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A split SUSY model from SUSY GUT

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Abstract: We propose to split the sparticle spectrum from the hierarchy between the GUT scale and the Planck scale. A split supersymmetric model, which gives non-universal gaugino masses, is built with proper high dimensional operators in the framework of SO(10) GUT. Based on a calculation of two-loop beta functions for gauge couplings (taking into account all weak scale threshold corrections), we check the gauge coupling unification and dark matter constraints (relic density and direct detections). We find that our scenario can achieve the gauge coupling unification and satisfy the dark matter constraints in some part of parameter space. We also examine the sensitivity of the future XENON1T experiment and find that the currently allowed parameter space in our scenario can be covered for a neutralino dark matter below about 1.0 TeV.

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JHEP03(2015)050

Contents

1 Introduction

The discovery of a 125 GeV Higgs boson [1, 2] by both the ATLAS and CMS collaborations has completed the Standard Model (SM). The experimental data reported so far by LHC agree quite well with the SM predictions. On the other hand, the lack of significant hints in electroweak precision tests and the absence of evidences for new particle contents at the LHC challenge many proposals of new physics beyond the SM, including weak scale supersymmetry (SUSY).

SUSY, which was regarded for a long time as one of the most appealing extensions of the SM, has many desirable features. For example, the observed 125 GeV Higgs boson falls within the narrow $115 - 135$ GeV window predicted by the minimal supersymmetric standard model (MSSM). Besides, the genuine unification of gauge couplings [3–9], which can hardly be achieved in the SM, can be successfully realized in the framework of low energy SUSY. Also, the puzzle of cosmic dark matter can naturally be explained in SUSY.

Although SUSY is appealing, null search results of sparticles at LHC suggest that either low energy SUSY needs to be tuned or sparticles are well above the weak scale. In fact, the LHC data has already set a limit [10–13] $m_{\tilde{g}} > 1.5$ TeV for $m_{\tilde{q}} \sim m_{\tilde{g}}$ and $m_{\tilde{g}} \gtrsim 1$ TeV for $m_{\tilde{q}} \gg m_{\tilde{g}}$ for certain popular CMSSM models. On the other hand, the observed mass of the Higgs boson requires rather large loop effects of top squarks in the MSSM and CMSSM, which implies some extent of fine-tuning [14, 15]. So naturalness in SUSY may be realized in a more involved way even though it was initially proposed to solve the hierarchy problem.

Split SUSY, proposed in [16–18], gives up the naturalness criterion while keeps the other two main advantages: the gauge coupling unification and viable dark matter candidates. This scenario assumes a very high scalar mass scale M_S and the low energy

JHEP03(2015)050

spectrum contains only the gauginos and higgsinos as well as a fine-tuned Higgs boson in addition to the SM sector. The SUSY CP and flavor problems can naturally be solved in this scenario due to very heavy sfermions. The latest results of the Higgs mass from the LHC measurement, together with the requirement of gauge coupling unification, suggest a scalar superpartner mass scale roughly of order $M_S \sim 100 - 1000\,\text{TeV}$ [19–21], which indicates a moderately split spectrum for split SUSY. The hierarchy of sparticle spectrum is argued to be determined by the gauge loop factor [20]. We propose in this paper to split the sparticle spectrum from the hierarchy between the GUT scale and the Planck scale.

Actually, there are various ways to split the sparticle spectrum. In this work we propose to use the generalized gravity mediation [22–33] with non-renormalizable Kähler potential and superpotential. Because GUT is one of the retained motivations of split SUSY, certain high-representation Higgs fields of GUT group could appear in the non-renormalizable Kähler potential and superpotential. Such high-representation fields could not only amend the gauge coupling unification condition at the GUT scale, but also establish new relations among various theory inputs. In our scenario we have non-universal gaugino masses at the GUT scale (for other scenarios giving non-universal gaugino masses, see [34–44]). Such non-universal gaugino masses can naturally appear with non-renormalizable Kähler potential and superpotential involving various high-representation Higgs fields.

This paper is organized as follows. In section 2 we present our model. In section 3 we check the phenomenology of our model. Based on a calculation of two-loop beta functions for gauge couplings (taking into account all weak scale threshold corrections), we check the gauge coupling unification and the dark matter constraints. The future XENON1T sensitivity to our scenario is also examined. Finally, in section 4 we give our conclusions.

2 A split SUSY model from SUSY GUT

There are many possible ways to mediate the SUSY breaking effects from the hidden sector to the visible sector. A very interesting and predictive possibility is the gravity mediation. With certain non-renormalizable terms, proper soft SUSY breaking parameters can be generated. In many popular gravity mediation scenarios, the Kähler potential is assumed to be minimal. However, a general Kähler potential seems to be more natural. When certain high-representation chiral fields for the GUT group are involved in the non-renormalizable Kähler potential, the kinetic terms of superfields could have other contributions after the GUT symmetry breaking. New non-renormalizable terms in the superpotential involving high-representation fields could also be important.

In general, the non-vanishing F-term VEVs of certain fields which break SUSY could be either gauge singlets or non-singlets. In this section, we propose that SUSY breaking is triggered by the GUT group non-singlet F-term VEVs. The gaugino and sfermion masses will be generated by some non-renormalizable operators which could arise from integrating out certain gravitational effects. In order to get more compact spectrum and simplify the relevant expressions, we adopt the SO(10) GUT group with the Georgi-Glashw SU(5) GUT group as an intermediate stage in the symmetry breaking chain:

$$\mathrm{SO}(10)\xrightarrow[\mathbf{16},\overline{\mathbf{16}}]{}\mathrm{SU}(5)\times\mathrm{U}(1)_X\xrightarrow[\mathbf{54}(\mathbf{45})]{}\mathrm{SU}(3)_c\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y.$$

JHEP03(2015)050

Here $U(1)_X$ is broken by possible ν_L^c component VEVs of $H_{\mathbf{16}}$ and $\bar{H}_{\overline{\mathbf{16}}}$ Higgs fields. Therefore, most of the results in this paper are also valid in SU(5) GUT. The possibility of SO(10) to Pati-Salam will be considered elsewhere.

