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Flavor-tuned 125 GeV supersymmetric Higgs boson at the LHC: Test of minimal and natural supersymmetric models

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We show that an enhanced two-photon signal of the Higgs boson, h, observed with 125 GeV mass by the ATLAS and CMS collaborations, can be obtained if it is identified principally with the neutral H_u^0 of the two Higgs doublets of minimal supersymmetry. We focus on sparticles and the pseudoscalar Higgs A with TeV masses. The off-diagonal element of the (H_u^0, H_d^0) mass matrix in the flavor basis must be suppressed, and this requires a large Higgsino mass parameter, $\mu \sim$ TeV, and large $\tan\beta$. A minimal supersymmetric Standard Model sum rule is derived that relates $\gamma\gamma$ and $b\bar{b}$ rates, and a $\gamma\gamma$ enhancement relative to the SM predicts $b\bar{b}$ reduction. On the contrary, natural supersymmetry requires $|\mu| <$ ~0.5 TeV, for which $\gamma\gamma$ is reduced and $b\bar{b}$ is enhanced. This conclusion is independent of the m_A value and the supersymmetry quantum correction Δ_b. Relative $\tau\bar{\tau}$ to $b\bar{b}$ rates are sensitive to Δ_b.

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A $\gamma\gamma$ enhancement of the 125 GeV Higgs boson signal relative to the Standard Model (SM) expectation has been reported by the ATLAS and CMS experiments at the LHC [1,2]. We investigate this in the minimal supersymmetric Standard Model (MSSM) in the region of large $m_A \sim$ TeV by flavor-tuning of the mixing angle α between two neutral charge conjugation parity even (CP-even) Higgs flavor states H_u^0 and H_d^0, with the 125 GeV Higgs signal identified principally with H_u^0. Then, the $b\bar{b}$ decay, which is predicted to be the dominant decay of the SM Higgs boson, is reduced. The production cross sections of other channels are correspondingly enhanced, except possibly $\tau\tau$. We relate the cross-section enhancements/suppressions in $\gamma\gamma/b\bar{b}/\tau\tau$ channels compared with those of the SM Higgs boson. We also consider the consequences for natural supersymmetry (SUSY) [3]. Our focus is on a heavy pseudoscalar A and large $\tan\beta \equiv \langle H_u^0\rangle/\langle H_d^0\rangle$, a region that has not yet been constrained by LHC experiments [4]. Light stau [5,6] and light stop quark [7] scenarios that have been considered are outside of our purview.

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 [8–10] 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 pseudoscalar 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

$$\frac{h}{\sqrt{2}} = c_\alpha H_u^0 - s_\alpha H_d^0, \qquad \frac{H}{\sqrt{2}} = s_\alpha H_u^0 + c_\alpha H_d^0, \quad (1)$$

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 $t_\alpha = \tan\alpha$. 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 $r_{PP}^{h,H} (\equiv g_{h,HP\bar{P}}/g_{h_{SM}P\bar{P}})$, are given by

$$\begin{aligned}
&r_{VV}^h = s_{\beta-\alpha}, \qquad r_{tt}^h = r_{cc}^h = \frac{c_\alpha}{s_\beta}, \qquad r_{\tau\tau}^h = \frac{-s_\alpha}{c_\beta},\\
&r_{bb}^h = \frac{-s_\alpha}{c_\beta}\left[1 - \frac{\Delta_b}{1+\Delta_b}\left(1 + \frac{1}{t_\alpha t_\beta}\right)\right],\\
&r_{VV}^H = c_{\beta-\alpha}, \qquad r_{tt}^H = r_{cc}^H = \frac{s_\alpha}{s_\beta}, \qquad r_{\tau\tau}^H = \frac{c_\alpha}{c_\beta},\\
&r_{bb}^H = \frac{c_\alpha}{c_\beta}\left[1 - \frac{\Delta_b}{1+\Delta_b}\left(1 - \frac{t_\alpha}{t_\beta}\right)\right]
\end{aligned} \quad (2)$$

where we include the 1-loop contribution Δ_b to the $b\bar{b}$ coupling. Δ_b is the b-quark mass correction factor [11,12], which may be sizable, especially if both μ and $\tan\beta$ are large.

