3 \left[1 - \left(\frac{\lambda_2}{2\lambda_1}\right)^2\right]. \tag{22}$$

With the relation given in eq. (5), the triple coupling is the same as in SM.

4.3 *Production and decay of the heavy Higgs at LHC*

The decay channels of the heavy Higgs in the SLRM model are given below. We discuss them one by one.

(1) $H \to hh$: for the scalar channel, the LO coupling $m_{Hhh}Hhh$ is given by $m_{Hhh} \simeq \frac{1}{2}\lambda_2 v_R$, with the exact value

$$m_{Hhh} = \frac{1}{2}\varepsilon(6\lambda_1 + (\varepsilon^2 - 2)\lambda_2)v_L$$
$$+\frac{1}{2}(6\varepsilon^2\lambda_1 + (1 - 2\varepsilon^2)\lambda_2)v_R, \tag{23}$$

where $\varepsilon = (\lambda_2/2\lambda_1)(v_L/v_R)$ is the mixing of 'left-handed' and 'right-handed' scalars. The decay width is then given by

$$\Gamma(H \to hh) = \frac{1}{8\pi}\frac{m_{Hhh}^2}{M_H}\left(1 - \frac{4m_h^2}{M_H^2}\right)^{1/2}. \tag{24}$$

(2) $H \to t\bar{t}$: for the fermion channel, we assume that the heavy Higgs boson is not heavy enough to decay into the vector-like fermion pairs but decays only into the SM fermions (this corresponds to a large region in the parameter space and there is no fine-tuning for the assumption). Amongst the couplings to the SM fermions, the top quark is expected to be the largest one. We start with the original Lagrangian given below:

$$-\mathcal{L} = \frac{1}{\sqrt{2}}\bar{t}_L Y_t h_L T_R + \frac{1}{\sqrt{2}}\bar{T}_L Y_t h_R t_R + \text{h.c.}$$
$$\Rightarrow \frac{1}{\sqrt{2}}\bar{t}_L^m H t_R^m \cdot Y_t(\varepsilon\cos\alpha_L^t \sin\alpha_R^t + \sin\alpha_L^t \cos\alpha_R^t)$$
$$\simeq \frac{1}{\sqrt{2}}\bar{t}_L^m H t_R^m \cdot Y_t(\varepsilon\sin\alpha_R^t + \sin\alpha_L^t\cos\alpha_R^t). \tag{25}$$

Here $T \equiv P_3$ is the top quark partner and t^m is the mass eigenstate. For a large top Yukawa coupling, the left-handed mixing α_L^t is generally very small, but the right-handed one α_R^t is always very large (generally of order one), as $Y_t v_R \sim M_F$. Denoting the Yukawa coupling $y_{Ht\bar{t}} = Y_t(\varepsilon\sin\alpha_R^t + \sin\alpha_L^t\cos\alpha_R^t)$ which is suppressed by the scalar mixing ε or the left-handed mixing α_L^t, the decay width is given by

$$\Gamma(H \to t\bar{t}) = \frac{3}{16\pi} \cdot y_{Ht\bar{t}}^2 M_H \left(1 - \frac{4m_t^2}{M_H^2}\right)^{3/2}. \tag{26}$$

(3) $H \to WW, ZZ$: In the SLRM, the gauge bosons W_L and W_R do not mix at tree level, but the scalars do; thus we can get the suppressed coupling $m_{HWW} = 2\varepsilon M_W^2/v$ with $v = v_L$ being the SM electroweak scale. For the decay width, we get

$$\Gamma(H \to WW) = \frac{1}{8\pi} \frac{m_{HWW}^2}{M_H} \left(1 - \frac{4m_W^2}{M_H^2}\right)^{1/2} \times \left[1 + \frac{1}{2}\left(1 - \frac{M_H^2}{2m_W^2}\right)^2\right]. \tag{27}$$

The width for the ZZ boson channel is similar (through the neutral gauge bosons Z and Z' mix at tree level but the mixing is suppressed by $(v_L/v_R)^2$),

$$\Gamma(H \to ZZ) = \frac{1}{16\pi} \frac{m_{HZZ}^2}{M_H} \left(1 - \frac{4m_Z^2}{M_H^2}\right)^{1/2} \left[1 + \frac{1}{2}\left(1 - \frac{M_H^2}{2m_Z^2}\right)^2\right], \tag{28}$$

with $m_{HZZ} = 2\varepsilon M_Z^2/v$.

