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Perturbative λ-supersymmetry and small κ-phenomenology

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Abstract For minimal λ-supersymmetry, the theory stays perturbative to the GUT scale for $\lambda \leq 0.7$. This upper bound is relaxed when one either takes the criterion that all couplings are close to $\sim 4\pi$ for non-perturbation or one allows for new fields at the intermediate scale between the weak and GUT scale. We show that a hidden $U(1)_X$ gauge sector with spontaneously broken scale ~ 10 TeV improves this bound as $\lambda \leq 1.23$ instead. This may induce significant effects on the Higgs physics, such as decreasing fine tuning involving the Higgs scalar mass, as well as on the small κ-phenomenology.

1 Introduction

The standard model (SM)-like Higgs boson with mass around 126 GeV [1,2] reported at the LHC strongly favors an extension beyond the minimal supersymmetric model (MSSM). A simple idea is adding a SM singlet S to the MSSM, with superpotential

$$W = \lambda \, SH_u H_d + \frac{\kappa}{3} S^3. \tag{1}$$

Due to the Yukawa coupling λ in Eq. (1), the Higgs boson mass can be naturally increased to the observed value for a moderate value $\lambda \geq 0.5$ at the electroweak (EW) scale. This specific extension is referred as the next-to-minimal supersymmetric model (NMSSM) [3] for $\lambda \leq 0.7$, or λ-supersymmetry (λ-SUSY) [4] for $0.7 \leq \lambda \leq 2$. Either the NMSSM or λ-SUSY is very attractive from the viewpoint of the naturalness argument [5,6], as the stringent tension on the stop masses required by the Higgs mass in the MSSM is dramatically reduced.

The origin of the upper bound on λ above can be seen from the beta function β_λ for λ. Since β_λ is dominated by the top Yukawa coupling y_t, the sign of which is positive, it imposes an upper bound on the EW value of λ in order to be still in the perturbative region at the high energy scale. Given this scale

near the grand unification (GUT) scale $\simeq 1.0 \times 10^{15}$ GeV, the critical value was found to be ~ 0.7 [7] in terms of the one-loop beta function,[1]

$$\beta_\lambda = \frac{\lambda}{16\pi^2} \left[4\lambda^2 + 3y_t^2 + 2\kappa^2 - 3g_2^2 - \frac{3}{5} g_1^2 \right], \tag{2}$$

where $t \equiv \ln(\mu/\mu_0)$, μ being the running renormalization scale and $\mu_0 \equiv 1$ TeV. On the other hand, the observed Higgs mass favors a large value of λ [9].

The intuition for solving the tension between the observed Higgs mass and perturbativity is obvious in the context of SUSY. The main observation is that the Yukawa couplings in the superpotential receive only wave-function induced renormalizations due to the protection of SUSY, in contrast to non-supersymmetric case.[2] By the fact that the beta functions of the Yukawa couplings are related to the anomalous dimensions of the chiral fields [8], and the anomalous dimensions are proportional to quadratic Yukawa couplings with positive coefficients, the sign of the contribution to the β function due to the Yukawa interactions is thus always positive. A way to alleviate this is introducing hidden super-confining gauge dynamics [14,15], from which the Yukawa couplings are asymptotically free, and S is a composite other than a fundamental state, with large λ at low energy as a result of this asymptotical freedom. In this context the Higgs doublets can be either fundamental [14] or composite states [15].

In this paper, we will explore λ-SUSY that stays perturbative up to the GUT scale in an alternative way. In this framework we will obtain a well-defined λ-SUSY in the realm of

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[1] In this paper, we follow the convention in [8], and present our β functions in the $\overline{DR}$ scheme.

[2] The phenomenon of Yukawa coupling being asymptotically free has been exposed in $\lambda\phi^4$ scalar field theory with dimension lower than four by Wilson and Fisher [10], and in the well-known nonlinear sigma model [11] together with an attempt to obtain a version of four-dimensional supersymmetry [12]. A similar phenomenon was also addressed by the authors of [13] in four-dimensional scalar field theory with nonpolynomial potentials.

perturbative analysis, and operators such as Higgs doublets and singlet S are all fundamental rather than composite states. This is the main motivation for this study.

