Higgs Mass Corrections in the SUSY $B-L$ Model with Inverse Seesaw

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In the context of the Supersymmetric (SUSY) $B-L$ (Baryon minus Lepton number) model with inverse seesaw mechanism, we calculate the one-loop radiative corrections due to right-handed (s)neutrinos to the mass of the lightest Higgs boson when the latter is Standard Model (SM)-like. We show that such effects can be as large as $\mathcal{O}(100)$ GeV, thereby giving an absolute upper limit on such a mass around 180 GeV. The importance of this result from a phenomenological point of view is twofold. On the one hand, this enhancement greatly reconciles theory and experiment, by alleviating the so-called 'little hierarchy problem' of the minimal SUSY realization, whereby the current experimental limit on the SM-like Higgs mass is very near its absolute upper limit predicted theoretically, of 130 GeV. On the other hand, a SM-like Higgs boson with mass below 180 GeV is still well within the reach of the Large Hadron Collider (LHC), so that the SUSY realisation discussed here is just as testable as the minimal version.

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The Higgs boson is the last missing particle in the SM. Higgs boson discovery at the LHC is, therefore, crucial for its validity as a low energy approximation of a new physics scenario valid to high energy scales. A possibility for the latter emerges in SUSY theories, wherein the Higgs mechanism is retained for mass generation and multiple Higgs bosons appear in order to cancel anomalies. In addition, the stabilization of the Higgs mass against loop corrections (gauge hierarchy problem) is possibly the strongest motivation for a SUSY theory of nature. Hence, Higgs boson discovery at the LHC is also crucial for SUSY as a whole. A consequence of a SUSY Higgs sector is the existence of a stringent upper bound on the mass of the lightest SUSY Higgs boson, h, when the latter is SM-like. In the Minimal Supersymmetric Standard Model (MSSM), this value is $m_h \lesssim 130$ GeV. Therefore, non-observing at the LHC a SM-like Higgs boson lighter than 130 GeV would rule out the MSSM.

In detail, in the MSSM, the mass of the lightest Higgs state can be approximated, at the one-loop level, as [1]

$$m_h^2 \leq M_Z^2 + \frac{3g^2}{16\pi^2 M_W^2}\frac{m_t^4}{\sin^2\beta}\log\left(\frac{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}{m_t^4}\right), \qquad (1)$$

where g is the $SU(2)$ gauge coupling and $m_{\tilde{t}_{1,2}}$ are the two stop physical masses. The ratio of the Electro-Weak (EW) Vacuum Expectation Values (VEVs) is given by $\tan\beta = v_2/v_1$. Note that the factor 3 in the above top-stop correction is due to color. If one assumes that the stop masses are of order TeV, then the one-loop effect leads to a correction of order $\mathcal{O}(100)$ GeV, which implies that

$$m_h^{\rm MSSM} \lesssim \sqrt{(90\ {\rm GeV})^2 + (100\ {\rm GeV})^2} \simeq 135\ {\rm GeV}. \qquad (2)$$

It is worth mentioning that the two-loop corrections reduce this upper bound by a few GeVs, to the aforementioned 130 GeV or so value [2].

Experimental evidence now exists for physics beyond the SM, in the form of neutrino oscillations, which imply neutrino masses [3]. In turn, the latter imply new physics beyond not only the SM, but also the MSSM. Right-handed neutrino superfields are usually introduced in order to implement the seesaw mechanism, which provides an elegant solution for the smallness of the left-handed neutrino masses. Right-handed neutrinos, which are heavy, can naturally be implemented in the SUSY $B-L$ extension of the SM (hereafter, the 'SUSY $B-L$ model' for short), which is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$. In this scenario, the scale of $B-L$ symmetry breaking is related to the soft SUSY breaking scale [4]. Thus, the right-handed neutrino masses are naturally of order TeV and the Dirac neutrino masses must be less than 10^{-4} GeV (*i.e.*, they are of order the electron mass) [5]. Nevertheless, due to the smallness of Dirac neutrino Yukawa couplings, the right-handed neutrino sector has very suppressed interactions with the SM particles. Therefore, the predictions of such a SUSY $B-L$ model (*i.e.*, with standard seesaw mechanism) remain close to the MSSM ones. In particular, the discussed MSSM prediction for the lightest Higgs boson mass upper bound remains intact. Same conclusion is obtained in the context of the minimal Supersymmetric seesaw model, where the right-handed neutrino masses are of order 10^{13} GeV [6].

