125 GeV HIGGS BOSON AND RADIATIVE NATURAL SUSY* **

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The observed mass 125 GeV of the Higgs boson at the LHC requires the large stop mass scale $M_{\tilde{t}} \simeq M_{\rm SUSY} \gtrsim 1$ TeV. The allowed region of stop parameters is investigated in MSSM of 3-loop accuracy. There is a sum rule between the MSSM Higgs $\gamma\gamma$-, $b\bar{b}$-production cross section ratios to the SM Higgs boson. Radiative natural SUSY (RNS), satisfying the weak-scale naturalness in MSSM, predicts the small Higgsio mass less than 500 GeV. Correspondingly, the $\gamma\gamma(b\bar{b})$ cross section ratio is reduced (enhanced) in RNS.

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1. Introduction

The discovery of the Higgs boson by the LHC has strong impacts on the particle physics. In the minimal supersymmetric standard model (MSSM), $M_h < M_Z$ is predicted in the tree-level, and thus, its large mass, $M_h \simeq$ 125 GeV, requires a large quantum correction, which is explained by large stop masses $m_{\tilde{t}_{1,2}}$ in the loop. They give the lower limit of the SUSY breaking scale $M_{\rm SUSY}(=\sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}) \gtrsim 1$ TeV [1]. Applying the naturalness condition in the unified scale predicts the stop mass scale less than 700 GeV [2]. However, this possibility is almost excluded from experiment, since CMS suggests no light stop [3] with the mass $m_{\tilde{t}_1} < 0.65$ TeV. Recent analyses from ATLAS and CMS, in the context of the minimal supergravity (mSUGRA or CMSSM) model [4], require $m_{\tilde{g}} \gtrsim 1.4$ TeV for $m_{\tilde{q}} \sim m_{\tilde{g}}$ and $m_{\tilde{g}} \gtrsim 1$ TeV for $m_{\tilde{g}} \ll m_{\tilde{q}}$. The SUSY breaking scale is considered plausibly to be more than ~ 1 TeV.

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Applying naturalness only at the electroweak scale permits the stop mass scale above 1 TeV. A small fine-tuning less than 10% level in reproducing M_h^2 and M_Z^2 requires

$$
\begin{aligned}
|\mu| &\lesssim 300\ \text{GeV}\,, \qquad & m_{\tilde{g}} &= 1 \sim 4\ \text{TeV}\,, \\
m_{\tilde{t}_1} &= 1 \sim 2\ \text{TeV}\,, \qquad & m_{\tilde{t}_2} &= 2 \sim 5\ \text{TeV}\,,
\end{aligned} \tag{1}
$$

where we have considered a SUSY-GUT type model with the non-universal Higgs mass to reproduce the electroweak symmetry breaking radiatively by RGE from GUT scale down to the weak scale. This model is called radiative natural SUSY (RNS) [5–7]. The small Higgsino mass $|\mu|$ is expected directly from naturalness, while the 1st-, 2nd-generation squarks and sleptons have masses $\sim$ 1–8 or 20–30 TeV. There is no SUSY CP problem because of these large masses and only a tiny non-standard contribution to the $(g-2)_\mu$ anomaly are expected. These spectra are consistent with the present data of $b \to s\gamma$, $B_s \to \mu^+\mu^-$. RNS is a promising candidate beyond the SM.

We will investigate the allowed region of stop parameters from $M_h \simeq$ 125 GeV in RNS of the three loop accuracy in the next section.