The study of well-defined λ-SUSY was initially addressed in Ref. [7], in which it was found that for $\tan\beta$ smaller than ~ 10, λ is the first Yukawa coupling running into the non-perturbative region at the high energy scale. It was also understood that either introducing new matter fields at the intermediate scale between the weak and GUT scale or adopting a smaller EW value $\kappa(\mu_0)$ can decrease the evolution rate for λ. Given an initial EW value $\kappa(\mu_0)$, one can derive the upper bound on $\lambda(\mu_0)$, or vice versa, once new matters which appear at the intermediate scale are identified explicitly. The minimal content of the new fields only includes the messenger sector which communicates the SUSY-breaking effect to the visible sector, namely the λ-SUSY.

Together with the messengers a spontaneously broken $U(1)_X$ gauge group will be considered as the new fields. We consider a class of models different from those studied in the literature [16], in which SM fermions and sfermions except the Higgs doublets and S are all charged under the hidden $U(1)_X$. The abelian gauge coupling g_X and the $U(1)_X$-breaking scale M_X enter into the parameter space as two new free parameters. In comparison with the traditional NMSSM or λ-SUSY, in our model the beta function for top Yukawa coupling β_t receives a new and negative contribution due to the hidden $U(1)_X$ sector, the magnitude of which is determined by the hidden gauge coupling g_X and scale M_X. As long as g_X is large enough, but still valid for perturbative analysis, these new effects decrease the slope of λ as a function of the renormalization scale μ above the scale M_X, and therefore leads to a larger critical value $\lambda(\mu_0)$.

The paper is organized as follows. In Sect. 2, in terms of one-loop renormalization group equations (RGEs) for the relevant coupling constants we estimate the critical values of $\kappa(\mu_0)$ and $\lambda(\mu_0)$ due to the hidden $U(1)_X$ effect. In Sect. 3, we revise the phenomenology for large λ but small κ. Finally we conclude in Sect. 4. In Appendix A, we show the details of the hidden sector and briefly review collider constraints on the model parameters.

2 Perturbative λ-SUSY

As mentioned in Sect. 1, the beta function for the Yukawa couplings of the superpotential can easily be estimated by using the non-renormalization of the superpotential. Firstly, in the case without hidden $U(1)_X$ gauge group the one-loop beta functions for the Yukawa couplings in λ-SUSY are given by[3]

$$\beta_\lambda = \frac{\lambda}{16\pi^2}\left[4\lambda^2 + 2\kappa^2 + 3y_t^2 - 3g_2^2 - \frac{3}{5}g_1^2\right],$$

$$\beta_{y_t} = \frac{y_t}{16\pi^2}\left[6y_t^2 + \lambda^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{15}g_1^2\right],$$

$$\beta_\kappa = \frac{\kappa}{16\pi^2}[6\lambda^2 + 6\kappa^2]. \tag{3}$$

The one-loop beta functions for the SM gauge couplings are the same as the MSSM,

$$\frac{\partial}{\partial t}\alpha_i^{-1} = -\frac{b_i}{2\pi}, \tag{4}$$

where $(b_1, b_2, b_3) = (33/5, 1, -3)$.

In the case with a hidden $U(1)_X$ gauge group, we study the model in which SM fermions and sfermions all carry a hidden $U(1)_X$ charge, with the spontaneously broken scale $M_X \simeq 10$ TeV. The Higgs doublets, however, are singlets of this $U(1)_X$ symmetry. Anomaly free conditions require three hidden matters $X_{1,2,3}$ with the same $U(1)_X$ charge added to the model. For details as regards the matter representations, see Appendix A. The modifications to the RGEs in Eq. (3) due to the hidden $U(1)_X$ effect are given by

$$\delta\beta_\lambda = 0,$$

$$\delta\beta_{y_t} = -\frac{y_t}{16\pi^2}\left(\frac{12}{5}a^2 \cdot g_X^2\right),$$

$$\delta\beta_\kappa = 0, \tag{5}$$

together with

$$\beta_{g_X} = \frac{g_X^3}{16\pi^2}\left[\frac{3}{5}a^2 \cdot \left(\frac{39}{2}n_g + 3\right)\right], \tag{6}$$

where $n_g = 3$ denotes the number of SM fermion generations. It is obvious that the large EW value $g_X(\mu_0)$ favors small $|a|$.

