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Light staus and enhanced Higgs diphoton rate with non-universal gaugino masses and SO(10) Yukawa unification

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Abstract: It is shown that substantially enhanced Higgs to diphoton rate induced by light staus with large left-right mixing in MSSM requires at the GUT scale non-universal gaugino masses with bino and/or wino lighter than gluino. The possibility of such enhancement is investigated in MSSM models with arbitrary gaugino masses at the GUT scale with additional restriction of top-bottom-tau Yukawa unification, as predicted by minimal SO(10) GUTs. Many patterns of gaugino masses leading to enhanced Higgs to diphoton rate and the Yukawa unification are identified. Some of these patterns can be accommodated in a well-motivated scenarios such as mirage mediation or SUSY breaking F-terms being a non-singlet of SO(10). Phenomenological implications of a scenario with non-universal gaugino masses generated by a mixture of the singlet F-term and the F-term in a 24-dimensional representation of $SU(5) \subset SO(10)$ are studied in detail. Possible non-universalities of other soft terms generated by such F-terms are discussed. The enhancement of Higgs to diphoton rate up to 30% can be obtained in agreement with all phenomenological constraints, including vacuum metastability bounds. The lightest sbottom and pseudoscalar Higgs are within easy reach of the 14 TeV LHC. The LSP can be either bino-like or wino-like. The thermal relic abundance in the former case may be in agreement with the cosmological data thanks to efficient stau coannihilation.

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Contents

1 Introduction

Discovery of a Higgs boson with mass around 125 GeV is now firmly established [1, 2]. On the other hand, it remains unclear whether the discovered boson is the Standard Model (SM) Higgs boson. Even though the measured values of the Higgs signal, μ_i, in most decay channels are within 1σ from the SM prediction, the errors are still rather large, of about 20–30%, even in the best-measured channels such as $\gamma\gamma$, WW^* and ZZ^* [3–7]. Moreover, there are some anomalies in the LHC data. Particularly interesting is an excess of events in the $\gamma\gamma$ channel observed by ATLAS. The fitted number of signal events in this channel is 2σ above the SM prediction [3]. Similar excess was observed before also in the CMS data [8] but it disappeared after analysing all data collected in the 7 TeV and 8 TeV LHC runs [5].

In general, there are many ways to enhance the Higgs signal strength in the $\gamma\gamma$ channel. One possibility is to have the Higgs coupling to $b\bar{b}$ smaller than in the SM because this leads to a reduction of the total Higgs decay width and, as a result, increases the Higgs branching ratios into other states. Since the $b\bar{b}$ channel dominates decays of the 125 GeV Higgs in the SM even small decrease of the $hb\bar{b}$ coupling gives non-negligible enhancement in other channels. Such effect is possible, for example, in the minimal supersymmetric standard model (MSSM) due to mixing between the the two CP-even Higgs bosons [9].[1] However, this effect enhances the Higgs signal strength in the WW^* and ZZ^* by the same amount as in the $\gamma\gamma$ channel unless partial Higgs decay widths are non-universally modified. There are no hints in the LHC data for any correlation between the Higgs signal strength in the WW^*/ZZ^* channels and the $\gamma\gamma$ channel so it seems more likely that the enhancement in the $\gamma\gamma$ channel is due to enhanced $\Gamma(h \to \gamma\gamma)$. That is why in this paper we study the possibility that the Higgs signal strength is enhanced in the $\gamma\gamma$ channel while other channels are SM-like.

Since in the SM the Higgs decays into $\gamma\gamma$ only at loop level, substantial corrections to $\Gamma(h \to \gamma\gamma)$ are possible due to new electromagnetically charged states with sizeable couplings to the Higgs [14]. Many models have recently appeared in the literature in which the $h \to \gamma\gamma$ rate is enhanced due to new charged scalars, gauge bosons or vector-like fermions. For representative examples of such scenarios see e.g. refs. [15–27]. In this paper we focus on the possibility that such new states are supersymmetric and study enhanced $h \to \gamma\gamma$ rate in the MSSM. Such possibility is very limited in the MSSM since only the third-generation sfermions [10, 11] and charginos [28] may couple to the Higgs strongly enough to have non-negligible impact on $\Gamma(h \to \gamma\gamma)$. The most attractive possibility is that the $h \to \gamma\gamma$ rate is enhanced by light staus with large left-right mixing [10, 11].[2]

Effects of staus on $h \to \gamma\gamma$ rate have been studied so far from a low-energy perspective. The purpose of the present work is to show that SUSY spectrum with light staus enhancing $h \to \gamma\gamma$ rate may emerge from a well-motivated high-energy scenario. In particular, we point out that GUT-scale boundary conditions for the MSSM soft terms that may lead to enhanced $h \to \gamma\gamma$ rate have to include non-universal gaugino masses with bino and/or wino lighter than gluino.

Strong left-right mixing in the stau sector, as required by enhanced $h \to \gamma\gamma$ rate, strongly prefers models with large $\tan\beta$. Large values of $\tan\beta$ can explain the observed hierarchy between the top and bottom masses [30] and are predicted in minimal SO(10) models since they impose unification of top, bottom and tau Yukawa couplings at the GUT scale ([31, 32]; for some more recent works see e.g. [33–39]). Therefore, these models are perfect candidates to accommodate enhanced $h \to \gamma\gamma$ rate. Unification of the Yukawa couplings by itself is very sensitive to low-energy SUSY threshold corrections, mainly because large values of $\tan\beta$ induce big threshold correction to the bottom mass [40, 41], so it gives constraints for soft SUSY breaking terms at the GUT scale, as well as for low-energy SUSY spectrum.[3] In particular, it has been known for a long time that Yukawa unification requires to go beyond universal soft SUSY breaking terms at the GUT scale [43]. The recent

[1]In the MSSM the Higgs coupling to $b\bar{b}$ can be reduced only if loop corrections to off-diagonal entries of the Higgs mass matrix from the-third-generation sfermion sector are significant [10, 11]. On the other hand, in the NMSSM this effect can be present already at the tree level because of the mixing of the two Higgs doublets with the singlet [12, 13].

[2]The impact of light staus on the $h \to \gamma\gamma$ rate was also investigated in extensions of the MSSM [29].

[3]The gauge coupling unification also has some sensitivity to the MSSM spectrum but as long as gaugino and higgsino masses are not far above $\mathcal{O}(10\text{ TeV})$ the gauge couplings unify within a few percent, see e.g. ref. [42]. On the other hand, for large $\tan\beta$ the correction to the bottom mass varies between a few and several tens of percent depending on some hierarchies in the SUSY spectrum.

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measurement of the Higgs mass of about 125 GeV gives additional constraints on the SUSY spectrum [44, 45]. Implications of the 125 GeV Higgs for SO(10) Yukawa unification were recently reviewed in ref. [46]. The most important effect of the measured Higgs mass is that most of the SUSY spectrum is pushed above 1 TeV, and typically sparticle masses are in a few TeV range. Given the above constraints it is not obvious from the start whether sufficiently light and strongly-mixed staus leading to enhanced $h \to \gamma\gamma$ rate are possible in SO(10) models. In this paper we show that enhanced $h \to \gamma\gamma$ rate can be obtained in agreement with top-bottom-tau Yukawa unification. Assuming D-term splitting of scalar masses, which generically arises in spontaneously broken SO(10) models [47] and is needed to account for proper radiative electroweak symmetry breaking (REWSB) [48], we identify patterns of gaugino masses that may allow for enhanced $h \to \gamma\gamma$ rate. We point out that appropriate patterns of gaugino masses can be accommodated in mirage mediation which often appears in low-energy limits of string theories [49]–[55]. Enhanced $h \to \gamma\gamma$ rate is possible also if SUSY breaking F-term that contributes to gaugino masses is a combination of a singlet and a non-singlet representation of SO(10). We investigate the latter scenario in more detail focusing on the most attractive case with the gaugino masses generated by a mixture of the singlet F-term and the F-term in $\mathbf{24} \subset \mathbf{54}$ of SU(5) $\subset$ SO(10).

The rest of the paper is organized as follows. In section 2 we review the Higgs to diphoton decays with a special emphasis on effects from light staus and argue that at the GUT scale the bino or wino mass should be smaller than the gluino mass in order to substantially enhance $h \to \gamma\gamma$ rate. In section 3 we review necessary conditions for top-bottom-tau Yukawa unification and identify patterns of gaugino masses that can accommodate enhanced $h \to \gamma\gamma$ rate in this class of models. In section 4 we investigate in detail a model with the gaugino masses generated by the singlet F-term and the F-term in $\mathbf{24} \subset \mathbf{54}$ of SU(5) $\subset$ SO(10). First we carefully discuss possible patterns of other soft terms - scalar masses and trilinear terms. Later we concentrate on the predictions of the model for the MSSM spectrum. Finally, we summarize our findings in section 5.

2 Enhanced Higgs to diphoton rate in MSSM

In the SM, $h \to \gamma\gamma$ is a loop mediated process with a dominant contribution from $W^\pm$ and a smaller, but non-negligible, contribution from the top quark [56]:

$$\Gamma^{\rm SM}(h \to \gamma\gamma) \approx \frac{G_\mu \alpha^2 m_h^3}{128\sqrt{2}\pi^3}\left| A_1\left(\frac{m_h^2}{4m_W^2}\right) + \frac{4}{3}A_{1/2}\left(\frac{m_h^2}{4m_t^2}\right)\right|^2 , \tag{2.1}$$

where the amplitudes for the spin-1 and spin-1/2 particle contributions are given at the leading order by:

$$A_{1/2}(x) = 2[x + (x-1)\arcsin^2\sqrt{x}]\, x^{-2} , \tag{2.2}$$

$$A_1(x) = -[2x^2 + 3x + 3(2x-1)\arcsin^2\sqrt{x}]\, x^{-2} , \tag{2.3}$$

with $x_i = m_h^2/4m_i^2$ and m_i denoting the mass of the particle running in the loop. In the MSSM, non-negligible contributions can originate from very light sfermions provided that their Yukawa couplings are large [9]. Thus, the candidates for the modification of the Higgs diphoton rate are stops, and at large $\tan\beta$ also sbottoms and staus. Light stops and sbottoms with large left-right mixing can enhance $\Gamma(h \to \gamma\gamma)$ but at the cost of even larger reduction of the Higgs production cross-section via gluon fusion [57].[4] On the other hand, light staus with large left-right mixing can enhance $\Gamma(h \to \gamma\gamma)$ without affecting the Higgs production rate [10, 11]. In the following subsection we describe the effects of light staus on the Higgs decay rate to two photons in some more detail.