2. SUSY Higgs boson mass

We analyze mass of the SUSY Higgs boson in the 3-loop accuracy by using the H3m package [8]. The input parameters are selected from a natural SUSY benchmark line (NS3): $(m_{\tilde{t}_1,B}, m_{\tilde{t}_2,B}) = (812.5, 1623.2)$ GeV which corresponds to $M_{\text{SUSY}} = 1212.9$ GeV. It is obtained by varying the third generation scalar mass m_0 [7] at the unification scale. The A_t dependence of H3m points are shown by solid circles in Fig. 1 (left). They are implemented

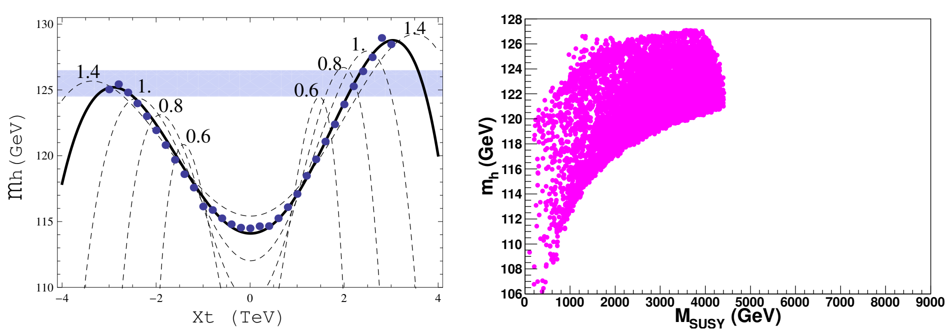

Fig. 1. (left) $A_t(M_{\text{SUSY}})$ dependence of M_h in 3-loop calculation by H3m (solid circles). The solid line implements the H3m results by using the 2LL formula. Based on this result, the dashed lines are obtained with different $M'_{\text{SUSY}}(=\frac{m_{\tilde{t}_1}+m_{\tilde{t}_2}}{2}) =$ 0.6, 0.8, 1.0, 1.4 TeV. $M_h = 125.5 \pm 1$ GeV is shown by gray/blue band. (right) M_{SUSY} dependence of M_h in natural SUSY points in Ref. [7]. See Ref. [9] for details.

by using the 2-loop leading log (2LL) formula in the effective field theory (EFT) approach. The dashed lines correspond to different $M_{\rm SUSY}$ values, 0.6, 0.8, 1.0, 1.4 TeV, see Ref. [9] for details. The peak value of M_h gradually increases with ln $M_{\rm SUSY}$. The Higgs mass constraint $M_h > 124.5$ GeV requires a SUSY breaking scale $M_{\rm SUSY} \gtrsim 0.6$ TeV.

The $M_{\rm SUSY}$ dependence in the radiative natural SUSY [7] is shown in Fig. 1 (right). The points indicate the $\ln M_{\rm SUSY}$ dependence, and in order to explain $M_h > 124.5$ GeV, it is indeed plausible that $M_{\rm SUSY} \gtrsim 1$ TeV.

3. Couplings of Higgs boson

3.1. Ratios of the SUSY Higgs couplings to those of the SM Higgs

The SUSY Higgs mechanism is based on the two Higgs doublet model of type II with the H_u doublet coupled to up-type quarks and the H_d doublet coupled to down-type quarks. After spontaneous symmetry breaking, the physical Higgs states are two CP-even neutral Higgs h, H, one CP-odd neutral pseudo-scalar A and the charged Higgs $H^{\pm}$. We focus on the CP-even neutral Higgs boson h and H, which are related to the flavor eigenstates H_u^0 and H_d^0 by

$$\frac{h}{\sqrt{2}} = c_\alpha H_u^0 - s_\alpha H_d^0 \,, \qquad \frac{H}{\sqrt{2}} = s_\alpha H_u^0 + c_\alpha H_d^0 \,, \tag{2}$$

where $H_{u,d}^0$ is the shorthand for the real part of $H_{u,d}^0 - \langle H_{u,d}^0 \rangle$. We use the notation $s_\alpha = \sin\alpha$, $c_\alpha = \cos\alpha$, and $t_\alpha = \tan\alpha$. $\tan\beta = v_u/v_d$ is defined in terms of the Higgs vacuum expectation values. Our interest is in large $\tan\beta$, $\tan\beta \gtrsim 20$, and in the decoupling regime with large m_A for which $\alpha \simeq \beta - \frac{\pi}{2}$.