$$\Delta_b = \bar{\mu} t_\beta\left[\frac{2\alpha_s}{3\pi}\hat{m}_{\tilde{g}} I(\hat{m}_{\tilde{g}}^2, \hat{m}_{\tilde{b}_1}^2, \hat{m}_{\tilde{b}_2}^2) + \frac{h_t^2}{16\pi^2} a_t I(\bar{\mu}^2, \hat{m}_{\tilde{t}_1}^2, \hat{m}_{\tilde{t}_2}^2)\right], \quad (3)$$

$$\begin{aligned}
&I(x,y,z) = -\frac{xy\ln x/y + yz\ln y/z + zx\ln z/x}{(x-y)(y-z)(z-x)},\\
&I(x,y,z=y) = -\left[x - y + x\log\frac{y}{x}\right]\Big/(x-y)^2,\\
&I(x,x,x) = \frac{1}{2x}.
\end{aligned} \quad (4)$$

The first (second) term of Δ_b is due to the sbottom-gluino (stop quark—chargino) loop. We nominally take $M_{\rm susy} =$ 1 TeV and express sparticle masses $\hat{m}$ in units of $M_{\rm susy}$. The top Yukawa coupling is $h_t = \bar{m}_t/v_u = \bar{m}_t/(v s_\beta)$ and

$\bar{m}_t = m_t(\bar{m}_t) = 163.5$ GeV is the running top quark mass [13]. We consider $m_Q = m_U = m_D = M_{\rm susy}$ for the squark masses in the third generation.

The off-diagonal element of the stop quark squared mass matrix is $\bar{m}_t X_t$ where the stop quark mixing parameter X_t is given by $X_t = A_t - \mu/t_\beta$. The quantities A_t, μ and X_t are also defined in units of $M_{\rm susy}$ as $a_t \equiv A_t/M_{\rm susy}$, $\bar{\mu} \equiv \mu/M_{\rm susy}$, and $x_t \equiv X_t/M_{\rm susy} = a_t - \bar{\mu}/t_\beta$. Our sign convention for μ and A_t is the same as [14], opposite to the sign convention of Ref. [15]. We fix $\hat{m}_{\tilde{g}} = 2$, well above the current LHC reach, $\hat{m}_{\tilde{b}_1} = \hat{m}_{\tilde{b}_2} = 1$, and $\hat{m}_{\tilde{t}_1} = 0.8$, $\hat{m}_{\tilde{t}_2} = 1.2$. A stop quark mass difference $m_{\tilde{t}_2} - m_{\tilde{t}_1} \geq 0.4$ TeV is chosen in accord with the natural SUSY prediction [16]. Then Δ_b is well approximated numerically by

$$\Delta_b \simeq \bar{\mu}\frac{t_\beta}{20}\left[0.26 + \left(\frac{0.09}{|\bar{\mu}| + 0.6} - 0.003\right)a_t\right], \tag{5}$$

where the first and the second terms in the square bracket are the values of the gluino and the chargino contributions, respectively.

The chargino and neutralino masses have no special role except possibly in $b \to s\gamma$ decay, but consistency with natural SUSY has been found there [17]. Large m_A implies a large charged Higgs H^+ mass and this suppresses the H^+ loop contribution to $b \to s\gamma$.

The gg, $\gamma\gamma$ coupling ratios $r^\phi_{gg,\gamma\gamma}$ for $\phi = h$, H, A relative to those of $h_{\rm SM}$ are [18]

$$r^\phi_{gg} = \frac{I^\phi_{tt} r^h_{tt} + I^\phi_{bb} r^h_{bb}}{I^\phi_{tt} + I^\phi_{bb}},$$
$$r^\phi_{\gamma\gamma} = \frac{\frac{7}{4}I^\phi_{WW} r^h_{VV} - \frac{4}{9}I^\phi_{tt} r^h_{tt} - \frac{1}{9}I^\phi_{bb} r^h_{bb}}{\frac{7}{4}I^\phi_{WW} - \frac{4}{9}I^\phi_{tt} - \frac{1}{9}I^\phi_{bb}}, \tag{6}$$

where $I^\phi_{WW,tt,bb}$ represent the triangle-loop contributions to the amplitudes normalized to the $m_h \to 0$ limit [19–21].