2.1 Effects of light staus

The possibility of the $h \to \gamma\gamma$ rate enhancement by light staus was first emphasized in ref. [10, 11]. In order to better understand the origin of this effect we collect below the most relevant formulae.

For the 125 GeV Higgs, the MSSM Higgs diphoton rate (normalized to its SM value) including the stau effects is given by:

$$\frac{\Gamma(h \to \gamma\gamma)}{\Gamma^{\rm SM}(h \to \gamma\gamma)} \approx \left|1.28 c_V - 0.28 c_t - 0.15 c_\gamma^{\rm stau}\right|^2 , \tag{2.4}$$

where c_V and c_t are the Higgs couplings to W boson and to top quark (normalized to their SM values), respectively, and the coefficients in front of them were obtained using A_i functions defined in eqs. (2.2)–(2.3). The contribution from staus is given by [14]:

$$c_\gamma^{\rm stau} = \sum_{i=1,2} g_{h\tilde{\tau}_i\tilde{\tau}_i} \frac{M_Z^2}{m_{\tilde{\tau}_i}^2} A_0\left(\frac{m_h^2}{4m_{\tilde{\tau}_i}^2}\right) \tag{2.5}$$

with[5]

$$g_{h\tilde{\tau}_1\tilde{\tau}_1} = \cos 2\beta \left(-\frac{1}{2}\cos^2\theta_{\tilde{\tau}} + \sin^2\theta_{\rm W}\cos 2\theta_{\tilde{\tau}}\right) + \frac{m_\tau^2}{M_Z^2} - \frac{m_\tau X_\tau}{2M_Z^2}\sin 2\theta_{\tilde{\tau}} , \tag{2.6}$$

$$g_{h\tilde{\tau}_2\tilde{\tau}_2} = \cos 2\beta \left(-\frac{1}{2}\cos^2\theta_{\tilde{\tau}} - \sin^2\theta_{\rm W}\cos 2\theta_{\tilde{\tau}}\right) + \frac{m_\tau^2}{M_Z^2} + \frac{m_\tau X_\tau}{2M_Z^2}\sin 2\theta_{\tilde{\tau}} , \tag{2.7}$$

where

$$X_\tau = A_\tau - \mu\tan\beta \tag{2.8}$$

and the form factor is given (for $x < 1$ i.e. $m_h < 2m_{\tilde{\tau}_i}$) by [9]

$$A_0(x) = -[x - \arcsin^2\sqrt{x}]\, x^{-2} . \tag{2.9}$$

[4] In principle, simultaneous enhancement of $\Gamma(h \to \gamma\gamma)$ and the production rate is possible for very light stops with left-right mixing large enough to make the stop contribution to the amplitude of the gluon fusion process more than two times larger than the corresponding contribution from the top quark. However, for stop masses consistent with the experimental constraints the required stop mixing gives negative contribution to the Higgs mass so obtaining 125 GeV for this mass would be very problematic.

[5] Different formulae for the Higgs-stau-stau couplings appear in the literature. The couplings we use in this paper agree, after taking into account the opposite sign convention for the stau mixing angle and an apparent misprint, with those from ref. [9]. They agree also with general results given in [14]. However, the formulae used in e.g. [58] and [59] differ from ours (and one from the other) even if rewritten using the same sign convention.

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The stau mixing angle, $\theta_{\tilde{\tau}}$, can be determined in terms of the soft, m_L and m_E, and physical, $m_{\tilde{\tau}_1}$ and $m_{\tilde{\tau}_2}$, stau masses using the following relations:

$$\cos 2\theta_{\tilde{\tau}} = \frac{m_L^2 - m_E^2}{m_{\tilde{\tau}_1}^2 - m_{\tilde{\tau}_2}^2}, \qquad \sin 2\theta_{\tilde{\tau}} = -\frac{2m_\tau X_\tau}{m_{\tilde{\tau}_1}^2 - m_{\tilde{\tau}_2}^2}. \tag{2.10}$$

We also note that the splitting between the stau masses equals

$$m_{\tilde{\tau}_2}^2 - m_{\tilde{\tau}_1}^2 = \sqrt{(m_L^2 - m_E^2)^2 + (2m_\tau X_\tau)^2} \tag{2.11}$$

and is the smallest for $m_L^2 = m_E^2$.

It is clear that the largest enhancement can be obtained for the smallest possible stau masses. Lower limits on the lightest stau vary between 82 and 94 GeV depending on the stau mixing angle and on the mass of the LSP (if it is not strongly degenerate with stau) [60–65].[6] After taking into account these LEP constraints, the modification of $\Gamma(h \to \gamma\gamma)$ coming from the first two terms in each of eqs. (2.6) and (2.7) is at most $\mathcal{O}(5\%)$. Substantial enhancement is possible only if the last terms in these equations are large which requires that the staus are strongly mixed. It is useful to note that the modification of the Higgs decay rate to two photons is well approximated by:

$$\frac{\Gamma(h \to \gamma\gamma)}{\Gamma^{\rm SM}(h \to \gamma\gamma)} \approx \left| 1 + 0.15 A_0\left(\frac{m_h^2}{4m_{\tilde{\tau}_1}^2}\right) \frac{m_\tau^2 X_\tau^2}{m_{\tilde{\tau}_1}^2 m_{\tilde{\tau}_2}^2} \right|^2 \tag{2.12}$$

where we assume that the Higgs couplings to W bosons and top quark are the same as in the SM, which is a very good approximation in the decoupling limit, $M_A \gg M_Z$. In the limit $x \to 0$ the function $A_0(x) \to 1/3$ but for a very light stau it gives additional enhancement e.g. $A_0(\frac{m_h^2}{4m_{\tilde{\tau}_1}^2}) \approx 1.3/3$ for $m_{\tilde{\tau}_1} = 100\,\text{GeV}$. Eq. (2.12) clearly demonstrates that a significant $\gamma\gamma$ rate enhancement is possible if both staus are very light and the value of X_τ is very large in order to compensate the suppression by the tau mass. For instance, the enhancement by 20% requires $X_\tau \approx 70\, m_{\tilde{\tau}_2}$ for $m_{\tilde{\tau}_1} = 100\,\text{GeV}$. This implies $|X_\tau| \gtrsim 20\,\text{TeV}$. Therefore, large values of μ and $\tan\beta$ are necessary to obtain a substantial enhancement of the $h \to \gamma\gamma$ rate.

The enhancement of $\Gamma(h \to \gamma\gamma)$ cannot be, however, arbitrarily large because for too large values of $\mu\tan\beta$ the electroweak vacuum becomes metastable [66–69]. The range of $\mu\tan\beta$ for which the vacuum is stable (or metastable with the life-time longer than the age of the Universe) can be estimated using the following phenomenological formula [69]:

$$\begin{aligned} |\mu\tan\beta| < {} & 56.9\sqrt{m_L m_E} + 57.1\,(m_L + 1.03 m_E) - 1.28\times 10^4\text{GeV} \\ & + \frac{1.67\times 10^6\text{GeV}^2}{m_L + m_E} - 6.41\times 10^7\text{GeV}^3\left(\frac{1}{m_L^2} + \frac{0.983}{m_E^2}\right). \end{aligned} \tag{2.13}$$

which gives good approximation for $m_L,\ m_E \le 2\,\text{TeV}$.[7] It was shown in ref. [68] that $\Gamma(h \to \gamma\gamma)$ can be enhanced by up to 50% without violating the metastability bound.

[6] Those limits can be substantially weaker if the mass splitting between the lighter stau and the LSP is below a few GeV.

[7] Similar formula, valid only for smaller values of the soft stau masses, was given earlier in ref. [67].

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2.2 Enhanced $h \to \gamma\gamma$ rate from GUT-scale perspective

Even though strongly-mixed light staus can lead to substantial enhancement of the $h \to \gamma\gamma$ rate it has not been demonstrated so far that such pattern of SUSY spectrum can be obtained from a well-motivated high-energy model. In the following we argue that, under reasonable assumptions, the $h \to \gamma\gamma$ rate cannot be substantially enhanced, say by 20% or more, if gaugino masses are universal at the GUT scale.

