

Hefty MSSM-like light Higgs in extended gauge models

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Abstract

It is well known that in the MSSM the lightest neutral Higgs h^0 must be, at the tree level, lighter than the Z boson and that the loop corrections shift this stringent upper bound up to about 130 GeV. Extending the MSSM gauge group in a suitable way, the new Higgs sector dynamics can push the tree-level mass of h^0 well above the tree-level MSSM limit if it couples to the new gauge sector. This effect is further pronounced at the loop level and h^0 masses in the 140 GeV ballpark can be reached easily. We exemplify this for a sample setting with a low-scale $U(1)_R \times U(1)_{B-L}$ gauge symmetry in which neutrino masses can be implemented via the inverse seesaw mechanism.

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I. INTRODUCTION

With the recent start of the LHC the hunt for the Higgs has recommenced and new limits for the Standard Model (SM) Higgs bosons have been obtained excluding various mass ranges between 145 GeV and 470 GeV [1, 2]. At the same time, a slight excess of events in the mass window at around 140 GeV has been observed. Unlike in the SM, there is a stringent tree-level upper bound on the mass of the lightest CP-even Higgs state in the Minimal Supersymmetric extension of the Standard Model (MSSM) $m_{h^0} \leq \text{Min}(m_A, m_Z)$ where m_A denotes the mass of the pseudoscalar Higgs and m_Z is the mass of the SM Z -boson. This comes from the fact that supersymmetry links the MSSM Higgs self-interactions to the gauge couplings (via the so-called D -terms), thus making the scalar potential of the theory rather rigid. It is well known [3–6] that radiative corrections are important in the Higgs sector of the MSSM. However, even at 2-loop and 3-loop level one obtains an upper bound of $m_{h^0} \lesssim 130$ GeV [7–18] for SUSY particles below roughly $m_{SUSY} \sim 2$ TeV and a top quark mass of about 173 GeV. If the above hint for a Higgs in the range of 140 GeV turns out to be correct, this would essentially rule out the MSSM, except perhaps for split-SUSY-like scenarios [19] where the squarks and sleptons are typically pushed far above the TeV scale. One popular way to resolve this issue is to add additional Higgs fields, e.g., the NMSSM-like singlet(s) [20], so that extra F -term contributions to the scalar potential lift the Higgs mass. Another option is to enhance the D -terms by employing extended gauge symmetries [21–25].

Besides the Higgs mass puzzle (assuming the Higgs boson exists at all) there is yet another eminent mass-related riddle in the particle physics, namely, why neutrinos are so much lighter than all other matter particles. In the “standard” seesaw picture [26–29] this is attributed to a new very-high-energy dynamics (often in the vicinity of the GUT scale) which, at low energies, exhibits itself as a dimension-five effective operator giving the SM neutrinos a lepton-number violating Majorana mass [30, 31]. On the other hand, there is nothing really fundamental about such high-energy realizations of the seesaw mechanism as, in principle, variants of seesaw can be implemented at virtually any mass scale; inverse seesaw proposed in [32] or the linear seesaw of [33] are just two examples. Such schemes are naturally realized in the class of left-right symmetric extensions of the SM [34, 35] based on the popular $SO(10)$ breaking chains

$$SO(10) \rightarrow SU(4)_C \times SU(2)_R \times SU(2)_L \rightarrow SU(3)_c \times SU(2)_L \times U(1)_Y \quad (1)$$

$$SO(10) \rightarrow SU(3)_c \times SU(2)_R \times SU(2)_L \times U(1)_{B-L} \rightarrow SU(3)_c \times SU(2)_L \times U(1)_Y \quad (2)$$

which, however, often call for further extension of the matter sector in order to maintain the near-perfect gauge coupling unification of the MSSM [36]. A potential problem in this context is that the gauge couplings can easily become non-perturbative well below the GUT scale as matter particles always yield a positive contribution to the corresponding beta functions [37]. Nevertheless, one can still devise models where an extended gauge symmetry can remain unbroken down to almost the electroweak scale and, at the same time, a perturbative unification of the gauge couplings at around 10^{16} GeV is retained [38].

In such extended models the MSSM Higgs bosons are often charged also under the additional group factor(s). Hence, extra D -terms contributing to the masses of the neutral Higgs bosons are naturally supplied. This implies that the Higgs boson, which is mainly the MSSM h^0 , can have a mass above the Z -boson mass already at the tree level. Moreover, the additional Higgs fields needed to break the extended gauge symmetry mix with the usual two Higgs doublets of the MSSM, thus affecting the phenomenology of the Higgs sector. In particular, the couplings of the mainly $SU(2)_L$ -doublet-like Higgs bosons to the SM particles can get reduced and one can find parameter regions where the lightest MSSM-like Higgs boson h^0 decays into two of the additional Higgs bosons, as they can be very light without violating existing experimental bounds [39].