As the heavy Higgs boson is expected to be close to the right-handed scale, which is much larger than the electroweak scale, we can approximate the decay widths and see what happens in the massive limit $v_R \to \infty$. In this limit, the fermion channel is suppressed by $(v_L/v_R)^2$ as long as $M_F \sim v_R$, while the expression for other channels are very simple, determined only by the parameters v_R, λ_1 and λ_2,

$$\begin{aligned}
\Gamma(H \to hh) &= \frac{1}{8\pi} \frac{\lambda_2^2}{4\sqrt{2\lambda_1}} v_R, \\
\Gamma(H \to WW) &= \frac{1}{8\pi} \frac{\lambda_2^2}{2\sqrt{2\lambda_1}} v_R, \\
\Gamma(H \to ZZ) &= \frac{1}{8\pi} \frac{\lambda_2^2}{4\sqrt{2\lambda_1}} v_R.
\end{aligned} \tag{29}$$

The suppression factor ε for the gauge boson channels is cancelled by the large enhancement factor $M_H^4/M_{W(Z)}^4$ from the interaction with the longitudinal components of gauge bosons. Ultimately, it is from the scalar interaction and is therefore not suppressed as these Goldstone bosons are 'eaten' by the gauge bosons. In this limit, we find a relation among these different decay widths, which we call 'the quartering rule' of heavy Higgs decay, whose origin lies in the coupling of H with the four components of χ_L before electroweak symmetry breaking. This is explicitly presented in figure 4. This extraordinary feature could be a smoking gun signal of the SLRM. The diphoton channel of the heavy Higgs decay $H \to \gamma\gamma$ is predominately mediated by the right-handed W boson, the top

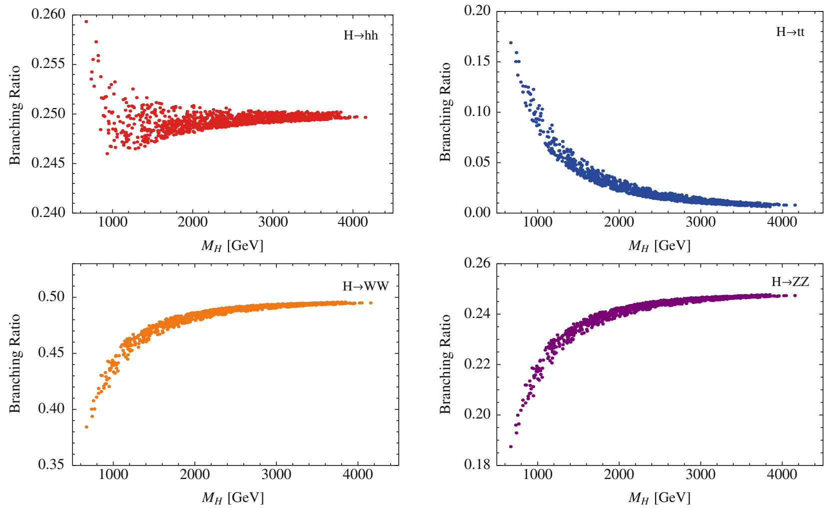

Figure 4. Branching ratios of heavy Higgs decay. In these plots we do not include the cases in which the heavy fermion pair channel(s) is kinematically allowed.

quark and its heavy partner. Numerical calculation reveals that the branching ratio of this channel is generally of order 10^{-5}. Even if the heavy Higgs is observed at colliders, it will be challenging to detect it in this specific channel.

For the heavy Higgs production at LHC, the dominant channel is the gluon fusion process via the top partner loop. The Yukawa coupling involved is approximately $Y_t \sin\alpha_R^t$; as stated above, this right-handed fermion mixing angle is generally very large, of order one and therefore this production process is not suppressed whereas the top loop is relatively suppressed by the scalar mixing angle ε or the left-handed fermion mixing angle α_L^t, as shown in eq. (25). The scatter plot of the production cross-section is depicted in figure 5. For a heavy Higgs with a mass of 1 TeV, with 100 fb^{-1} of 14 TeV data, we can

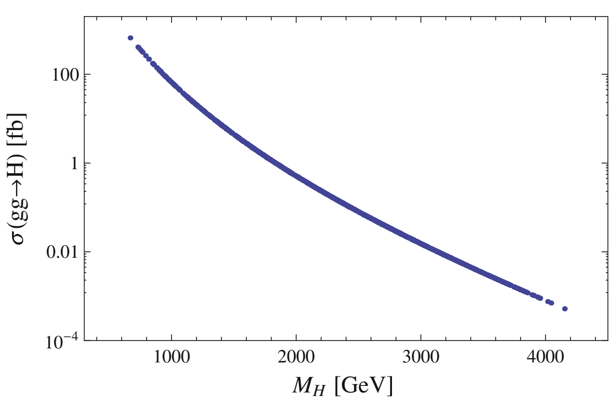

Figure 5. Heavy Higgs production cross-section $\sigma(gg \to H)$ at LHC with a centre-of-mass energy of 14 TeV, as a function of H mass.

expect thousands of heavy Higgs to be produced at LHC. For heavier H, the cross-section drops rapidly.