The SUSY $B-L$ model with inverse seesaw, where Dirac neutrino Yukawa couplings are of order 1, has recently been considered [7]. The superpotential of the leptonic sector associated to this model is given by

$$W = Y_E L H_1 E^c + Y_\nu L H_2 N^c + Y_S N^c \chi_1 S_2 + \mu H_1 H_2 + \mu' \chi_1 \chi_2, \tag{3}$$

where $\chi_{1,2}$ are SM singlet superfields with $B-L$ charges $+1$ and -1, respectively. Therefore, $U(1)_{B-L}$ is spontaneously broken by the VEVs of the scalar components of these superfields. N_i are three SM singlet chiral superfields with $B-L$ charge $= -1$, introduced to cancel the $U(1)_{B-L}$ anomaly. The fermion components of N_i account for right-handed neutrinos. Finally, chiral SM singlet superfields $S_{1,2}$ with $B-L$ charge $= +2, -2$ are considered to implement the inverse seesaw mechanism. Note that a $\mathbb{Z}_2$ symmetry is assumed in order to prevent the interactions between the field S_1 and any other field.

After $B-L$ and EW symmetry breaking, the neutrino Yukawa interaction terms lead to the following expression:

$$\mathcal{L}_m^\nu = m_D \bar{\nu}_L N^c + M_N \bar{N}^c S_2, \tag{4}$$

where $m_D = Y_\nu v \sin\beta$ and $M_N = Y_S v' \sin\theta$, with $\tan\beta = \frac{v_2}{v_1} = \frac{\langle H_2 \rangle}{\langle H_1 \rangle}$, $\tan\theta = \frac{v_1'}{v_2'} = \frac{\langle \chi_1 \rangle}{\langle \chi_2 \rangle}$, $v^2 = v_1^2 + v_2^2$, and $v'^2 = v_1'^2 + v_2'^2$. Light neutrino masses are related to a small mass term $\mu_S \bar{S}_2^c S_2$, with $\mu_S \lesssim \mathcal{O}(1)$ KeV, which can emerge at the $B-L$ scale from a non-renormalizable term in the superpotential, $\frac{\chi_1^4 S_2^2}{M_I^3}$, with M_I an intermediate scale of order $\mathcal{O}(10^7)$ GeV. Note that the non-renormalizable scale M_I can be related to a more fundamental scale and couplings of the S_2 and χ fields with integrated out fields and a suppression factor. In this case one can write for instance $1/M_I^3 \sim \lambda^4/M_*^3$, therefore, if $\lambda \sim \mathcal{O}(0.01)$, then M_* is of order 10^{12} GeV. Thus, the Lagrangian of neutrino masses, in the flavor basis, is given by:

$$\mathcal{L}_m^\nu = m_D \bar{\nu}_L N^c + M_N \bar{N}^c S_2 + \mu_S \bar{S}_2^c S_2. \tag{5}$$

In the basis $\{\nu_L, N^c, S_2\}$, the 3×3 neutrino mass matrix of one generation takes the form:

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & M_N \\ 0 & M_N & \mu_S \end{pmatrix}. \tag{6}$$

Thus, the physical light and heavy neutrino masses are given by

$$m_{\nu_\ell} = \frac{m_D^2 \mu_S}{M_N^2 + m_D^2}, \tag{7}$$

$$m_{\nu_{H,H'}} = \pm\sqrt{M_N^2 + m_D^2} + \frac{1}{2} \frac{M_N^2 \mu_S}{M_N^2 + m_D^2}. \tag{8}$$

Therefore, neglecting μ_S one finds

$$m_{\nu_\ell}^2 \simeq 0, \qquad m_{\nu_{H,H'}}^2 \simeq m_D^2 + M_N^2. \tag{9}$$