The ratios of the h and H couplings to those of the SM Higgs $h_{\rm SM}$, denoted as $r^{h,H}_{PP} (\equiv g_{h,H\ P\bar{P}}/g_{h_{\rm SM}P\bar{P}})$, are given by [10]

$$\begin{aligned}
r^h_{VV} &= s_{\beta-\alpha} \,, \qquad r^h_{tt} = r^h_{cc} = \frac{c_\alpha}{s_\beta} \,, \qquad r^h_{\tau\tau} = \frac{-s_\alpha}{c_\beta} \,, \\
r^h_{bb} &= \frac{-s_\alpha}{c_\beta}\left[1 - \frac{\Delta_b}{1+\Delta_b}\left(1 + \frac{1}{t_\alpha t_\beta}\right)\right] , \\
r^H_{VV} &= c_{\beta-\alpha} \,, \qquad r^H_{tt} = r^H_{cc} = \frac{s_\alpha}{s_\beta} \,, \qquad r^H_{\tau\tau} = \frac{c_\alpha}{c_\beta} \,, \\
r^H_{bb} &= \frac{c_\alpha}{c_\beta}\left[1 - \frac{\Delta_b}{1+\Delta_b}\left(1 - \frac{t_\alpha}{t_\beta}\right)\right] .
\end{aligned} \tag{3}$$

Here, we include the 1-loop contribution Δ_b to the $b\bar{b}$ coupling, which is negligibly small for the small μ value in RNS.

The $gg, \gamma\gamma$ coupling ratios $r^{\phi}_{gg,\gamma\gamma}$ for $\phi = h, H, A$ relative to $h_{\rm SM}$ are [11]

$$r^{\phi}_{gg} = \frac{I^{\phi}_{tt} r^{h}_{tt} + I^{\phi}_{bb} r^{h}_{bb}}{I^{\phi}_{tt} + I^{\phi}_{bb}}\,, \qquad r^{\phi}_{\gamma\gamma} = \frac{\frac{7}{4} I^{\phi}_{WW}\; r^{h}_{VV} - \frac{4}{9} I^{\phi}_{tt} r^{h}_{tt} - \frac{1}{9} I^{\phi}_{bb} r^{h}_{bb}}{\frac{7}{4} I^{\phi}_{WW} - \frac{4}{9} I^{\phi}_{tt} - \frac{1}{9} I^{\phi}_{bb}}\,, \tag{4}$$

where $I^{\phi}_{WW,tt,bb}$ represent the triangle-loop contributions to the amplitudes normalized to the $M_h \to 0$ limit [12–14].

$XX \to h \to PP$ cross section ratios [11] relative to $h_{\rm SM}$ are given by

$$\begin{aligned} \sigma_P &\equiv \frac{\sigma_{PP}}{\sigma_{\rm SM}} = \frac{\sigma_{XX\to PP}}{\sigma_{XX\to h_{\rm SM}\to PP}} = \frac{\left|r^{h}_{XX} r^{h}_{PP}\right|^2}{R^h}\,, \\ R^h &= \frac{\mathit{\Gamma}^{h}_{\rm tot}}{\mathit{\Gamma}^{h_{\rm SM}}_{\rm tot}} = 0.57\left|r^{h}_{bb}\right|^2 + 0.06\left|r^{h}_{\tau\tau}\right|^2 + 0.25\left|r^{h}_{VV}\right|^2 + 0.09\left|r^{h}_{gg}\right|^2 + 0.03\left|r^{h}_{cc}\right|^2\,, \end{aligned} \tag{5}$$

where R^h is the ratio of the h total width to that of $h_{\rm SM}$, $\mathit{\Gamma}^{\rm tot}_{h_{\rm SM}} = 4.14$ MeV [15] for $M_h = 125.5$ GeV. The coefficients in RHS of R^h are the SM Higgs branching fractions. Here, we have assumed no appreciable decays to dark matter.