In this letter we exemplify this basic mechanism at the Higgs sector of a simple extension of the SM featuring a $U(1)_R \times U(1)_{B-L}$ gauge symmetry which gets broken to $U(1)_Y$ at energies close to the electroweak scale. The $U(1)_R$ factor can be viewed as, e.g., a remnant of a complete gauged $SU(2)_R$ symmetry that can be restored at higher energies, thus facilitating a possible embedding into a full GUT based, for instance, on an $SO(10)$ gauge symmetry [34]. To this end, we complement the existing literature in several aspects; namely, by performing a complete one-loop analysis of the light Higgs sector and by checking that the light Higgs phenomenology is fully consistent with the current data by inspecting carefully all the relevant constraints from colliders (in particular, those coming from the LEP and LHC searches) and lepton flavour violation ($\mu \rightarrow e\gamma$). In particular, we consider not only shifts in the masses of the lightest Higgs CP-even eigenstates but also the changes in their character, i.e., the amount of the $SU(2)_L$ doublet components within, and their implications.

In the next section we present the details of the model focusing namely on its extended Higgs sector. Working out the relevant mass matrices we argue that the CP-even mass eigenstate most similar to the lightest MSSM Higgs boson h^0 can have a mass above 100 GeV already at the tree level. We also briefly discuss the one-loop corrections to the tree-level situation. In Section III we present results of a dedicated numerical analysis where the complete one-loop corrections to the Higgs sector were taken into account. Finally, we draw our conclusions in Section IV.

II. THE MODEL AND ITS HIGGS SECTOR

We shall consider a sample model based on the $SU(3)_c \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$ gauge group which can emerge, e.g., in a class of $SO(10)$ GUTs broken along the “minimal” left-right symmetric chain

$$SO(10) \rightarrow SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \rightarrow SU(3)_c \times SU(2)_L \times U(1)_R \times U(1)_{B-L},$$

as advocated, for instance, in [34]. The main virtue of this setting is that, even in the minimally fine-tuned version, an MSSM-like gauge unification is perfectly compatible with a $U(1)_R \times U(1)_{B-L}$ stage stretching down to a few TeV. Renormalization group evaluation usually lead to $U(1)$ mixing effects [40] if more than one $U(1)$ factor is present which not

	Superfield	$SU(3)_c \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$	Generations
Matter	\hat{Q}	$(\mathbf{3}, \mathbf{2}, 0, +\frac{1}{6})$	3
	\hat{d}^c	$(\bar{\mathbf{3}}, \mathbf{1}, +\frac{1}{2}, -\frac{1}{6})$	3
	\hat{u}^c	$(\bar{\mathbf{3}}, \mathbf{1}, -\frac{1}{2}, -\frac{1}{6})$	3
	\hat{L}	$(\mathbf{1}, \mathbf{2}, 0, -\frac{1}{2})$	3
	\hat{e}^c	$(\mathbf{1}, \mathbf{1}, +\frac{1}{2}, +\frac{1}{2})$	3
	$\hat{\nu}^c$	$(\mathbf{1}, \mathbf{1}, -\frac{1}{2}, +\frac{1}{2})$	3
	\hat{S}	$(\mathbf{1}, \mathbf{1}, 0, 0)$	3
Higgs	\hat{H}_u	$(\mathbf{1}, \mathbf{2}, +\frac{1}{2}, 0)$	1
	\hat{H}_d	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2}, 0)$	1
	$\hat{\chi}_R$	$(\mathbf{1}, \mathbf{1}, +\frac{1}{2}, -\frac{1}{2})$	1
	$\hat{\bar{\chi}}_R$	$(\mathbf{1}, \mathbf{1}, -\frac{1}{2}, +\frac{1}{2})$	1

TABLE I: The Matter and Higgs sector field content of the $U(1)_R$ model under consideration. Matter generation indices have been suppressed. The \hat{S} superfields are included to generate neutrino masses via the inverse seesaw mechanism. Under matter parity, the matter fields are odd while the Higgses are even. The additive charges were conveniently normalized in such a way that $Y = T_R + B - L$ and $Q = T_L^3 + Y$.

only introduces additional couplings in the gauge sector but also affects the evolution of the soft SUSY breaking parameters [41]. However, it turns out that for the model under study the corresponding effects are small [42]. Thus, we neglect these additional couplings in this letter.