In addition, the gauge eigenstates $\{\nu_L, N^c, S_2\}$ can be written in terms of the physical eigenstates as follows [10]

$$\begin{array}{rcl} \nu_L & \simeq & \nu_\ell + \alpha\, a_2\, \nu_H + \alpha\, a_2\, \nu_{H'}, \\ N^c & \simeq & a_1\, \nu_\ell - \alpha\, \nu_H + \alpha\, \nu_{H'}, \\ S_2 & \simeq & -a_2\, \nu_\ell + \alpha\, \nu_H + \alpha\, \nu_{H'}, \end{array} \tag{10}$$

where $\alpha \sim \sin(\pi/4)$, $a_1 \sim \mathcal{O}(10^{-10})$, and $a_2 \sim \mathcal{O}(0.1)$. In this respect, new couplings among the heavy neutrinos, the charged leptons and W bosons are generated. These couplings are of order the gauge coupling times the mixing angles between light and heavy neutrinos. Therefore, the heavy neutrinos of this class of models may have a clean signature at the LHC [7] and imply very interesting phenomenological implications [8].

The sneutrino mass matrix is obtained from the sneutrino scalar potential, which is given by

$$V_{\rm scalar} = V_F + V_D + V_{\rm soft}, \tag{11}$$

where V_F is defined as usual as $|\partial W/\partial \phi|^2$ and

$$V_D = \frac{M_Z^2 \cos 2\beta}{2} \tilde{\nu}_L^* \tilde{\nu}_L + M_{Z'}^2 \cos 2\theta \Big(\tilde{\nu}_L^* \tilde{\nu}_L - \tilde{\nu}_R^* \tilde{\nu}_R + 2 \tilde{S}_2^* \tilde{S}_2\Big),$$

where $M_{Z'}$ is the mass of the $B-L$ neutral gauge boson Z', given by $M_{Z'}^2 = g''^2 v'^2$. From the LEP II experimental limits, one finds $M_{Z'}/g'' > 6$ TeV [9]. Finally, $V_{\rm soft}$ is defined as

$$V_{\rm soft} = m_0^2 \sum_\phi |\phi|^2 + \frac{1}{2} M_{1/2} \sum_i \widetilde{\lambda}_i \widetilde{\lambda}_i + \Big[A_0\Big(Y_\nu \widetilde{N}^c \widetilde{L} H_2 + Y_e \widetilde{E}^c \widetilde{L} H_1 + Y_S \widetilde{N}^c \widetilde{S}_2 \chi_1\Big) + B_0\Big(\mu H_1 H_2 + \mu' \chi_1 \chi_2\Big) + h.c.\Big]. \tag{12}$$

Here, the sum in the first term runs over $\phi = H_1, H_2, \chi_1, \chi_2, \tilde{L}, \tilde{E}^c, \tilde{N}^c, \tilde{S}_1, \tilde{S}_2$ and the sum in the second term runs over the gauginos: $\lambda_i = \widetilde{B}, \widetilde{W}^a, \widetilde{g}^a, \widetilde{Z}'$. Note that in the above expression, the trilinear coupling is defined, as usual, by $A_{ij} = A_0 Y_{ij}$.

In general, one finds that the sneutrino mass matrix, for one generation, in the basis $(\tilde{\nu}_L, \tilde{\nu}_L^*, \tilde{\nu}_R, \tilde{\nu}_R^*, \tilde{S}_2, \tilde{S}_2^*)$, can be written as a 3×3 matrix, with entries multiplied by the identity 2×2 matrix [10], *i.e.*, with one generation, one obtains two left-handed sneutrinos $\tilde{\nu}_{L_{1,2}}$ and four right-handed sneutrinos $\tilde{\nu}_{H_{3,4,5,6}}$:

$$\mathcal{M}^2_{\tilde{\nu}} = \begin{pmatrix} m^2_{\tilde{L}} + m^2_D + \frac{M^2_Z}{2}\cos 2\beta + M^2_{Z'}\cos 2\theta & m_D(A_\nu + \mu\cot\beta) & m_D M_N \\ m_D(A_\nu + \mu\cot\beta) & m^2_{\tilde{N}} + m^2_D + M^2_N - M^2_{Z'}\cos 2\theta & M_N(A_S + \mu'\cot\theta) \\ m_D M_N & M_N(A_S + \mu'\cot\theta) & m^2_{\tilde{S}} + M^2_N + 2M^2_{Z'}\cos 2\theta \end{pmatrix}. \quad (13)$$