In order to accommodate the non-minimal Kähler potential, we need to know the group products of various SO(10) representations [45]. The spinor representation (for the matter part) can be decomposed as

$$\overline{\mathbf{16}} \otimes \mathbf{16} = \mathbf{1} \oplus \mathbf{45} \oplus \mathbf{210}, \tag{2.1}$$

while the adjoint representation for gaugino is

$$(\mathbf{45} \otimes \mathbf{45})_{\text{symmetric}} = \mathbf{1} \oplus \mathbf{54} \oplus \mathbf{210} \oplus \mathbf{770}. \tag{2.2}$$

and the fundamental representation of Higgs is

$$\mathbf{10} \otimes \mathbf{10} = \mathbf{1} \oplus \mathbf{45} \oplus \mathbf{54}. \tag{2.3}$$

We assume the non-minimal gauge kinetic term for vector supermultiplets

$$\mathcal{L} = \int d^2\theta W^\alpha \left(\delta_{\alpha\beta} + \eta \frac{\Phi_{\alpha\beta}}{M_*}\right) W^\beta \ . \tag{2.4}$$

with M_* being the reduced Planck scale. So the kinetic part for gauge field with non-minimal Kähler potential is given by

$$\mathcal{L} = -\frac{1}{4k} Tr[F_{\mu\nu}F^{\mu\nu}] - \frac{\eta}{4kM_*} Tr[F_{\mu\nu}\Phi F^{\mu\nu}] \,, \tag{2.5}$$

with k being the normalization factor for various representations according to $Tr(T^aT^b) = k\delta^{ab}$. After the GUT non-singlet develops a VEV, $\langle\Phi\rangle = v + F_\Phi\theta^2$ with $v^2 \gg F_\Phi$, the unification condition turns into

$$g_1^2(M_X)\left(1 + \frac{\eta v}{M_*}\delta_1\right) = g_2^2(M_X)\left(1 + \frac{\eta v}{M_*}\delta_2\right) = g_3^2(M_X)\left(1 + \frac{\eta v}{M_*}\delta_3\right) \tag{2.6}$$

with $\delta_{3,2,1}$ being the appropriate group factors for $SU(3)_c, SU(2)_L, U(1)_Y$, respectively. If the hierarchy between the (first step Georgi-Glashow SU(5)) GUT scale and the Planck scale is not small, the previous GUT conditions turn into two independent new GUT conditions:

$$F_1 \equiv \frac{\frac{g_1^2}{g_2^2}(M_X) - 1}{\frac{g_1^2}{g_3^2}(M_X) - 1} = \frac{\delta_2 - \delta_1}{\delta_3 - \delta_1} = -\frac{2}{3} \,, \tag{2.7}$$

$$F_2 \equiv \frac{\frac{g_2^2}{g_1^2}(M_X) - 1}{\frac{g_2^2}{g_3^2}(M_X) - 1} = \frac{\delta_1 - \delta_2}{\delta_3 - \delta_2} = \frac{2}{5} \ . \tag{2.8}$$

For non-singlet Φ, the F-term F_Φ can be decomposed as $(F_\Phi)_{ab} = F_U \cdot A_{ab}$ with A_{ab} being the group factor and F_U the universal part. The review of the group structure can be found in [45].

The group structure of the **24** component F-term VEV of **54** representation Higgs can be written in terms of 10×10 matrix

$$\langle F_{\mathbf{54}} \rangle_{ab} = F^U_{\mathbf{54}} A_{ab} = F^U_{\mathbf{54}} \sqrt{\frac{3}{5}} \left(\frac{1}{3}, \frac{1}{3}, \frac{1}{3}, -\frac{1}{2}, -\frac{1}{2}, \frac{1}{3}, \frac{1}{3}, \frac{1}{3}, -\frac{1}{2}, -\frac{1}{2} \right). \tag{2.9}$$

The gaugino will get contribution from

$$\mathcal{L} \supseteq \frac{< F_\Phi >_{\alpha\beta}}{M_*} \lambda^\alpha \lambda^\beta \tag{2.10}$$

with the F-term VEVs $(F_{\mathbf{54}})_{ab} = F^U_{\mathbf{54}}.A_{ab}$. Here the universal part $F^U_{\mathbf{54}}$ is independent of the group structure. Then the non-universal gaugino masses are given by

$$M_1 = -\frac{1}{6}\sqrt{\frac{3}{5}} m_{1/2}, \quad M_2 = -\frac{1}{2}\sqrt{\frac{3}{5}} m_{1/2}, \quad M_3 = \frac{1}{3}\sqrt{\frac{3}{5}} m_{1/2} \tag{2.11}$$

with

$$m_{1/2} = \frac{F^U_{\mathbf{54}}}{M_*}. \tag{2.12}$$

The sfermion masses and kinetic term will be generated by the following non-renormalizable Kähler potential

$$K = \frac{1}{M_*^2} \phi_a^\dagger (\Phi^\dagger \Phi)_{ab} \phi_b \tag{2.13}$$

with proper F-term and lowest component VEVs of Φ, respectively. Here we assume that the universal part of the kinetic terms $\Phi^\dagger \Phi$ is approximately canceled by a similar high dimensional operators with a lowest component VEV of singlet.

We know from the group theory that possible contributions to Kähler potential for matter content (filled in **16** representation of SO(10)) can arise from the following type of Higgs fields

$$\mathbf{1} \oplus \mathbf{45} \supset \mathbf{54} \otimes \mathbf{54}\,, \qquad \mathbf{1} \oplus \mathbf{210} \supset \overline{\mathbf{16}} \otimes \mathbf{16}\ . \tag{2.14}$$

