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Phys. Rev. D **86**, 115010 — Published 3 December 2012

DOI: 10.1103/PhysRevD.86.115010

BI-TP 2012/31
UAB-FT-714

Large diphoton Higgs rates from supersymmetric triplets

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Abstract

Recent results on Higgs searches at the LHC point towards the existence of a Higgs boson with mass of about 126 GeV whose diphoton decay rate tends to be larger than in the Standard Model. These results are in tension with natural MSSM scenarios: such a Higgs mass requires heavy (third-generation) squarks which reintroduce some amount of fine-tuning and in general the Higgs diphoton decay rate tends to follow the Standard Model result. In this paper we prove that these problems can be alleviated by introducing an extra supersymmetric triplet coupled to the Higgs in the superpotential. This superfield generates a sizeable tree-level correction to the Higgs mass so that the third generation is no longer required to be heavy, and its charged component enhances the diphoton Higgs decay rates by as much as 70% with respect to the Standard Model values. We also show that such a scenario is compatible with present electroweak precision observables.

1. Introduction The ATLAS and CMS collaborations at CERN have recently reported [1–6] excesses in several channels compatible with a Higgs with mass $m_h \simeq 126$ GeV. This value of the Higgs mass puts a strong tension on the minimal supersymmetric extension of the Standard Model (MSSM) as very heavy third generation squarks and large stop mixing are required in order to reproduce it [7]. Such mass values in the stop sector are in conflict with the MSSM as a natural solution to the hierarchy problem and create a *little hierarchy* problem.

In fact in the MSSM the couplings of the Higgs sector are a prediction of the model so that in the decoupling limit the SM-like Higgs mass turns out to be [8]

$$\begin{aligned} m_h^2 &= m_Z^2 \cos^2 2\beta \left(1 - \frac{3}{8\pi^2}\frac{m_t^2}{v^2}\, t\right) \\ &+ \frac{3}{4\pi^2}\frac{m_t^4}{v^2}\left[\frac{1}{2}X_t + t + \frac{1}{16\pi^2}\left(\frac{3}{2}\frac{m_t^2}{v^2} - 32\pi\alpha_3\right)\left(X_t t + t^2\right)\right] , \end{aligned} \tag{1}$$

with

$$X_t = \frac{2\tilde{A}_t^2}{m_Q^2}\left(1 - \frac{\tilde{A}_t^2}{12 m_Q^2}\right) , \tag{2}$$

$$t = \log\left(\frac{m_Q^2}{m_t^2}\right) , \qquad \tilde{A}_t = A_t - \mu/\tan\beta , \tag{3}$$

where m_t and m_Z are the top and Z masses, α_3 is the QCD coupling, m_Q is the (common) supersymmetry breaking mass of the third generation squarks [1] and μ is the holomorphic Higgsino mass. Moreover, we use the notation $v = \sqrt{v_1^2 + v_2^2} = 174\,\text{GeV}$ and $\tan\beta = v_2/v_1$ where v_1 (v_2) is the vacuum expectation value (VEV) of the Higgs H_1 (H_2) coupled to down (up) quarks. As it was pointed out in Ref. [8], Eq. (1) for values of $m_Q \lesssim 1.5$ TeV provides a good approximation (within 2 GeV error) to more sophisticated numerical results. Actually one easily obtains from Eq. (1) that quite heavy third generation squarks with $m_Q = \mathcal{O}(1\,\text{TeV})$, large $\tan\beta$ and a sizeable mixing X_t are needed [7] in order to reproduce a SM-like Higgs mass of about 126 GeV. These heavy squarks however induce large radiative contributions to the electroweak breaking mechanism and then tend to rise the electroweak scale far away from the observed one (say the Z mass) unless a fine-tuning of around one per mille is done.

In short, by taking naturalness as a guiding criterion for physics beyond the Standard Model (SM), there is a tension in the MSSM between the actual values of the Higgs and Z boson masses. To alleviate this tension a simple supersymmetric possibility (without enlarging the SM gauge group) is to extend the MSSM with some extra multiplets which trigger extra tree level contributions to the Higgs quartic coupling, in such a way that one

[1] For simplicity we will consider in this paper degenerate supersymmetry breaking masses m_Q for up and down-type third generation squarks. Notice that perturbative problems can spoil the approximations used in Eq. (1) when m_Q is very large in which case resumming logarithms is required.

could reproduce the experimental value of the Higgs mass without the need of heavy third-generation squarks. Several proposals have been made in the literature [9] [2]. As the extra multiplet has to be coupled to the Higgs sector in the superpotential by renormalizable couplings, the number of possible extra multiplets is reduced: either $SU(2)_L$ singlets or triplets with hypercharge $Y = 0$ or ± 1 can play this role.