If one chooses the A-terms such that elements 12 and 23 vanish, then the sneutrino masses can be written as

$$\begin{aligned} m_1^2 &= d, \\ m_2^2 &= \frac{1}{2}\left[(a+f) + \sqrt{(a-f)^2 + 4c^2}\right], \\ m_3^2 &= \frac{1}{2}\left[(a+f) - \sqrt{(a-f)^2 + 4c^2}\right], \end{aligned} \quad (14)$$

$$\begin{aligned} \text{where} \quad a &= m^2_{\tilde{L}} + m^2_D + \frac{M^2_Z}{2}\cos 2\beta + M^2_{Z'}\cos 2\theta, \\ c &= m_D M_N, \\ d &= m^2_{\tilde{N}} + m^2_D + M^2_N - M^2_{Z'}\cos 2\theta, \\ f &= m^2_{\tilde{S}} + M^2_N + 2M^2_{Z'}\cos 2\theta. \end{aligned} \quad (15)$$

If one assumes that $m_{\tilde{L}} = m_{\tilde{N}} = m_{\tilde{S}} = \tilde{m}$ and neglects the D-term, then the sneutrino masses can be written as

$$m^2_{\tilde{\nu}_{L_{1,2}}} = \tilde{m}^2, \qquad m^2_{\tilde{\nu}_{H_{3,4,5,6}}} = m^2_D + M^2_N + \tilde{m}^2. \quad (16)$$

It is important to note that, unlike the squark sector, where only the third generation (stops) has a large Yukawa coupling with the Higgs boson, hence giving the relevant contribution to the Higgs mass correction, all three generations of the (s)neutrino sector may lead to important effects since the neutrino Yukawa couplings are generally not hierarchical. Also, due to the large mixing between the right-handed neutrinos N_i and S_{2_j}, all the right-handed sneutrinos $\tilde{\nu}_H$ are coupled to the Higgs boson H_2, hence they can give significant contribution to the Higgs mass correction. In this respect, it is useful to note that the stop effect is due to the running of 12 degrees of freedom (3 colors times 2 charges times 2 for Left and right stops) in the Higgs mass loop corrections. Also, in the case of right-handed sneutrinos, there are 12 degrees of freedom (3 generations times 4 eigenvalues).

As example of a generic 3×3 neutrino Yukawa coupling, Y_ν, we consider $Y_\nu = m_D/v_2$, with the Dirac neutrino mass matrix m_D [8]

$$m_D = U_{\rm MNS}\sqrt{m^{\rm phys}_{\nu_\ell}}\, R\, \sqrt{\mu_S^{-1}}\, M_N, \quad (17)$$

where R is an arbitrary orthogonal matrix and $U_{\rm MNS}$ is the light neutrino mixing matrix. If we assume that $R = I_{3\times 3}$ and $\sqrt{m^{\rm phys}_{\nu_\ell}/\mu_S} \sim \mathcal{O}(0.1)$, then we find $Y_\nu \simeq U_{\rm MNS}$. Note that here we assume a hierarchical μ_s in order to account for a possible hierarchy between light neutrino masses. For simplicity, we also assume universal Majorana neutrino masses, $M_N = {\rm diag}\{M, M, M\}$.