The RG running for g_X with different values at μ_0 is plotted in Fig. 1. In this figure the non-perturbative region lies below the red line. It shows the critical value $g_X(\mu_0) \leq \{2.0, 1.18, 0.82\}$ for $a = \{\frac{1}{6}, \frac{1}{3}, \frac{1}{2}\}$, respectively. Given the upper bound on $g_X(\mu_0)$, the upper bound on the Yukawa coupling λ can be estimated in terms of RGEs in Eqs. (5) and (3). The input parameters related to the critical value $\lambda(\mu_0)$ are composed of

$$\{y_t(\mu_0), \kappa(\mu_0), g_X(\mu_0), M, n_{5\bar{5}}\}, \tag{7}$$

where M denotes the messenger mass scale, $n_{5\bar{5}}$ represents the number of $\bar{5} + 5$ representations of $SU(5)$ for messengers. We will use the updated pole mass of top quark $m_t = 174$ GeV [18] instead of $m_t = 180$ GeV in [7] for our analysis, and we take the criterion that the theory enters

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[3] The value of $\tan\beta$ defined as the ratio $\langle H_u^0\rangle/\langle H_d^0\rangle$ is required to be smaller than ~ 10 in light of the 126 Higgs mass in the NMSSM or

Footnote 3 continued

λ-SUSY. In [17] it has been shown that experimental constraints require $\tan\beta \leq 10$ for $\lambda \geq 1.0$. Thus, it is a good approximation to ignore the bottom Yukawa coupling in comparison with the top Yukawa coupling.
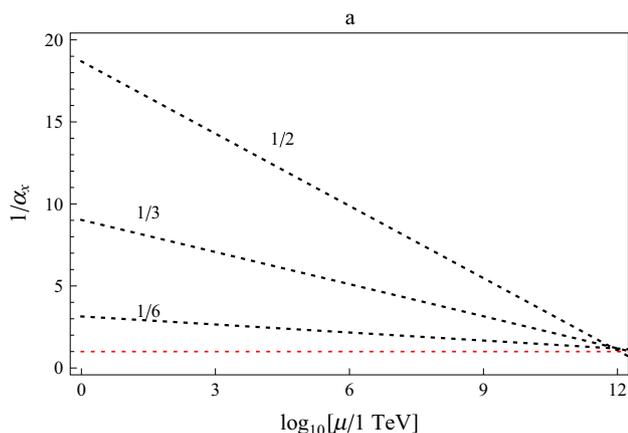
Fig. 1 Critical value $g_X(\mu_0)$ for $a = \left\{\frac{1}{6}, \frac{1}{3}, \frac{1}{2}\right\}$, respectively

into the non-perturbative region when any of the coupling constants in the theory is larger than $\sim 4\pi$.

We want to emphasize that (i) below the scale M_X, the modifications to RGEs arising from $U(1)_X$ can be ignored, and (ii) above the scale of M, the coefficients b_i in the beta functions for the SM gauge couplings should change as $b_i \rightarrow b_i + n_{5\bar{5}}$. Figure 2 shows the modifications to the upper bound on $\kappa(\mu_0)$ due to the hidden $U(1)_X$ sector, given fixed $\lambda(\mu_0)$. We choose the initial EW value $\lambda(\mu_0) = 1.0$, $g_X(\mu_0) = 0.82$ for $a = 1/2$, and the messenger parameters $M = 10^7$ GeV and $n_{5\bar{5}} = \{1, 4\}$ for illustration. Given the critical value $\kappa(\mu_0)$, the solid line in gray (green) represents the RG running for λ^{-1} without (with) $U(1)_X$ effect. The RG runnings of κ^{-1} (dotted) and y_t^{-1} (dot-dashed) shown in the figure imply that the upper bound $\kappa_0 \leq 0.8$ is improved to $\kappa_0 \leq 0.85 \sim 0.86$ when the hidden $U(1)_X$ effect is taken into account. Alternatively, given the same $\kappa(\mu_0)$, the upper bound on $\lambda(\mu_0)$ can be improved by the $U(1)_X$ effect. Combination of the left and

right panel shows that the deviation due to the change of $n_{5\bar{5}}$ is small, in comparison with the $U(1)_X$ effect. Note that the plots in Fig. 1 are actually critical lines, because the change from the perturbative to the non-perturbative region is rather abrupt.