So the F-term VEV of the **54** representation chiral superfield will contribute to non-universal sfermion masses of order $(F_{\mathbf{54}}/M_*)^2$ in addition to possible contributions from $F_{\mathbf{16}}$. We assume $F_{\mathbf{16}} \sim F_{\mathbf{54}}$ and we always have $F_{\overline{\mathbf{16}}} = F_{\mathbf{16}}$. Since $v_{\mathbf{16}} \gg v_{\mathbf{54}}$, the **54** representation field will give sub-leading contributions and we will not include them explicitly in the following expressions. From the group structure there are several possible contractions for matter fields ϕ with the form of Kähler potential

$$\begin{aligned} K \supset\ & \frac{1}{M_*^2} \sum_{a=1}^{3} (\phi^\dagger_{a,\mathbf{16}} \otimes \phi_{a,\mathbf{16}})^{\mathbf{1}} [(d_{\mathbf{1}} \Phi^\dagger_{H_{\mathbf{16}}} + \tilde{d}_{\mathbf{1}} \Phi_{\bar{H}_{\overline{\mathbf{16}}}}) \otimes (f_{\mathbf{1}} \Phi_{H_{\mathbf{16}}} + \tilde{f}_{\mathbf{1}} \Phi^\dagger_{\bar{H}_{\overline{\mathbf{16}}}})]^{\mathbf{1}} \\ & + \frac{1}{M_*^2} \sum_{a=1}^{3} (\phi^\dagger_{a,\mathbf{16}} \otimes \phi_{a,\mathbf{16}})^{\mathbf{45}}_{mn} [(d_{\mathbf{45}} \Phi^\dagger_{H_{\mathbf{16}}} + \tilde{d}_{\mathbf{45}} \Phi_{\bar{H}_{\overline{\mathbf{16}}}}) \otimes (f_{\mathbf{45}} \Phi_{H_{\mathbf{16}}} + \tilde{f}_{\mathbf{45}} \Phi^\dagger_{\bar{H}_{\overline{\mathbf{16}}}})]^{\mathbf{45}}_{mn} \\ & + \frac{1}{M_*^2} \sum_{a=1}^{3} (\phi^\dagger_{a,\mathbf{16}} \otimes \phi_{a,\mathbf{16}})^{\mathbf{210}}_{mnlp} [(d_{\mathbf{210}} \Phi^\dagger_{H_{\mathbf{16}}} + \tilde{d}_{\mathbf{210}} \Phi_{\bar{H}_{\overline{\mathbf{16}}}}) \otimes (f_{\mathbf{210}} \Phi_{H_{\mathbf{16}}} + \tilde{f}_{\mathbf{210}} \Phi^\dagger_{\bar{H}_{\overline{\mathbf{16}}}})]^{\mathbf{210}}_{mnlp} \end{aligned} \tag{2.15}$$

where $d, f, \tilde{d}, \tilde{f}$ denote the corresponding combination coefficients and a is the family index.

It can be checked that only the first term in eq. (2.15) contributes in this scenario. Then the soft sfermion masses are given by

$$\epsilon^2 \tilde{m}^2_{\mathbf{16}} = d_{\mathbf{1}} f_{\mathbf{1}} \frac{F^*_{\mathbf{16}} F_{\mathbf{16}}}{M_*^2} + \tilde{d}_{\mathbf{1}} \tilde{f}_{\mathbf{1}} \frac{F^*_{\overline{\mathbf{16}}} F_{\overline{\mathbf{16}}}}{M_*^2} + d_{\mathbf{1}} \tilde{f}_{\mathbf{1}} \frac{F^*_{\mathbf{16}} F_{\overline{\mathbf{16}}}}{M_*^2} + \tilde{d}_{\mathbf{1}} f_{\mathbf{1}} \frac{F^*_{\overline{\mathbf{16}}} F_{\mathbf{16}}}{M_*^2} \tag{2.16}$$

after we take into account the normalization factor $\epsilon^2 = v^2_{\mathbf{16}}/M_*^2$ of the kinetic term. So we have

$$\tilde{m}^2_{\mathbf{16}} \sim \frac{F^*_{\mathbf{16}} F_{\mathbf{16}}}{M_*^2 \epsilon^2} \sim \frac{F^*_{\mathbf{16}} F_{\mathbf{16}}}{v^2_{\mathbf{16}}}. \tag{2.17}$$

We can see that the typical gaugino mass scale $m_{1/2}$ is suppressed by a factor ϵ relative to the sfermion mass scale.

The soft SUSY breaking masses for the Higgs potential can be similarly obtained

$$\begin{aligned} K \supset\ & \frac{1}{M_*^2} (\phi^\dagger_{\mathbf{10}} \otimes \phi_{\mathbf{10}})^{\mathbf{1}} [(g_{\mathbf{1}} \Phi^\dagger_{H_{\mathbf{16}}} + \tilde{g}_{\mathbf{1}} \Phi_{\bar{H}_{\overline{\mathbf{16}}}}) \otimes (h_{\mathbf{1}} \Phi_{H_{\mathbf{16}}} + \tilde{h}_{\mathbf{1}} \Phi^\dagger_{\bar{H}_{\overline{\mathbf{16}}}})]^{\mathbf{1}} \\ & + \frac{1}{M_*^2} (\phi^\dagger_{\mathbf{10}} \otimes \phi_{\mathbf{10}})^{\mathbf{45}}_{mn} [(g_{\mathbf{45}} \Phi^\dagger_{H_{\mathbf{16}}} + \tilde{g}_{\mathbf{45}} \Phi_{\bar{H}_{\overline{\mathbf{16}}}}) \otimes (h_{\mathbf{45}} \Phi_{H_{\mathbf{16}}} + \tilde{h}_{\mathbf{45}} \Phi^\dagger_{\bar{H}_{\overline{\mathbf{16}}}})]^{\mathbf{45}}_{mn}, \end{aligned} \tag{2.18}$$

with also contributions from $F_{\mathbf{54}}$ after we add similar terms involving the **54** representation Higgs fields. Both contributions are at the same order and the soft SUSY breaking Higgs masses are given by $m^2_{H_{u,d}} \sim (m_{1/2})^2$. Thus we can see that the soft SUSY breaking Higgs mass parameters can be at the same order as the gaugino.