The one-loop radiative correction to the effective potential is given by the relation

$$\Delta V = \frac{1}{64\pi^2}\,{\rm STr}\left[M^4\left(\log\frac{M^2}{Q^2} - \frac{3}{2}\right)\right], \quad (18)$$

where M^2 is the field dependent generalized mass matrix and Q is the renormalization scale. The complete effective potential is Q independent. However, the effective potential at one-loop level contains implicit dependence on this scale. The renormalization scale Q is usually chosen such that the large logarithms, appearing in higher order corrections, are suppressed. The supertrace in Eq.(18) is defined by

$${\rm STr} f(M^2) = \sum_i (-1)^{2J_i}(2J_i + 1) f(m_i^2). \quad (19)$$

Here m_i^2 is a field dependent squared mass eigenvalue of a particle with spin J_i. Therefore ΔV, due to one generation of neutrinos and sneutrinos, is given by

$$\Delta V_{\nu,\tilde{\nu}} = \frac{1}{64\pi^2}\Big[\sum_{i=1}^{6} m^4_{\tilde{\nu}_i}\Big(\log\frac{m^2_{\tilde{\nu}_i}}{Q^2} - \frac{3}{2}\Big) - 2\sum_{i=1}^{3} m^4_{\nu_i}\Big(\log\frac{m^2_{\nu_i}}{Q^2} - \frac{3}{2}\Big)\Big]. \quad (20)$$

The first sum runs over the sneutrino mass eigenvalues, while the second sum runs over the neutrino masses (with vanishing m_{ν_1}). In case of our above example, where $Y_\nu \sim U_{\rm MNS}$, one finds that the total ΔV is given by three times the value of ΔV for one generation. This factor then compensates the color factor of (s)top contributions.

The genuine $B-L$ correction to the CP-even Higgs mass matrix, due to the (s)neutrinos, at any renormalization scale Q, is given by

$$\Delta M^2_{ij} = \frac{1}{2}\left[\frac{\partial^2(\Delta V)_{\nu,\tilde{\nu}}}{\partial H_i \partial H_j} - \frac{\delta_{ij}}{2H_i}\frac{\partial \Delta V_{\nu\tilde{\nu}}}{\partial H_i}\right]_{H_i=v_i}. \quad (21)$$

From the (s)neutrino masses, given in Eqs. (9) and (16), one can easily show that, for one generation, we have

$$\delta M^2_{11} \;=\; \delta M^2_{12} = \delta M^2_{21} = 0, \quad (22)$$

$$\delta M^2_{22} = \frac{1}{16\pi^2}\left[\left(\frac{\partial m^2_{\tilde{\nu}_H}}{\partial v_2}\right)^2 \log\frac{m^2_{\tilde{\nu}_H}}{\hat{Q}^2} - \left(\frac{\partial m^2_{\nu_H}}{\partial v_2}\right)^2 \log\frac{m^2_{\nu_H}}{\hat{Q}^2}\right] = \frac{m_D^4}{4\pi^2 v_2^2}\,\log\frac{m^2_{\tilde{\nu}_H}}{m^2_{\nu_H}}. \quad (23)$$

Therefore, the complete one-loop squared-mass matrix of CP-even Higgs bosons will be given by $M^2_{\rm tree} + \Delta M^2$, with $\Delta M^2 = \begin{pmatrix} 0 & 0 \\ 0 & \delta_t^2 + \delta_\nu^2 \end{pmatrix}$, where δ_t^2 refers to the (s)top contribution presented in Eq. (1) and δ_ν^2 is the (s)neutrino correction given in Eq. (23). In this case, the lightest Higgs bosons mass is given by

$$m_h^2 = \frac{1}{2}(M_A^2 + M_Z^2 + \delta_t^2 + \delta_\nu^2)\left[1 - \sqrt{1 - 4\frac{M_Z^2 M_A^2 \cos^2 2\beta + (\delta_t^2+\delta_\nu^2)(M_A^2\sin^2\beta + M_Z^2\cos^2\beta)}{(M_A^2+M_Z^2+\delta_t^2+\delta_\nu^2)^2}}\right] \quad (24)$$