The $XX \to h \to PP$ cross section ratios [18] relative to $h_{\rm SM}$ are obtained from

$$\sigma_P \equiv \frac{\sigma_{PP}}{\sigma_{\rm SM}} = \frac{\sigma_{XX\to PP}}{\sigma_{XX\to h_{\rm SM}\to PP}} = \frac{|r^h_{XX} r^h_{PP}|^2}{R^h}, \tag{7}$$

$$R^h = \frac{\Gamma^h_{\rm tot}}{\Gamma^{h_{\rm SM}}_{\rm tot}} = 0.57|r^h_{bb}|^2 + 0.06|r^h_{\tau\tau}|^2 + 0.25|r^h_{VV}|^2 + 0.09|r^h_{gg}|^2 + 0.03|r^h_{cc}|^2, \tag{8}$$

where R^h is the ratio of the h total width to that of $h_{\rm SM}$, $\Gamma^{\rm tot}_{h_{\rm SM}} = 4.14$ MeV [22] for $m_h = 125.5$ GeV. The coefficients in Eq. (8) are the SM Higgs branching fractions. Here we have assumed no appreciable h decays to dark matter.

Sum rule of cross-section ratios.—In the large m_A region close to the decoupling limit, α takes a value

$$\alpha = \beta - \frac{\pi}{2} + \epsilon \tag{9}$$

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

$$r^h_{VV} = 1, \qquad r^h_{tt,cc} = 1 + \epsilon/t_\beta,$$
$$r^h_{\tau\tau} \simeq 1 - \epsilon t_\beta, \qquad r^h_{bb} \simeq 1 - \frac{1}{1+\Delta_b}\epsilon t_\beta \tag{10}$$

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

$$r^h_{gg} \simeq r^h_{\gamma\gamma} \simeq 1, \tag{11}$$

since the bottom triangle loop function I^h_{bb} is negligible in Eq. (6). Only r^h_{bb}, $r^h_{\tau\tau}$ can deviate sizably from unity for large m_A and large $\tan\beta$. Following Eqs. (7) and (8), 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

$$\sigma_\gamma = \sigma_W = \sigma_Z = \frac{1}{0.6(r^h_{bb})^2 + 0.4}, \tag{12}$$

and

$$0.4\sigma_\gamma + 0.6\sigma_b = 1, \tag{13}$$

where the SM $b\bar{b}$ branching fraction [23] is approximated as 60%. Equation (12) holds independently of the production process. Enhanced σ_γ implies reduced σ_b, as well as enhanced σ_W and σ_Z.

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 nonzero is necessary to obtain a significant variation of r^h_{bb} from unity. Positive (negative) ϵ gives bb reduction (enhancement).

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

$$(M^2_{ij})^{\rm tree} = M^2_Z s^2_\beta + m^2_A c^2_\beta;\ M^2_Z c^2_\beta + m^2_A s^2_\beta;\ -(M^2_Z + m^2_A)s_\beta c_\beta, \tag{14}$$

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

In the 2-loop leading-log (LL) approximation the ΔM^2_{ij} are given [14,24] by

$$M^2_{ij} = (M^2_{ij})^{\rm tree} + \Delta M^2_{ij}, \tag{15}$$

where

$$\Delta M_{11}^2 = F_3 \frac{3\bar{m}_t^4}{4\pi^2 v^2 s_\beta^2}\Big[t\Big(1 - G_{\frac{15}{2}}t\Big) + a_t x_t\Big(1 - \frac{a_t x_t}{12}\Big)\Big(1 - 2G_{\frac{9}{2}}t\Big)\Big] - M_Z^2 s_\beta^2(1 - F_3),$$

$$\Delta M_{22}^2 = -F_{\frac{3}{2}} \frac{\bar{m}_t^4}{16\pi^2 v^2 s_\beta^2}\Big[\Big(1 - 2G_{\frac{9}{2}}t\Big)(x_t\bar{\mu})^2\Big],$$

$$\Delta M_{12}^2 = -F_{\frac{9}{4}} \frac{3\bar{m}_t^4}{8\pi^2 v^2 s_\beta^2}\Big[\Big(1 - 2G_{\frac{9}{2}}t\Big)(x_t\bar{\mu})\Big(1 - \frac{a_t x_t}{6}\Big)\Big] + M_Z^2 s_\beta c_\beta\Big(1 - F_{\frac{3}{2}}\Big), \quad (16)$$