In Table 1, we show the improvement on the upper bound on $\lambda(\mu_0)$ due to the hidden $U(1)_X$. Without $U(1)_X$ sector, the model is not well defined up to the GUT scale when $\lambda > 0.80$. Instead, $\lambda \leq 1.23$ stays perturbative up to the GUT scale when the hidden $U(1)_X$ is added to the model. The critical value on λ may be modified due to different choices of the parameters $\{n_{5\bar{5}}, M, a\}$. As shown in Fig. 2, the deviation to $\lambda(\mu_0)$ due to messenger parameters M and $n_{5\bar{5}}$ is rather small. A similar conclusion holds when one tunes a. This is because an a larger than $1/2$ leads to a larger contribution to $\delta\beta_{y_t}$ but unfortunately a smaller $g_X(\mu_0)$, and vice versa.

In summary, in terms of introducing a hidden $U(1)_X$ sector with gauge symmetry broken scale $\simeq 10$ TeV a well-defined λ-SUSY up to the GUT scale is allowed for $\lambda \leq 1.23$. In this class of models all fields including the singlet S and the Higgs doublets are fundamental. The implication for $\lambda \sim 1$–2 has been addressed, e.g., in [9]. In the next section, we revise the phenomenological implication for such large λ together with small κ.

Table 1 Upper bound on $\lambda(\mu_0)$ for $g_X(\mu_0) = 0.82$ for $a = 1/2$, messenger parameters $M = 10^7$ GeV and $n_{5\bar{5}} = 1$, and different initial values of $\kappa(\mu_0)$. In contrast to the result $\lambda(\mu_0) \leq 0.8$ for $\kappa(\mu_0) \simeq 0$ in [7], we have $\lambda(\mu_0) \leq 1.23$ instead for the case with the hidden $U(1)_X$

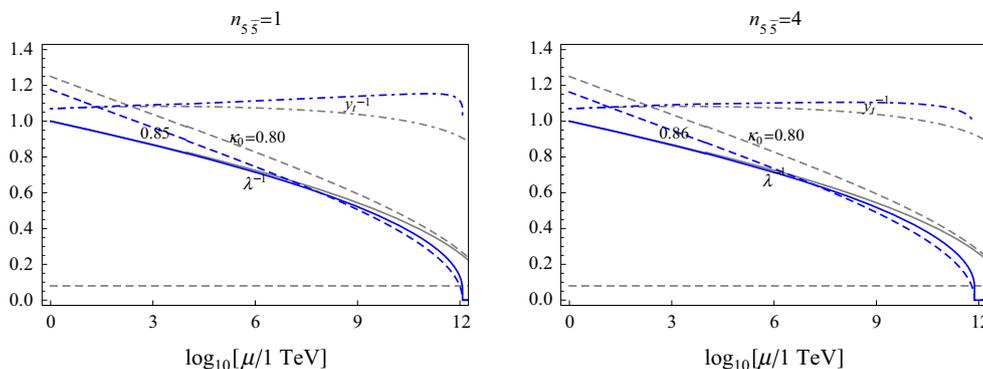| $\lambda(1$ TeV$)$	Without $U(1)_X$	With $U(1)_X$
1.0	$\kappa(\mu_0) \leq 0.80$	$\kappa(\mu_0) \leq 0.85$
1.10	$\kappa(\mu_0) = 0$	$\kappa(\mu_0) \leq 0.81$
1.23	Not well defined	$\kappa(\mu_0) = 0$

Fig. 2 RG running of λ^{-1} (*solid line*) for initial EW value $\lambda(\mu_0) = 1.0$ and $g_X(\mu_0) = 0.82$ for $a = 1/2$, and messenger parameters $M = 10^7$ GeV and $n_{5\bar{5}} = 1$ (*left*) and $n_{5\bar{5}} = 4$ (*right*), respectively. The *gray* and *green colors* correspond to λ-SUSY without and with $U(1)_X$, respectively. The RG runnings of κ^{-1} (*dotted*) and y_t^{-1} (*dot dashed*)

shown in the figure indicate that the upper bound $\kappa_0 \leq 0.8$ is improved to $\kappa_0 \leq 0.85 \sim 0.86$ when the hidden $U(1)_X$ effect is taken into account. The *horizontal line* refers to the critical point between the perturbative and the non-perturbative region