The transformation properties of all the matter and Higgs superfields are summarized in Table I. The relevant R -parity¹ conserving superpotential is given by

$$W = Y_u \hat{u}^c \hat{Q} \hat{H}_u - Y_d \hat{d}^c \hat{Q} \hat{H}_d + Y_\nu \hat{\nu}^c \hat{L} \hat{H}_u - Y_e \hat{e}^c \hat{L} \hat{H}_d + \mu \hat{H}_u \hat{H}_d - \mu_R \hat{\bar{\chi}}_R \hat{\chi}_R + Y_s \hat{\nu}^c \hat{\chi}_R \hat{S} \quad (3)$$

where Y_e , Y_d and Y_u are the usual MSSM Yukawa couplings for the charged leptons and the quarks. In addition there are the neutrino Yukawa couplings Y_ν and Y_s ; the latter mixes the ν^c fields with the S fields giving rise to an inverse seesaw mechanism for neutrino masses. For completeness we note that for realistic neutrino masses and mixing angles one needs also a $\mu_S \hat{S} \hat{S}$ term with a small parameter μ_S which, however, hardly affects the Higgs sector and, thus, is omitted here for simplicity. Its effect for the phenomenology will be discussed elsewhere [42]. Note that, besides the role it plays in neutrino physics, the Y_s coupling is relevant also for the Higgs phenomenology at the loop level as it enters the mixing of χ_R and $\bar{\chi}_R$ Higgs fields with the $SU(2)_L$ Higgs doublets. The fields χ_R and $\bar{\chi}_R$ can be viewed as the (electric charge neutral) remnants of $SU(2)_R$ doublets, which remain light in the spectrum

¹ More precisely, an effective R -parity is implemented by means of an extra Z_2 matter parity.

when the $SU(2)_R$ gauge factor is broken by the VEV of a $B-L$ neutral triplet down to the $U(1)_R$.

Following the notation and conventions of [43] the soft SUSY breaking Lagrangian reads

$$V_{soft} = \sum_a M_a \tilde{G}_a \tilde{G}_a + \sum_{ij} m_{ij}^2 \phi_i^* \phi_j + T_u \tilde{u}_R^* \tilde{Q} H_u - T_d \tilde{d}_R^* \tilde{Q} H_d + T_\nu \tilde{\nu}_R^* \tilde{L} H_u - T_e \tilde{e}_R^* \tilde{L} H_d + B_\mu H_u H_d - B_{\mu_R} \bar{\chi}_R \chi_R + T_s \tilde{\nu}_R^* \chi_R \tilde{S}. \quad (4)$$

The first sum runs over all gauginos for the different gauge groups and the second one contains the scalar masses squared.

The $U(1)_R \times U(1)_{B-L}$ gauge symmetry is spontaneously broken to the hypercharge $U(1)_Y$ by the VEVs v_{χ_R} and $v_{\bar{\chi}_R}$ of the scalar components of the $\hat{\chi}_R$ and $\hat{\bar{\chi}}_R$ superfields while the subsequent $SU(2)_L \otimes U(1)_Y \rightarrow U(1)_Q$ is governed by the VEVs v_d and v_u of the neutral scalar components of the $SU(2)_L$ Higgs doublets \hat{H}_d and \hat{H}_u . Thus, one can write

$$\chi_R = \frac{1}{\sqrt{2}} (\sigma_R + i\varphi_R + v_{\chi_R}), \quad \bar{\chi}_R = \frac{1}{\sqrt{2}} (\bar{\sigma}_R + i\bar{\varphi}_R + v_{\bar{\chi}_R}), \quad (5)$$

$$H_d^0 = \frac{1}{\sqrt{2}} (\sigma_d + i\varphi_d + v_d), \quad H_u^0 = \frac{1}{\sqrt{2}} (\sigma_u + i\varphi_u + v_u). \quad (6)$$

where the generic symbols σ and φ denote the CP-even and CP-odd components of the relevant fields, respectively.

Let us mention at this point that in order to avoid a decoupling of the beyond-MSSM gauge and Higgs sectors, one has to assume that the $U(1)_R \times U(1)_{B-L}$ breaking VEVs v_{χ_R} and $v_{\bar{\chi}_R}$ are not very far from the electroweak scale. This, however, facilitates the simplified approach to the Higgs sector analysis where the desired $SU(2)_L \times U(1)_R \times U(1)_{B-L} \rightarrow U(1)_{QED}$ transition is treated as a one-step breaking.