Since enhancing the $\gamma\gamma$ rate requires the lightest stau mass to be around 100 GeV, avoiding a charged LSP implies $|\mu|$, $|M_1|$ or $|M_2|$ at the EW scale to be at most $\sim \mathcal{O}(100\ \text{GeV})$. If the LSP is higgsino-like, then $X_\tau \gtrsim 20$ TeV (which is required for enhancing the $h \to \gamma\gamma$ by at least 20%) would require values of $\tan\beta \gtrsim 200$ (unless $|A_\tau| \gtrsim 20$ TeV) leading to non-perturbative bottom and tau Yukawa couplings. Universal gaugino masses at the GUT scale are not compatible with a gaugino-like LSP with mass of about 100 GeV. The reason is that universal gaugino masses lead to the relation $M_1 : M_2 : M_3 \approx 1 : 2 : 6$ at the EW scale, as a consequence of the one-loop RGEs,[8] which would require gluino mass to be smaller than about $\mathcal{O}(600\ \text{GeV})$ ($\mathcal{O}(300\ \text{GeV})$) for a bino (wino) LSP. Such light gluinos are excluded by the LHC.[9] The lower limit on the gluino mass varies depending on the features of the SUSY spectrum. Nevertheless, gluino mass below about 1.2 TeV is generically excluded even if the first-generation squarks are much heavier than the gluino [72, 73], and the bound gets stronger with smaller masses of the first-generation squarks. The gluino mass above 1.2 TeV together with a gaugino-like LSP with mass below about 100 GeV imply the following condition for the gaugino masses at the GUT scale:

$$|c_1| \equiv \left|\frac{M_1}{M_3}\right| \lesssim \frac{1}{2} \qquad \text{or} \qquad |c_2| \equiv \left|\frac{M_2}{M_3}\right| \lesssim \frac{1}{4}\,, \tag{2.14}$$

where we introduced parameters c_1 and c_2 defined as the GUT scale ratios of the bino and wino masses to the gluino mass. We should emphasize that the condition (2.14) was obtained without assuming Yukawa unification so it is valid in any (R-parity conserving) MSSM model and should be valid also in many MSSM extensions such as NMSSM. A more detailed discussion of ranges of c_i values leading to enhanced $h \to \gamma\gamma$ rates and compatible with the experimental constraints and unification of the Yukawa couplings is presented in the next section.[10]

3 Top-bottom-tau Yukawa unification and enhanced Higgs diphoton rate

We now would like to go one step further and ask the question whether it is possible to obtain enhanced $h \to \gamma\gamma$ rates in SO(10) models predicting top-bottom-tau Yukawa unification. Large values of $\tan\beta \sim \mathcal{O}(50)$ necessary for substantial stau mixing are predicted in

[8]The pattern of gaugino masses at the EW scale can be altered in the presence of extreme hierarchy between the A-terms and gaugino masses, $|A_0/M_{1/2}| > \mathcal{O}(100)$ because in such a case two-loop effects in the RGEs for gaugino masses are non-negligible. See e.g. ref. [70] for an example of such a scenario.

[9]Gluino could be so light only if the SUSY spectrum is extremely compressed [71] which is not the case with the LSP mass around 100 GeV.

[10]It was shown very recently [59] that an enhanced $h \to \gamma\gamma$ rate can be accommodated for a specific case of $c_1 = c_2 = 1/10$, however without imposing the Yukawa coupling unification. More general scenarios have not been studied yet.

such models. On top of the condition (2.14) for enhanced $h \to \gamma\gamma$ rate, there are also conditions for the MSSM parameters coming from the assumption of SO(10) GUT symmetry group and top-bottom-tau Yukawa unification which we discuss in the following.

3.1 Top-bottom-tau Yukawa unification and REWSB

It is well known that in models with top-bottom-tau Yukawa unification proper REWSB is endangered because of the RGE effects of large bottom and tau Yukawa couplings which tend to make M_A^2 tachyonic [46]. In particular, REWSB is incompatible with top-bottom-tau Yukawa unification in CMSSM [41]. In order to solve this problem some non-universalities in the soft scalar masses at the GUT scale need to be introduced [43]. The pattern of soft scalar masses consistent with SO(10) gauge symmetry is given by:

$$\begin{aligned} m_{H_d}^2 &= m_{10}^2 + 2D\,, \\ m_{H_u}^2 &= m_{10}^2 - 2D\,, \\ m_{Q,U,E}^2 &= m_{16}^2 + D\,, \\ m_{D,L}^2 &= m_{16}^2 - 3D\,, \end{aligned} \tag{3.1}$$

where D parametrizes the size of a U(1) D-term contribution to soft scalar masses which generically arises in an effective theory below the GUT scale when SO(10) gauge symmetry (which has rank bigger than that of the SM gauge group) is spontaneously broken to its SM subgroup (or other subgroup with the rank smaller than that of SO(10)) [47]. The coefficients in front of D are fixed by charges under the broken U(1). It has been shown that for $D > 0$, $m_{10} > m_{16}$ and universal other soft terms proper REWSB is consistent with top-bottom-tau Yukawa unification [48].

The crucial role in top-bottom-tau Yukawa unification is played by the sign of μ because it controls the sign of the dominant finite SUSY threshold corrections to the bottom mass [40, 41, 74]:

$$\left(\frac{\delta m_b}{m_b}\right)^{\text{finite}} \approx \frac{g_3^2}{6\pi^2}\mu m_{\tilde{g}} \tan\beta\, I(m_{\tilde{b}_1}^2, m_{\tilde{b}_2}^2, m_{\tilde{g}}^2) + \frac{h_t^2}{16\pi^2}\mu A_t \tan\beta\, I(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2, \mu^2)\,, \tag{3.2}$$

where the loop integral $I(x, y, z)$ is defined e.g. in the appendix of [40]. Bottom-tau Yukawa unification requires the finite correction to be negative with the magnitude between 10 and 20% [75]. The first term in eq. (3.2) comes from the gluino-sbottom contribution and is the dominant one over the most part of the parameter space so Yukawa unification strongly prefers $\mu < 0$. Top-bottom-tau Yukawa unification is also possible for $\mu > 0$ [76] but this requires very heavy scalars with m_{16} exceeding $\mathcal{O}(20\text{ TeV})$, and even larger A-terms, [77, 78] so corresponding fine-tuning of electroweak symmetry breaking is very big, much bigger than the fine-tuning imposed on the MSSM by the measured Higgs mass of 125 GeV. Moreover, staus in models with positive μ also tend to be heavy. The reason is that in those models gluinos have to be light and sbottoms very heavy in order to suppress the gluino-sbottom correction to the bottom mass (which has the wrong sign for positive μ). The masses of staus cannot be much smaller than the masses of sbottoms because they have common value at the GUT scale and their RG evolution (which is dominated by the RGE terms proportional to the Yukawa couplings) is similar. It was found in ref. [78] that for positive μ staus are even heavier than sbottoms and that the properties of the 125 GeV Higgs are almost the same as those of the SM Higgs. In contrast, for negative μ top-bottom-tau Yukawa unification is possible also for heavy gluino. As a result, staus can be much lighter than stops and sbottoms because RG contribution from gauginos to squarks may be large while the one to sleptons small. For the above reasons in this paper we consider only $\mu < 0$.

3.2 Non-universal gaugino masses

As explained in section 2, non-universal gaugino masses are necessary for enhancing $h \to \gamma\gamma$ rate because the experimental data from the LHC set lower bounds on the gluino mass. There is also another reason why gluinos should be rather heavy (what implies necessity of non-universal gaugino masses at the GUT scale). It is directly related to the requirement of Yukawa unification and enhanced $\Gamma(h \to \gamma\gamma)$. The condition of b-τ Yukawa unification introduces rather strong correlation between μ and the gluino mass. Since b-τ Yukawa unification requires the finite threshold correction to the bottom mass to be below about 20%, one can set the following bound on the ratio of the gluino mass and μ [46]:

$$\frac{m_{\tilde{g}}}{|\mu|} \gtrsim 2.5\,, \tag{3.3}$$

where we used eq. (3.2) with the chargino-stop contribution neglected, $\tan\beta \approx 50$ and assumed that gluino is heavier than the heaviest sbottom which is usually the case unless $M_3 \ll m_{16}$. Big enhancement of $h \to \gamma\gamma$ rate requires big $|\mu|$, e.g. for $m_{\tilde{\tau}_1} = 100$ GeV and $\tan\beta \approx 50$ the enhancement by 20% requires $|\mu| \gtrsim 500$ GeV so the bound (3.3) implies that the gluino mass has to be at least about 1.2 TeV. Accidentally, this lower bound on the gluino mass coincides with the experimental bound so it leads to the same constraint (2.14) on the gaugino masses at the GUT scale.

Even if gaugino masses at the GUT scale are such that the LSP is neutral when the lightest stau mass is around 100 GeV it is not guaranteed that the $\gamma\gamma$ rate is substantially enhanced. As explained in section 2, also the second stau should be relatively light because the stau mixing should be close to the maximal (i.e. $\theta_{\tilde{\tau}} \approx \pi/4$). For a fixed value of the off-diagonal term in the stau mass matrix, X_τ, the stau mixing is maximized for $m_L \approx m_E$ at the EW scale. Due to the D-term splitting of the scalar masses (3.1) m_L is smaller than m_E at tree level. On the other hand, RG running of m_L and m_E is significantly different. In particular, the negative RGE contribution proportional to a large τ Yukawa coupling is two times larger for m_E^2 than the corresponding contribution to m_L^2 so this effect can diminish the initial splitting between m_L and m_E. On top of that, the electroweak gaugino masses also contribute differently to m_L and m_E and this effect strongly depends on the values of c_1 and c_2. Therefore, the analysis of the $\gamma\gamma$ rate requires a careful numerical treatment.

In the following we focus on a numerical scan of the parameter space of a model with D-term splitting of scalar masses (3.1), universal trilinear coupling A_0 (justification for these assumptions will be discussed in the next section), and arbitrary gaugino masses parametrized by M_3, c_1 and c_2. There are eight free parameters altogether so a homogeneous scan of the parameter space would not be very efficient. Therefore, we use Markov Chain Monte Carlo techniques to sample the parameter space. More precisely, we adopt a similar procedure to that proposed in ref. [70] which makes use of the Metropolis-Hastings algorithm [79, 80]. We consider the following ranges of parameters:

$$\begin{aligned} m_{16} &\in (0, 10\ \text{TeV})\ , \\ M_3 &\in (0, 10\ \text{TeV})\ , \\ \frac{m_{10}}{m_{16}} &\in (0, 10)\ , \\ \frac{A_0}{m_{16}} &\in (-3, 3)\ , \\ \frac{D}{m_{16}^2} &\in \left(0, \frac{1}{3}\right)\ . \end{aligned} \tag{3.4}$$

We also reject values of D that lead to a negative value of any of the soft scalar squared masses (3.1) at the GUT scale. On the other hand, for c_1, c_2 and $\tan\beta$ we do not specify ranges over which they are scanned.