3.2. Sum rule of cross-section ratios

In the large m_A region close to the decoupling limit, α takes a value

$$\alpha = \beta - \frac{\pi}{2} + \epsilon \tag{6}$$

with $|\epsilon| < \frac{\pi}{2} - \beta$. Then, the r^{h}_{XX} of Eq. (3) are well approximated by

$$r^{h}_{VV} = 1\,, \quad r^{h}_{tt,cc} = 1 + \epsilon/t_\beta\,, \quad r^{h}_{\tau\tau} \simeq 1 - \epsilon t_\beta\,, \quad r^{h}_{bb} \simeq 1 - \frac{1}{1+\Delta_b}\epsilon t_\beta \tag{7}$$

through first order in ϵ. The $r^{h}_{tt,cc}$ are close to unity because those deviations from SM are t_β suppressed. Thus,

$$r^{h}_{gg} \simeq r^{h}_{\gamma\gamma} \simeq 1\,, \tag{8}$$

since the bottom triangle loop function I^{h}_{bb} is negligible in Eq. (4). Only $r^{h}_{bb}, r^{h}_{\tau\tau}$ can deviate sizably from unity for large m_A and large $\tan\beta$. Following Eq. (5), the $\sigma_P \equiv \sigma_{PP}/\sigma_{\rm SM}$ of the other channels are commonly reduced (enhanced) in correspondence with $r^h_{bb} > 1$ ($r^h_{bb} < 1$). We predict the cross sections relative to their individual SM expectations

$$\sigma_\gamma = \sigma_W = \sigma_Z = \frac{1}{0.6\left(r^h_{bb}\right)^2 + 0.4}\,, \tag{9}$$

and

$$0.4\sigma_\gamma + 0.6\sigma_b = 1\,, \tag{10}$$

where the SM $b\bar{b}$ branching fraction is approximated as 60%. Equation (10) holds independently of the production process. Enhanced σ_γ implies reduced σ_b, as well as enhanced σ_W and σ_Z. Or, reduced σ_γ implies enhanced σ_b, as well as reduced σ_W and σ_Z.

3.3. Flavor-tuning of mixing angle α

Note that $r^h_{bb,\tau\tau} = 1$ in the exact decoupling limit $m_A \to \infty$ for which $\epsilon = 0$. Flavor-tuning of ϵ to be small but non-zero is necessary to obtain a significant variation of r^h_{bb} from unity. Positive (negative) ϵ gives bb-reduction (enhancement) from Eq. (7).

The mixing angle α is obtained by diagonalizing the squared-mass matrix of the neutral Higgs in the u, d basis. Their elements at tree-level are

$$\left(M^2_{ij}\right)^{\text{tree}} = M_Z^2 s_\beta^2 + m_A^2 c_\beta^2\,; \quad M_Z^2 c_\beta^2 + m_A^2 s_\beta^2\,; \quad -\left(M_Z^2 + m_A^2\right) s_\beta c_\beta \tag{11}$$

for $ij = 11; 22; 12$, respectively, which gives $\epsilon < 0$ in all region of m_A. Thus, in order to get $b\bar{b}$-reduction, it is necessary to cancel $(M^2_{12})^{\text{tree}}$ by higher order terms ΔM^2_{ij}.