3 PQ symmetry and small κ-phenomenology

This section is devoted to the study of the phenomenology in λ-SUSY with small κ. Although it is a consequence of taking a large λ for well-defined λ-SUSY, the study of small κ-phenomenology can be considered as an independent subject from the viewpoint of phenomenology. The smallness of κ can be understood in terms of a broken $U(1)$ global symmetry, i.e., the Peccei–Quinn (PQ) symmetry. Without κ, the model is invariant under the following $U(1)$ symmetry transformation:

$$H_{\rm u} \to H_{\rm u} \exp(i\phi), \;\; H_{\rm d} \to H_{\rm d} \exp(i\phi), \;\; S \to S \exp(-2i\phi). \tag{8}$$

Terms like $\delta V = m^2 S^2 + B_\mu H_{\rm u} H_{\rm d}$ explicitly break this symmetry. For this breaking to be small, we have a pseudo-Goldstone boson with small mass.

The soft terms in the scalar potential for small κ-phenomenology is given by

$$V_{\rm soft} = m_{\rm S}^2 \mid S \mid^2 + m_{H_{\rm u}}^2 H_{\rm u}^2 + m_{H_{\rm d}}^2 H_{\rm d}^2 + (A_\lambda \lambda S H_{\rm u} H_{\rm d} + {\rm H.c}), \tag{9}$$

where the κ relevant terms are ignored. The EW symmetry-breaking vacuum can be determined from Eq. (9). For details of the analysis of the EW-breaking vacuum, see, e.g., [3,19]. There are five free parameters

$$\{m_{H_{\rm u}}^2, \; m_{H_{\rm d}}^2, m_{\rm S}^2, \; A_\lambda, \; \lambda\}, \tag{10}$$

which define the small κ-phenomenology. The first two can be traded for $\upsilon = 174$ GeV and $\tan\beta$. In what follows we explore the constraints on these parameters, the Higgs scalar spectrum, and their couplings to SM particles.

The 3×3 squared mass matrix of CP-even neutral scalars reads[4]

LHC with $\mathcal{M}_{22}^2 = (126\,{\rm GeV})^2$, and s decouples from the SM gauge bosons with mass $\lambda\upsilon \sim 209$ GeV for $\lambda = 1.2$. Such a scalar s easily escapes searches performed by the LEP2 and Run-I of LHC.

The precision measurement of h including its couplings to SM gauge bosons and fermions powerfully constrains the magnitude of the mixing effect. After mixing, we define the mass eigenstates as h_1 and h_2, where $h_1 = \cos\theta h - \sin\theta s$, $h_2 = \cos\theta s + \sin\theta h$. Normalized to the SM Higgs boson couplings, h_1 and h_2 couple to the SM particles as

$$\xi_{h_1 VV} = \frac{g_{h_1 VV}^2}{g_{hVV}^2} = \cos^2\theta, \;\; \xi_{h_2 VV} = \frac{g_{h_2 VV}^2}{g_{hVV}^2} = \sin^2\theta, \tag{12}$$

where $V = \{W, Z, t, b, \ldots\}$. In the present status, LHC data suggests that $0.96 \leq \cos^2\theta \leq 1$ [20]. Alternatively we have $\sin^2\theta \leq 0.04$.

In Fig. 3 we show the parameter space in the two-parameter plane of $m_{\rm S}$ and A_λ, for $\lambda = 1.2$ and $\tan\beta = \{4.5, 5, 5.5\}$, respectively. For each $\tan\beta$ region below the color contour is excluded by the condition of stability of the potential and the mass bound on the chargino mass $m_\chi > 103$ GeV [21–23]. The region in the right up corner is excluded by the precision measurement of Higgs coupling presently. In particular, $m_{\rm S}$ values heavier than $\sim$90, 105, and 110 GeV are excluded for $\tan\beta = 4.5$, 5, and 5.5, respectively. In comparison with the choice $\lambda = 0.7$ discussed in [19], the main difference is that for $\lambda = 1.2$ it allows for larger $m_{\rm S}$.