The trilinear terms will also get contributions from both **16** and $\overline{\mathbf{16}}$ representation Higgs fields. The relevant non-renormalizable superpotential has the form

$$\begin{aligned} W \supset\ & \frac{y_{\text{Yukawa}}}{M_*^2} (\Phi_{\overline{H}_{\overline{\mathbf{16}}}} \otimes \Phi_{H_{\mathbf{16}}})^{\mathbf{1}} [C^{\mathbf{1}} (\mathbf{16_i} \otimes \mathbf{16_i})^m_{\mathbf{10}} \mathbf{10}^m] \\ & + \frac{y_{\text{Yukawa}}}{M_*^2} (\Phi_{\overline{H}_{\overline{\mathbf{16}}}} \otimes \Phi_{H_{\mathbf{16}}})^{\mathbf{45}}_{mn} [C^{\mathbf{45}} (\mathbf{16_i} \otimes \mathbf{16_i})^{mnl}_{\mathbf{120}} \mathbf{10}^l] \\ & + \frac{y_{\text{Yukawa}}}{M_*^2} (\Phi_{\overline{H}_{\overline{\mathbf{16}}}} \otimes \Phi_{H_{\mathbf{16}}})^{\mathbf{210}}_{mnlp} [C^{\mathbf{210}} (\mathbf{16_i} \otimes \mathbf{16_i})^{mnlpq}_{\mathbf{126}} \mathbf{10}^q] \ . \end{aligned} \tag{2.19}$$

Again, here only the first term contributes. From the lowest component VEV and F-term VEV of $H_{\mathbf{16}}$ and $H_{\overline{\mathbf{16}}}$, we can obtain the trilinear coupling as

$$\epsilon^2 A_y = C^{\mathbf{1}} \left[\frac{v_{\mathbf{16}} F_{\overline{\mathbf{16}}}}{M_*^2} + \frac{v_{\overline{\mathbf{16}}} F_{\mathbf{16}}}{M_*^2} \right] y_{\text{Yukawa}} \ . \tag{2.20}$$

after we normalize the kinetic term for the **16** representation matter contents. Thus, we see that

$$A_0 = \frac{A_y}{y_{\text{Yukawa}}} \sim \frac{F_{\overline{\mathbf{16}}}}{v_{\mathbf{16}}} \tag{2.21}$$

is typically of the same scale as sfermion masses.

The B_μ-term and μ term are given by

$$W \supset (M + \Phi_{\mathbf{45}}) \phi_{\mathbf{10}} \phi_{\mathbf{10}} + k_1 (M + \Phi_{\mathbf{45}}) \frac{(\Phi_{\mathbf{54}})}{M_*} \phi_{\mathbf{10}} \phi_{\mathbf{10}} \ . \tag{2.22}$$

JHEP03(2015)050

We will not discuss in detail the doublet-triplet splitting problem in GUT and just use the fact that the combination $M + \langle \Phi_{\mathbf{45}} \rangle$ will lead to (by possible tuning or some mechanism) proper low energy effective μ term.

The B_μ-term will be generated after $\Phi_{\mathbf{54}}$ acquires F-term VEVs

$$B_\mu = -k_1 \sqrt{\frac{3}{5}} \frac{F^U_{\mathbf{54}}}{M_*} \mu \ . \tag{2.23}$$

The electroweak symmetry breaking condition can relate different parameters as

$$|\mu|^2 = -\frac{M_Z^2}{2} + \frac{1}{\tan^2\beta - 1}(m^2_{H_d} - \tan^2\beta m^2_{H_u}) \, , \tag{2.24}$$

$$2B_\mu = \sin 2\beta (m^2_{H_d} + m^2_{H_u} + 2|\mu|^2) \ . \tag{2.25}$$

Although there are many possibilities for the choice of μ from electroweak symmetry breaking condition, the relation eq. (2.23) will further constrain the choice of μ. Combining eq. (2.23) with the (RGE modified) approximation $m^2_{H_d} \gtrsim m^2_{H_u}$, we can see that the symmetry breaking condition requires $|\mu|^2 \sim -m^2_{H_u} \sim B_\mu \sim m^2_{1/2}$ for a negative $m^2_{H_u}$. For a positive $m^2_{H_u}$, it is difficult to reconcile a large $\tan\beta$ and $m^2_{H_d} \gtrsim m^2_{H_u}$ in case of $\mu \ll m^2_{h_u}$. Only if the RGE running can affect greatly the GUT scale relation $m^2_{H_d} = m^2_{H_u}$ can a relatively heavy μ be possible. The cases of $\tan\beta \sim 1$ which indicate $|\mu|^2 \gg |m^2_{h_d}|$ for both signs of $m^2_{H_u}$ can hardly be compatible with eq. (2.23).

So we see that in our scenario the hierarchy between the GUT scale and the Planck scale is used in splitting the SUSY soft spectrum in contrast to the gauge loop factor appeared in simply unnatural supersymmetry [20].

3 Gauge coupling unification and dark matter constraints

3.1 Inputs of our model

In this section we check the gauge coupling unification and dark matter constraints in our scenario. From the previous section, we can see the inputs of our scenario at the GUT scale:

- The gaugino masses with the raito

$$M_1 : M_2 : M_3 = -1 : -3 : \ 2$$

- The hierarchy between the GUT scale (SO(10) breaking scale) and the Planck scale ϵ.
- The higgsino mass μ which should be at the same sacle as gaugino mass.
- The parameter $\tan\beta$. We scan it in the range $1 \sim 50$.

The sfermion mass m_0, determined by $m_{1/2}/\epsilon$, and the trilinear term $A_0 \simeq m_0$ are not independent parameters in our model. This scenario predicts different parameter values in comparison to the universal gaugino mass scenario [21]. It is well known that the ratios of

JHEP03(2015)050

gaugino masses and the corresponding gauge couplings in SUSY are RGE-invariant up to one-loop level

$$\frac{d}{d\ln\mu}\left(\frac{M_i}{g_i^2}\right) = 0. \tag{3.1}$$

If we assume universal gaugino masses $M_1 = M_2 = M_3 \equiv M_U$ at the GUT scale, from the RGE-invariant ratio

$$\frac{M_1}{g_1^2} = \frac{M_2}{g_2^2} = \frac{M_3}{g_3^2} = \frac{M_U}{g_U^2} \tag{3.2}$$

and the electroweak scale inputs we can have a mass relation given by

$$M_1 : M_2 : M_3 \simeq 1 : 2 : 6 \tag{3.3}$$

at electroweak scale. On the other hand, if we assume non-universal gaugino masses at the GUT scale, the above gaugino mass ratios at the electroweak scale will be approximately changed to

$$M_1 : M_2 : M_3 \simeq 1 : 6 : -12\,. \tag{3.4}$$

We should note again that the formula eq. (3.1) can no longer be valid below M_S where heavy scalars decouple. So the exact gaugino mass ratio in split SUSY should be obtained by the subsequent RGE running from M_S to EW scale.

