1. Introduction

One of the main missing links of the otherwise immensely successful Standard Model (SM) is the Higgs boson which plays the crucial role in giving masses to all elementary particles in nature. It is therefore rightly the focus of a great deal of theoretical [1] and experimental enquiries. Even though the Higgs boson mass in SM is arbitrary, some ideas about how heavy the Higgs boson can be gained in the context of different plausible extensions of SM as well as from other considerations [2,3]. Typical upper limits from, say unitarity considerations [2] is in the TeV range. This bound is however considerably strengthened in one of the most widely discussed possibility for TeV scale physics, supersymmetry. Specifically in the minimal supersymmetric SM (MSSM), the upper bound on the Higgs boson mass is $M_{h}^{\text{max}} \leq 135$ GeV [4] when one and two loop radiative corrections are included. Present collider searches provide a lower bound on the SM Higgs mass [5] of 114 GeV leaving a narrow region which need to be probed to test MSSM. If the Higgs mass is found to be above this upper limit, does it mean that supersymmetry is not relevant for physics at TeV scale? The answer is of course “No” since there exist simple and well motivated extensions of MSSM, e.g. the next-to-MSSM, which extends the MSSM only by the addition of a singlet field [6] where there is a relaxation of this bound to about 142 GeV or so [7]. There are also other examples in literature [8] where simple modifications of the post-MSSM physics can provide additional room for Higgs mass.

In this paper we discuss an alternative scenario motivated by neutrino mass as well as understanding of the origin of parity violation [9] where the upper limit on the light Higgs mass is relaxed compared to MSSM. The model is the supersymmetric left-right model (SUSYLR) [10,11] with TeV scale parity restoration (or TeV right-handed gauge boson mass $W_{R}$). The change in the Higgs mass upper limit comes from the contribution of the D-terms and satisfies the decoupling theorem i.e. as the $W_{R}$ mass goes to infinity, the Higgs mass upper bound coincides with that for MSSM. This effect is to be expected on general grounds [13] in gauge extensions of MSSM.

2. Basics of the SUSYLR model

The gauge group of this model is $SU(2)_{L} \times SU(2)_{R} \times U(1)_{B-L} \times SU(3)_{c}$. The chiral left-handed and right-handed quark superfields are denoted by $Q \equiv (u, d)$ and $Q^{c} \equiv (w, c, \nu_{c})$ respectively and similarly the lepton superfields are given by $L \equiv (\nu, e)$ and $L^{c} \equiv (\nu^{c}, e^{c})$. The $Q$ and $L$ transform as left-handed doublets with the obvious values for the $B - L$ and the $Q^{c}$ and $L^{c}$ transform as the right-handed doublets with opposite $B - L$ values. The symmetry breaking is achieved by the following set of Higgs superfields: $\Phi_{a}(2, 2, 0, 1) (a = 1, 2); \Delta(3, 1, +2, 1); \Delta(3, 1, −2, 1); \Delta^{c}(1, 3, −2, 1)$ and $\Delta^{c}(1, 3, +2, 1)$. We include a gauge singlet superfield $S$ to facilitate the right handed symmetry breaking. The symmetry breaking can also be carried out by $B - L = 1$ doublet fields, for which our results also apply. A virtue of using triplet Higgs fields is that they lead to the see-saw mechanism for small neutrino masses using only renormalizable couplings. In addition, as was noted many years ago [11], low scale $W_{R}$ requires that R-parity must break spontaneously. This leads to many interesting phenomenological implications that we do not address here.

The superpotential for the model is given by:

$$\begin{align*}
W &= h_{Q}^{a}Q^{T}r_{2}\Phi_{a}r_{2}Q^{c} + h_{L}^{a}L^{T}r_{2}\Phi_{a}r_{2}L^{c} \\
&+ \frac{1}{2}f \left( L^{T}r_{2}\Delta L + L^{cT}r_{2}\Delta^{c} L^{c} \right) + \mu_{ab}\text{Tr} \left( \Phi_{a}^{T}r_{2}\Phi_{b}r_{2} \right) \\
&+ S \left[ \text{Tr} \left( \Delta \Delta^{c} + \Delta^{c}\Delta \right) - \nu_{\Delta}^{2} \right]
\end{align*}$$

(1)