On the other hand, a further source of tension comes from the LHC measurements of the Higgs decay rates. From the experimental results on the Higgs decay into ZZ and WW channels one may infer that in this sector no dominant contributions beyond the SM ones do exist in Nature, as well as in the Higgs production through either gluon or weak vector boson fusion. However new physics beyond the SM will arise from the diphoton Higgs decay rate if LHC keeps on showing a significant excess with respect to the SM prediction when more LHC data will be collected. Assuming SM-like Higgs production, the ratio between the diphoton rate observed at LHC and the one expected in the SM is $R_{\gamma\gamma} = 1.8 \pm 0.5$ for ATLAS ($m_h = 126$ GeV) [5] and $R_{\gamma\gamma} = 1.6 \pm 0.4$ for CMS ($m_h = 125$ GeV) [6]. The central value of the combination of these measurements, which roughly corresponds to an enhancement of ~ 1.7 with respect to the SM prediction (with an error around ± 0.3) is hard to reproduce in the MSSM [7] and, if this central value were confirmed with more statistics, it would represent a strong tension between the MSSM and LHC data. A supersymmetric extension of the SM solving this problem should then extend the MSSM by some electrically charged (extra) states coupled to the Higgs that should contribute to $R_{\gamma\gamma}$ at one-loop. Notice that these extra states, possibly charged under both $SU(2)_L$ and $U(1)_Y$, would generate a subleading (one-loop) correction to the SM (tree-level) weak vector fusion production. Moreover they should be colorless in order to not modify at leading order the gluon-fusion Higgs production arising at one-loop in the SM.

In the present paper we analyze minimal MSSM extensions where extra states can relax the little hierarchy problem in the presence of a 126 GeV Higgs mass as well as reproduce the diphoton excess in the Higgs production rate.

2. The model As the singlet is electrically neutral only triplets, which contain charged states and can thus contribute to the $h \to \gamma\gamma$ decay width, are natural candidates to the MSSM extension. In particular we consider the effect of a supersymmetric $Y = 0$ triplet [3]

$$\Sigma = \begin{pmatrix} \xi^0/\sqrt{2} & -\xi_2^+ \\ \xi_1^- & -\xi^0/\sqrt{2} \end{pmatrix} \tag{4}$$

on both issues: the Higgs mass generation and the diphoton rate.

The most general renormalizable coupling of the triplet Σ to the Higgs sector is provided by the superpotential

$$\Delta W = \lambda H_1 \cdot \Sigma H_2 + \frac{1}{2}\mu_\Sigma \operatorname{tr} \Sigma^2 \tag{5}$$

[2]In this paper we focus on low-energy extensions of the MSSM. For non-minimal ultraviolet completions with the extra content at the multi-TeV scale, see for instance [10].

[3]Considering supersymmetric triplets with $Y = \pm 1$ should lead to similar results as those founds in the present paper.

that gets added to the MSSM one. Notice that $\mathrm{tr}\,\Sigma^3 \equiv 0$ due to its own structure. The new interaction modifies the Higgs potential and, in the decoupling limit, the tree-level mass of the SM-like Higgs is given by

$$m^2_{h,\mathrm{tree}} = m_Z^2 \cos^2 2\beta + \frac{\lambda^2}{2} v^2 \sin^2 2\beta \ . \tag{6}$$

We see that for moderate values of λ the tree-level mass can be lifted so that no large contributions from loop corrections are required to reproduce $m_h \simeq 126\,\mathrm{GeV}$. In particular, stops can be light, thus reducing the fine tuning for the electroweak scale. Notice also that the new contribution is relevant mostly when $\tan\beta \simeq 1$ while it vanishes when $\tan\beta \to \infty$. In this way it will be possible to cope with the LHC Higgs mass for moderate values of $\tan\beta$ without large stop mixing, contrarily to what happens in the MSSM.

One strong constraint on models with triplets comes from the electroweak precision tests (EWPT), in particular from the ρ-parameter constraint [11]. If the scalar component of the triplet acquires a VEV $\langle \xi^0 \rangle$ it will give a tree-level contribution to the ρ parameter as $\rho = 1 + 2\langle \xi^0 \rangle^2/v^2$ [11] which will easily be in conflict with experimental data. On the other hand a triplet VEV will always exist in a theory with a superpotential given by Eq. (5) after electroweak breaking. Moreover once supersymmetry is broken, one expects the soft-breaking term $\lambda A_\lambda H_1 \cdot \Sigma H_2$ to appear in the scalar potential, depending on the particular mechanism of supersymmetry breaking, where now all symbols denote just the scalar components of the chiral superfields. So once the Higgs gets a VEV it will generate a tadpole for the neutral component ξ^0 of Σ and the previous terms will induce the VEV

$$\langle \xi^0 \rangle \simeq \sqrt{2}\lambda \left[\mu + \frac{1}{2}\left(\mu_\Sigma + \frac{A_\lambda}{2}\right) \sin 2\beta \right] \frac{v^2}{m_\Sigma^2 + \mu_\Sigma^2} \ , \tag{7}$$

where m_Σ is the supersymmetry breaking mass for the triplet and A_λ is the trilinear supersymmetry breaking parameter associated to the superpotential (5). The present bound on the ρ parameter, $\rho = 1.0004^{+0.0003}_{-0.0004}$ [11], imposes the constraint $\langle \xi^0 \rangle \lesssim 4$ GeV at 95% CL. In order to be consistent with the experimental value of the ρ parameter we are going to suppose that there is a hierarchy between the trilinear coupling A_λ and supersymmetric masses in the superpotential (5), and the soft mass for the Σ-scalar [4], i.e. A_λ, μ, $\mu_\Sigma \ll m_\Sigma$. Moreover to cope with the experimental value of the ρ parameter we will consider that the scalar component is in the TeV range and thus its presence is negligible in this discussion. We will then hereafter neglect $\langle \xi^0 \rangle$ and put it to zero.