For $M_A \gg M_Z$ and $\cos 2\beta \simeq 1$, one finds that

$$m_h^2 \simeq M_Z^2 + \delta_t^2 + \delta_\nu^2. \quad (25)$$

If $\tilde{m} \simeq \mathcal{O}(1)$ TeV, $Y_\nu \simeq \mathcal{O}(1)$ and $M_N \simeq \mathcal{O}(500)$ GeV, one finds that $\delta_\nu^2 \simeq \mathcal{O}(100)^2$, thus the Higgs mass will be of order $\sqrt{(90)^2 + \mathcal{O}(100)^2 + \mathcal{O}(100)^2} \simeq 170$ GeV. It is worth mentioning that the Yukawa coupling Y_ν at the $B-L$ scale, M_{B-L}, can be derived from the renormalization group equations [13]:

$$\begin{aligned}
16\pi^2\frac{dY_\nu}{dt} &= 3Y_\nu Y_\nu^\dagger Y_\nu + Y_\nu\left[{\rm Tr}\Big(Y_\nu Y_\nu^\dagger\Big) + 3{\rm Tr}\Big(Y_u Y_u^\dagger\Big) - \frac{3}{2}g_B^2 - \frac{3}{5}g_1^2 - 3g_2^2\right] + Y_\nu Y_S^* Y_S^T + Y_e^T Y_e^* Y_\nu && (26)\\
16\pi^2\frac{dY_u}{dt} &= 3Y_u Y_u^\dagger Y_u + Y_u\left[{\rm Tr}\Big(Y_\nu Y_\nu^\dagger\Big) + 3{\rm Tr}\Big(Y_u Y_u^\dagger\Big) - \frac{1}{6}g_B^2 - \frac{13}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2\right] + Y_u Y_d^\dagger Y_d && (27)\\
16\pi^2\frac{dY_S}{dt} &= 2Y_S Y_S^\dagger Y_S + Y_S\left[{\rm Tr}\Big(Y_S Y_S^\dagger\Big) - \frac{9}{2}g_B^2\right] + 2Y_\nu^T Y_\nu^* Y_S && (28)
\end{aligned}$$

By ignoring the small Yukawa couplings: Y_d and Y_e, one can solve these equations numerically and verify that for $Y_t \simeq \mathcal{O}(1)$ the neutrino Yukawa coupling, Y_ν, at TeV scale, can be of order one as well. As shown in Fig. 1, these values of Y_ν are sufficient for enhancing the Higgs mass significantly. In this figure we present the Higgs mass, m_h, as a function of the sneutrino mass, $m_{\tilde{\nu}}$, for $M_N = {\rm diag}\{300, 400, 500\}$ GeV and Y_ν couplings are given by: $Y_\nu = {\rm diag}\{0, 0, 0.85\}$, $Y_\nu = 0.85\,{\rm I}_{3\times 3}$, and $Y_\nu = U_{\rm MNS}$. As can be seen, the neutrino Yukawa couplings indeed play a crucial rule in enhancing the lightest Higgs mass, m_h. Also the large mixing in $U_{\rm MNS}$ is favored by a large Higgs mass.

Finally, we consider the impact of the trilinear couplings A_N and A_S, which contribute to the off-diagonal elements of the sneutrino mass matrix (13). For sim-

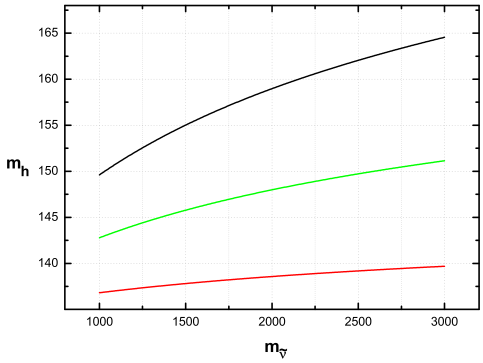

FIG. 1: Lightest Higgs boson mass versus the lightest sneutrino mass for $M_N = \mathrm{diag}\{300, 400, 500\}$, $Y_\nu = \mathrm{diag}\{0, 0, 0.85\}$, $0.85\, \mathrm{I}_{3\times3}$, and $U_{\rm MNS}$ (for curves from bottom to up respectively).