As for the Higgs mass constraint, the discrepancy between $\mathcal{M}_{22}^2$ and 126 GeV is compensated by the stop induced loop correction. The stop mass for the zero mixing effect (i.e., $A_{\rm t} \sim 0$) is $\sim$340 GeV for $\tan\beta = 4.5$, $\sim$550 GeV for $\tan\beta = 5$, and $\sim$800 GeV for $\tan\beta = 5.5$. The stop mass below 1 TeV is favored by naturalness. Nevertheless, there is tension for

$$\mathcal{M}^2 = \begin{pmatrix} \frac{A_\lambda^2}{1+x} + (M_Z^2 - \lambda^2\upsilon^2)\sin^2 2\beta & -\frac{1}{2}(M_Z^2 - \lambda^2\upsilon^2)\sin 4\beta & -A_\lambda \lambda\upsilon \cos 2\beta \\ * & M_Z^2 \cos^2 2\beta + \lambda^2\upsilon^2 \sin^2 2\beta & -A_\lambda \lambda\upsilon \sin 2\beta \frac{x}{1+x} \\ * & * & \lambda^2\upsilon^2(1+x) \end{pmatrix} \tag{11}$$

in the basis (H, h, s), where $x \equiv m_{\rm S}^2/(\lambda\upsilon)^2$, $H = \cos\beta h_2 - \sin\beta h_1$, and $h = \cos\beta h_1 + \sin\beta h_2$ under the compositions $H_{\rm u}^0 = \frac{1}{\sqrt{2}}(h_1 + i\pi_1)$, $H_{\rm d}^0 = \frac{1}{\sqrt{2}}(h_2 + i\pi_2)$, and $S = \frac{1}{\sqrt{2}}(s + i\pi_{\rm s})$. The state h mixes with the other two states H and s via $\mathcal{M}_{12}^2$ and $\mathcal{M}_{23}^2$, respectively. Due to the smallness of $\mathcal{M}_{12}^2$, state h mixes dominantly with s. Without such mixing (i.e., $\mathcal{M}_{23}^2 \to 0$), h couples to the SM gauge bosons and fermions exactly as the SM Higgs. This implies that in the parameter space of small $m_{\rm S}$, h is the SM-like scalar discovered at the

such a light stop with present LHC data. This problem can be resolved in some situation, which we will not discuss here.

We show in Fig. 4 the ratio $\xi_{h_2 VV}$ defined in Eq. (12), which determines the production rate for scalar h_2. Its magnitude increases slowly as $m_{\rm S}$ becomes larger. $\xi_{h_2 VV}$ reaches $\sim$0.03 at most when $m_{\rm S}$ closes in to its upper bound (suggested by Fig. 3). For such a strength of the coupling and the mass $\sim$200 GeV, h_2 is easily out of reach of Run-I at the LHC and earlier attempts at the LEP2. A small ratio of the strength of the coupling similarly holds for the heaviest CP-even neutral state H. In this sense, it is probably more

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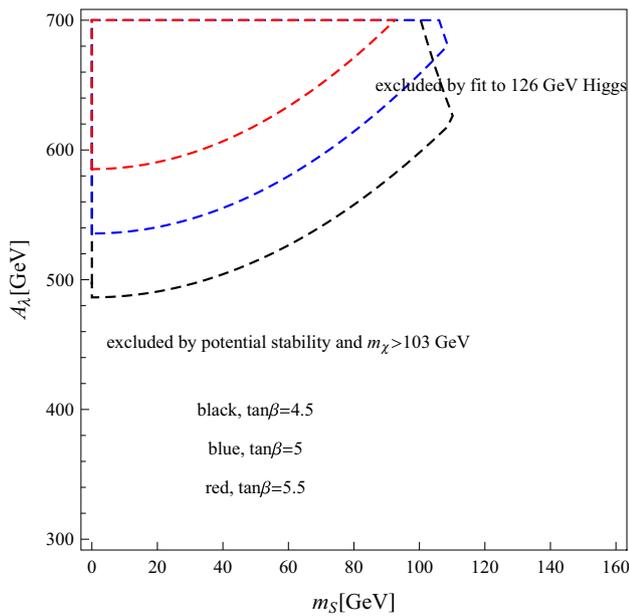[4] Here we follow the conventions and notation in Ref. [19].