In the fermion sector $\tilde{\xi}^0$ mixes with the MSSM neutralinos $(\widetilde{W}^3, \widetilde{W}^0, \widetilde{H}_1^0, \widetilde{H}_2^0)$ while $\tilde{\xi}_1^-$ and $\tilde{\xi}_2^+$ mix with the MSSM charginos $(\widetilde{W}^-, \widetilde{W}^+, \widetilde{H}_1^-, \widetilde{H}_2^+)$. As the relevant states for the ratio $R_{\gamma\gamma}$ are the charged ones we will concentrate on charginos. Their mass matrix is given by

$$\left(\widetilde{W}^-,\ \widetilde{H}_1^-,\ \tilde{\xi}_1^-\right) \mathcal{M}_{ch} \begin{pmatrix} \widetilde{W}^+ \\ \widetilde{H}_2^+ \\ \tilde{\xi}_2^+ \end{pmatrix}, \quad \mathcal{M}_{ch} = \begin{pmatrix} M_2 & gv\sin\beta & 0 \\ gv\cos\beta & \mu & \lambda v \sin\beta \\ 0 & \lambda v\cos\beta & \mu_\Sigma \end{pmatrix} \ , \tag{8}$$

[4]Notice that the required hierarchy $A_\lambda \ll m_\Sigma$ naturally arises if the supersymmetry breaking mechanism is driven by gauge interactions where both m_Σ^2 and A_λ are obtained at two-loop.

where we have put for simplicity $\langle \xi^0 \rangle = 0$ while for definiteness the parameters M_2, μ and μ_Σ will be assumed to be positive in the rest of the paper. In the limit where $M_2 \gg m_W$, the three charged eigenstates $(\widetilde{\chi}_1^+, \widetilde{\chi}_2^+, \widetilde{\chi}_3^+)$ have respectively masses (m_-, m_+, M_2) where

$$\begin{aligned} m_\pm &= \mu_+ \pm \sqrt{v_1 v_2 \lambda^2 + \mu_-^2}\ , \\ \mu_\pm &= \frac{\mu \pm \mu_\Sigma}{2}\ . \end{aligned} \tag{9}$$

Therefore the parameter $m \equiv m_-$ is a good estimate for the mass of the lightest chargino $m_{\widetilde{\chi}_1^\pm}$, which we will impose to be heavier than 94 GeV [11], and the whole chargino sector can be expressed as function of $m, \lambda, \tan\beta, \mu_-$ and M_2. In all cases we will choose the value of the mass parameter m such that, in the corresponding region of the parameter space, the experimental bounds on the chargino masses are fulfilled.

3. The ratio $h \to \gamma\gamma$ In the limit where $m_h^2 \ll 4m^2_{\widetilde{\chi}_i^+}$ and by taking into account the main contributions due to charginos, W boson and top quark t, the diphoton Higgs decay rate with respect to the SM value $R_{\gamma\gamma}$ turns out to be [12–14]

$$R_{\gamma\gamma} = \left| 1 + \frac{\frac{4}{3}\frac{\partial}{\partial \log v} \log\det \mathcal{M}_{ch}(v)}{A_1(\tau_W) + \frac{4}{3} A_{1/2}(\tau_t)} \right|^2 , \tag{10}$$

where $\tau_i = m_h^2/4m_i^2$. The functions $A_1(\tau)$ and $A_{1/2}(\tau)$ are given in Ref. [15] and their numerical values for the W and top fields are $A_1(\tau_W) \simeq -8.3$ and $A_{1/2}(\tau_t) \simeq 1.4$ so that the denominator in Eq. (10) is negative. The numerator in Eq. (10) is given by

$$\frac{\partial}{\partial \log v} \log\det \mathcal{M}_{ch}(v) = -\frac{\sin 2\beta v^2(\lambda^2 M_2 + g^2 \mu_\Sigma)}{M_2 \mu \mu_\Sigma - \frac{1}{2}\sin 2\beta \lambda^2 v^2 (\lambda^2 M_2 + g^2 \mu_\Sigma)} , \tag{11}$$

and its sign depends on the specific values that one can choose for the parameters [5]. In the present paper we are interested in cases where the r.h.s. of Eq. (11) is positive, which leads to an enhancement of the diphoton Higgs decay rate as suggested by present LHC data. This will be studied by imposing $m_h = 126$ GeV. In particular as we do not need very heavy third-generation squarks to achieve such a Higgs mass, in the following illustrative examples we will fix $m_Q \simeq 700$, $X_t = 0$ (i.e. $m_{\tilde{t}_1} = m_{\tilde{t}_2} = 700$ GeV) and $X_t = 4$ (i.e. $m_{\tilde{t}_1} = 545$ GeV and $m_{\tilde{t}_2} = 828$ GeV).

In Fig. 1 we present the contour lines of $m_h = 126$ GeV for $X_t = 0$ [thick solid (red) curves] and $X_t = 4$ [thin solid (red) curves] and the contour lines of $R_{\gamma\gamma}$ (dashed black curves) in the plane $(\lambda, \tan\beta)$ where we fix $\mu = \mu_\Sigma$. In the left panel, where we take $M_2 = 750$ GeV and $m = 100$ GeV, the chargino masses vary as $m_{\widetilde{\chi}_1^+} \in [95, 110]$ GeV, $m_{\widetilde{\chi}_2^+} \in [217, 302]$ GeV and $m_{\widetilde{\chi}_3^+} \in [754, 762]$ GeV. As we can see the considered model cannot reach $R_{\gamma\gamma} \simeq 1.45$ for mixing $X_t = 4$ but it can for $X_t = 0$, $\lambda \simeq 0.85$ and $\tan\beta \simeq 1$. Of course, larger

[5] Notice that forbidding massless charginos imply finiteness of the numerator in Eq. (10).