In order to quantify the goodness of top-bottom-tau Yukawa unification we introduce the following quantity:

$$R \equiv \left.\frac{\max\left(h_t, h_b, h_\tau\right)}{\min\left(h_t, h_b, h_\tau\right)}\right|_{\text{GUT}}\ . \tag{3.5}$$

We use SOFTSUSY [81] to solve the 2-loop renormalization group equations and calculate the MSSM spectrum. For every randomly generated point in the parameter space we demand proper REWSB and one of the neutralinos being the LSP. We also reject points that do not satisfy the vacuum metastability condition (2.13). We calculate the thermal relic abundance of the lightest neutralino, as well as $\text{BR}(b \to s\gamma)$, $\text{BR}(B_s \to \mu^+\mu^-)$ and $(g-2)_\mu$ using MicrOmegas [82]. We apply the following constraints: [83–89]

$$\begin{aligned} 2.52 \cdot 10^{-4} < \text{BR}(b \to s\gamma) < 4.34 \cdot 10^{-4}\,, \\ 1.5 \cdot 10^{-9} < \text{BR}(B_s \to \mu^+\mu^-) < 4.3 \cdot 10^{-9}\,, \\ \Omega_{\text{DM}} h^2 < 0.13\,. \end{aligned} \tag{3.6}$$

For $\text{BR}(b \to s\gamma)$ we use the 2σ experimental constraint combined in quadrature with the theoretical uncertainty of $4 \cdot 10^{-5}$. Note that the computation of $\text{BR}(b \to s\gamma)$ in the MSSM is completed at the NLO [85], while in the SM at the NNLO [84]. Moreover, the NNLO corrections shift the NLO result in the SM [90] by about $4 \cdot 10^{-5}$ so we use the value of this shift as an estimate of the theoretical error in MSSM. The theoretical uncertainty for $\text{BR}(B_s \to \mu^+\mu^-)$ [87] is still much smaller than the experimental one so we use for this observable the 2σ limit obtained by combining [86] the CMS [91] and LCHb [92] results. We demand that only the upper bound on $\Omega_{\text{DM}} h^2$ is satisfied but, as we shall discuss later, also the lower bound can be satisfied in some circumstances. The relevant lower mass limits on the MSSM particles [1, 2, 60–65, 93]:

$$\begin{aligned} 90\,\mathrm{GeV} &< m_{\tilde{\tau}}\ , \\ 103.5\,\mathrm{GeV} &< m_{\tilde{\chi}^{\pm}}\ , \\ 750\,\mathrm{GeV} &< m_A\ , \\ 123\,\mathrm{GeV} &< m_h < 128\,\mathrm{GeV} \end{aligned} \tag{3.7}$$

are also applied. Experimental lower mass limit on m_A depends on $\tan\beta$ but we fix it to a constant value since the assumption of top-bottom-tau Yukawa unification constrains $\tan\beta$ to be between 40 and 50, where the limit of 750 GeV is a good approximation. The theoretical uncertainty in the prediction of the Higgs mass, calculated by SOFTSUSY at two-loop level, is about 3 GeV [94] so we assume that the Higgs mass between 123 and 128 GeV is consistent with the experimentally measured value of about 125.5 GeV. In practice, we found that only the lower bound on the Higgs mass constrains the parameter space. It was argued in ref. [95–97] that dominant three-loop corrections to the Higgs mass are positive with the magnitude up to 3 GeV. These results strengthen our assumption that points for which SOFTSUSY gives the Higgs mass of 123 GeV are compatible with the experimental data. We found that other lower limits on sparticle masses from direct LHC searches do not impose any additional constraints on the model.

We also use the quantity $R_{\gamma\gamma}$ defined as the predicted signal strength in the $\gamma\gamma$ channel normalized to the corresponding SM prediction:

$$R_{\gamma\gamma} \equiv \frac{\sigma(gg\to h)\times \mathrm{BR}(h\to\gamma\gamma)}{\sigma(gg\to h)^{\mathrm{SM}}\times \mathrm{BR}(h\to\gamma\gamma)^{\mathrm{SM}}}\,. \tag{3.8}$$

The strong LHC constraints on m_A push the model to the decoupling region of the MSSM so the fermion and gauge boson couplings are almost the same as in the SM. In consequence, both the production cross-section and the total decay width of the Higgs are practically the same as in the SM. In principle, light stops or sbottoms could modify the Higgs production cross-section and decay width into $\gamma\gamma$ but we found in our numerical analysis that stops and sbottoms are relatively heavy and such effects are negligible. Therefore, any deviations of $R_{\gamma\gamma}$ from one in this model are due to light and strongly-mixed staus. We numerically compute $R_{\gamma\gamma}$ using the formulae collected in subsection 2.1.

In figure 1 we present the results of our scan in the c_1–c_2 plane. Only points that correspond to $R < 1.1$ (Yukawa unification better than 10%) and $R_{\gamma\gamma} > 1.1$ (enhancement of the $\gamma\gamma$ rate by at least 10%) are shown. Black points satisfy all the constraints mentioned before, including the metastability bound (2.13). For blue (red) points $b\to s\gamma$ ($B_s\to\mu^+\mu^-$) constraint is relaxed, while yellow points violate both $b\to s\gamma$ and $B_s\to\mu^+\mu^-$. Note that all points have either c_1 or c_2 small, in agreement with our qualitative constraint (2.14). However, there are very few points in the quadrant with $c_1 > 0$ and $c_2 < 0$. In this quadrant of the c_1–c_2 plane the Yukawa-unified solutions with enhanced $h\to\gamma\gamma$ rate are characterized by a light pseudoscalar Higgs with mass typically below the experimental lower bound [93]. The lightness of the CP-odd Higgs in this region of parameter space follows from the assumption of Yukawa unification. Without imposing this assumption enhanced $h\to\gamma\gamma$ rate can be easily obtained also in the quadrant with $c_1 > 0$ and $c_2 < 0$.

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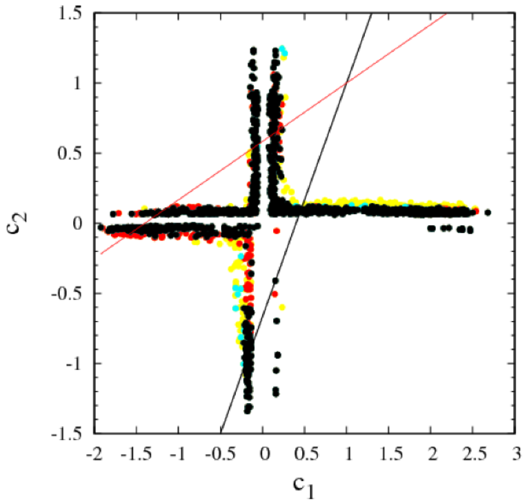

Figure 1. Points with $R_{\gamma\gamma} > 1.1$ and $R < 1.1$ in the c_1–c_2 plane. Black points satisfy all the constraints in (3.6) and (3.7). Blue (red) points violate $b \to s\gamma$ ($B_s \to \mu^+\mu^-$). Yellow points violate both $b \to s\gamma$ and $B_s \to \mu^+\mu^-$. Black (red) line corresponds to gaugino masses generated by a mixture of the singlet and **24** F-term (in mirage mediation).

The scan with arbitrary gaugino masses is very useful for illustrative purposes but it is more interesting to focus on specific models of SUSY breaking that predict some patterns of gaugino masses. One example of such pattern is mirage mediation [49]–[55] which predicts

$$M_a = M(\rho + b_a g_a^2)\,, \tag{3.9}$$

where M and ρ are free parameters, $b_a = (33/5, 1, -3)$ for $a = 1, 2, 3$ and g_a are the gauge coupling constants (which in our model are assumed to unify at the GUT scale). In this case the relation between c_1 and c_2 is fixed:

$$c_2 = \frac{5c_1 + 7}{12} \tag{3.10}$$

and corresponds to the red line in figure 1. The line intersects regions in which top-bottom-tau Yukawa unification and enhanced $h \to \gamma\gamma$ rate can be obtained. This indicates that Yukawa unification with enhanced $\gamma\gamma$ rate can be realized in some mirage mediation scenario. It is an indication and not a proof because the points shown in figure 1 do not correspond exactly to any mirage mediation model. Those points were obtained assuming universal trilinear soft terms while mirage mediation models predict trilinear terms with non-universalities correlated with the non-universalities of the gaugino masses (3.9). Thus, to draw any firm conclusions about such models it is necessary to perform separate calculations dedicated to each of them. Our preliminary results show that top-bottom-tau Yukawa unification and enhanced $h \to \gamma\gamma$ rate indeed can be obtained in some mirage mediation models.[11] The full results will be presented elsewhere. In the present paper we concentrate on another class of models.

[11] Yukawa unification in "effective" mirage mediation was discussed in a recent ref. [98], however, neglecting the necessary non-universalities of the A-terms and without any discussion of $\Gamma(h \to \gamma\gamma)$.

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Non-universal gaugino masses are also possible in GUT models with pure gravity mediation provided that the SUSY breaking F-term belongs to an appropriate non-singlet representation of the unifying gauge group. In order to see this, notice first that the gaugino masses in supergravity can arise from the following dimension five operator:

$$\mathcal{L} \supset -\frac{F^{ab}}{2M_{\rm Planck}}\lambda^a\lambda^b + \text{c.c.}\,, \tag{3.11}$$

where λ^a are the gaugino fields and the resulting gaugino mass matrix is $\frac{\langle F^{ab}\rangle}{M_{\rm Planck}}$. Non-zero gaugino masses require that the vacuum expectation value of the relevant F-term, $\langle F^{ab}\rangle$, transforms as the singlet of the SM gauge group and, in order to make the term (3.11) invariant under the GUT group, as any of the representations present in the symmetric part of the direct product of the two adjoint representations, which for SO(10) are:

$$(\mathbf{45}\times\mathbf{45})_{\rm S} = \mathbf{1}+\mathbf{54}+\mathbf{210}+\mathbf{770}\,. \tag{3.12}$$

If SUSY is broken by an F-term transforming as a non-singlet representation of SO(10), gaugino masses are not universal. Complete classification of non-universal gaugino masses for SO(10) and its subgroups can be found in ref. [99].