Fig. 3 Parameter space in the two-parameter plane of m_S and A_λ, for $\lambda = 1.2$ and $\tan\beta = 4.5$ (*black*), 5 (*blue*), 5.5 (*red*), respectively. For each $\tan\beta$ region below the *color contour* is excluded by the condition of stability of potential and mass bound on chargino mass $m_\chi > 103$ GeV. The region in the *right up corner* is excluded by the fit to the Higgs couplings presently
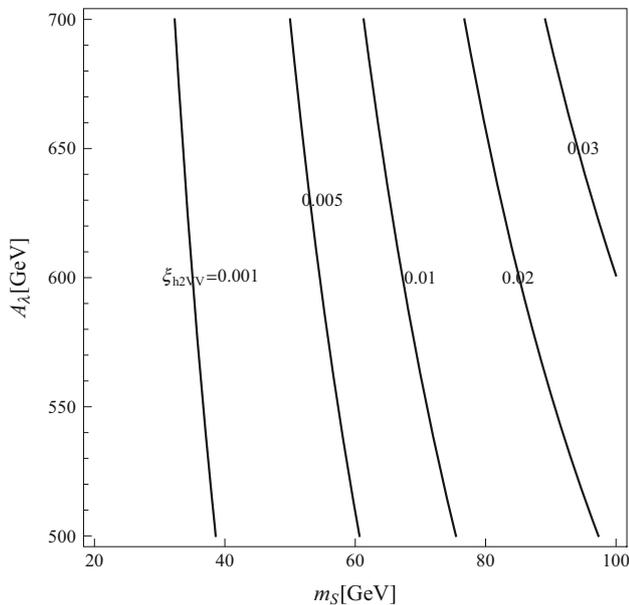

Fig. 4 The ratio of the coupling strength for h_2 shown in the plane of m_S–A_λ for $\lambda = 1.2$ and $\tan\beta = 4.5$

efficient to probe a charged Higgs scalar $H^\pm$, a CP-odd scalar A, or a light pseudo-boson G. Studies along this line can be found in, e.g., [24].

As for the other scalar masses: we show them in Table 2 for three sets of A_λ and m_S, which are chosen in the parameter space shown in Fig. 3. It is shown that for each case the

Table 2 Mass spectrum for three sets of choices, which are subtracted from Fig. 3

(A_λ, m_S) (GeV)	$H^\pm$ (GeV)	H (GeV)	A (GeV)
(500, 40)	463	482	535
(560, 50)	527	530	585
(620, 50)	590	586	640

mass of the charged Higgs boson exceeds its experimental bound $\sim$300 GeV, and A scalar is always the heaviest with mass around 600 GeV. In comparison with the scalar mass spectrum in the fat Higgs model [15], the spectrum in Table 2 is similar to it, although their high energy completions are rather different. As a final note we want to mention that the smallness of m_S compared with A_λ can be achieved in model building, e.g., in gauge mediation. This is because the singlet S only couples to messengers indirectly through Higgs doublets.

4 Conclusions

In this paper, we have studied λ-SUSY, which stays perturbative up to the GUT scale. We find that the bound $\lambda \leq 0.7 \sim 0.8$ in the minimal model is relaxed to $\lambda \leq 1.23$ if a hidden $U(1)_X$ gauge theory is introduced above the scale $\sim$10 TeV and a small $\kappa(\mu_0)$ is assumed at the same time. This improvement gives rise to several interesting consequences in phenomenology. For example, the fine tuning related to the 126 Higgs mass can be reduced, and the light stop below 1 TeV can be allowed. In the light of such a single $U(1)_X$, one may introduce multiple $U(1)$s or other gauge sectors at the intermediate scale, and further increase the bound on λ.

In the second part of the paper, we have revised the small κ-phenomenology for large λ. In comparison with the fat Higgs model [15], the spectrum in the small κ-phenomenology is similar, although their high energy completions are different. The null result for signals of the other two CP-even neutral scalars h_2 and H is due to the perfect match between the scalar discovered at the LHC (here it is referred to as h_1) and the SM Higgs. Because the perfect fit dramatically reduces the mixing effect between h_1 and the others, which results in a tiny strength of the coupling for h_2 and H relative to the SM expectation. The studies on the signals of the charged scalar $H^\pm$, the CP-odd scalar A, and the pseudo-boson G will shed light on this type of model.