At the tree level we find that in the $(\varphi_d, \varphi_u, \bar{\varphi}_R, \varphi_R)$ basis the pseudoscalar sector has a block-diagonal form²

$$M_{AA}^2 = \begin{pmatrix} M_{AA,L}^2 & 0 \\ 0 & M_{AA,R}^2 \end{pmatrix} \quad (7)$$

with

$$M_{AA,L}^2 = B_\mu \begin{pmatrix} \tan \beta & 1 \\ 1 & \cot \beta \end{pmatrix}, \quad M_{AA,R}^2 = B_{\mu_R} \begin{pmatrix} \tan \beta_R & 1 \\ 1 & \cot \beta_R \end{pmatrix}, \quad (8)$$

$\tan \beta = v_u/v_d$ and $\tan \beta_R = v_{\chi_R}/v_{\bar{\chi}_R}$. From these four states two are Goldstone bosons which become the longitudinal parts of the massive neutral vector bosons Z and a Z' . In the physical spectrum there are two pseudoscalars A^0 and A_R^0 with masses

$$m_A^2 = B_\mu (\tan \beta + 1/\tan \beta), \quad m_{A_R}^2 = B_{\mu_R} (\tan \beta_R + 1/\tan \beta_R) \quad (9)$$

² Note that this remains to be the case even if the kinetic mixing effects are turned on.

where the first formula is identical to the MSSM case. For later convenience we define

$$v_R^2 = v_{\chi_R}^2 + v_{\bar{\chi}_R}^2, \quad v^2 = v_d^2 + v_u^2. \quad (10)$$

The tree-level CP-even Higgs mass matrix in the $(\sigma_d, \sigma_u, \bar{\sigma}_R, \sigma_R)$ basis reads

$$M_{hh}^2 = \begin{pmatrix} m_{LL}^2 & m_{LR}^2 \\ m_{LR}^{2,T} & m_{RR}^2 \end{pmatrix}, \quad (11)$$

where

$$m_{LL}^2 = \begin{pmatrix} g_Z^2 v^2 c_\beta^2 + m_A^2 s_\beta^2 & -\frac{1}{2} (m_A^2 + g_Z^2 v^2) s_{2\beta} \\ -\frac{1}{2} (m_A^2 + g_Z^2 v^2) s_{2\beta} & g_Z^2 v^2 s_\beta^2 + m_A^2 c_\beta^2 \end{pmatrix}, \quad (12)$$

$$m_{LR}^2 = \begin{pmatrix} g_R^2 v v_R c_\beta c_{\beta_R} & -g_R^2 v v_R c_\beta s_{\beta_R} \\ -g_R^2 v v_R s_\beta c_{\beta_R} & g_R^2 v v_R s_\beta s_{\beta_R} \end{pmatrix}, \quad (13)$$

$$m_{RR}^2 = \begin{pmatrix} g_{Z_R}^2 v_R^2 c_{\beta_R}^2 + m_{A_R}^2 s_{\beta_R}^2 & -\frac{1}{2} (m_{A_R}^2 + g_{Z_R}^2 v_R^2) s_{2\beta_R} \\ -\frac{1}{2} (m_{A_R}^2 + g_{Z_R}^2 v_R^2) s_{2\beta_R} & g_{Z_R}^2 v_R^2 s_{\beta_R}^2 + m_{A_R}^2 c_{\beta_R}^2 \end{pmatrix}, \quad (14)$$

$s_x = \sin(x)$, $c_x = \cos(x)$ ($x = \beta, \beta_R, 2\beta, 2\beta_R$), $g_Z^2 = (g_L^2 + g_R^2)/4$, $g_{Z_R}^2 = (g_{BL}^2 + g_R^2)/4$ and g_L , g_R and g_{BL} are gauge couplings associated to the $SU(2)_L$, $U(1)_R$ and $U(1)_{B-L}$ gauge factors, respectively. The matrix m_{LL}^2 contains the standard MSSM doublet mass matrix³, m_{RR}^2 corresponds to the $U(1)_R \times U(1)_{B-L}$ Higgs bosons and m_{LR}^2 provides the essential mixing among the two sectors. It is in particular this sector which gives rise to the increase of the mass of the MSSM-like lighter Higgs boson already at tree-level overcoming the stringent MSSM bound. In what follows we shall denote the orthogonal matrix diagonalizing the mass matrix in eq. (11) by R and the eigenvalues/eigenstates will be ordered in such a way that $m_i \leq m_j$ for $i < j$. In a full analogy to the MSSM, the entire Higgs spectrum can be parametrized in terms of the pseudoscalar masses m_A and m_{A_R} and the relevant mixings encoded by $\tan \beta$ and $\tan \beta_R$.