In order to analyze the Higgs mass spectrum, we write down the Higgs potential for the model including the soft SUSY-breaking terms:

$$V = V_{F} + V_{S} + V_{D}$$

(2)

where $V_{F}$ and $V_{D}$ are the standard F-term and D-term potential and $V_{S}$ is the soft-SUSY-breaking terms which can be found in the literature [11,12]. Minimization of the Higgs potential leads to the following vacuum configuration for the $\Delta^{c}$ and $\nu^{c}$ Higgs fields [11]:

$$\langle \Delta^{c} \rangle = \left( \begin{array}{c}
\nu^{c} \\
0
\end{array} \right), \quad \langle \nu^{c} \rangle = \left( \begin{array}{c}
\nu_{\Delta}^{c} \\
0
\end{array} \right), \quad \langle \Delta^{c} \rangle = \left( \begin{array}{c}
0 \\
\nu_{\Delta}^{c}
\end{array} \right)$$

Note that in the SUSY limit $\nu_{\Delta} = 0$ and $\langle \nu^{c} \rangle = 0$. In the presence of supersymmetry breaking terms however, $\langle \nu^{c} \rangle$ is nonzero. On the other hand, if this model is extended to include a $B-L=0$ right handed triplet with nonzero vev, there appears a global minimum of the potential which has $\langle \nu^{c} \rangle = 0$ [14] even in the presence of susy breaking terms. Since the vev of $\nu^{c}$ is not relevant to our discussion, we will work with
B-L=0 triplet model and set \( \langle \tilde{\nu} \rangle = 0 \) henceforth. The SM symmetry remains unbroken at this stage and is broken by the vevs of the \( \Phi \) fields. We can write these fields in terms of their MSSM Higgs content:

\[
\Phi_i = \left( \begin{array}{c} \phi_{iL}^0 \\ \phi_{iL}^\pm \end{array} \right) \equiv (H_{di}, H_{ui}),
\]

with vevs \( \langle H_{di} \rangle = \kappa_i \), \( \langle H_{ui} \rangle = \kappa'_i \).

Before proceeding to discuss upper bound on the light neutral Higgs mass in this model, we wish to make a few comments on the implications of the TeV scale \( W_R \) models for neutrino masses. First, in the non-SUSY left-right model where neutrino mass has both type I and type II seesaw contributions, having a TeV scale \( W_R \) is unnatural since the type II seesaw contribution then becomes extremely large. This is due to the presence of non-zero couplings of type \( Tr[\phi^3 \Delta \phi^\dagger \Delta^c] \), which are allowed by the symmetries of the theory. On the other hand, in the SUSYLR model, this coupling is absent due to supersymmetry and therefore there is no type II seesaw contribution to neutrino mass. As far as the type I contribution is concerned, if we choose \( h_\nu \sim h_\tau \), where \( h_\nu \) is the Dirac neutrino Yukawa coupling, then we can have a few TeV \( W_R \) and neutrino masses of order of eV. Thus as far as neutrino masses go, low scale \( W_R \) is a realistic model.

3. Light Higgs mass bound: single bi-doublet case

We proceed to consider the bound on the light neutral Higgs mass in the SUSYLR model. We work in the limit where \( v_R \) and \( \bar{v}_R \) are much bigger than the SM scale. In this limit, we search for additional contributions to the MSSM Higgs potential, which will be at the heart of the change in the upper limit of the Higgs boson mass.

We first illustrate this in a one bi-doublet model. This simple model leads to vanishing CKM angles at the tree level, which can be fixed in one of two ways: (i) by including radiative correction effects from squark mixings \( [15] \) or (ii) by including a second bi-doublet which decouples from the low energy sector but it has a tadpole induced vev that generates the correct CKM angles. We discuss the case (ii) toward the end of the paper. The interesting point is that neither of these affects the Higgs mass upper bound that we derive. For the model under consideration, we first show that in the SUSY limit the low energy Higgs potential recovers that of MSSM in the same limit. Then soft SUSY breaking terms are taken into account, there appear new contributions to the MSSM Higgs potential, serving to raise the upper limit on the light Higgs mass, which can be significant for TeV scale \( W_R \).