In 2LL approximation, the ΔM^2_{ij} are given by [16, 17]

$$\begin{aligned} M^2_{ij} &= \left(M^2_{ij}\right)^{\text{tree}} + \Delta M^2_{ij}\,, \qquad (12)\\ \Delta M^2_{11} &= F_3 \frac{3\bar{m}_t^4}{4\pi^2 v^2 s_\beta^2}\left[t\left(1 - G_{\frac{15}{2}} t\right) + a_t x_t \left(1 - \frac{a_t x_t}{12}\right)\left(1 - 2G_{\frac{9}{2}} t\right)\right] \\ &\quad - M_Z^2 s_\beta^2 \left(1 - F_3\right)\,, \\ \Delta M^2_{22} &= -F_{\frac{3}{2}} \frac{\bar{m}_t^4}{16\pi^2 v^2 s_\beta^2}\left[\left(1 - 2G_{\frac{9}{2}} t\right)\left(x_t \bar{\mu}\right)^2\right]\,, \\ \Delta M^2_{12} &= -F_{\frac{9}{4}} \frac{3\bar{m}_t^4}{8\pi^2 v^2 s_\beta^2}\left[\left(1 - 2G_{\frac{9}{2}} t\right)\left(x_t \bar{\mu}\right)\left(1 - \frac{a_t x_t}{6}\right)\right] + M_Z^2 s_\beta c_\beta \left(1 - F_{\frac{3}{2}}\right)\,, \qquad (13) \end{aligned}$$

where $F_l = 1/(1 + l\frac{h_t^2}{8\pi^2} t)$ with $l = 3, \frac{3}{2}, \frac{9}{4}$ and $G_l = -\frac{1}{16\pi^2}(l h_t^2 - 32\pi\alpha_s)$ with $l = \frac{15}{2}, \frac{9}{2}$. The F_l are due to the wave function (WF) renormalization of the

H_u field and the index l is related to numbers of H_u^0 fields in the effective potential of the two Higgs doublet model. $F_3\xi^4 \simeq F_{\frac{9}{4}}\xi^3 \simeq F_{\frac{3}{2}}\xi^2 \simeq 1$, where ξ is defined by $H_u(M_s) = H_u(\bar{m}_t)\xi$, where $\xi = F_{\frac{3}{4}}^{-1}$. We get

$$\epsilon = -\frac{2M_Z^2 + \Delta M_{11}^2 - \Delta M_{22}^2 - \Delta M_{12}^2 \tan\beta}{m_A^2 \tan\beta} \,. \tag{14}$$

In order to get $\gamma\gamma$ enhancement, $\sigma_\gamma > 1$ ($\epsilon > 0$), flavor-tuning (FT), a cancellation of $(M_{12}^2)^{\text{tree}}$ by the loop-level ΔM_{12}^2 contribution is required. This is possible for rather large values of $\bar{\mu}$ and $\tan\beta$ [10, 18].

Here, $a_t, x_t, \bar{\mu}$ have scale $Q = M_{\text{SUSY}}$, while the $\tan\beta = v_u/v_d$ is defined at the weak scale $Q = \bar{m}_t \simeq 163.5$ GeV. The relation $\cot\beta = \cot\beta(\bar{m}_t) = \cot\beta(M_s)\ \xi^{-1}$ will be used in the following calculation.

Numerically $\alpha_{\text{s}} = \alpha_{\text{s}}(\bar{m}_t) = 0.109$ giving $-32\pi\alpha_{\text{s}} = -10.9$, while $h_t = \bar{m}_t/v = 0.939$ is small. $G_{\frac{15}{2},\frac{9}{2}} = 0.0274, 0.0442$ and $t = \log(\frac{1\ \text{TeV}}{\bar{m}_t})^2 = 3.62$; thus, $G_{\frac{15}{2}}t = 0.099$ and $2G_{\frac{9}{2}}t = 0.320$, and $F_3 = 0.892$.