The similarity to the MSSM makes it also clear that the loop corrections can be potentially large and, thus, very important. In particular, top and stop loops affect the $SU(2)_L$ -doublet part of the mass matrix in the usual manner. Furthermore, in certain parts of the parameter space also the neutrino/sneutrino loops can be large. Technically, we have been using the **SARAH** package [44, 45] to obtain the relevant $SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L}$ generalizations of the basic MSSM formulae given in [46]; in this respect, let us stress that this accounts for the full one-loop structure of the Higgs sector. For further details an interested reader should refer to a dedicated work [42].

Finally, besides direct bounds from various Higgs boson searches, there is an important constraint on the parametric space of the model associated to the heavy Z' . In the

³ To see this explicitly one has to integrate out the additional Higgs fields in the $v_R \rightarrow \infty$ limit which yields a shift in the gauge couplings such the the MSSM limit is achieved.

(W_3^0, B, B') basis (corresponding to the electrically neutral generators of $SU(2)_L$, $U(1)_R$ and $U(1)_{B-L}$, respectively) the relevant vector boson mass matrix reads

$$M_{VV}^2 = \frac{1}{4} \begin{pmatrix} g_L^2 v^2 & 0 & -g_L g_R v^2 \\ 0 & g_{BL}^2 v_R^2 & -g_{BL} g_R v_R^2 \\ -g_L g_R v^2 & -g_{BL} g_R v_R^2 & g_R^2 (v^2 + v_R^2) \end{pmatrix}, \quad (15)$$

from where the masses of the photon, Z and Z' are readily identified

$$m_\gamma = 0, \quad m_Z^2 = \frac{1}{8} \left(A - \sqrt{A^2 - 4B} \right), \quad m_{Z'}^2 = \frac{1}{8} \left(A + \sqrt{A^2 - 4B} \right), \quad (16)$$

where

$$A = (g_L^2 + g_R^2) v^2 + (g_{BL}^2 + g_R^2) v_R^2, \quad B = [g_L^2 (g_R^2 + g_{BL}^2) + g_{BL}^2 g_R^2] v^2 v_R^2. \quad (17)$$

In particular, the product $g_{Z'R}^2 v_R^2 [= m_{Z'}^2 + O(v^2/v_R^2)]$ is constrained by the Z' searches at LEP and at Tevatron as well as from the precision measurements [47, 48].

III. NUMERICAL RESULTS

The numerical results given below have been calculated in **SPheno** [49, 50] for which the necessary subroutines and input files were generated by the relevant extension of **SARAH** [51]. Hence, the complete one-loop corrections in the extended Higgs sector have been included [42]. We will concentrate the discussion on the lightest two mass eigenstates, since here the changes with respect to the MSSM are expected to be most important for the choice of m_A and m_{A_R} used below. We always check that we are at the minimum of the potential, by solving the (1-loop improved) tadpole equations for the soft Higgs masses.

Throughout the numerical analysis we have adopted a CMSSM-like configuration specified by $M_{1/2} = 600$ GeV, $m_0 = 120$ GeV, $A_0 = 0$ and $\tan \beta = 10$. The stop-sector soft masses in (4) were chosen as $m_{\tilde{Q}_3} = m_{\tilde{U}_3} = 2$ TeV, $T_{u33} = 3$ TeV and the top quark mass has been fixed to $m_t = 172.9$ GeV. In addition we have assumed $v_R = 5$ TeV, $\mu = 800$ GeV, $m_A = 800$ GeV, $\mu_\chi = -500$ GeV, $m_{A_R} = 2$ TeV and $\tan \beta_R = 1.1$ unless specified otherwise⁴. For the sake of completeness⁵ we have taken $g_{BL} = 0.46$ and $g_R = 0.48$.

Some further remarks concerning the parameters of the extended Higgs sector are in order here. The experimental constraints on the Z' mass yield a lower bound on v_R of about 2.5 TeV [48]⁶ for the assumed gauge couplings. This VEV, however, also enters the sfermion

⁴ For this choice of parameters m_{h^0} in the MSSM limit is about 125 GeV (1-loop), while for $m_{\tilde{Q}_3} = 1.1$ TeV, $m_{\tilde{U}_3} = 0.96$ TeV and $T_{u33} = 1.1$ TeV corresponding to the RGE solutions for these CMSSM one finds $m_{h^0} = 111$ GeV at 1-loop.

⁵ It is perhaps worth mentioning that from the effective theory point of view the specific values of g_{BL} and g_R do not matter as long as they yield the correct MSSM hypercharge coupling. Indeed, we have verified that different choices lead to results very similar to those quoted in the text.