Appendix A: The hidden $U(1)_X$ sector

Realistic model. The model that we are going to study for the case with hidden $U(1)_X$ gauge symmetry is presented in Table 3. The $U(1)_X$ charges in this table must satisfy the gauge anomaly free conditions and are consistent with the superpotential of the visible sector, $W \sim y_u Q_i \bar{u}_i H_u + y_d Q_i \bar{d}_i H_d + y_e L_i \bar{e}_i H_d + \mu H_u H_d$. We restrict ourselves to the case in which the Higgs doublets $H_{u,d}$ are singlets of $U(1)_X$. Without the three hidden matters $X_{1,2,3}$ with the same $U(1)_X$ charge anomaly free conditions such as $U(1)_X$–graviton–graviton and $U(1)_X$–$U(1)_X$–$U(1)_X$ cannot be satisfied.[5]

Symmetry breaking. There is no signal of Z' from the broken $U(1)_X$ gauge group yet, so it should be spontaneously broken above the weak scale for g_X of the order of the SM gauge coupling. This can be achieved in various ways. Here we simply take the gauge mediation as an example. If the hidden $U(1)_X$ gauge group communicates the D-type of SUSY-breaking effects into Xs, the sign of the soft mass squared m_X^2 would be negative [25]. For an earlier application of such a property in gauge mediation, see, e.g., [26]. Below the SUSY-breaking scale for the potential for X_i we have

$$V = -m_{X_i}^2 X_i^\dagger X_i + (X_i^\dagger X_i)^2, \tag{13}$$

which spontaneously breaks $U(1)_X$, with $U(1)_X$-breaking scale $M_X \sim m_X$. Note that the magnitude of m_X can be either larger or smaller than the soft breaking masses in the visible sector, which depends on the ratio of the D term relative to the F term and also the ratio of g_X relative to SM gauge couplings. With a D term which gives rise to effects an order of magnitude larger than the soft-breaking masses and g_X of the same order as the SM gauge coupling, one can obtain $M_X \simeq \mathcal{O}(10)$ and $m_{\tilde{f}_i} \sim \mathcal{O}(1)$ TeV in the visible sector.

Limit on $g_X(\mu_0)$. The experimental constraint on M_X and the gauge coupling g_X is obtained from direct production of Z' at colliders through leptonic decays; and also indirect searches from flavor violation and electroweak precision tests. For a review of the status of hidden $U(1)_X$, see e.g., [27,28]. The parameter M_X is constrained to be above

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[5] We thank the referee for reminding us that the $U(1)_X$ charges in this model are identical to the $B - L$ quantum numbers of the SM fields, and the charges of the "hidden" matter, $X_{1,2,3}$ are the same as those of the right-handed neutrino fields.

Table 3 Charge assignments for the matter superfields in the visible and hidden sectors. Q_i, $\bar{u}_i$, $\bar{d}_i$, L_i, and $\bar{e}_i$ denote the three-generation matters of MSSM. S is a singlet of both SM gauge group and the hidden $U(1)_X$. The $X_{1,2,3}$ are a set of hidden matters charged under the hidden $U(1)_X$, which are responsible for the $U(1)_X$ gauge symmetry breaking. a is a real number

	$SU(3) \times SU(2) \times U(1)_Y$	$U(1)_X$
Q_i	$(\mathbf{3}, \mathbf{2}, \frac{1}{6})$	a
$\bar{u}_i$	$(\bar{\mathbf{3}}, \mathbf{1}, -\frac{2}{3})$	$-a$
$\bar{d}_i$	$(\bar{\mathbf{3}}, \mathbf{1}, \frac{1}{3})$	$-a$
L_i	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$	$-3a$
$\bar{e}_i$	$(\mathbf{1}, \mathbf{1}, 1)$	$3a$
H_u	$(\mathbf{1}, \mathbf{2}, \frac{1}{2})$	0
H_d	$(\mathbf{1}, \mathbf{2}, -\frac{1}{2})$	0
$X_{1,2,3}$	$(\mathbf{1}, \mathbf{1}, 0)$	$3a$
S	$(\mathbf{1}, \mathbf{1}, 0)$	0

$1 \sim 2$ TeV for the coupling $g_X \sim 1$–2. For $M_X \simeq 10$ TeV, adopted in this note, we have the experimental limit $g_X \leq M_X/(\mid a \mid \cdot \text{a few TeV}) \sim \mid a \mid^{-1}$ [29,30].

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