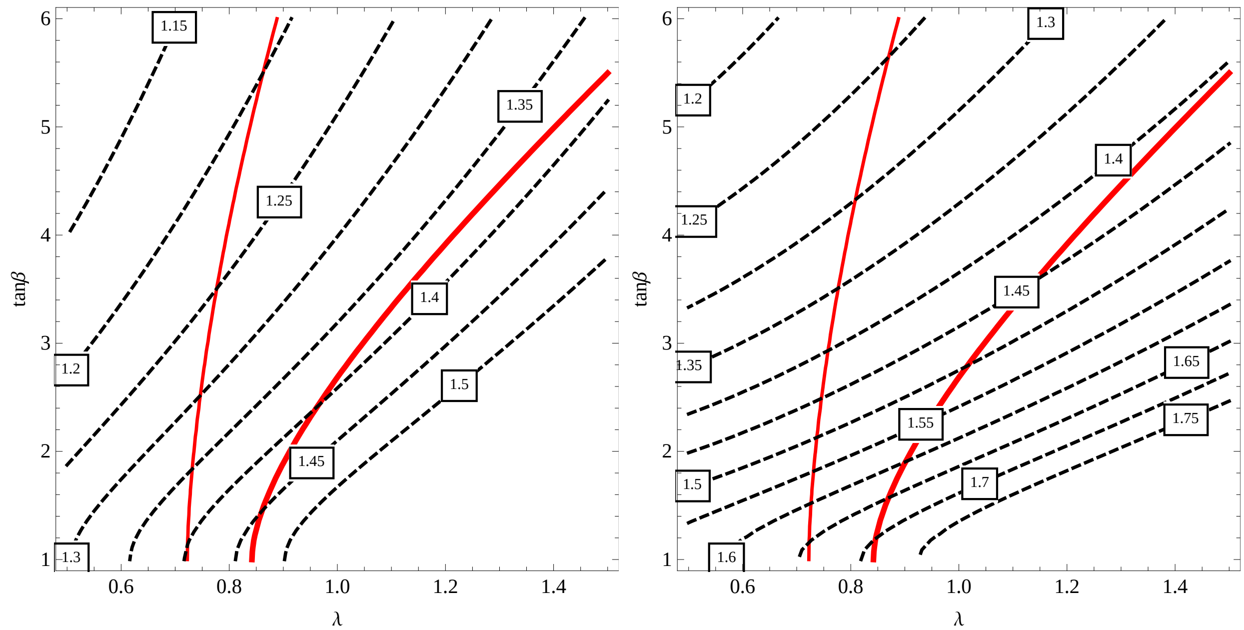

Figure 1: *Left panel: Contour plots of $R_{\gamma\gamma}$ (dashed lines) for $M_2 = 750$ GeV, $m = 100$ GeV and $\mu_- = 0$. Contour plots of $m_h = 126$ GeV for $X_t = 0$ [thick solid (red) line] and for $X_t = 4$ [thin solid (red) line]. Right panel: The same as in the left panel but for $M_2 = 250$ GeV and $m = 117$ GeV.*

values of $R_{\gamma\gamma}$ can be achieved by decreasing M_2 a possibility shown in the right panel of Fig. 1 where $M_2 = 250\,\text{GeV}$ and $m = 117\,\text{GeV}$ are fixed. In such a case, the model can roughly reproduce the combined ATLAS and CMS central value $R_{\gamma\gamma} \simeq 1.7$ for $X_t = 0$ in correspondence with $\tan\beta \simeq 1$ and $\lambda \simeq 0.85$. In the region $(\lambda, \tan\beta)$ shown in the right panel of Fig. 1 the masses of the charginos vary as $m_{\widetilde{\chi}_1^+} \in [95, 110]\,\text{GeV}$, $m_{\widetilde{\chi}_2^+} \in [215, 240]\,\text{GeV}$ and $m_{\widetilde{\chi}_3^+} \in [280, 360]\,\text{GeV}$.

Up to now we have presented results on $R_{\gamma\gamma}$ for $\mu_- = 0$. The variation with μ_- for $\tan\beta = 1$ is shown in Fig. 2 where we vary μ_- for $M_2 = 750$ GeV, $m = 100$ GeV (left panel) and $M_2 = 250$ GeV, $m = 117$ GeV (right panel). We can see that increasing μ_- (starting from $\mu_- = 0$) does not enhance $R_{\gamma\gamma}$ for a given Higgs mass curve (thin red line for $X_t = 4$ and thick red line for $X_t = 0$). Instead, for $\mu_- < 0$ one might naively extrapolate from the figure that negative values of μ_- lead to better results, but for the shown choices of parameters such a possibility would actually yield too small lightest chargino masses.