3.2 The gauge coupling unification requirement

We study the gauge coupling unification with the two-loop RGE running of gauge couplings, taking into account all weak scale threshold corrections. The relevant analytical results for the two-loop beta functions are given in the appendix. It is well known that the two-loop RGE running for gauge couplings are scheme independent [46], so we use the $\overline{MS}$ couplings in our study of the gauge coupling unification. Lacking the full knowledge responsible for GUT symmetry breaking, especially the mechanism used to solve the doublet-triplet splitting problem which is model dependent, we therefore neglect such GUT scale threshold corrections in our study. In order to make our calculation reliable, the first step SU(5) GUT scale must be much lower than the Planck scale so that the gravitational effects can be neglected. Besides, the GUT scale can not be too low, otherwise it will lead to fast proton decay. The constraint is [21, 47–50]

$$1.0 \times 10^{19}\text{GeV} > M_{\text{GUT}} > \sqrt{35\alpha_{\text{GUT}}}\left(6.9 \times 10^{15}\right)\text{GeV}\ . \tag{3.5}$$

In our numerical study we input the central values of g_1 and g_2 while for g_3 we require it in the 3σ at the electroweak scale. Other inputs at the electroweak scale, for example, the top Yukawa coupling y_t, are extracted from the SM taking into account the threshold corrections. Relevant details can be seen in the appendix of [21, 51, 52]. Because of the uncertainty of the GUT scale threshold contributions, we adopt the criteria that the gauge coupling unification is satisfied when the three couplings differ within the range < 0.005.

JHEP03(2015)050

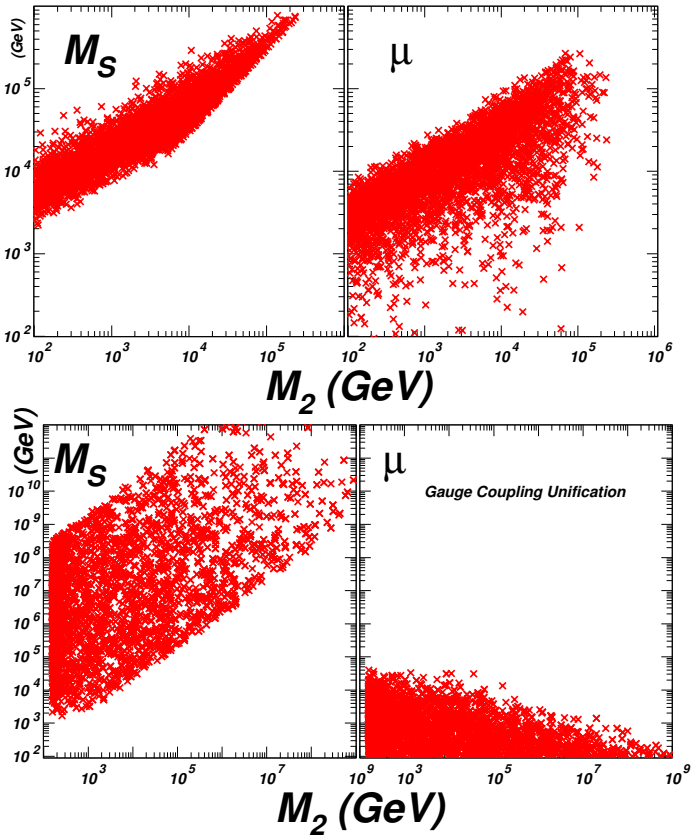

Figure 1. The scatter plots of the parameter space with gauge coupling unification in case of universal gaugino input (upper) and non-universal gaugino input (lower) at the GUT scale.

Figure 1 shows the result of the parameter space (M_2, M_S, μ) with successful gauge coupling unification in case of universal condition eq. (3.3) and non-universal condition eq. (3.4), respectively. We can see that the gaugino unification gives a stringent constraint on the parameter space. From the upper-left panel we can find an upper bound for M_S, which is about 10^6 GeV. Since split SUSY requires $M_S \gg M_{\tilde{g}_i}$, we can also obtain an upper bound for M_2 correspondingly. From the upper-right panel we can find upper limits for μ and M_2, which are around 100 TeV, independent of the M_S value. However, the constraints for the non-universal gaugino scenario are rather mild. From the lower panels, we can see that M_S can be as high as 10^{12} GeV while M_2 can be 10^9 GeV. The μ parameter, which plays an important role in gauge coupling unification, also has an upper bound around 10^5 GeV for both non-universal and universal case.

3.3 Dark matter constraints

Now we study the dark matter constraints in our scenario, using the latest dark matter relic density data from Planck [53], WMAP [54] and the direct detection limits from XENON100 [55] and the LUX [56]. The package DarkSUSY [57] is used to scan the parameter space of our scenario. In addition to the inputs defined above, we also require $1 < \tan\beta < 50$. In order to calculate the relic density of dark matter, we use the fact

JHEP03(2015)050

that the effects of heavy sfermions and heavy Higgs fields almost entirely decouple when $M_S = M_A > 5$TeV [58]. So in our numerical study, we keep the samples which satisfy the GUT constraints and then set $M_S = M_A = 10$ TeV in DarkSUSY to carry out the dark matter related calculations.

For illustration we set $M_S = 30,\ 50,\ 100,\ 200$ TeV and show the allowed ranges for the parameters $M_2, \mu, \tan\beta$. In our scan, we take into account the current dark matter and collider constraints (the details can be found in our previous work [21]):

(1) We require that the relic density of the neutralino dark matter satisfies the Planck result $\Omega_{DM} = 0.1199 \pm 0.0027$ [53] (in combination with the WMAP data [54]).

(2) The LEP lower bounds on neutralino and charginos, including the invisible decay of Z-boson. For LEP experiments, the most stringent constraints come from the chargino mass and the invisible Z-boson decay. We require $m_{\tilde{\chi}^\pm} > 103$GeV and the invisible decay width $\Gamma(Z \to \tilde{\chi}_0\tilde{\chi}_0) < 1.71$ MeV, which is consistent with the 2σ precision EW measurement $\Gamma_{inv}^{non-SM} < 2.0$ MeV.