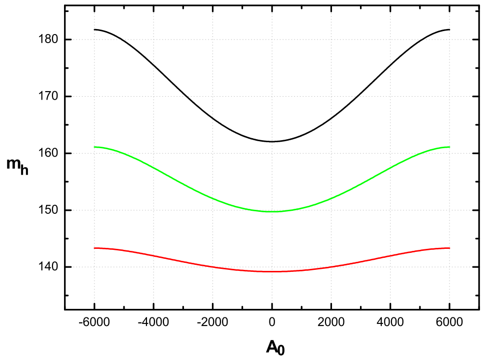

FIG. 2: Lightest Higgs boson mass as a function of the triliniar coupling A_0 for $M_N = \mathrm{diag}\{300, 400, 500\}$, $\tilde{m} = 2500\ \mathrm{I}_{3\times3}$, and $Y_\nu = \mathrm{diag}\{0, 0, 0.85\}$, $0.85\, \mathrm{I}_{3\times3}$, and $U_{\rm MNS}$ (for curves from bottom to up respectively).

plicity, we assume $A_N = A_S = A_0$. In Fig. 2 we display the dependence of m_h on A_0 for the above mentioned three examples of Y_ν, $\tilde{m} = 2.5\ \mathrm{I}_{3\times3}$ TeV, and $M_N = \mathrm{diag}\{300, 400, 500\}$ GeV. From this figure, it can be noted that a large value of A_0 may enhance the value of the Higgs mass by about 20 GeV. This can be understood as follows: for non-vanishing A_0, one may approximate the sneutrino masses as

$$m^2_{\tilde{\nu}} \ \simeq \ m^2_{\tilde{\nu}}\Big|_{A_0=0} + \alpha\, m_D\, A_0, \tag{29}$$

where the coefficient α can be fixed by fitting the sneutrino masses with this expression. Here we kept the dependence on the trilinear coupling to be consistent with the definition of the A-term that led to the expressions in the sneutrino squared-mass matrix (13), namely $A_{ij} = A_0 Y_{ij}$. In this case, Eq. (23) can be written as

$$\delta M_{22} \ \simeq \ \frac{m_D^4}{4\pi^2 v_2^2}\left[\log\frac{m^2_{\tilde{\nu}_H}}{m^2_{\nu_H}} + \log\left(1 + \frac{\alpha\, m_D\, A_0}{m^2_{\tilde{\nu}_H}}\right)\right]. \tag{30}$$

Thus, for A_0 of order TeV, one finds that the Higgs mass is enhanced by few GeVs.

For such large A_ν-term, one should be careful with possible $B-L$ symmetry breaking through a non-vanishing vev of the sneutrino, which also breaks R-parity and makes the model quite involved. In order to avoid this minimum one has to satisfy the following constraint, which is very similar to the usual one imposed in the MSSM to avoid the electric-charge and color symmetry breaking minimum [11], namely

$$|A_\nu|^2 \leq 3\left(m_L^2 + m_N^2 + m^2_{H_2}\right). \tag{31}$$

Since we have m_L and m_N of order $|A_\nu|$, it is clear that this minimum can be safely avoided.

It is interesting to note that recent results from the CMS and ATLAS experiments may indicate the existence of a Higgs boson with mass around 125 GeV [12]. In this case, it is remarkable that in the model described herein the required loop corrections to the Higgs mass can be obtained easily from a combination of the well established MSSM ones and those specific to the $B-L$ sector with stop and sneutrino masses that are smaller than 1 TeV, hence promptly testable at the LHC, while retaining a nature for the light Higgs state which is rather SM-like.

In conclusion, we have calculated the one-loop radiative corrections to the lightest SM-like Higgs boson mass, due to the right-handed (s)neutrinos in a SUSY $B-L$ extension of the SM with inverse seesaw mechanism. We have shown that the upper bound on the Higgs mass can be enhanced to be around 180 GeV. This enhancement may alleviate a possible conflict between the experimental limits from Higgs searches at the LHC and the absolute upper limit predicted in MSSM theoretically, of 130 GeV. It is remarkable that our result remains valid for any model beyond the MSSM with TeV scale right-handed neutrinos (including Left-Right, Pati-Salam and other models derived from $SO(10)$).

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