where $F_l = 1/(1 + l\frac{h_t^2}{8\pi^2}t)$ with $l = 3, \frac{3}{2}, \frac{9}{4}$ and $G_l = -\frac{1}{16\pi^2}(lh_t^2 - 32\pi\alpha_s)$ with $l = \frac{15}{2}, \frac{9}{2}$. The F_l are due to the wave function renormalization of the H_u field and the index l is the 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 $\xi = F_{\frac{3}{4}}^{-1}$. The formulas of Refs. [14,24] are based on the expansion $F_l = 1 - l\frac{h_t^2}{8\pi^2}t$, but our formula of F_l is more exact and has better approximation at large t. The parameter $\tan\beta = v_u/v_d$ is defined in terms of the Higgs vacuum expectation values $v_{u,d} = \langle H_{u,d}^0\rangle$ at the minimum of the 1-loop effective potential at the weak scale $\mu = \bar{m}_t$ and $v = \sqrt{v_u^2 + v_d^2} \simeq 174$ GeV, while a_t, x_t, $\bar{\mu}$ have scale $\mu = M_{\rm susy}$. The relation $\cot\beta(\bar{m}_t) = \cot\beta(M_s)\xi^{-1}$ will be used in the following calculation.

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

In the large m_A limit, the m_h^2 expression is

$$m_h^2 = M_Z^2 c_{2\beta}^2 + F_3 \frac{3\bar{m}_t^4}{4\pi^2 v^2}\Big[t\Big(1 - G_{\frac{15}{2}}t\Big) + \Big(1 - 2G_{\frac{9}{2}}t\Big)\Big(x_t^2 - \frac{x_t^4}{12}\Big)\Big] - M_Z^2\Big[s_\beta^4(1 - F_3) - 2s_\beta^2 c_\beta^2\Big(1 - F_{\frac{3}{2}}\Big)\Big], \quad (17)$$

where the Higgs wave function renormalization factor ξ is retained in the denominator of F_3. In the usual expansion of the F_3 denominator $G_{\frac{15}{2}}$ and $G_{\frac{9}{2}}$ are replaced by $G_{\frac{3}{2}}$: $m_h^2 = M_Z^2 c_{2\beta}^2 + \frac{3\bar{m}_t^4}{4\pi^2 v^2}[t(1 - G_3 t) + (1 - 2G_3 t) \times (x_t^2 - \frac{x_t^4}{12})] - M_Z^2 s_\beta^4 \frac{3h_t^2}{8\pi^2}t$. However, numerically Eq. (17) significantly increases m_h at large $M_{\rm susy}$ as shown in Fig. 1. Equation (17) gives increasing m_h as $M_{\rm susy}$ increases up to ~7 TeV, while the usual formula with the expansion approximation of F_3 gives decreasing m_h when $M_{\rm susy} >$ 1.3 TeV and it is not applicable at large $M_{\rm susy}$.

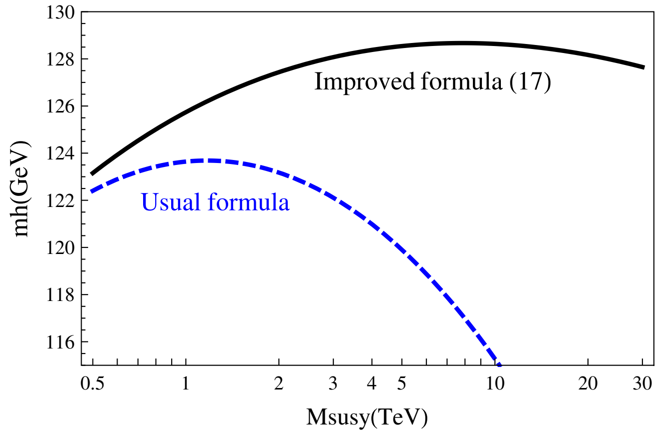

FIG. 1 (color online). $M_{\rm susy}$ dependence of Higgs mass m_h by the improved formula Eq. (17) (solid black) in comparison with the one by the usual 2LL approximation (dashed blue) with a linear expansion in t of F_3. In this illustration x_t is taken to be $\sqrt{6}$ following the "maximal-mixing" condition, and $\tan\beta = 20$.

The experimental m_h determinations from the LHC experiments are [1,2]

$$m_h = 125.3 \pm 0.4 \pm 0.4, \quad 126.0 \pm 0.4 \pm 0.4\,\text{GeV}. \quad (18)$$

It seems unlikely that the central m_h determination will change much with larger statistics because of the excellent mass resolution in the $\gamma\gamma$ channel. The experimental m_h value is near the maximum possible value of m_h in Eq. (17) and this constrains the value of x_t to $|x_t| \simeq \sqrt{6}$, to maximize the term $x_t^2 - \frac{x_t^4}{12}$. This is known as "maximal-mixing" in the stop quark mass matrix [16]. In Eq. (17) we require $m_h \geq 124$ GeV. This implies

$$1.95(\equiv x_{\rm tmin}) < |x_t| < 2.86(\equiv x_{\rm tmax}), \quad (19)$$

where we should note that the positive x_t branch is favored by the SUSY renormalization group prediction [16].