(3) The precision electroweak measurements. We require the oblique parameters [59, 60] to be compatible with the LEP/SLD data at 2σ confidence level [61].

(4) The combined mass range for the Higgs boson: $123\text{GeV} < M_h < 127\text{GeV}$ from ATLAS and CMS collaborations of LHC [1, 2]. In split SUSY, due to large M_S, $\log(m_{\tilde{f}}^2/m_t^2) \gg 1$ will spoil the convergence of the traditional loop expansion in evaluating the SUSY effects of Higgs boson self-energy. So in order to calculate mass of the SM-like Higgs boson, we use the RGE improved effective potential [62] which is employed in the NMSSMTools package [63] after we set $\lambda = \kappa \to 0$.

Note that the spin-independent (SI) dark matter-nucleon scattering rate is calculated with relevant parameters chosen as [64–67]: $f_{T_u}^{(p)} = 0.023, f_{T_d}^{(p)} = 0.032, f_{T_u}^{(n)} = 0.017, f_{T_d}^{(n)} = 0.041$ and $f_{T_s}^{(p)} = f_{T_s}^{(n)} = 0.020$. We take into account all the contributions known so far (including QCD corrections) in our calculation of the scattering rate. The value of f_{T_s} is taken from the recent lattice simulation results [68–70].

It is instructive to compare the dark matter constraints for non-universal and universal gaugino scenarios. Results for the universal gaugino scenario are taken from our previous work [21]. In figures 2 and 3 we show the samples surviving the constraints (1-4), where the green '$\times$' and red '$\triangle$' denote respectively the samples allowed and excluded by the gauge coupling unification requirement. For these results we have the following discussions:

- We can see from figure 2 that increasing M_S tends to slightly relax the gauge coupling unification constraints in non-universal gaugino scenario, in contrast to universal gaugino scenario where increasing M_S tends to spoil the gauge coupling unification.

- For both scenarios, a strip corresponding to the higgsino dark matter with mass range from 1.0 to 1.3 TeV can always survive the combined constraints of gauge coupling unification and dark matter direct detection (except $M_S > 200$TeV in universal gaugino scenario which is not preferred by gauge coupling unification requirement).

JHEP03(2015)050

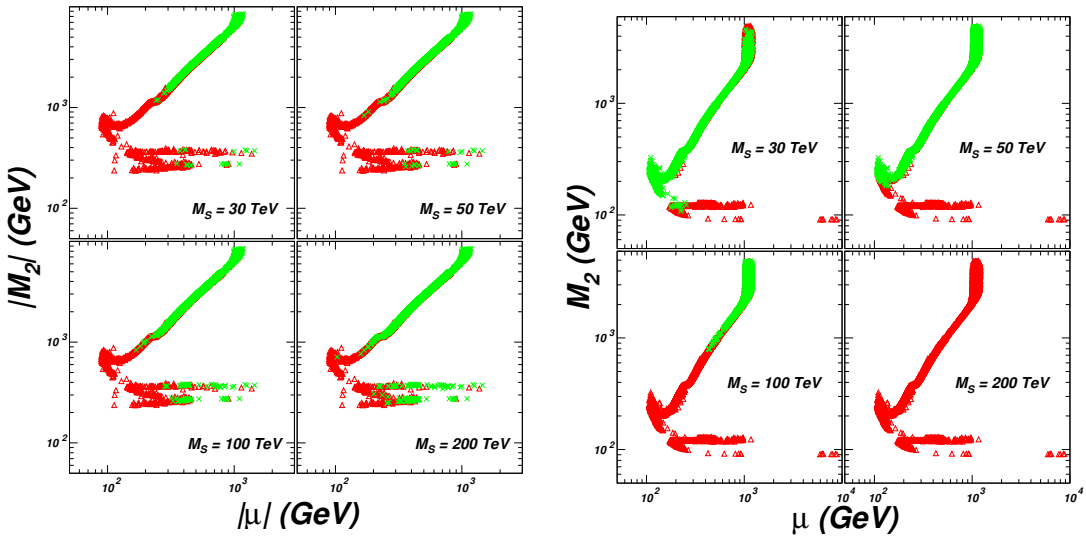

Figure 2. The scatter plots of the (μ, M_2) parameter space satisfying constraints (1-4) including dark matter relic density. The green '×' (red '△') can (cannot) achieve the gauge coupling unification. The left panel is for the non-universal gaugino scenario proposed in this work while the right panel is for the universal gaugino scenario studied in our previous work [21].

This is the well known fact that higgsino at about 1.2 TeV can be a viable dark matter candidate.

- Outside the strip of higgsino dark matter, almost all the survived points will be covered by XEON1T in both scenarios. An interesting exception occurs in non-universal gaugino scenario, where a tiny strip at about 50 GeV cannot be covered by XEON1T and such a strip enlarges as M_S increases. However, a careful analysis indicates that this strip corresponds to a bino dark matter. Although the gaugino mass ratio $M_1 : M_2 : M_3 = 1 : 6 : 12$ is no longer valid at the weak scale in split SUSY, the RGE running in general will not change significantly the mass ratio for a not too large M_S. So we can estimate that this strip corresponds to a gluino below 700 GeV. The current preliminary limits on gluino mass using 20 fb^{-1} of 8 TeV data are $M_{\tilde{g}} = 1350$ GeV ([74]) and $M_{\tilde{g}} = 1200$ GeV ([75]) assuming a massless neutralino for mini-split SUSY [76]. So this tiny strip should have been ruled out by the LHC data.

So we can see that a neutralino dark matter below about 1.0 TeV will be fully covered by XENON1T for both non-universal and universal gaugino scenarios in split SUSY. We checked that our numerical calculation results are not sensitive to the sign of μ, except that for the minus sign a very small part of parameter space (called blind spots) can survive all the constraints including the future XENON1T limits [73].