4. *Electroweak observables* An important question is how the zero-hypercharge supersymmetric triplet modifies the electroweak observables, especially the T parameter [16] which is particularly sensitive to the presence of the triplet. Indeed, as already mentioned above, the triplet contributes to the T parameter at tree-level as the previously considered ρ-parameter and T are related by $\rho - 1 = \alpha T$, where α is the electromagnetic constant at the m_Z scale. As we already mentioned electroweak breaking produces at tree-level a tadpole in its neutral

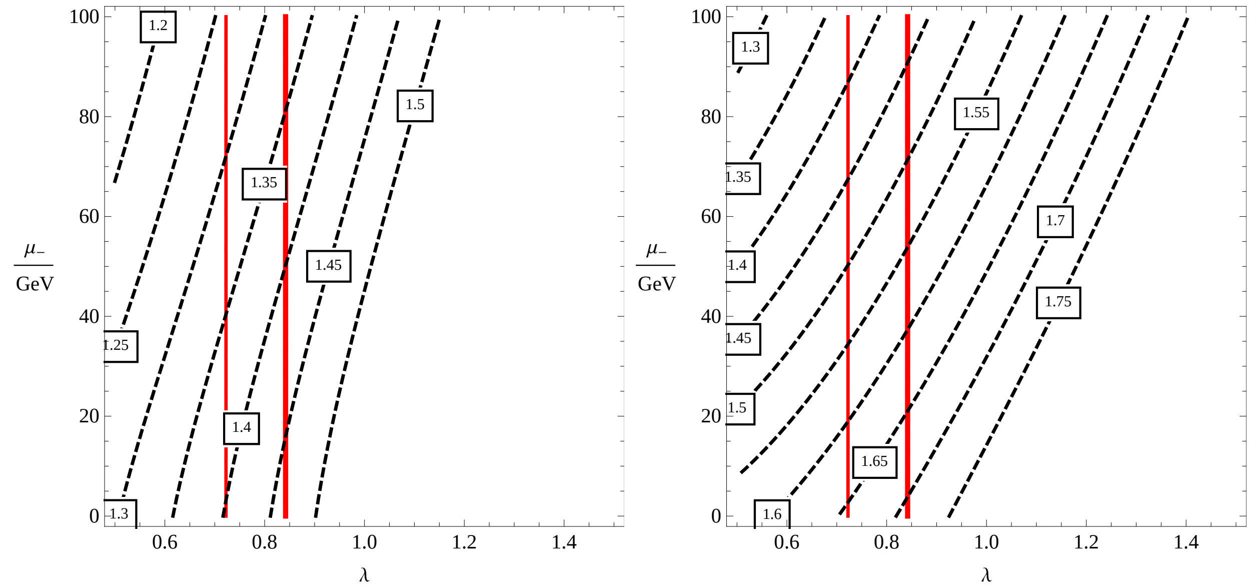

Figure 2: *Left panel: Contour plots of $R_{\gamma\gamma}$ (dashed lines) for $M_2 = 750$ GeV, $m = 100$ GeV and $\tan\beta = 1$ in the plane $(\lambda, \mu_-/GeV)$ and $m_h = 126$ GeV for $X_t = 0$ [thick solid (red) line] and $X_t = 4$ [thin solid (red) line]. Right panel: Same as in the left panel but for $M_2 = 250$ GeV and $m = 117$ GeV.*

component ξ^0 and the experimental constraint on this contribution then requires $\langle \xi^0 \rangle \lesssim 4$ at 95% CL.

Moreover at one-loop the supersymmetric triplet contributes to the electroweak observables through its coupling to the Higgs sector and, for $\mu = \mu_\Sigma$, the oblique S and T parameters [16] get modified compared to the MSSM as

$$
\begin{aligned}
\alpha S &= \frac{s_W^2 \lambda^2}{10\pi^2} \frac{m_W^2}{\mu^2} \left[1 + \frac{19}{24} \sin 2\beta \right] + \mathcal{O}(g^4) \ , \\
\alpha T &= \frac{3\lambda^2}{128\pi^2} \frac{m_W^2}{\mu^2} \cos^2 2\beta + \mathcal{O}(g^4) \ , \qquad (12)
\end{aligned}
$$

where the $\mathcal{O}(g^4)$ correction is coming from the (MSSM) Wino-Higgsino mixing. These corrections are based on an expansion at $\mathcal{O}(s_W^2)$ [17] and turn out to be small in the considered region as compared with the experimental values [11]

$$
S = 0.04 \pm 0.09, \quad T = 0.07 \pm 0.08 \qquad \text{(88\% correlation)} \ . \qquad (13)
$$

We plot in Fig. 3 contour lines for the predicted values of S (left panel) and T (right panel) for $m = 117\,\mathrm{GeV}$, $M_2 = 250$ GeV and $\mu = \mu_\Sigma$. We have included on top of the triplet contributions from Eq. (12) the $\mathcal{O}(g^4)$ contribution coming from the MSSM gaugino-Higgsino mixing. The predicted values of the S and T parameters are in all cases in agreement