In large m_A limit, the M_h^2 is expressed by

$$\begin{aligned} M_h^2 \;=\; & M_Z^2 c_{2\beta}^2 + F_3 \frac{3\bar{m}_t^4}{4\pi^2 v^2}\left[t\left(1 - G_{\frac{15}{2}}t\right) + \left(1 - 2G_{\frac{9}{2}}t\right)\left(x_t^2 - \frac{x_t^4}{12}\right)\right] \\ & -M_Z^2\left[s_\beta^4\left(1 - F_3\right) - 2s_\beta^2 c_\beta^2\left(1 - F_{\frac{3}{2}}\right)\right] , \end{aligned} \tag{15}$$

where the Higgs WF renormalization factor ξ is retained in the denominator of F_3. This F_3 factor is usually expanded to the numerator in 2LL approximation, and correspondingly $G_{\frac{15}{2}}$ and $G_{\frac{9}{2}}$ are replaced by $G_{\frac{3}{2}}$: $M_h^2 = M_Z^2 c_{2\beta}^2 + \frac{3\bar{m}_t^4}{4\pi^2 v^2}[t(1 - G_{\frac{3}{2}}t) + (1 - 2G_{\frac{3}{2}}t)(x_t^2 - \frac{x_t^4}{12})] - M_Z^2 s_\beta^4 \frac{3h_t^2}{8\pi^2}t$. However, numerically Eq. (15) significantly increases M_h at large M_{SUSY}. Equation (15) gives increasing M_h as M_{SUSY} increases up to ~ 7 TeV, while the usual formula with the expansion approximated for F_3 gives decreasing M_h when $M_{\text{SUSY}} > 1.3$ TeV and is not applicable at large M_{SUSY}.

In Eq. (15) we require $m_h \geq 124$ GeV. This implies

$$1.95(\equiv x_{\text{tmin}}) < |x_t| < 2.86(\equiv x_{\text{tmax}}) \,, \tag{16}$$

where we should note that the positive x_t branch is favored by the SUSY renormalization group prediction [9].

Correspondingly, the μ dependence of σ_γ, σ_b and σ_τ are given respectively by the two curves in Fig. 2, where we take $\tan\beta = 20$.