4 Conclusion

In this work we proposed to achieve the hierarchy of sparticle spectrum in split SUSY from the gap between the GUT scale and the Planck scale. We built a split SUSY model

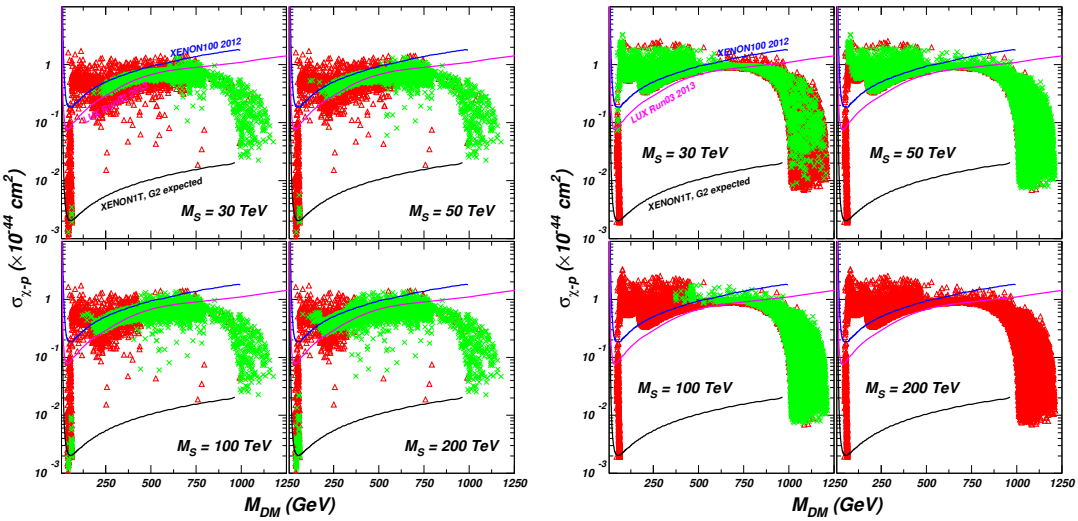

Figure 3. Same as figure 2, but showing the spin-independent cross section of dark matter scattering off the nucleon. The curves denote the limits from XENON100(2012) and LUX as well as the future XENON1T sensitivity.

(which gives non-universal gaugino masses) with proper high dimensional operators in the framework of SO(10) GUT. Based on a calculation of two-loop beta functions for gauge couplings (taking into account all weak scale threshold corrections), we checked the gauge coupling unification and dark matter constraints. We found that our scenario can achieve the gauge coupling unification and satisfy the dark matter constraints in some part of parameter space. We also examined the sensitivity of the future XENON1T experiment and found that the currently allowed parameter space in our scenario can be covered for a neutralino dark matter below about 1.0 TeV.

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A The two-loop beta function for gauge couplings

In this work, we adopt the method in [77–79] to calculate the two-loop beta functions for three gauge couplings in the split SUSY, taking into account the weak scale threshold corrections. Our result is in agreement with [17, 19]. The notations for the two-loop RGE can be found in the appendix of our previous work [21]. The one-loop beta functions for the three gauge couplings are given by

$$b_3 = -7\theta(\mu - M_Z) + 2\theta(\mu - M_{\tilde{g}}) + 2\theta(\mu - M_S)\,, \tag{A.1}$$

$$b_2 = -\frac{19}{6}\theta(\mu - M_Z) + \frac{4}{3}\theta(\mu - M_{\tilde{W}}) + \frac{2}{3}\theta(\mu - M_{\tilde{H}}) + \frac{13}{6}\theta(\mu - M_S)\,, \tag{A.2}$$

$$b_1 = \frac{41}{10}\theta(\mu - M_Z) + \frac{2}{5}\theta(\mu - M_{\tilde{H}}) + \frac{21}{10}\theta(\mu - M_S), \tag{A.3}$$

with the step function defined as

$$\theta(x) = \begin{cases} 1, & x \geq 0; \\ 0, & x < 0. \end{cases} \tag{A.4}$$

The 2-loop RGE for $\mathrm{SU}(3)_c, \mathrm{SU}(2)_L, \mathrm{U}(1)_Y$ gauge couplings (g_3, g_2, g_1, respectively) are given by

$$\frac{d}{d\ln E} g_i = \frac{b_i}{(4\pi)^2} g_i^3 + \frac{g_i^3}{(4\pi)^4} \left[\sum_j B_{ij} g_j^2 - \sum_{a=u,d,e} d_i^a Tr(h^{a\dagger} h^a) - d_W(\tilde{g}_u^2 + \tilde{g}_d^2) - d_B(\tilde{g}_u'^2 + \tilde{g}_d'^2) \right] \tag{A.5}$$

with the $\mathrm{U}(1)_Y$ normalization $g_1^2 = \frac{5}{3}(g_Y)^2$.

The two-loop beta functions for gauge couplings are given by

$$\begin{aligned} B_{ij} = {} & \theta(\mu - M_Z) \begin{pmatrix} \frac{199}{50} & \frac{27}{10} & \frac{44}{5} \\ \frac{9}{10} & \frac{35}{6} & 12 \\ \frac{11}{10} & \frac{9}{2} & -26 \end{pmatrix} + \theta(\mu - M_2) \begin{pmatrix} 0 & 0 & 0 \\ 0 & \frac{64}{3} & 0 \\ 0 & 0 & 0 \end{pmatrix} + \theta(\mu - M_3) \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 48 \end{pmatrix} \\ & + \theta(\mu - M_{\tilde{H}}) \begin{pmatrix} \frac{9}{50} & \frac{9}{10} & 0 \\ \frac{3}{10} & \frac{49}{6} & 0 \\ 0 & 0 & 0 \end{pmatrix} + \theta(\mu - M_S) \begin{pmatrix} \frac{19}{5} & \frac{9}{5} & \frac{44}{5} \\ \frac{3}{5} & -\frac{31}{3} & 12 \\ \frac{11}{10} & \frac{9}{2} & -8 \end{pmatrix}. \end{aligned} \tag{A.6}$$