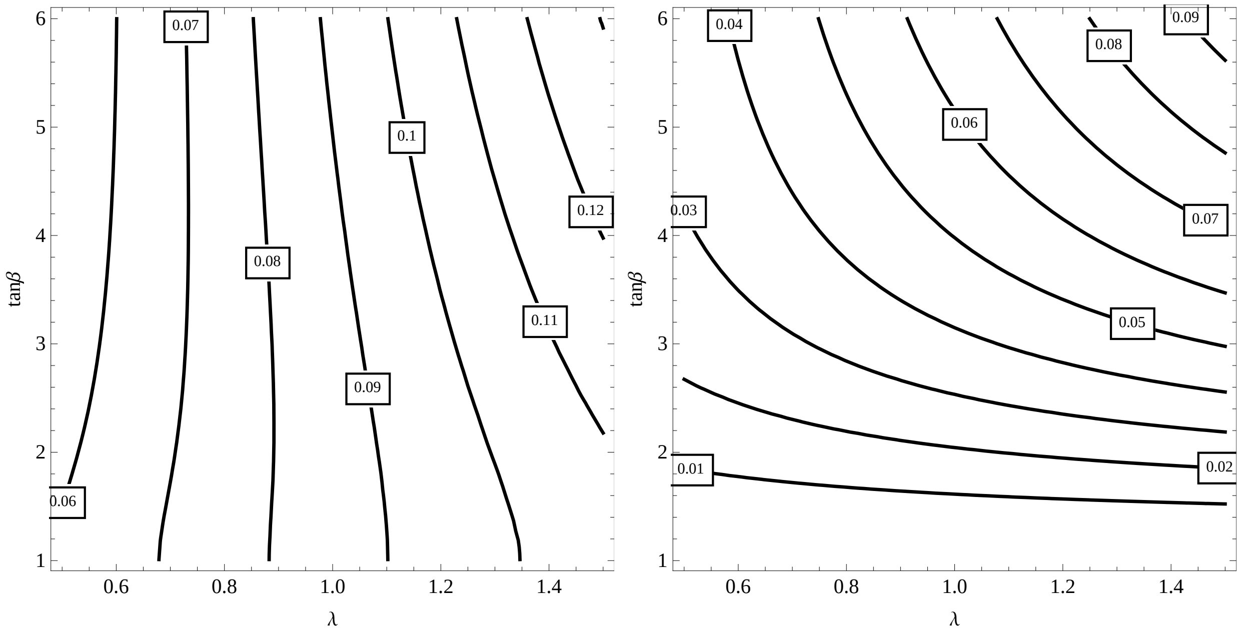

Figure 3: *Contour plots of the S (left panel) and T (right panel) parameters in the plane* $(\lambda, \tan\beta)$ *for $m = 117$ GeV, $M_2 = 250$ GeV and $\mu = \mu_\Sigma$.*

with the experimental values within 1σ. Moreover, in the region where $m_h = 126\,\mathrm{GeV}$ (for $m_Q = 700\,\mathrm{GeV}$) and the diphoton rate is maximally enhanced (i.e. $\tan\beta \simeq 1$ and $\lambda \simeq 0.85$) we obtain $S \simeq 0.08$ and $T \simeq 0$. The technical reasons why triplet corrections to the electroweak observables are generally small in the considered region are:

- For $\tan\beta = 1$ (and zero triplet VEV) the custodial symmetry is unbroken by the triplet and thus $T = 0$.
- Apart from the loop factor there is the extra suppression m_W^2/μ^2 and it turns out that this ratio is small. For instance with $m = 117$ GeV, $\lambda = 0.85$ and $\tan\beta = 1$ we get $\mu = 221$ GeV for $\tan\beta = 1$ and $\mu = 177$ GeV for $\tan\beta = 6$.

5. Perturbativity A final issue which has to be considered is perturbativity of coupling constants. In fact the evolution with the scale of the couplings λ and h_t are given by the renormalization group equations (RGE) [9]

$$\begin{aligned} 8\pi^2\dot{\lambda} &= \left(-\frac{7}{2}g^2 - \frac{1}{2}g'^2 + 2\lambda^2 + \frac{3}{2}h_t^2\right)\lambda\ , \\ 8\pi^2\dot{h}_t &= \left(-\frac{3}{2}g^2 - \frac{13}{18}g'^2 - \frac{8}{3}g_3^2 + \frac{3}{4}\lambda^2 + 3h_t^2\right)h_t\ , \\ 16\pi^2\dot{g} &= 3g^2\ , \quad 16\pi^2\dot{g}' = 11g'^2\ , \quad 16\pi^2\dot{g}_3 = -3g_3^3\ , \end{aligned} \tag{14}$$

where the dots stand for d/dt and $t = \log(\mathcal{Q}/\mathrm{GeV})$. We can see from the first equality in Eq. (14) that for large enough initial values of $\lambda \equiv \lambda(m_t)$, the running coupling $\lambda(\mathcal{Q})$

is driven to larger values at high scales and eventually it reaches non-perturbative values ($\lambda(\mathcal{Q}) \simeq 4\pi$) in the ultraviolet (UV) at some scale $\mathcal{Q} \simeq \Lambda$, the cutoff of the theory, near its Landau pole. This means that the theory becomes non-perturbative, unless it is UV completed at some scale smaller than Λ. It is thus clear that the MSSM with an extra zero-hypercharge triplet does not unify perturbatively [6]. However if the theory is still valid near its Landau pole all couplings will feel the singularity through the higher loop RGE and they can unify through the so-called non-perturbative unification [18] [7]. With respect to conventional unification, non-perturbative unification has the attractive feature that low energy couplings are less sensitive to high energy physics. Fig. 4 (left panel) shows the value of the cutoff Λ (in GeV) in the plane $(\lambda, \tan\beta)$ for the MSSM model completed with the supersymmetric $Y = 0$ triplet Σ. We can see that when $\lambda = 0.8$ the cutoff Λ ranges from 10^8