We start with a review of the well known symmetry breaking of the model by the triplet Higgs fields \( \Delta^c \) and \( \Delta^c \) in the SUSY case. The gauge bosons get mass from the kinetic terms of triplets and after symmetry breaking, the massless gauge boson and gaugino corresponding to \( U(1)_Y \) is the combination \( B = g_R^2 v_R^2 + g_{BL}^2 v_{BL}^2 \), \( W_{3R} = g_R^2 v_R^2 v_{3R}^2 + g_{BL}^2 v_{BL}^2 \), with the hypercharge gauge coupling given by \( \frac{g_Y^2}{g_R^2} = \frac{1}{g_R^2} + \frac{1}{g_{BL}^2} \). The heavy \( Z' \)-boson has mass squared \( M_{Z'}^2 = 2(g_R^2 + g_{BL}^2)v_R^2 \). There is a factor 2 compared with the charged \( W_R \) boson mass \( M_{W_R}^2 = g_R^2 v_R^2 \) because the triplet vev breaks custodial symmetry for the right-handed sector.

Since there is no coupling between the bidoublet Higgs and the triplet Higgs fields responsible for parity breaking in Eq. (1), any change in the effective MSSM doublet Higgs potential below the \( v_R \) scale must originate from the D-terms,

\[
V_D = \frac{g_R^2}{8} \left| \text{Tr}[2\Delta^c v_m \Delta^c + 2\Delta^c v_m \Delta^c + \Phi_T \Phi^*] \right|^2
\]

where \( \delta V_D \) is the contribution to the neutral Higgs fields couplings is,

\[
V_D^{\text{neut.}} = \frac{g_R^2}{8} \left| \text{Tr}[\Phi_T \Phi^*] \right|^2 + 4(g_R^2 + g_{BL}^2)v_R^2 \left| \frac{\Delta^c - \Delta^c}{\sqrt{2}} \right|^2
\]

The coupling is linear in the field \( \text{Re}[\Delta^c - \Delta^c] \), which will we call \( \sigma_- \). As \( \sigma_- \) field becomes heavy, its coupling to \( [\Phi_T \Phi^*] \) will generate new quartic term in the MSSM doublet field potential, which in turn will lead to new contributions to Higgs mass upper bound. Collecting this new effect, we get for the Higgs quartic term:

\[
\delta V(\Phi) = \frac{1}{8} \left( g_R^2 - \frac{g_{BL}^2 v_{BL}^2}{M_{\sigma_-}^2} \right) \left| \text{Tr}[\Phi_T \Phi^*] \right|^2
\]

To evaluate this new contribution, we need to know \( M_{\sigma_-} \). This has two potential contributions: (i) from the D-term and (ii) from the F-term contribution to the Higgs potential. It turns out that in the SUSY limit, the only contribution to \( M_{\sigma_-} \) is from the D-terms and we have \( M_{\sigma_-}^2 = 2(g_R^2 + g_{BL}^2)v_R^2 \). This follows not only from actual calculations but also from the fact that \( \sigma_- \) is a member of the Goldstone supermultiplet, all members of which must have the same mass as \( Z' \) in the SUSY limit. This result would hold even if the superpotential had a term of the form \( \mu \Delta^c \Delta^c \). Using this in Eq. (6), it is easy to see that the net contribution to the quartic term in the Higgs superpotential becomes \( \frac{g_R^2}{8} \left( |H_u|^2 - |H_d|^2 \right)^2 \). This is nothing but the \( D_Y \) contribution to quartic Higgs doublet term in MSSM. Since in the decoupling limit, we get MSSM, as expected from the decoupling theorem.

Let us now switch on the supersymmetry breaking terms. In their presence, the \( \sigma_- \) field has aditional contributions which lead to a shift in the Higgs masses. To see this we introduce soft mass term \( m_{H_u}^2 S^T S \) as well as SUSY breaking mass terms for \( \Delta^c \) and \( \Delta^c \). Taking the same SUSY breaking terms for all the fields gives different value for the \( \sigma_- \) field mass and we get for the contribution to the quartic Higgs term

\[
\frac{g_R^2 v_R^2}{8} \left( |H_u|^2 - |H_d|^2 \right)^2
\]

where

\[
g_N^2 = g_R^2 - \frac{g_R^2}{g_R^2 + g_{BL}^2 + \frac{m_d^2}{2v_R^2}} = \frac{g_R^2 g_{BL}^2 + g_{BL}^2}{g_R^2 + g_{BL}^2 + \frac{m_d^2}{2v_R^2}}
\]
where $m_0$ in the above equation is a generic soft mass term for sparticles that breaks supersymmetry. This leads to an enhancement of the Higgs mass upper bound since $g_{V_SM}^2 > \theta_r^2$. To get an idea about how large the change in the upper bound is likely to be, we take $m_0 = 1 \text{ TeV}$, $v_R = 2 \text{ TeV}$, then from $\theta_r^2 \approx 0.42$ and $g_{V_SM}^2 \approx 0.13$, we get the ratio $r = \frac{\theta_r^2 + g_{V_SM}^2}{\theta_r^2} \approx 1.1$, which will give 10% increase of the tree level upper bound on the lightest Higgs mass, i.e., it increases from $M_Z = 90 \text{ GeV}$ to 100 GeV.