By using Eq. (15) the Higgs mixing angle α is determined from

$$t_{2\alpha} = \frac{2M_{12}^2}{M_{22}^2 - M_{11}^2} \simeq \frac{(m_A^2 + M_Z^2)s_{2\beta} - 2\Delta M_{12}^2}{(m_A^2 - M_Z^2)c_{2\beta} + (\Delta M_{11}^2 - \Delta M_{22}^2)}, \quad (20)$$

$$\Delta M_{12}^2 \simeq -\frac{\bar{\mu}}{s_\beta^2}x_t\Big(1 - \frac{x_t^2}{6}\Big)558\ \text{GeV}^2 + 24 \cdot \frac{20}{\tan\beta}\ \text{GeV}^2. \quad (21)$$

Defining $z(\equiv M_Z^2/m_A^2)$, $\delta(\equiv \Delta M_{12}^2/m_A^2)$, and $\eta(\equiv \frac{1}{2} \times (\Delta M_{11}^2 - \Delta M_{22}^2)/m_A^2)$, ϵ is simply given in the first order of z, δ, and η by

$$\epsilon = -2\frac{z + \eta}{\tan\beta} + \delta. \quad (22)$$

We note that r_{bb}^h is related to ϵ through Eq. (10). With the x_t constraint in Eq. (19), we can derive the allowed