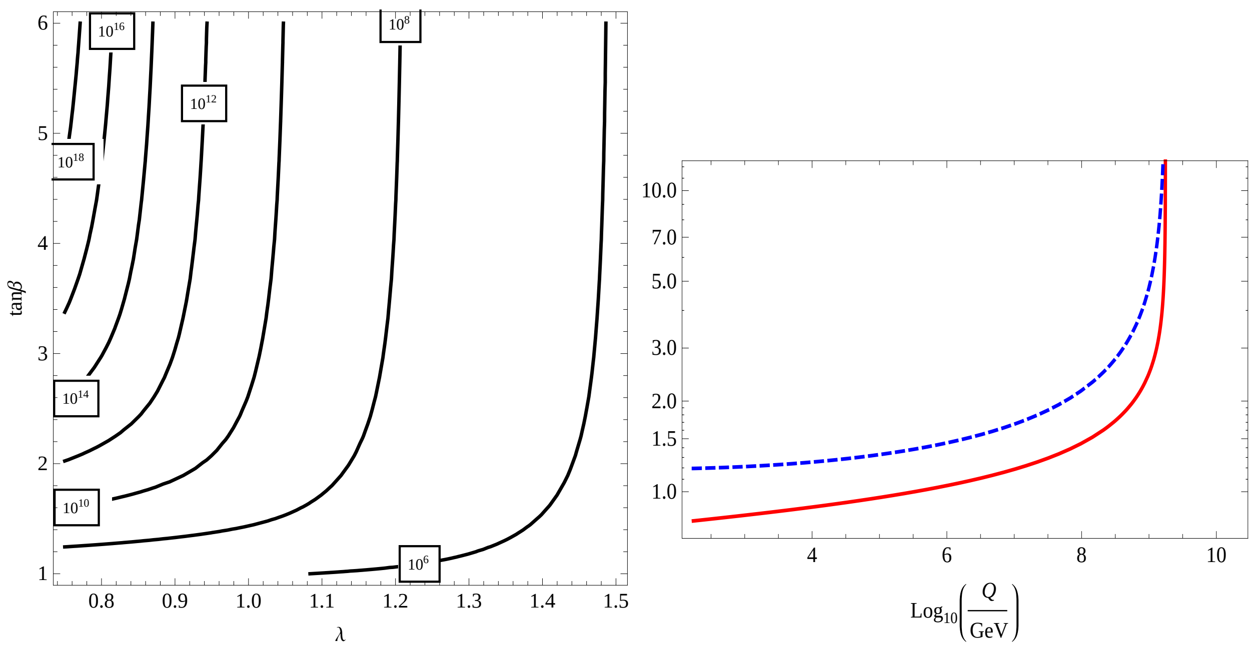

Figure 4: *Left panel: Contour lines for constant values of the cutoff Λ (in GeV) in the plane $(\lambda, \tan\beta)$. Right panel: Plot of $\lambda(t)$ [solid (red) line] and $h_t(t)$ [dashed (blue) line] for* $\tan\beta = 1.5$ *and* $\lambda = 0.8$.

GeV to 10^{16} GeV for $\tan\beta \simeq 1$ to $\tan\beta \simeq 6$, respectively. In particular, in the parameter region considered in the right panels of Figs. 1 and 2, for $R_{\gamma\gamma} \simeq 1.7$ and $m_h = 126\,\mathrm{GeV}$ it

[6]Indeed, even in the regime where λ is small so that the theory remains perturbative until the Planck scale, unification cannot be achieved because the extra triplet modifies the MSSM beta functions of the gauge coupling in an incomplete way.

[7]See Ref. [9] for examples of non-perturbative unification applied to the extensions of the MSSM with additional matter.

turns out that $\Lambda \simeq 2\times 10^9$ GeV. This can be also seen in the right panel of Fig. 4 where the evolution of the couplings $\lambda(t)$ and $h_t(t)$ is presented.

6. Conclusion In conclusion if one interprets the excess discovered at LHC as a Higgs boson with a mass about 126 GeV, a little hierarchy problem emerges in the MSSM. In addition a further tension between the MSSM and experimental results would arise if the actual tendency of ATLAS and CMS data, on Higgs production and decay, towards deviations from the SM results only in the Higgs diphoton decay channel, would be confirmed with better precision in the present (and forthcoming) LHC run. In this paper we have proven that both problems can be naturally overcome by minimally extending the MSSM by a colorless zero-hypercharge $SU(2)_L$-triplet with a coupling to the Higgs superfields of order one, in fact $\lesssim h_t$, the top quark Yukawa coupling. In such a case, even for small $\tan\beta$ and no mixing between the third-generation squarks, it is possible to reach the experimental value of the Higgs mass $m_h = 126$ GeV and at the same time $\sim 70\%$ enhancement in the Higgs diphoton decay rate. In the considered parameter range the theory is consistent with electroweak precision tests and hints towards a non-perturbative unification at the scale $\sim 10^9$ GeV. Finally the experimental signals of this triplet could come from pair production of neutralinos (charginos) which then decay mainly into Higgs plus missing energy.

Acknowledgments

AD was partly supported by the National Science Foundation under grants PHY-0905383-ARRA and PHY-1215979. MQ was supported by the Spanish Consolider-Ingenio 2010 Programme CPAN (CSD2007-00042) and by CICYT-FEDER-FPA2008-01430 and FPA2011-25948.

References

[1] F. Gianotti, CERN Seminar, "Update on the Standard Model Higgs searches in ATLAS", July, 4 2012. ATLAS-CONF-2012-093

[2] J. Incandela, CERN Seminar, "Update on the Standard Model Higgs searches in CMS", July, 4 2012.

[3] G. Aad *et al.* [ATLAS Collaboration], "Search for the Standard Model Higgs boson in the diphoton decay channel with 4.9 fb-1 of pp collisions at sqrt(s)=7 TeV with ATLAS," Phys. Rev. Lett. **108**, 111803 (2012) [arXiv:1202.1414 [hep-ex]].