Particularly interesting is the case when the SUSY breaking F-term transforms as $\mathbf{24}$ of SU(5) $\subset$ SO(10) for which the gaugino masses have the following pattern:

$$M_1 : M_2 : M_3 = -\frac{1}{2} : -\frac{3}{2} : 1\,. \tag{3.13}$$

The negative value of M_2 (with respect to M_3) is preferred from the phenomenological point of view because for $\mu < 0$ such values make the SUSY contribution to $(g-2)_\mu$ positive. Moreover, for $\mu < 0$ and $M_2 < 0$ the chargino-stop contribution to $b \to s\gamma$ is smaller than for $M_2 > 0$ so the tension with this observable is relaxed. Top-bottom-tau Yukawa unification in a model with gaugino masses generated by the SUSY breaking F-term in $\mathbf{24}$ of SU(5) $\subset$ SO(10) and the soft scalar masses (3.1) with D-term splitting was investigated in ref. [38] with a special emphasize on the constraints from $(g-2)_\mu$ and $b \to s\gamma$, while the LHC constraints on that model from the Higgs and SUSY searches were studied in refs. [39, 46].

The gaugino mass pattern (3.13) is interesting but still not suitable to allow for substantial enhancement of $h \to \gamma\gamma$ rate. It gives $c_1 = -0.5$ and $c_2 = -1.5$ which are too big values, as can be seen in figure 1. This problem may be solved when supersymmetry is broken by more than just one F-term. The simplest possibility is to consider two such F-terms, one transforming as $\mathbf{24}$ of SU(5) $\subset$ SO(10) and second transforming as the gauge singlet. If both such F-terms have non-zero VEVs the gaugino masses can be parametrized in the following way:

$$\begin{aligned} M_1 &= M_3^{(1)} - \frac{1}{2}M_3^{(24)} \equiv \frac{1-\frac{1}{2}c_{24}}{1+c_{24}}M_3\,, \\ M_2 &= M_3^{(1)} - \frac{3}{2}M_3^{(24)} \equiv \frac{1-\frac{3}{2}c_{24}}{1+c_{24}}M_3\,, \end{aligned} \tag{3.14}$$

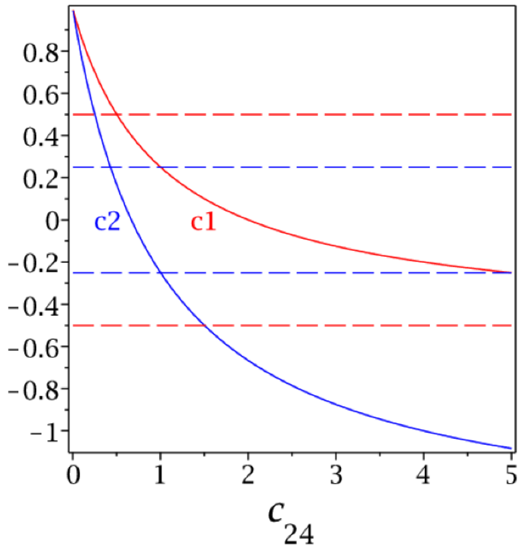

Figure 2. c_1 (red) and c_2 (blue) as functions of c_{24}. The dashed lines correspond to $|c_1| = 1/2$ and $|c_2| = 1/4$.

where $c_{24} \equiv M_3^{(24)}/M_3^{(1)}$ is the ratio of the **24** and singlet F-term contributions to the gluino mass. In this case the relation between c_1 and c_2 is given by:

$$c_2 = \frac{5c_1 - 2}{3} \tag{3.15}$$

and corresponds to the black line in figure 1. This line intersects several regions with good Yukawa unification and enhanced $h \to \gamma\gamma$ rate. A more detailed analysis of these regions will be presented in the next section.

In order to get a feeling how the ratios of the gaugino masses, M_1/M_3 and M_2/M_3, depend on c_{24} we plot them in figure 2. Notice that c_{24} has to be positive and larger than about 0.4 in order to get $|c_1| \lesssim 1/2$ or $|c_2| \lesssim 1/4$. Moreover, for c_{24} below (above) about 0.9 the LSP is dominated by the wino (bino). Notice also that for $c_{24} < 2/3$, $M_2 > 0$ so the contribution to $(g-2)_\mu$ from the chargino-sneutrino loop, which is typically a dominant SUSY contribution [100, 101], is negative and the discrepancy between the theoretical prediction and experimental result becomes even larger than in the SM.

4 Model with SUSY broken by F-terms in 1 and 24 representations of SU(5)

Phenomenological implications resulting from non-universal gaugino masses generated by a mixture of the singlet and **24** F-term were investigated before [102–104]. However, neither the impact of that assumption on Yukawa unification nor the $h \to \gamma\gamma$ rate has been considered so far. In this section we study in detail implications of these assumptions. First of all, we should check how the assumed structure of the F-terms influences all the soft SUSY breaking terms. In the previous section we concentrated on the gaugino masses

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assuming for simplicity that the trilinear terms are universal and the structure of the soft scalar masses is as given in eq. (3.1). We show in the following subsection whether and when such simplifying assumptions can be justified.

4.1 (Non)universalities of other soft terms

24 of SU(5) appears in each of the three non-singlet representations of SO(10) in the symmetric part of the product $\mathbf{45} \times \mathbf{45}$ given in eq. (3.12). The pattern of the gaugino masses is the same for each of these SO(10) representations. This is not true for the soft scalar masses and trilinear terms.

Soft scalar masses terms arising in supergravity from dimension six operators have the following structure

$$\frac{\langle F\, F\rangle^{ij}}{M^2_{\rm Planck}} \phi_i^\dagger \phi_j \,, \tag{4.1}$$

where $\phi_{i,j}$ are scalar fields. In order to contribute to the scalar masses without breaking the SM gauge symmetry, the VEV of the product of F-terms in the above formula must transform as a singlet of the SM and some representation present in the product

$$\overline{\mathbf{16}} \times \mathbf{16} = \mathbf{1} + \mathbf{45} + \mathbf{210}\,. \tag{4.2}$$

Any F-term transforming as a tensor representation $\mathbf{R}$ of SO(10) gives some universal contribution to the scalar masses because the singlet is always present in the product $(\mathbf{R} \times \mathbf{R})_{\rm S}$. In addition, some of them may give also non-universal contributions. This happens if the symmetric part of the product $\mathbf{R} \times \mathbf{R}$ contains **45** or **210**. This is the case for the representations **210** and **770** but not for **54**.[12] Let us also note that for the representation **210** additional contribution to the soft scalar masses may arise from the mixed product of the singlet and non-singlet F-terms because **210** is present in both products, (3.12) and (4.2).

The soft trilinear terms, generated by dimension five operators, have the form

$$\frac{\langle F\rangle^{ijk}}{M_{\rm Planck}} \phi_i \phi_j \phi_k \,, \tag{4.3}$$

and may arise for F-terms transforming as any of the representations appearing in the product

$$(\mathbf{16} \times \mathbf{16})_{\rm S} \times \mathbf{10} = \mathbf{1} + \mathbf{45} + \mathbf{54} + \mathbf{210} + \mathbf{1050}\,. \tag{4.4}$$

Let us now discuss in turn each of the non-singlet representations present in the r.h.s. of eq. (3.12). Representation **770** is absent in the r.h.s. of eq. (4.4) so the corresponding F-term does not generate any trilinear terms. F-terms transforming as **54** and **210** do generate soft trilinear terms. In fact, **210** leads to two kinds of such terms: universal and non-universal. This follows from the fact that there are two independent singlets in the product $(\mathbf{16} \times \mathbf{16})_{\rm S} \times \mathbf{10} \times \mathbf{210}$. There is only one contribution to the trilinear terms coming

[12] The products $\mathbf{54} \times \mathbf{54}$ and $\mathbf{210} \times \mathbf{210}$ are given in ref. [105]. In the case of $\mathbf{770} \times \mathbf{770}$ one can prove that it contains **45** and **210** using the Young tableaux technique.

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from F-terms transforming as **54**. A more detailed analysis shows that this contribution is universal.

The last results may be understood using the following argument. In the only singlet in the product $(\mathbf{16}\times\mathbf{16})_S\times\mathbf{10}\times\mathbf{54}$ the part containing the Higgs multiplet and the F-term multiplet, $\mathbf{10}\times\mathbf{54}$, transforms as **10**. Representation **10** may be understood as a vector of SO(10) while **54** as a symmetric matrix. Their product transforming as **10** is obtained by multiplying this vector with this matrix. $\langle F\rangle$ must be a SM singlet so it is represented by a diagonal matrix. Moreover, this matrix is proportional to unit matrices in the subspaces corresponding to SU(3) and SU(2) subgroups of SO(10). So, the only non-universality generated by **54** is that between trilinear terms involving Higgs doublets versus trilinear terms involving Higgs triplets. We neglect the latter assuming some mechanism for the Higgs doublet-triplet splitting. Trilinear terms for the Higgs doublets alone are universal.

Let us summarize the above results. There are three possibilities in models with SUSY broken by two F-terms, one transforming as **1** and one as **24** of SU(5). Each case leads to the same pattern of the non-universalities in the gaugino masses, the one discussed in the previous subsection. The patterns of other soft terms depends on the representation of SO(10) in which **24** of SU(5) is embedded. In the case of **54** of SO(10) all other soft terms are universal at the GUT scale. Representation **210** leads to non-universality in soft scalar masses and trilinear terms. Finally, **770** gives non-universal scalar masses but universal trilinear terms.

4.2 Numerical results

The discussion in the previous subsection shows that embedding **24** of SU(5) in **54** of SO(10) is the most attractive possibility. We have shown that in such a case all soft trilinear terms have one common value, A_0, at the GUT scale. As for the soft scalar masses we assume the pattern given by (3.1). The F-term transforming as **54** does not give non-universalities in the scalar masses but the difference between m_{10} and m_{16} may be generated by the RGE running between the Planck scale and the GUT scale.[13] Our choice is also the simplest because **54** is the smallest representation of SO(10) leading to non-universal gaugino masses.