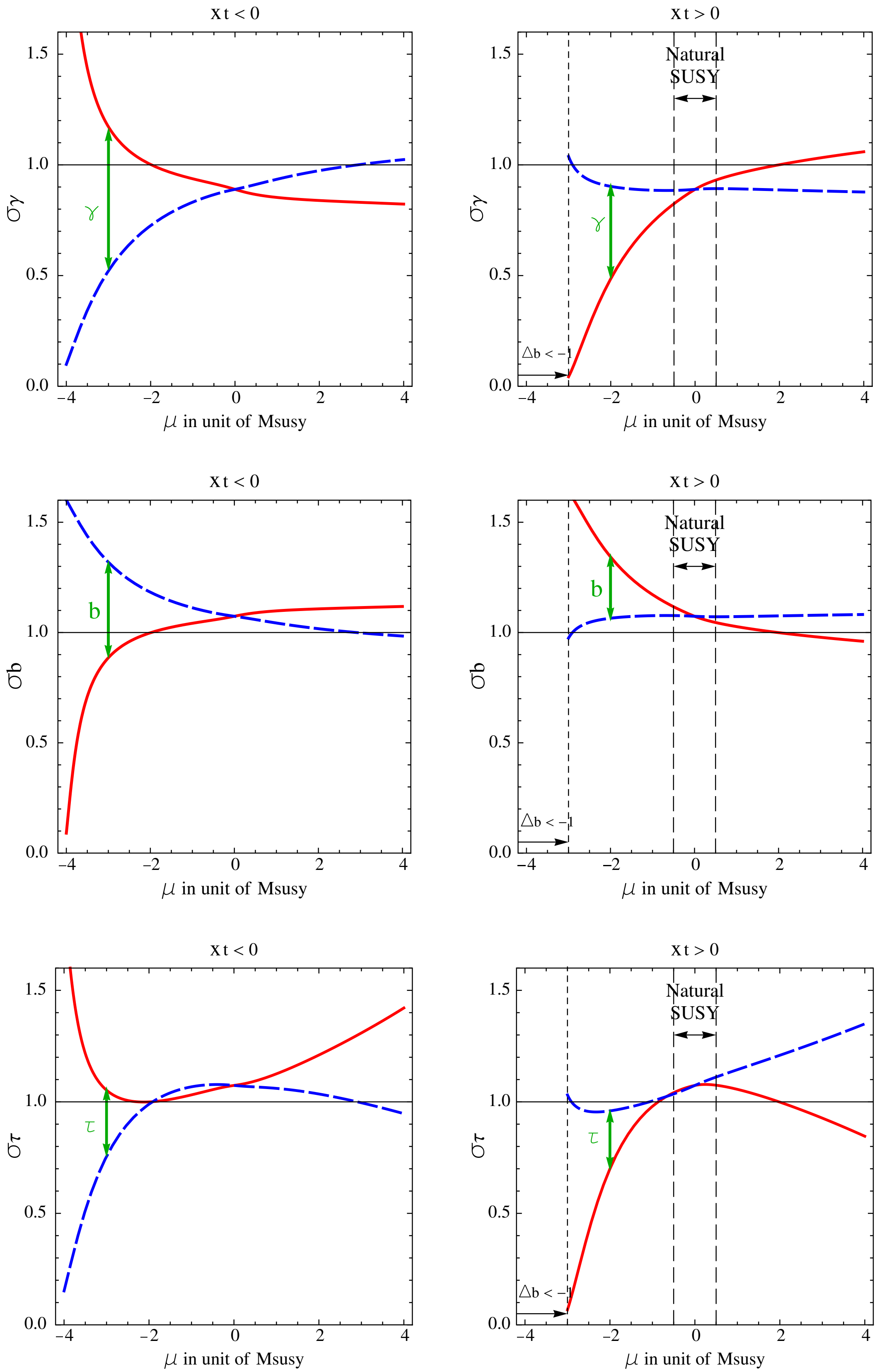

FIG. 2 (color online). $\bar{\mu}$ dependence of $\sigma_\gamma = \sigma_{\gamma\gamma}/\sigma_{\rm SM}$ (upper panels), $\sigma_b = \sigma_{b\bar{b}}/\sigma_{\rm SM}$ (middle panels), and $\sigma_b = \sigma_{b\bar{b}}/\sigma_{\rm SM}$ (lower panels) for $m_A = 500$ GeV. Their allowed values are between 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.

region of r^h_{bb} for each $\bar{\mu}$ value. Correspondingly, the allowed regions of $\sigma_\gamma(=\sigma_{\gamma\gamma}/\sigma_{\rm SM}=\frac{1}{0.6(r^h_{bb})^2+0.4})$, $\sigma_b(=\sigma_{b\bar{b}}/\sigma_{\rm SM}=\frac{(r^h_{bb})^2}{0.6(r^h_{bb})^2+0.4})$ and $\sigma_\tau(=\sigma_{\tau\tau}/\sigma_{\rm SM}=\frac{(r^h_{\tau\tau})^2}{0.6(r^h_{bb})^2+0.4})$ are given respectively by the two curves in Fig. 2 where we take $\tan\beta=50$.

The condition $r^h_{bb}=1$, or equivalently $\epsilon=0$, $t_{2\alpha}=t_{2\beta}$, defines the boundary that separates $\gamma\gamma$ enhancement and suppression in the parameter space:

$$r^h_{bb}=1\Leftrightarrow\epsilon=0\Leftrightarrow\Delta M^2_{12}=M^2_Z s_{2\beta}-\frac{\Delta M^2_{11}-\Delta M^2_{22}}{2}t_{2\beta}. \tag{23}$$

This condition is independent of m_A and the quantum correction Δ_b. $\Delta M^2_{12}>M^2_Z s_{2\beta}-\frac{\Delta M^2_{11}-\Delta M^2_{22}}{2}t_{2\beta}$ gives $b\bar{b}$ reduction. Flavor-tuning (FT) with small α requires a cancellation of $(M^2_{12})^{\rm tree}$ by the loop-level ΔM^2_{12} contribution, which requires rather large values of $\bar{\mu}$ and $\tan\beta$. This possibility was raised in Ref. [25].

The region of $\gamma\gamma$ enhancement does not overlap with the region $|\bar{\mu}|<0.5$ of natural SUSY for any value of $\tan\beta$ from 20 to 60. For $\tan\beta=20$, $|\bar{\mu}|\gtrsim 2$ is necessary for $\gamma\gamma$ enhancement.

We give a benchmark point of the FT model in the MSSM (FT1) and two benchmark points of natural SUSY (NFT1, NFT2).

	$\bar{\mu}$	$\tan\beta$	x_t	m_h(GeV)	σ_γ	σ_b	σ_τ
FT1	−3	20	−2.86	124	1.17	0.89	1.05
NFT1	−0.5	20	2.70	125	0.84	1.11	1.04
NFT2	−0.15	20	2.70	125	0.87	1.08	1.07

(24)

where $M_{\rm susy}=1$ TeV and $m_A=0.5$ TeV. The relevant sparticle masses are taken commonly with the values given above Eq. (5). The m_h value is predicted by Eq. (17). We also note that the predicted values of $B_s\to\mu^+\mu^-$ branching fraction of these benchmark points are consistent with the experimental measurement [26], $(3.2^{+1.4}_{-1.2\,\rm stat}{}^{+0.5}_{-0.3\,\rm syst})\times 10^{-9}$ within 2σ.