[4] S. Chatrchyan *et al.* [CMS Collaboration], "Search for the standard model Higgs boson decaying into two photons in pp collisions at sqrt(s)=7 TeV," Phys. Lett. B **710**, 403 (2012) [arXiv:1202.1487 [hep-ex]].

[5] G. Aad *et al.* [ATLAS Collaboration], "Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC," Phys. Lett. B [arXiv:1207.7214 [hep-ex]].

[6] S. Chatrchyan *et al.* [CMS Collaboration], "Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC," Phys. Lett. B [arXiv:1207.7235 [hep-ex]].

[7] M. Carena, S. Gori, N. R. Shah and C. E. M. Wagner, "A 125 GeV SM-like Higgs in the MSSM and the $\gamma\gamma$ rate," JHEP **1203**, 014 (2012) [arXiv:1112.3336 [hep-ph]].

[8] M. S. Carena, J. R. Espinosa, M. Quiros and C. E. M. Wagner, "Analytical expressions for radiatively corrected Higgs masses and couplings in the MSSM," Phys. Lett. B **355**, 209 (1995) [hep-ph/9504316]; M. S. Carena, M. Quiros and C. E. M. Wagner, "Effective potential methods and the Higgs mass spectrum in the MSSM," Nucl. Phys. B **461**, 407 (1996) [hep-ph/9508343].

[9] J. R. Espinosa and M. Quiros, "On Higgs boson masses in nonminimal supersymmetric standard models," Phys. Lett. B **279**, 92 (1992); J. R. Espinosa and M. Quiros, "Upper bounds on the lightest Higgs boson mass in general supersymmetric Standard Models," Phys. Lett. B **302**, 51 (1993) [hep-ph/9212305]; G. L. Kane, C. F. Kolda and J. D. Wells, "Calculable upper limit on the mass of the lightest Higgs boson in any perturbatively valid supersymmetric theory," Phys. Rev. Lett. **70**, 2686 (1993) [hep-ph/9210242]; J. R. Espinosa and M. Quiros, "Gauge unification and the supersymmetric light Higgs mass," Phys. Rev. Lett. **81**, 516 (1998) [hep-ph/9804235].

[10] I. Antoniadis, E. Dudas, D. M. Ghilencea and P. Tziveloglou, "MSSM with Dimension-five Operators (MSSM(5))," Nucl. Phys. B **808** (2009) 155 [arXiv:0806.3778 [hep-ph]]; M. Carena, K. Kong, E. Ponton and J. Zurita, "Supersymmetric Higgs Bosons and Beyond," Phys. Rev. D **81**, 015001 (2010) [arXiv:0909.5434 [hep-ph]]; I. Antoniadis, E. Dudas, D. M. Ghilencea and P. Tziveloglou, "MSSM Higgs with dimension-six operators," Nucl. Phys. B **831**, 133 (2010) [arXiv:0910.1100 [hep-ph]]; A. Delgado, G. Nardini and M. Quiros, "The Light Stop Scenario from Gauge Mediation," JHEP **1204**, 137 (2012) [arXiv:1201.5164 [hep-ph]]; P. Tziveloglou, "Aspects of Effective Supersymmetric Theories," arXiv:1201.6040 [hep-th]; K. Blum, R. T. D'Agnolo and J. Fan, "Natural SUSY Predicts: Higgs Couplings," arXiv:1206.5303 [hep-ph].

[11] J. Beringer *et al.* [Particle Data Group Collaboration], "Review of particle physics," Phys. Rev. D **86**, 0100001 (2012).

[12] J. R. Ellis, M. K. Gaillard and D. V. Nanopoulos, "A Phenomenological Profile of the Higgs Boson," Nucl. Phys. B **106** (1976) 292.

[13] M. A. Shifman, A. I. Vainshtein, M. B. Voloshin and V. I. Zakharov, "Low-Energy Theorems for Higgs Boson Couplings to Photons," Sov. J. Nucl. Phys. **30** (1979) 711 [Yad. Fiz. **30** (1979) 1368].

[14] M. Carena, I. Low and C. E. M. Wagner, "Implications of a Modified Higgs to Diphoton Decay Width," arXiv:1206.1082 [hep-ph].

[15] A. Djouadi, "The Anatomy of electro-weak symmetry breaking. II. The Higgs bosons in the minimal supersymmetric model," Phys. Rept. **459** (2008) 1 [hep-ph/0503173].

[16] M. E. Peskin and T. Takeuchi, "Estimation of oblique electroweak corrections," Phys. Rev. D **46**, 381 (1992).

[17] G. Marandella, C. Schappacher and A. Strumia, "Supersymmetry and precision data after LEP2," Nucl. Phys. B **715** (2005) 173 [hep-ph/0502095].

[18] L. Maiani, G. Parisi and R. Petronzio, "Bounds on the Number and Masses of Quarks and Leptons," Nucl. Phys. B **136** (1978) 115; L. Maiani and R. Petronzio, "Low-energy Gauge Couplings And The Mass Gap Of N=1 Supersymmetry," Phys. Lett. B **176** (1986) 120 [Erratum-ibid. **178B** (1986) 457]; N. Cabibbo and G. R. Farrar, "An Alternative To Perturbative Grand Unification: How Asymptotically Nonfree Theories Can Successfully Predict Low-energy Gauge Couplings," Phys. Lett. B **110** (1982) 107.