It is also worth pointing out that as the scale of parity violation goes to infinity, this new contribution goes to zero and one recovers the MSSM result. This is an important consistency check on our result [16].

4. One-loop radiative corrections and numerical result

In the following, we discuss the radiative corrections to the Higgs boson mass for the model above. It is well known that

$$\Delta V_1 = \frac{1}{64\pi^2} \text{Str} \mathcal{M}^2 \left( \log \frac{\mathcal{M}^2}{Q^2} - \frac{3}{2} \right),$$

where the supertrace $\text{Str}$ means $\text{Str} f(M^2) = \sum_i (-1)^{2J_i}(2J_i + 1)f(m_i^2)$, where the mass $m_i$ is calculated in background fields, and the sum counts all the fermionic and bosonic degrees of freedom. $J_i$ is the spin for particle $i$ and $Q$ is the renormalization scale on the order of electroweak symmetry breaking. In MSSM, top and stop contributes dominantly to $\Delta V_1$ because of large Yukawa coupling $y_t$, while the bottom and sbottom contribution is only important for large $\tan \beta \gg 20$.

In SUSYLR model with one bidoublet, we have $\tan \beta = m_t/m_b \approx 40$, so the sbottom quark can couple to $H_u$ with a large coupling $y_t$. This is from the F-term of bidoublet field $\Phi_1$.

$$(F_{\Phi_1})_{ij} = h_1(Q_i^T \tau_2)_{i} (\tau_2 Q_1^T)_{j} + \mu_{11}(\tau_2 \Phi_1^T \tau_2)_{ij}$$

$$|F_{\Phi_1}|^2 = \mu_{11} m_t y_t \beta \cdot \tau_1^T \tau_R + \mu_{11} m_t y_t \tau_1 \tau_R + \cdots$$

Note the second term in the second line differs from MSSM by a factor $m_t/(m_t \tan \beta)$, since $H_u$ and $H_d$ are unified into the same bidoublet. This means the sbottom mass receives a large LR mixing proportional to the top Yukawa coupling, even for low $\tan \beta$ (in the presence of a second bidoublet in the realistic model below). Therefore now there are three fields that have to be taken into account: top, stop and sbottom. If we neglect the small couplings except for $y_t$, and also neglect the $A$-terms their masses can be approximated by

$$m_t^2 = y_t^2 \frac{|H_u|^2}{2},$$
$$m_{t_2}^2 \simeq m_{Q_2}^2 \cong y_t^2 |H_u|^2,$$
$$m_{b_1}^2 \simeq m_Q + y_{t11} |H_u| + \cdots,$$
$$m_{b_2}^2 \simeq m_Q - y_{t11} |H_u| + \cdots$$

The $\cdots$ represents dependence on $|H_d|$, but since the lightest Higgs boson is mainly made up of $H_u$ for large $\tan \beta$ with only a small $H_d$ component, this dependence can be neglected. We can choose a proper scale $Q$ so that the first derivative vanishes, which will be important in eliminating the explicit $Q$ dependence of Higgs mass, i.e., it can only depend on $Q$ through depending on other parameters. Second derivative gives radiative corrections to the lightest Higgs mass

$$\delta M_h^2 \approx \frac{1}{2} \frac{\partial^2 \Delta V_1}{\partial |H_u|^2} = \frac{3g_{V_SM}^2}{8\pi^2} \frac{m_t^4}{M_{W_L}^2} \log \frac{m_t^2}{m_0^2}$$
$$- \frac{3g_{V_SM}^2 m_{0} \mu_{11} m_t}{64\pi^2} \log \frac{m_t^2}{m_{0}^2} + \frac{3g_{V_SM}^2 \mu_{11} m_t^2}{32\pi^2} \frac{1}{M_{W_L}^2}$$