Natural SUSY predictions.—Natural SUSY always predicts $b\bar{b}$ enhancement and $\gamma\gamma$ reduction [27].

m_A	σ_γ	σ_b	σ_τ
$m_A\geq 500$ GeV	$0.82\sim 0.91$	$1.06\sim 1.12$	$1.04\sim 1.08$
$m_A\geq 1000$ GeV	$0.95\sim 0.98$	$1.01\sim 1.03$	$1.01\sim 1.02$

(25)

Here we have taken $|\mu|\leq 500$ GeV and the other parameters are fixed with the values given above Eq. (5).

Concluding remarks.—We have explored the $\gamma\gamma$, $b\bar{b}$ and $\tau\tau$ signals in the MSSM, relative to SM, and also in natural SUSY. In MSSM an enhancement in the diphoton signal of the 125 GeV Higgs boson relative to the SM Higgs can be obtained in a flavor-tuned model with $h=H^0_u$ provided that $|\mu|$ is large (TeV) and μ is negative. A $\gamma\gamma$ enhancement is principally due to the reduction of the $b\bar{b}$ decay width compared to $h_{\rm SM}$. The ratios of WW^* and ZZ^* to their SM values are predicted to be the same as that of $\gamma\gamma$. There is also a corresponding reduction of the h to $\tau\tau$ signal. The Tevatron evidence of a Higgs to $b\bar{b}$ signal in W + Higgs production [29] does not favor much $b\bar{b}$ reduction. The flavor-tuning of the neutral Higgs mixing angle α requires a large $\mu\sim$ TeV and large $\tan\beta$. For small $|\mu|<\sim$0.5 TeV of natural SUSY, $\gamma\gamma$ suppression relative to the SM is predicted. Thus, precision LHC measurements of the $\gamma\gamma$, WW^*, ZZ^* and $b\bar{b}$ signals of the 125 GeV Higgs boson can test MSSM models.

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[1] (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] M. Dine, A. Kagan, and S. Samuel, Phys. Lett. B **243**, 250 (1990); S. Kelley, J. Lopez, D. Nanopoulos, H. Pois, and K. Yuan, Nucl. Phys. B **398**, 3 (1993); V. Barger, M. Berger, and P. Ohmann, Phys. Rev. D **49**, 4908 (1994); S. Dimopoulos and G. Giudice, Phys. Lett. B **357**, 573 (1995); F. Gabbiani, E. Gabrielli, A. Masiero, and L. Silvestrini, Nucl. Phys. B **477**, 321 (1996); A. Cohen, D. Kaplan, and A. Nelson, Phys. Lett. B **388**, 588 (1996); K. Chan, U. Chattopadhyay, and P. Nath, Phys. Rev. D **58**, 096004 (1998); J. L. Feng, K. T. Matchev, and

T. Moroi, Phys. Rev. D **61**, 075005 (2000); Phys. Rev. Lett. **84**, 2322 (2000); R. Kitano and Y. Nomura, Phys. Lett. B **631**, 58 (2005); C Brust, A. Katz, S. Lawrence, and R. Sundrum, J. High Energy Phys. 03 (2012) 103; S. Akula, M. Liu, P. Nath, and G. Peim, Phys. Lett. B **709**, 192 (2012); R. Essig, E. Izaguirre, J. Kaplan, and J. G. Wacker, J. High Energy Phys. 01 (2012) 074; M. Papucci, J. T. Ruderman, and A. Weiler, J. High Energy Phys. 09 (2012) 035; L. Hall, D. Pinner, and J. Rudermn, J. High Energy Phys. 04 (2012) 131; S. King, M. Muhlleitner, and R. Nevzorov, Nucl. Phys. B **860**, 207 (2012); N. Arkani-Hamed, *WG2 meeting, Oct. 31, 2012*, CERN, Geneva.H. Baer, V. Barger, and P. Huang, J. High Energy Phys. 11 (2011) 031; H. Baer, V. Barger, P. Huang, and X. Tata, J. High Energy Phys. 05 (2012) 109; H. Baer, V. Barger, P. Huang, A. Mustafayev, and X. Tata, Phys. Rev. Lett. **109**, 161802 (2012); D. M. Ghilencea, H. M. Lee, and M. Park, J. High Energy Phys. 07 (2012) 046; J. Feng and D. Sanford, Phys. Rev. D **86**, 055015 (2012); L. Randall and M. Reece, arXiv:1206.6540.