⁶ Our Z' corresponds to the Z_χ in the notation of [48].

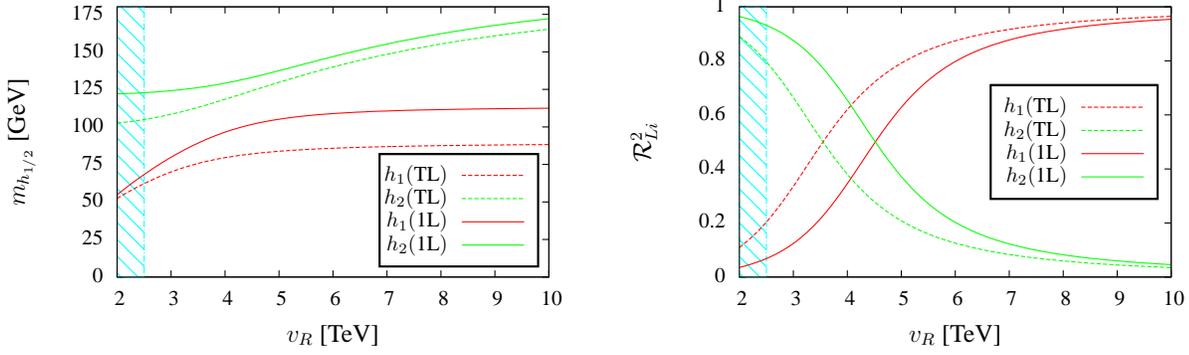


FIG. 1: The tree level and one-loop masses of the two lightest Higgs bosons $h_{1,2}$ (left) and \mathcal{R}_{Li}^2 (right) as a function of v_R ; at tree level (TL) in dashed and at one loop (1L) in solid lines. The values of all the other parameters are given in the text. The shaded area is excluded by the Z' searches.

mass matrices via the D -term contributions. Focusing, e.g., at the charged sleptons the relevant mass matrix reads

$$M_i^2 = \begin{pmatrix} M_L^2 + \frac{1}{8}M_{DL}^2 + m_f^2 & \frac{1}{\sqrt{2}}(v_d T_l - \mu Y_l v_u) \\ \frac{1}{\sqrt{2}}(v_d T_l - \mu Y_l v_u) & M_E^2 + \frac{1}{8}M_{DR}^2 + m_f^2 \end{pmatrix}, \quad (18)$$

where

$$M_{DL}^2 = g_{BL}^2(v_{\chi_R}^2 - v_{\tilde{\chi}_R}^2) + g_L^2(v_u^2 - v_d^2) \quad \text{and} \quad M_{DR}^2 = (g_R^2 - g_{BL}^2)(v_{\chi_R}^2 - v_{\tilde{\chi}_R}^2) + g_R^2(v_u^2 - v_d^2) \quad (19)$$

are just the D -terms, M_L^2 and M_E^2 are the soft SUSY masses for the L-type and R-type sleptons and all flavour indices have been suppressed in the above formula.

Since the (dominant) v_R^2 -parts of the two D -terms have opposite signs, the breaking of the extra gauge group must be nearly “ D -flat”, i.e., $\tan \beta_R \simeq 1$ as otherwise one of the sleptons would become tachyonic⁷. For completeness we also note that the cases of $\tan \beta = 1$ and $\tan \beta_R = 1$ lead to saddle points of the potential but not to minima which is a well known fact within the MSSM. In a complete analogy with the MSSM one can also show that for $\tan \beta_R \rightarrow 1$ one of the Higgs states gets massless at the tree level. Thus, since $\tan \beta_R$ has to be close to one, we generally expect two light Higgs bosons in the spectrum, which holds even at the one-loop level.

In Figure 1 we show the masses of the two lightest Higgs bosons together with

$$\mathcal{R}_{Li}^2 \equiv R_{i1}^2 + R_{i2}^2 \quad (20)$$

⁷ Let us note that this is indeed the case in all supersymmetric models featuring a spontaneously broken extended gauge symmetry well above the TeV scale (like, e.g., SUSY GUTs) and as such this requirement should be viewed as a phenomenological constraint rather than a fine-tuning.

as a function of v_R where $i = 1, 2$ labels the light Higgs scalars in the model. Note that the quantity $\mathcal{R}_{L_i}^2$, which reaches one in the MSSM limit, is a rough measure of how much the corresponding Higgs with index i resembles an MSSM Higgs boson. Roughly speaking, the smaller this quantities is, the smaller is the i -th Higgs coupling to the Z - and W -bosons, implying a reduced production cross sections at LEP, Tevatron and the LHC.