(11)

The first term is as the usual MSSM one. Now we have two new terms proportional to $\mu_{11}$. Their sum is an even function of $\mu_{11}$ since changing the sign also interchanges $\tilde{b}_1 \leftrightarrow \tilde{b}_2$. We find the sum of second and third term are negative definite for arbitrary $\mu_{11}$. Actually, for $m_0 \sim 1 \text{ TeV}$, one can expand with $\mu_{11} m_t$. The net contribution is non-vanishing only up to the third order, which is $-g_{V_SM}^2 \frac{1}{32\pi^2} \frac{m_{0} \mu_{11} m_t}{M_{W_L}^2}$, depending on $\mu_{11}$ very mildly. For $\mu_{11}$ ranging from 100 GeV to 300 GeV around EW scale, this negative contribution is smaller than 1 GeV.

Note that in this case, we did not have to discuss the details of EWSB since it is very similar to MSSM.

Let us present the numerical results for the Higgs mass upper bound for this scenario. In Fig. 1, we plot the difference in the prediction of upper bound on the lightest Higgs boson mass between SUSYLR model and MSSM: $\Delta M_h = m_h^{\text{SUSYLR}} - m_h^{\text{MSSM}}$. For the right-handed scale near 2-3 TeV, the upward shift of the Higgs mass bound can be of a few GeV, increasing as $v_R$ decreases. (For symmetry breaking using Higgs doublets, $M_{Z}^2 = (g_{V_SM}^2 + g_{B,L}^2)v_R^2$ and $g_{V_SM}^2 \neq 0$. In Eq. 4 gets increased to $g_{V_SM}^2 g_{B,L} + g_{V_SM}^2 y_{t11}^2 h_0^2$. Then $\Delta M_h$ can further increase by a factor of 2.) The bound also increases with the soft mass scale $m_0$ as expected. From the discussions below Eq. (11), non-zero $\mu_{11}$ always gives small and negative contribution. In order not to violate the lower bound on chargino mass at LEP2, we choose $\mu_{11} \geq 100 \text{ GeV}$.

5. Extension to two bidoublet case

We can extend our discussion to the more realistic two bidoublet case as needed to generate the correct CKM angles at the tree level. There are two possible ways to do that; both these we discuss below.

Model A:

In this case, we identify $\Phi_1$ as the bidoublet of the previous section. We can diagonalize the corresponding Yukawa coupling matrix $h_Q$. We are then forced to have all elements of the second bi-doublet $\Phi_2$ Yukawa coupling, $h_Q^2 \neq 0$. Once the second $\Phi_2$ has vev, by appropriate choice of this matrix, we can generate the desired quark masses and mixings. We will see that even though there are four real neutral Higgs fields in this case, the upper limit on the light Higgs field remains the same as in the single bi-doublet case.

To see this, first we choose a basis in the $\Phi$ space such that the superpotential for $\Phi$'s has the form

$$W_{\text{extra}} = \mu_{11} \text{Tr}(\Phi_1 \Phi_1) + \mu_{22} \text{Tr}(\Phi_2 \Phi_2)$$

(12)
We assume that $\mu_{22} \gg \mu_{11}$. We have then no freedom to diagonalize the soft SUSY breaking terms. The sum of the $V_F + V_S$ can then in general be written as

$$V_F + V_S(\Phi_i) = m_{ij}^2 \text{Tr}(\phi_i^\dagger \phi_j) + b_{ij} \text{Tr}(\phi_i \phi_1) + b_{22} \text{Tr}(\phi_2 \phi_2) + \text{h.c.} \quad (13)$$