[4] (ATLAS Collaboration), Report No. ATLAS-CONF-2012-094.

[5] M. Carena, S. Gori, N. R. Shah, C. E. M. Wagner, and L.-T. Wang, J. High Energy Phys. 07 (2012) 175.

[6] K. Hagiwara, J. S. Lee, and J. Nakamura, J. High Energy Phys. 10 (2012) 002.

[7] D. Carmi, A. Falkowski, E. Kuflik, and T. Volansky, J. High Energy Phys. 07 (2012) 136.

[8] V. Barger, H. E. Logan, and G. Shaughnessy, Phys. Rev. D **79**, 115018 (2009).

[9] D. Chung, L. Everett, G. Kane, S. King, J. Lykken, and L. Wang, Phys. Rep. **407**, 1 (2005).

[10] J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, *The Higgs Hunter's Guide* (Perseus Books, Boulder, Colorado, 1990).

[11] L. Hall, R. Rattazzi, and U. Sarid, Phys. Rev. D **50**, 7048 (1994); M. Carena, M. Olechowski, S. Pokorski, and C. E. M. Wagner, Nucl. Phys. B **426**, 269 (1994); D. Pierce, J. Bagger, K. Matchev, and R. Zhang, Nucl. Phys. B **491**, 3 (1997).

[12] A. Djouadi, Phys. Rep. **459**, 1 (2008).

[13] P. Langacker, J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).

[14] M. Carena, S. Gori, N. R. Shah, C. E. M. Wagner, and L.-T. Wang, J. High Energy Phys. 07 (2012) 175; M. Carena, S. Gori, N. R. Shah, and C. E. M. Wagner, J. High Energy Phys. 03 (2012) 014.

[15] H. Baer and X. Tata, *Weak Scale Supersymmetry: From Superfields to Scattering Events* (Cambridge University Press, Cambridge, England, 2006).

[16] V. Barger, P. Huang, M. Ishida, and W.-Y. Keung, arXiv:1206.1777 [Phys. Lett. B (to be published)].

[17] H. Baer, V. Barger, and A. Mustafayev, Phys. Rev. D **85**, 075010 (2012).

[18] V. Barger, M. Ishida, and W.-Y. Keung, Phys. Rev. Lett. **108**, 261801 (2012).

[19] V. D. Barger and R. J. N. Phillips, *Collider Physics* (Perseus Books, Boulder, Colorado, 1997).

[20] V. Barger, M. Ishida, and W.-Y. Keung, Phys. Rev. D **85**, 015024 (2012).

[21] J. Chang, K. Cheung, P.-Y. Tseng, and T.-C. Yuan, J. High Energy Phys. 12 (2012) 058.

[22] A. Denner, S. Heinemeyer, I. Puljak, D. Rebuzzi, and M. Spira, Eur. Phys. J. C **71**, 1753 (2011).

[23] A. Djouadi, J. Kalinowski, and M. Spira, Comput. Phys. Commun. **108**, 56 (1998).

[24] H. E. Haber and R. Hempfling, Phys. Rev. D **48**, 4280 (1993).

[25] M. Carena, S. Mrenna, and C. E. M. Wagner, Phys. Rev. D **62**, 055008 (2000).

[26] R. Aaij *et al.* (LHCb Collaboration), [arXiv:1211.2674].

[27] Here we note that as $\bar{\mu} = 0$, then $r^h_{bb} = 1 + 2M_Z^2/m_A^2 \rightarrow \sigma_{\gamma\gamma}/\sigma_{\rm SM} = 1 - 2.4M_Z^2/m_A^2$, independently of $\tan\beta$. The smaller m_A gives the larger suppression of σ_γ [28]. Then, from the $\gamma\gamma$ deviation from unity, the *CP*-odd state mass m_A could be estimated.

[28] K. Blum and R. T. D'Agnolo, Phys. Lett. B **714**, 66 (2012).

[29] The TEVNPH Working Group for the CDF and (D0 Collaboration), Report No. FERMILAB-CONF-12-065-E.