As claimed above there are two light CP-even states $h_{1,2}$ which essentially correspond to an admixture of the “standard” MSSM-like doublet component h^0 and its counterpart h_R^0 spanning over the $\chi_R - \bar{\chi}_R$ sector; this can also be seen by noticing that $\mathcal{R}_{L_1}^2 + \mathcal{R}_{L_2}^2 \simeq 1$ as displayed on the right hand side of Figure 1. We stress that the state which mainly resembles the MSSM h^0 (i.e., the one with a large $\mathcal{R}_{L_i}^2$) has already a tree-level mass of around 110 GeV or larger and reaches up to 140 GeV once loop corrections are included⁸. The lighter state with a mass below 100 GeV hardly couples to the Z -boson and, thus, the LEP constraints from the Higgs searches do not apply for it. To this end, we have used the `HiggsBounds` package [52, 53] to check explicitly that all the configurations of our concern here are experimentally allowed. Note also that the large variation in $\mathcal{R}_{L_i}^2$ as seen on the right panel of the Figure 1 and, in particular, its high sensitivity to radiative corrections is expected because the parameters have been deliberately chosen close to a level-crossing region.

This can be also seen in Figure 2 where we display the $m_{h_{1,2}}$ dependence on $\tan\beta_R$ and m_{A_R} . All results shown in this figure are at the one-loop level. The upper bound on $\tan\beta_R$ is given by the requirement that for a given value of v_R all sfermions masses are consistent with existing data (which, however, depends also on the sfermion mass parameters). The observed dependence on $\tan\beta_R$ is, indeed, rather strong. Note also that very light h_1 can be obtained for⁹ $\tan\beta_R \lesssim 1.05$. As in this regime it is mainly a combination of $\bar{\chi}_R$ and χ_R (see the right panel) the usual bounds do not apply. However, the second lightest Higgs boson (similar to the MSSM h^0) can decay into a pair of these states with sizable branching ratio which in turn can change the Higgs phenomenology drastically [42]. For $1.2 \lesssim \tan\beta_R \lesssim 1.3$ the lightest state becomes mainly the MSSM h^0 with a mass close to 130 GeV which is a consequence of the stop-sector parameter choice. In this figure one also sees that there is still quite some mixing between the two lightest states even for $m_{A_R} = 5$ TeV; this implies a change in the phenomenology with respect to that of the MSSM (for a given set of the MSSM parameters).

We checked that in this model, in general, the loops due to third generation sfermions (in particular the stops) give the largest contribution. In reference [54] it has been shown that in inverse seesaw models also the sneutrino loops can give large contributions. Indeed, we find that there can be huge contributions if the neutrino Yukawa couplings are $O(1)$ or

⁸ Actually, even larger values can be obtained when varying the parameters, e.g. m_{A_R} .

⁹ The exact value as well as the others given below depend on the other parameters.