Having defined the model we can investigate its properties. We perform similar numerical scan to the one described in the previous section with the only difference that the gaugino masses are determined by (3.14) as a function of c_{24} and M_3 (i.e. we scan along the black line in figure 1).

In figure 3 we present points from our scan characterized by $R_{\gamma\gamma} > 1.1$ in the R–M_3 plane. It shows that $R_{\gamma\gamma} > 1.1$ may be consistent with the experimental constraints provided that the gluino is heavy enough. The lower bound on M_3 could be somewhat relaxed if the constraints from $b \to s\gamma$ and $B_s \to \mu^+\mu^-$ were not taken into account. Moreover, the bound on the gluino mass depends very weakly on R and even demanding the perfect top-bottom-tau Yukawa unification, i.e. $R \approx 1$, is not an obstacle to get enhanced $h \to \gamma\gamma$ rate.

[13] Using **210** or **770** instead of **54** would lead to much more complicated pattern of scalar masses, splitting e.g. masses of sfermions belonging to one **16** representation of SO(10).

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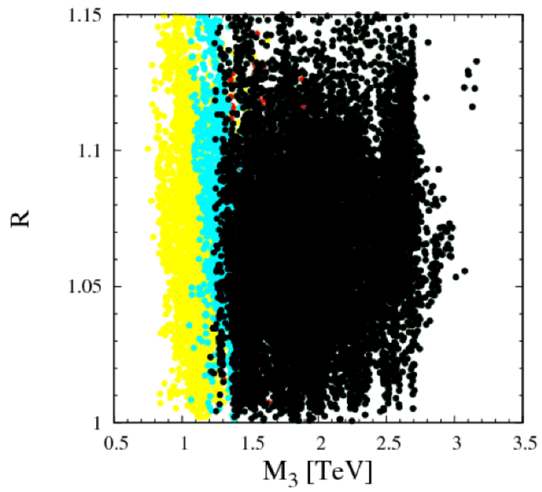

Figure 3. Points with $R_{\gamma\gamma} > 1.1$ in the R–M_3 plane. The colour coding is as in figure 1.

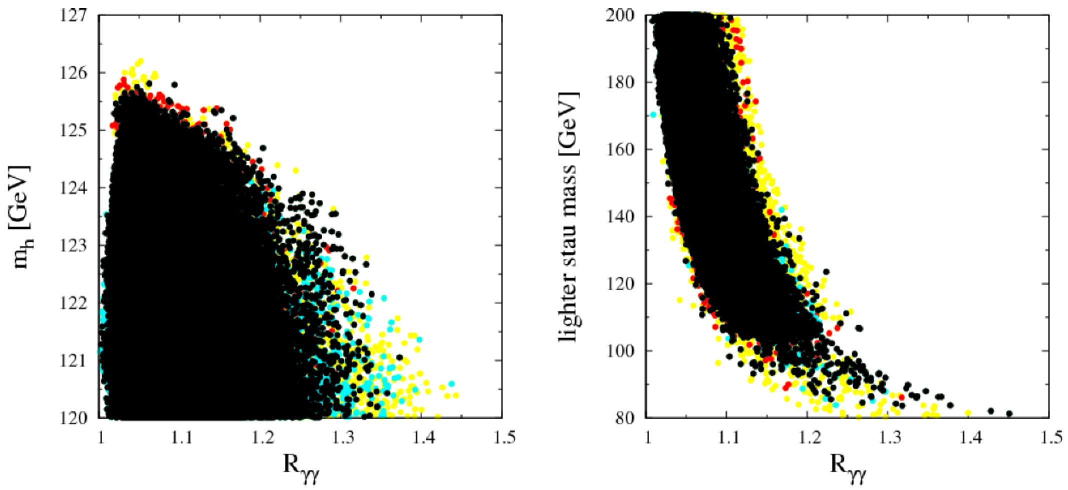

Figure 4. Scatter plots of m_h (left panel) and the lighter stau mass (right panel) versus $R_{\gamma\gamma}$ for the points with $R < 1.1$. Points in the left panel were obtained without imposing the bound (3.7) on the Higgs boson mass. The colour coding is the same as in figure 1.

In the left panel of figure 4 we plot the points with $R < 1.1$ in the $m_h - R_{\gamma\gamma}$ plane. It can be seen from this plot that $h \to \gamma\gamma$ rate can be enhanced even by 30%. Notice also that $R_{\gamma\gamma}$ is anti-correlated with the Higgs mass. This is because $R_{\gamma\gamma}$ grows when $X_\tau^2/(m_{\tilde{\tau}_1} m_{\tilde{\tau}_2})$ increases but this, in turn, implies that negative contribution to the Higgs mass from the stau sector grows as well [10, 11]. The anti-correlation is weaker for smaller values of m_h because too large values of $X_\tau^2/(m_{\tilde{\tau}_1} m_{\tilde{\tau}_2})$ are excluded by the constraints on the vacuum metastability.

The plot in the right panel of figure 4 confirms our expectations showing that large enhancement of the $\gamma\gamma$ rate requires the lighter stau to have mass around 100 GeV. Taking the most conservative lower limit on the stau mass of 82 GeV even 40% enhancement can be obtained. Note, however, that non-negligible enhancement do not require extremely

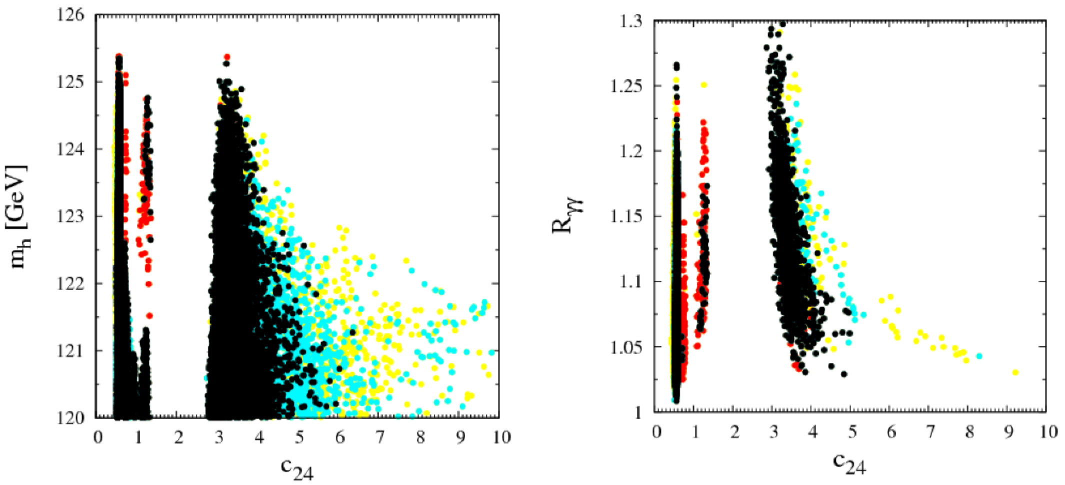

Figure 5. Scatter plots of m_h (left panel) and $R_{\gamma\gamma}$ (right panel) versus c_{24} for the points with $R < 1.1$. Points in the left panel have $R_{\gamma\gamma} > 1.1$ but they were obtained without imposing the bound (3.7) on the Higgs boson mass. The colour coding is the same as in figure 1.

light stau. For instance, $R_{\gamma\gamma} > 1.1$ can be obtained for the lightest stau mass as large as about 200 GeV.

In figure 5 we present the dependence of the Higgs mass and $R_{\gamma\gamma}$ on c_{24}. In accord with our qualitative discussion in subsection 3.2, $m_h > 123$ GeV and $R_{\gamma\gamma} > 1.1$ is possible only if c_{24} is positive and larger than about 0.5 but smaller than about 4 (5) when the B-physics constraints are (not) taken into account. However, not all values of c_{24} in this range are possible due to various phenomenological constraints. Since for $c_{24} = 2/3$ the wino mass vanishes, values of c_{24} very close to this value are excluded by the LEP constraint on the chargino mass. For $c_{24} = 2$ the bino mass vanishes so values of c_{24} close to this value imply very light bino-like LSP with too large thermal relic abundance which cannot be reduced by stau coannihilation because the mass splitting between LSP and stau is too large. On the other hand, moving away from $c_{24} = 2$ increases the bino mass so the mass splitting between LSP and stau gets smaller. For $c_{24} \lesssim 1.4$ and $c_{24} \gtrsim 3$ the mass splitting can be small enough to have the bino-like LSP relic abundance in agreement with observations. Finally, also some region around $c_{24} = 1$ is excluded because the pseudoscalar Higgs mass turns out to be below 750 GeV there. We should stress at this point that very light pseudoscalar Higgs in the region around $c_{24} = 1$ is a consequence of the assumption of Yukawa unification. Without this assumption enhanced $h \to \gamma\gamma$ rate can be obtained also in this region with the pseudoscalar Higgs mass satisfying the experimental constraints.

For large $\tan\beta$, constraints from flavour changing observables are very important. For most values of c_{24}, $b \to s\gamma$ is the most constraining observable. This can be seen from the right panel of figure 5. However, for c_{24} around one the main constraint comes from $B_s \to \mu^+\mu^-$ because in this region the CP-odd Higgs turns out to be light, with mass around 1 TeV at most.

It was argued in ref. [58] that stau-induced enhanced $h \to \gamma\gamma$ rate is correlated with SUSY contribution to $(g-2)_\mu$ and that enhancement of $h \to \gamma\gamma$ rate by a few tens of percent typically leads to $(g-2)_\mu$ within 2σ from the experimental central value [106, 107]. However, that statement was based on the assumption of the slepton mass universality which is strongly violated in our case by the RG effects of the τ Yukawa coupling. In consequence, we found that SUSY contribution to $(g-2)_\mu$ is quite small. For $c_{24} > 2/3$, this contribution is positive, as preferred by the experiment, but smaller than approximately 4×10^{-10} so about 3σ below the experimental value. For $c_{24} < 2/3$, the SUSY contribution to $(g-2)_\mu$ (dominated by the chargino-sneutrino loop which sign is given by the sign of μM_2 [100, 101]) is negative with the absolute value below about 10^{-9}. Therefore, $(g-2)_\mu$ slightly favours $c_{24} > 2/3$.