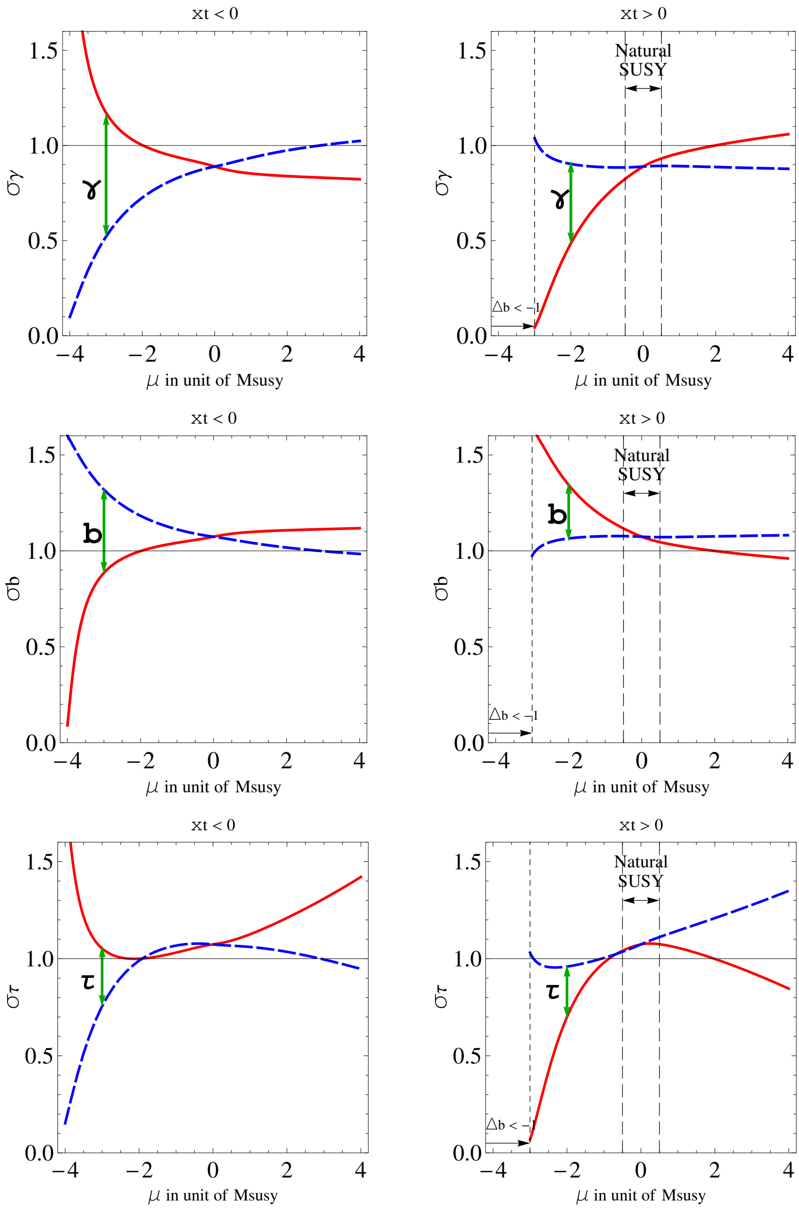

Fig. 2. $\bar{\mu}$ dependence of $\sigma_\gamma = \sigma_{\gamma\gamma}/\sigma_{\rm SM}$ (upper panel), $\sigma_b = \sigma_{b\bar{b}}/\sigma_{\rm SM}$ (middle panel), and $\sigma_b = \sigma_{b\bar{b}}/\sigma_{\rm SM}$ (lower panel) for $m_A = 500$ GeV: Their allowed values are between the solid/red curve (corresponding to $|x_t| = x_{\rm tmax}$) and the dashed/blue curve (corresponding to $|x_t| = x_{\rm tmin}$). Left (Right) panels show negative (positive) x_t region. Deviations from unity are enlarged for a large negative $\bar{\mu}$, but there the perturbative calculation is unreliable due to a large quantum correction.

3.4. Natural SUSY predictions

The allowed region in RNS is shown by a narrow range between two vertical lines in each panel of Fig. 2, where $|\mu| \leq 500$ GeV. The RNS always predicts the $b\bar{b}$-enhancement and $\gamma\gamma$ reduction [10].

TABLE I

Natural SUSY predictions of σ_γ, σ_b, and σ_τ for $m_A = 0.5, 1$ TeV.

m_A	σ_γ	σ_b	σ_τ
500 GeV	$0.82 \sim 0.91$	$1.06 \sim 1.12$	$1.04 \sim 1.08$
1 TeV	$0.95 \sim 0.98$	$1.01 \sim 1.03$	$1.01 \sim 1.02$

4. Concluding remarks

The mass 125 GeV of the Higgs boson is consistent with the SUSY breaking scale $M_{\rm SUSY} > 0.6$ TeV due to the three loop analysis of the SUSY Higgs boson. It is consistent with the stop masses anticipated in RNS. The Higgs couplings to the SM particles are predicted in RNS. The MSSM Higgs $\gamma\gamma$ cross section ratio to the SM Higgs satisfies the sum rule together with the $b\bar{b}$ cross section ratio. The flavor-tuning of the neutral Higgs mixing angle α needed to reproduce $\gamma\gamma$ enhancement requires a large μ ~TeV and large $\tan\beta$. For small $|\mu| \lesssim 0.5$ TeV in RNS, the $\gamma\gamma$-suppression relative to the SM is always predicted. Thus, the precision of LHC measurements of the $\gamma\gamma$, $W*W$, $Z*Z$ and $b\bar{b}$ signals of the 125 GeV Higgs boson can test MSSM and RNS models.

Finally, I remark that the relatively small wino and bino masses are also expected in RNS. The wino pair production signal $\tilde{W}_2^{\pm}\tilde{Z}_4 \to (W^{\pm}\tilde{Z}_{1,2}) + (W^{\pm}\tilde{W}_1^{\mp})$ occurs at substantial rates and the detection of the same-sign diboson signal is a promising method to check the RNS [19] at the LHC.

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