Note that if $m_{22}^2 \gg m_{12}^2$, then the mixed term in the $\phi$'s will induce a vev for the $\phi_2$ field which is small compared to that for the $\phi_1$ i.e. $\kappa_2, \kappa'_2 \ll \kappa_1, \kappa'_1$. To generate the correct mass and mixing pattern for the quarks, it is sufficient to have the $\phi_2$ vevs of order of a 100 MeV. For instance if $m_{12} \leq 10 \text{ GeV}$ and $m_{22} = 1 \text{ TeV}$, then we can estimate $\kappa'_2 = \kappa'_1 m_{12}/m_{22} = \kappa'_1/100 \sim 100 \text{ MeV}$, which is enough to generate the strange quark mass as well as other CKM angles. Note also that one should include the one loop effects coming from squark masses and mixings[13]. While we do not give a detailed fit here, it seems clear that this is a realistic model where the new Higgs mixing parameter $m_{12}$ is in the 1-10 GeV range. When it is close to one GeV, the effect on the Higgs mass upper bound is also about a GeV lower due to off diagonal contributions. We can also keep the $\Phi_2$ Yukawa couplings sufficiently small so that their radiative corrections do not affect the one loop result. This vacuum then is a perturbation around the vacuum of the single bi-doublet case and furthermore due to large $\mu_{22}$, the $H_{u,d}$ coming from the second bi-doublet will acquire heavy mass and decouple without affecting the light Higgs mass upper bound except perhaps a small one GeV or so shift. This case corresponds to large $\tan \beta \approx 40$.

**Model B:**

In this case, we choose two bi-doublets with the vev pattern given by:

$$\langle \Phi_1 \rangle = \begin{pmatrix} \kappa & 0 \\ 0 & 0 \end{pmatrix}, \quad \langle \Phi_2 \rangle = \begin{pmatrix} 0 & 0 \\ 0 & \kappa' \end{pmatrix} \quad (14)$$

Since the down quark masses in this case come from a second Yukawa coupling unlike the model A, we can have the value of $\tan \beta \approx 30$ much lower than 40 by appropriate choice of the second Yukawa coupling matrix. There are generally four electroweak scale Higgs doublets. Using the standard formula in Eq. (6), we have calculated the 1-loop radiative corrections to the 4×4 neutral Higgs mass matrix in the presence of SUSY breaking thresholds. As before, we neglected all other couplings but $y_t$ when calculating $\Delta V_1$. We also keep the effect of non-zero vev for the right-handed sneutrino. However due to small neutrino Dirac Yukawa couplings, the mixing effect between the Higgs field and the left-handed sneutrino caused by the right-handed sneutrino vev is very small and does not affect our result. In order to estimate the upper bound, we have done a numerical study to obtain the Higgs mass for random choice of parameters. The results are the scatter points in Fig. 2 below for a choice of the generic soft mass scale $m_0 = 1 \text{ TeV}$ and right-handed scale $v_R = 1.5 \text{ TeV}$. Each point in the scatter plot represents the lightest Higgs mass for a specific choice of parameters. The upper limit therefore corresponds to the topmost set of points in Fig. 2. In contrast, the MSSM Higgs mass upper bound is plotted as the yellow (lower) curve, which is at most 130 GeV after 1-loop radiative corrections with the same choice of $m_0$. The red (upper) curve is for Model A. We find Fig. 2 that in general SUSYLR model the upper bound can be as high as 140 GeV or even more especially in the regime $5 < \tan \beta < 10$. This is higher than the prediction of MSSM. Clearly, as the right-handed scale goes down, the upper bound increases.

**6. Comments and Conclusion:** Before concluding, we wish to make a few comments on the model:

(i) Low scale non-SUSY left-right models have strong constraints coming from the tree level Higgs contribution to flavor changing processes. In the SUSY version however, there are additional contributions to the same from squark and slepton sector which can be used to cancel this effect[17]. While
strictly this is not natural, from a phenomenological point of view, this makes the model consistent when both the $W_R$ and Higgs masses are in the few TeV range.

(ii) The second point is that unlike other models such as NMSSM where the light Higgs mass bound is changed by making additional assumptions about the Higgs couplings (e.g. not hitting the Landau pole at the GUT scale), in our model the increase in the bound is purely gauge coupling induced and is independent of the Higgs couplings.

(iii) It is also worth stressing the obvious point that observation of a Higgs with mass above the MSSM bound of 135 GeV is not necessarily an evidence for the SUSYLR model since there exist other models with which also relax this bound. One needs other direct evidences such as the mass of $W_R$ or $Z'$ produced at LHC which when combined with observed higher Higgs mass could provide evidence for low left-right seesaw

To conclude, we have pointed out that the upper bound on the light Higgs mass is higher if MSSM is assumed to be an effective low energy theory of a TeV scale SUSYLR model. The increase can be as much as 10 GeV or more depending on the scale of parity breaking. If the Higgs boson mass in the collider searches is found to exceed the MSSM upper limit of 135 GeV, one interpretation of that could be in terms of a TeV scale seesaw in the context of a SUSYLR model.

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