We found that Yukawa-unified solutions with enhanced $h \to \gamma\gamma$ rate have typically large and negative A-terms. A_0/m_{16} is between -3 and -2, except the region with $c_{24} \in (0.5, 0.6)$ where A_0/m_{16} between -2.5 and +0.5 is possible. Large negative A-terms at the GUT scale are generally needed to generate large enough stop mixing to account for the observed Higgs boson mass. Nevertheless, we expect that large negative values of A-terms at the GUT scale are strictly related to our assumption of intergenerational degeneracy of the soft sfermion masses because it was shown that if at the GUT scale the first two-generations sfermions are much heavier than the third-generation ones large stop mixing can be generated with small, or even vanishing, A-terms at the GUT scale [108]. It is beyond the scope of the present work to study in detail the intergenerational splitting of scalar masses in the context of top-bottom-tau Yukawa unification and we leave it for a future work.

4.3 Predictions for the MSSM spectrum

The MSSM spectrum is substantially different for different regions of c_{24}. Examples of such differences, for the pseudoscalar and for the lighter chargino masses, are shown in figure 6. In the following we discuss separately $c_{24} < 1$ and $c_{24} > 1$, corresponding to the dominant contribution to the gluino mass coming from the SUSY breaking F-term in the singlet and non-singlet representations, respectively.

4.3.1 Singlet F-term domination

For $c_{24} < 1$, the dominant contribution to the gluino mass comes for the singlet F-term. However, enhanced $h \to \gamma\gamma$ rate requires non-negligible contribution of the non-singlet F-term. There are two separate regions with $c_{24} \in (0.5, 0.6)$ and $c_{24} \in (0.7, 0.8)$. Partial cancellation between the contributions to M_2 from the singlet and non-singlet F-terms results in a wino-like LSP with mass above 100 GeV, due to the LEP constraint on the chargino mass which is almost degenerate with the LSP. In consequence, also the lightest stau mass is above 100 GeV. Nevertheless, significant enhancement of the $h \to \gamma\gamma$ rate, especially for $c_{24} \in (0.5, 0.6)$, is possible. Since the LSP in this region is wino-like its thermal relic abundance is much too small to explain the observed relic abundance of dark matter [109] unless non-standard cosmological history is assumed.

Even though the wino-like chargino is very light in this scenario it is very challenging to discover it at the LHC because it decays predominantly to the LSP and a very soft pion [110]. It is also rather difficult, but not impossible, to probe this scenario with

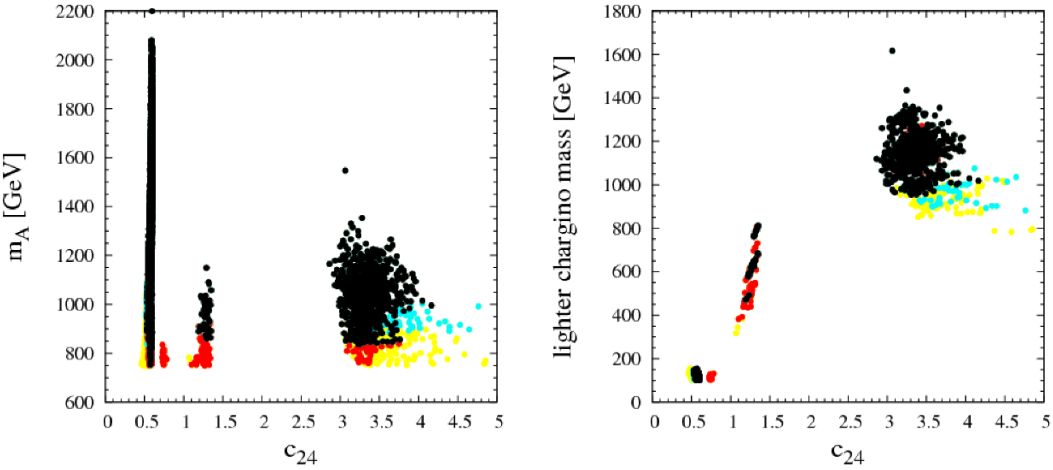

Figure 6. Scatter plots of m_A (left panel) and the lighter chargino mass (right panel) versus c_{24} for the points with $R < 1.1$ and $R_{\gamma\gamma} > 1.1$. The colour coding is the same as in figure 1.

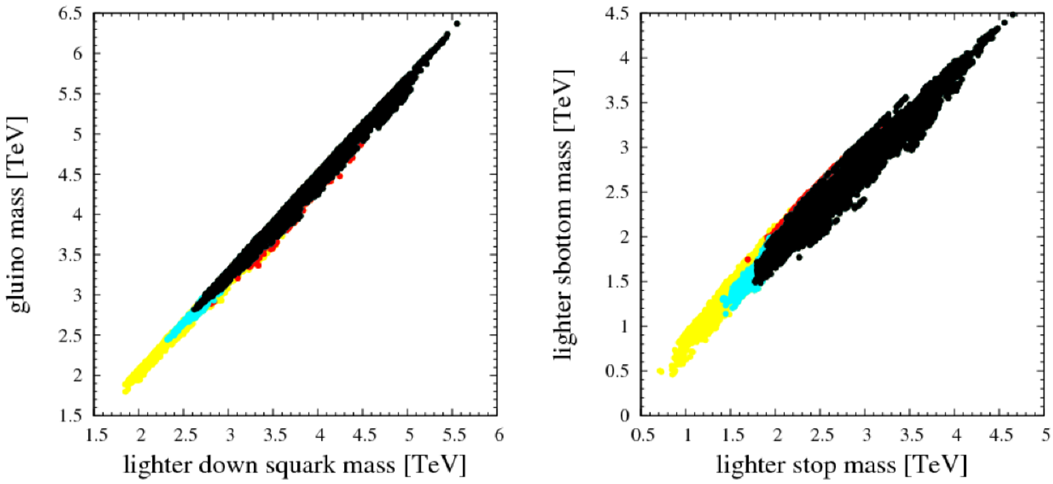

Figure 7. Gluino mass versus the right-handed down quark mass (left panel) and the lighter sbottom mass versus the lighter stop mass (right panel) for the points with $R < 1.1$, $R_{\gamma\gamma} > 1.1$ and $c_{24} < 1$. The colour coding is the same as in figure 1.

the LHC searches for coloured sparticles because the gluino mass is above 2.8 TeV while the right-handed down squark, which is the lightest first-generation squark, have mass at least 2.6 TeV, as can be seen from the left panel of figure 7. Such SUSY states could be discovered at the 14 TeV LHC but only if their masses are close to the lower bounds quoted above [111]. The third-generation squarks are lighter than those from the first two generations. Nevertheless, the mass of the lightest sbottom (stop) is above about 1.5 (1.8) TeV which will require very large statistics to discover it at the LHC.

It is worth pointing out that constraints from $b \to s\gamma$ should not be used for an ultimate exclusion of a given MSSM model. This is because the MSSM prediction for BR$(b \to s\gamma)$ can be easily affected by a flavour violating gluino contribution, see e.g. [112–114]. It

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was recently argued in [114] that even if BR($b \to s\gamma$) calculated assuming minimal flavour violation disagrees with the experimental result the flavour violating gluino contribution can brought it in agreement with the experimental data without large fine-tuning. Admitting such additional contributions the blue points in our figures should be considered allowed. In such a case the lower bound for the gluino (the right-handed down squark) mass is reduced to 2.4 (2.3) TeV so should be probed with $\mathcal{O}(100\ \mathrm{fb}^{-1})$ at the 14 TeV LHC. The third-generation squarks can also be lighter in such a case. The lower bound on the lightest sbottom (stop) reduces to 1.1 (1.4) TeV making them accessible in the early stage of the LHC run after the energy upgrade.

The most promising signature of the model at the LHC is a light pseudoscalar Higgs which can be arbitrarily close to the present experimental lower bound on m_A. Interestingly, one branch of solutions with $c_{24} \in (0.7, 0.8)$ have also quite strong upper bound on m_A of about 900 GeV which will be entirely probed in the very early stage of the 14 TeV LHC operation. However, this branch of solutions is incompatible with the recent measurement of BR($B_s \to \mu^+\mu^-$).

4.3.2 Non-singlet F-term domination

In the region of the parameter space with $c_{24} > 1$, corresponding to the gluino mass generated mainly by the non-singlet F-term, the LSP is bino-like which, in contrast to the wino-like LSP, can be below 100 GeV without violating the collider constraints. Therefore, in this case the lightest stau can also be below 100 GeV and the lower limit on its mass from LEP is about 82-90 GeV (depending on the stau mixing angle and the mass difference with the LSP). That is why the $h \to \gamma\gamma$ rate can be somewhat larger than for the $c_{24} < 1$.

The lightest chargino is also wino-like when $c_{24} > 1$ but it is significantly heavier than in the $c_{24} < 1$ case, as seen from the right panel of figure 6. For $c_{24} \in (3, 4)$ it is in the $1 \div 1.5$ TeV range, while for $c_{24} \in (1.2, 1.4)$ it is lighter (due to larger cancellation between the singlet and non-singlet F-term contributions), between about 400 and 800 GeV. The latter region could be, in principle, interesting from the point of view of the LHC phenomenology. This is because the lower mass limits for wino decaying via an on-shell stau, which is the dominant decay channel in our case, reach about 300 GeV in certain circumstances [115, 116]. However, if one demands that the upper bound on Ω_{LSP} is satisfied than the mass splitting between the stau and the LSP is below about 10 GeV so taus produced by decays of intermediate staus are soft and very hard to detect. In consequence, for mass splitting allowing efficient stau co-annihilation the LHC does not provide any constraints on the wino mass.