We also wish to note that if $M_H > 2M_F$, new decay modes open up. However, for a large range of parameters of the model, the mass of heavy Higgs boson is not large enough to produce heavy top partner pairs. On the other hand, the heavy bottom and tau partners, which are lighter than heavy top partner, could in principle be produced but these channels are suppressed by the small scalar or light–heavy fermion mixings. Therefore, the heavy fermion pair channels are always suppressed, with the branching ratio generally of order 10^{-3}. We also note that, the $Z - Z'$ mixing effects are suppressed by $M_Z^2/M_{Z'}^2$ and are therefore very small. We ignore these effects here.

5. Neutrinos

In this section, we briefly address the scale of neutrino masses in the universal see-saw models. The simplest option is to introduce the vector-like gauge singlet field $\mathcal{N}_{L,R}$ with both Dirac mass $M_{\mathcal{N}}$ and Majorana masses $M_{L,R}$ for $\mathcal{N}$ fields. The neutrino mass matrix in this case reads, on the basis of $(\nu, \mathcal{N}, \nu^C, \mathcal{N}^C)$ (where all fields are left handed) as

$$\begin{pmatrix} 0 & 0 & 0 & (1/\sqrt{2})Yv_L \\ 0 & M_L & (1/\sqrt{2})Y^T v_R & M_N \\ 0 & (1/\sqrt{2})Yv_R & 0 & 0 \\ (1/\sqrt{2})Y^T v_L & M_N & 0 & M_R \end{pmatrix}. \tag{30}$$

In the parameter regime where $M_R \sim M_L \gg M_N \gg Yv_R \gg Yv_L$, the light left-handed neutrino masses are given by $\mathcal{M}_\nu \sim -\frac{1}{2}v_L^2 Y M_R^{-1} Y^T$. For $M_N \leq M_{L,R}$, the formula is roughly

$$\mathcal{M}_\nu \simeq -\frac{1}{2}v_L^2 Y \left(M_R - M_N^T M_L^{-1} M_N\right)^{-1} Y^T \tag{31}$$

and for the right-handed neutrinos (ν^cs), replace $L \leftrightarrow R$ in the above formulae. Naively, one might think that in the Majorana alternative, the right-handed neutrino masses will be (v_R^2/v_L^2) times those of the left-handed neutrinos (roughly 100 times larger). However, this is true only if parity symmetry is exact. If we take the Majorana mass terms for $\mathcal{N}$, $\mathcal{N}^c$ to be different and therefore break parity softly, they could have very different forms i.e., mass scales as well as textures. Therefore, by adjusting these terms, one can make the right-handed (ν^c) mass terms in the 10–100 GeV range, and keep them in conformity with cosmology and low-energy weak constraints. As an example, consider the case where the magnitudes of all elements of M_R are in the range of 10^{10} GeV and those of M_L are in the TeV range. In this case, the light 'left-handed' neutrinos can have sub-eV masses as observed with right-handed neutrino masses being in the 100 GeV range. As the neutrinos in this case are Majorana fermions, they would give rise to neutrinoless double beta decay. Our goal in this paper is simply to demonstrate that getting small neutrino masses does not pose any challenge to the viability of these models.

6. Summary

We have discussed the question of vacuum stability in the left–right see-saw embedding of Standard Model with the universal see-saw implemented by TeV scale vector-like fermions. This model has only one extra scalar Higgs coupling beyond the Standard Model and it helps to stabilize the electroweak vacuum till GUT scale. This model has only two neutral Higgs bosons. Identifying the lighter of them with the 126 Higgs boson of Standard Model, the heavy Higgs mass is found to be below the v_R scale. For parity-breaking scale in the few TeV range, it can be accessible at the LHC. We discuss its collider phenomenology such as production cross-section and decay properties. We also find that the vector-like top, bottom and τ partners P_3, N_3, E_3 are below v_R making them LHC accessible. We also find an interesting relation between the three heavy Higgs boson decay modes: $H \to hh$, WW, ZZ, which can provide a test of this model once the heavy Higgs boson is discovered.

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