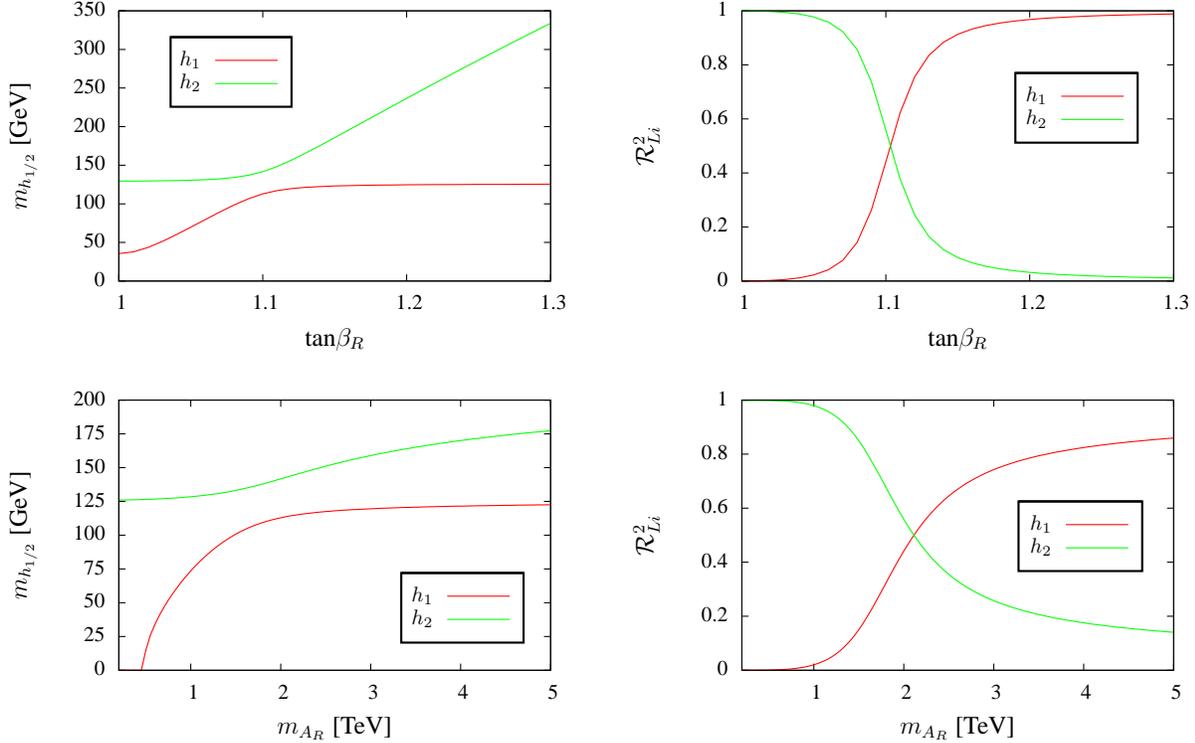


FIG. 2: One-loop masses of the two lightest Higgs bosons (left column) and \mathcal{R}_{Li}^2 (right column) as a function of $\tan\beta_R$ (upper row) and m_{A_R} (lower row). The values of all the other parameters are given in the text.

larger as can be seen in Figure 3. The neutrino Yukawa couplings are parametrized as

$$Y_\nu = f \begin{pmatrix} 0 & 0 & 0 \\ a & a & -a \\ 0 & 1 & 1 \end{pmatrix}, \quad (21)$$

with

$$a = (\Delta m_\odot^2 / \Delta m_A^2)^{\frac{1}{4}} \sim 0.4, \quad (22)$$

and the structure has been chosen such that one correctly accommodates the neutrino data. However, we find that the bound $\text{BR}(\mu \rightarrow e\gamma) \lesssim 2.4 \cdot 10^{-12}$ [55] severely constrains this option as, for large f , one gets a large contribution to $\mu \rightarrow e\gamma$ due to the chargino-sneutrino and W -neutrino loops.

There are, of course, several ways to tune the parameters such that this bound is avoided. For example, one can add a non-minimal flavour structure into the slepton sector [56] or tune the structure of the neutrino Yukawa couplings so that very specific values for θ_{13} , the reactor mixing angle, are obtained [57, 58]. This implies that, in principle, larger values for the neutrino Yukawa couplings are possible, hence rendering the corresponding loops more important. On the other hand, making them competitive even to the stop loops already requires quite some tuning [42].

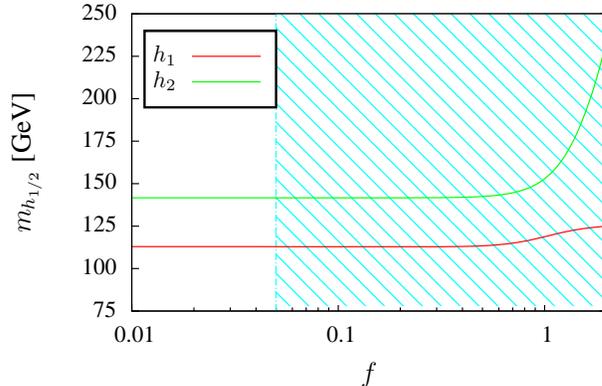


FIG. 3: One-loop masses of the two lightest Higgs bosons as functions of f . The values of all the other parameters are given in the text. The shaded area is excluded by $\mu \rightarrow e\gamma$.

IV. CONCLUSIONS

In this letter we have discussed the Higgs sector of a supersymmetric model where the SM gauge group has been extended to $SU(3)_c \times SU(2)_L \times U(1)_R \times U(1)_{B-L}$. In particular, we have shown that, already at the tree-level, the CP-even Higgs boson resembling the lightest neutral Higgs h^0 of the MSSM, can have a mass well above m_Z . At the one loop level, masses of 140 GeV and even above can easily be reached. In addition to such an h^0 -like Higgs, one can also have a second light state which, however, hardly couples to the SM vector bosons as it predominantly spans over the SM-neutral components. We have found regions where the h^0 -like Higgs can decay into two such states which, however, alters the standard search techniques at the LHC. Finally, we would like to stress that the general features discussed here also apply to other extensions of the SM gauge group, e.g., to full-featured left-right symmetric models, provided the MSSM Higgs doublets are charged with respect to the extended gauge symmetry.

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