Searches for staus are even more challenging because the production cross-section is much smaller than the wino production cross-section and taus resulting from stau decays are also very soft.

The squarks of the first two generations and the gluino in the $c_{24} > 1$ case are also rather heavy. The gluino mass and the right-handed down squark are heavier than about 2.7 TeV, as seen from the left panel of figure 8. However, the third-generation squarks are somewhat lighter. The lighter stop mass is above 1.4 TeV. The (mostly right-handed)

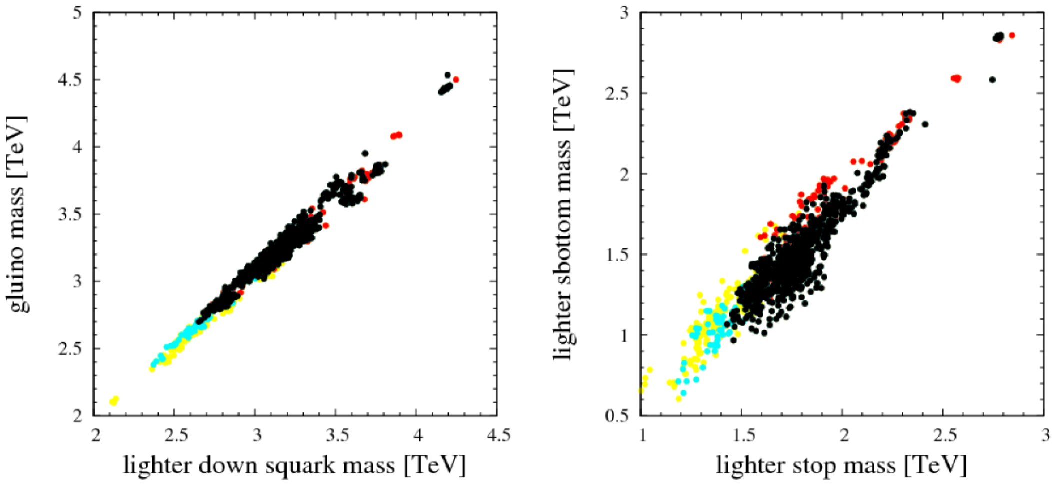

Figure 8. The same as in figure 7 but for $c_{24} > 1$.

sbottom can be as light as 1 TeV so not far away from the present experimental lower limit of about 600 GeV [117].

Relaxing the constraint from $b \to s\gamma$ reduces the lower bound on the gluino and right-handed down squark masses to 2.4 TeV. The lower limit on the stop mass is reduced in such a case to 1.2 TeV. What is the most interesting, the lower limit on the sbottom mass becomes about 600 GeV so the sbottom can be around the corner.

A very interesting prediction of the model with the non-singlet F-term domination is that the CP-odd Higgs is relatively light. It can be seen from the left panel of figure 6 that m_A is below about 1.6 TeV. It was recently shown in ref. [118] that majority of pMSSM points with $\tan\beta \sim 45$, which is a typical value for points with top-bottom-tau Yukawa unification, and m_A about 1.5 TeV can be excluded with 150 fb^{-1} of data at the 14 TeV LHC.[14] Points that could avoid exclusion are characterized by very large SUSY threshold correction to the bottom mass or by substantial branching ratio of the heavy Higgses to SUSY particles. The former case does not apply to our model since the condition of Yukawa unification implies that the SUSY threshold correction to the bottom mass can not be large (see subsection 3.1). We also checked with SUSYHIT [119–121] that in this scenario the branching ratio of H and A to stau pairs is always below ten percent (H/A decays to neutralinos are completely negligible). In addition, the 14 TeV LHC is expected to deliver much more luminosity than 150 fb^{-1} (used in the study of ref. [118]), see e.g. ref. [122]. Therefore, it is likely that the whole region of parameter space with the non-singlet F-term domination (or at least large part of it) can be ruled out by the heavy MSSM Higgs searches at the 14 TeV LHC.

It should be emphasized that the lightness of the CP-odd Higgs is tightly connected with the assumption of the enhanced $h \to \gamma\gamma$ rate. This is confirmed by the analysis performed in ref. [46] where the case when only **24** F-term contributes to the gaugino masses (which corresponds to the limiting case $c_{24} \to \infty$ of the present model) was considered and

[14] We would like to thank Ian Lewis for turning our attention to ref. [118].

enhancement of the $h \to \gamma\gamma$ rate was not required. It was shown that in such a case there is no sharp prediction for m_A which could be in the multi-TeV range, well outside of the LHC reach.

A qualitative argument for the prediction of small m_A is following. For large $\tan\beta$ and after imposing REWSB m_A^2 is well approximated by $m_{H_d}^2 - m_{H_u}^2$. RG contribution to $m_{H_d}^2 - m_{H_u}^2$ from gauginos is $\mathcal{O}(0.1)M_3^2$, while the contributions from a universal soft scalar masses and soft trilinear terms are negative, see e.g. ref. [46]. This imply $m_A \lesssim 0.1 m_{\tilde{g}}$ with the inequality saturated in the limit $M_3 \gg m_0, A_0$. Too light pseudoscalar Higgs can be avoided with the help of the D-term contribution, cf. eq. (3.1), since $D > 0$ gives positive contribution to $m_{H_d}^2 - m_{H_u}^2$. However, the requirement of light staus to enhance $h \to \gamma\gamma$ rate and rather heavy stops (to account for the observed Higgs mass) implies that the positive RGE contribution from gauginos to the stop masses should be large as compared to m_{16} and the negative RGE contribution proportional to the Yukawa couplings. This happens only if m_{16}/M_3 is not too large. Since $D/m_{16}^2 < 1/3$, in order to have positive soft mass squared at the GUT scale, the D-term contribution to $m_{H_d}^2 - m_{H_u}^2$, hence also to m_A^2, is also constrained from above by the requirement of large splitting between the stau and stop masses.

5 Conclusions

We studied enhanced $h \to \gamma\gamma$ rate induced by light staus with a strong left-right mixing. We found that the requirement of a substantial enhancement of the $h \to \gamma\gamma$ rate leads to strong constraints on the gaugino masses at the GUT scale: $|M_1/M_3| \lesssim 1/2$ and/or $|M_2/M_3| \lesssim 1/4$. This constraint follows from the requirement of a neutral LSP, the LHC limits on the gluino mass and the one-loop RGE prediction for the low-energy gaugino masses. Therefore, it is applicable not only in the (R-parity conserving) MSSM but also in many of its extensions such as the NMSSM.

We made a successful attempt to accommodate the MSSM spectrum with strongly-mixed staus inducing enhanced $h \to \gamma\gamma$ rate in SO(10) models predicting top-bottom-tau Yukawa unification. We argued that a substantial enhancement of the $h \to \gamma\gamma$ rate is possible only for the negative sign of μ, so future measurements of the $h \to \gamma\gamma$ rate may discriminate between models with different signs of μ. Assuming the D-term splitting of scalar masses, we identified patterns of the gaugino masses that allow for top-bottom-tau Yukawa unification and enhanced $h \to \gamma\gamma$ rate. These patterns can be accommodated in well-motivated models of SUSY breaking such as mirage mediation or gravity mediation with the SUSY breaking F-term that is a mixture of the singlet and non-singlet representations of SO(10).

We investigated in detail a particular scenario in which the gaugino masses are generated by a combination of the singlet F-term and the F-term in $\mathbf{24} \subset \mathbf{54}$ of $SU(5) \subset SO(10)$. We found that the $h \to \gamma\gamma$ rate in this scenario can be enhanced by more than 30% in agreement with the phenomenological constraints, including vacuum metastability bounds. In order to account for enhanced $h \to \gamma\gamma$ rate and top-bottom-tau Yukawa unification the singlet and non-singlet F-term contributions to the gluino mass should be of the same or-

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der. Nevertheless, these contributions can differ by a factor of a few and only some ratios of these contributions can be consistent with the experimental constraints.

There are some phenomenological differences between models depending on whether the gluino mass arises mainly from the singlet or the non-singlet F-term. If the singlet F-term dominates, the LSP is wino-like with too small thermal relic abundance to account for the observed energy density of dark matter. On the other hand, if the non-singlet F-term dominates the LSP is bino-like and its thermal relic abundance can be brought to cosmologically acceptable values due to efficient coannihilation with staus.

In any case the resulting spectrum of coloured sparticles is rather heavy. The lower bounds on most of them are only slightly below 3 TeV so it will require a lot of data at the 14 TeV LHC to start to probe them. The exception is the lightest sbottom which in the non-singlet F-term domination case may have mass around 1 TeV. The best prospects for testing the model is due to the pseudoscalar Higgs which generically has mass close to the present experimental lower bound.

It is interesting to note that in spite of the correlation between the soft masses for squarks and sleptons at the GUT scale dictated by the structure of SO(10) GUTs, in the low-energy spectrum very light stau, with mass $\mathcal{O}(100\text{ GeV})$ to account for enhanced $h \to \gamma\gamma$ rate, and heavy squarks can be simultaneously present. This is possible due to RG effect of gluino which contribute substantially to squark masses without affecting the slepton masses. The mass splitting between staus and the first two generation squarks is additionally enhanced due to large tau Yukawa coupling which reduces the stau masses via RGEs.

We also found that the $b \to s\gamma$ and $B_s \to \mu^+\mu^-$ constraints push the supersymmetric spectrum up in a significant way. If these are neglected, due to non-minimally-flavour-violating contributions, most of the coloured sparticles could be accessible in the early stage of the next LHC run.

The only observable that cannot be fitted in this model is $(g-2)_\mu$ since slepton masses of the second generation are much heavier than staus because only the latter acquire a negative RGE contribution proportional to large Yukawa coupling. In principle, the $(g-2)_\mu$ anomaly could be explained in SO(10) models in the framework of the NMSSM since at large $\tan\beta$ additional contributions to the Higgs mass due to mixing with the singlet [123] allow to lower the scale of superpartner masses. We plan to investigate this issue in the future.

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