Similarly, we have

$$d^u = \theta(\mu - M_Z)\left(\frac{17}{10}, \frac{3}{2}, 2\right) + \theta(\mu - M_S)\left(\frac{7}{2}, \frac{9}{2}, 2\right), \tag{A.7}$$

$$d^d = \theta(\mu - M_Z)\left(\frac{1}{2}, \frac{3}{2}, 2\right) + \theta(\mu - M_S)\left(\frac{23}{10}, \frac{9}{2}, 2\right), \tag{A.8}$$

$$d^e = \theta(\mu - M_Z)\left(\frac{3}{2}, \frac{1}{2}, 0\right) + \theta(\mu - M_S)\left(\frac{21}{10}, \frac{3}{2}, 0\right), \tag{A.9}$$

$$d^W = \left(\frac{9}{20}, \frac{11}{4}, 0\right)\theta(\mu - \max(M_2, M_{\tilde{H}})) + \theta(\mu - M_S)\left(-\frac{9}{20}, -\frac{11}{4}, 0\right), \tag{A.10}$$

$$d^B = \left(\frac{3}{20}, \frac{1}{4}, 0\right)\theta(\mu - \max(M_1, M_{\tilde{H}})) + \theta(\mu - M_S)\left(-\frac{9}{20}, -\frac{11}{4}, 0\right). \tag{A.11}$$

The one-loop renormalization group equations for Yukawa couplings below the M_S scale can be written as

$$16\pi^2 \frac{d}{dt} h^u = h^u \left[-3c_i^u g_i^2 + c_T^u T + c_{S_1}^u S_1 + c_{S_2}^u S_2 + \frac{3}{2}\left(h^{u\dagger} h^u - h^{d\dagger} h^d\right) \right], \tag{A.12}$$

$$16\pi^2 \frac{d}{dt} h^d = h^d \left[-3c_i^d g_i^2 + c_T^d T + c_{S_1}^d S_1 + c_{S_2}^d S_2 + \frac{3}{2}\left(h^{d\dagger} h^d - h^{u\dagger} h^u\right) \right], \tag{A.13}$$

$$16\pi^2 \frac{d}{dt} h^e = h^e \left[-3c_i^e g_i^2 + c_T^e TT + c_{S_1}^e S_1 + c_{S_2}^e S_2 + \frac{3}{2} h^{e\dagger} h^e \right], \tag{A.14}$$

with

$$T = Tr(3h^{u\dagger} h^u + 3h^{d\dagger} h^d + h^{e\dagger} h^e), \;\; S_1 = \frac{1}{2}\left[(\tilde{g}_u')^2 + (\tilde{g}_d')^2\right], \;\; S_2 = \frac{3}{2}\left(\tilde{g}_u^2 + \tilde{g}_d^2\right). \tag{A.15}$$

JHEP03(2015)050

Above M_S, we will recover the MSSM result and the one-loop RGE for Yukawa-type interactions in the superpotential are well known to be

$$16\pi^2 \frac{d}{dt}\lambda^u = \lambda^u \left[-2c_i^u g_i^2 + 3Tr(\lambda^{u\dagger}\lambda^u) + 3\lambda^{u\dagger}\lambda^u + \lambda^{d\dagger}\lambda^d\right], \tag{A.16}$$

$$16\pi^2 \frac{d}{dt}\lambda^d = \lambda^d \left[-2c_i^d g_i^2 + Tr(3\lambda^{d\dagger}\lambda^d + \lambda^{e\dagger}\lambda^e) + \lambda^{u\dagger}\lambda^u + 3\lambda^{d\dagger}\lambda^d\right], \tag{A.17}$$

$$16\pi^2 \frac{d}{dt}\lambda^e = \lambda^e \left[-2c_i^e g_i^2 + Tr(3\lambda^{d\dagger}\lambda^d + \lambda^{e\dagger}\lambda^e) + 3\lambda^{e\dagger}\lambda^e\right], \tag{A.18}$$

with

$$c_i^u = \left(\frac{13}{30}, \frac{3}{2}, \frac{8}{3}\right),\ c_i^d = \left(\frac{7}{30}, \frac{3}{2}, \frac{8}{3}\right),\ c_i^e = \left(\frac{9}{10}, \frac{3}{2}, 0\right). \tag{A.19}$$

The one-loop Yukawa couplings for c_i^u, c_i^d, c_i^e are calculated to be

$$\begin{pmatrix} c_i^u \\ c_i^d \\ c_i^e \end{pmatrix} = \theta(\mu - M_Z) \begin{pmatrix} \frac{17}{60} & \frac{3}{4} & \frac{8}{3} \\ \frac{1}{12} & \frac{3}{4} & \frac{8}{3} \\ \frac{3}{4} & \frac{3}{4} & 0 \end{pmatrix} + \theta(\mu - M_S) \begin{pmatrix} \frac{3}{20} & \frac{3}{4} & 0 \\ \frac{3}{20} & \frac{3}{4} & 0 \\ \frac{3}{20} & \frac{3}{4} & 0 \end{pmatrix} \tag{A.20}$$

All terms are set to zero above M_S. Besides, we have

$$\begin{pmatrix} c_T^u & c_{S_1}^u & c_{S_2}^u \\ c_T^d & c_{S_1}^d & c_{S_1}^d \\ c_T^e & c_{S_1}^e & c_{S_2}^e \end{pmatrix} = \theta(\mu - M_Z) \begin{pmatrix} 1 & 0 & 0 \\ 1 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix} + \theta(\mu - \max(M_2, M_{\tilde{H}})) \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 1 \\ 0 & 0 & 1 \end{pmatrix},$$
$$+\theta(\mu - \max(M_1, M_{\tilde{H}})) \begin{pmatrix} 0 & 1 & 0 \\ 0 & 1 & 0 \\ 0 & 1 & 0 \end{pmatrix} \tag{A.21}$$

and all set to zero above M_S.

One loop RGE for gaugino couplings $\tilde{g}_u, \tilde{g}'_u, \tilde{g}_d, \tilde{g}'_d$ below M_S are given as (with the gaugino relation $M_1 < M_2$)

- Between $[\max(M_2, M_{\tilde{H}}), M_S]$, the RGE for $\tilde{g}_{u,d}$ are given by

$$16\pi^2 \frac{d}{dt}\tilde{g}_u = -3\tilde{g}_u c_i^u g_i^2 + \frac{5}{4}\tilde{g}_u^3 - \frac{1}{2}\tilde{g}_u \tilde{g}_d^2$$
$$+\frac{1}{4}\tilde{g}_u \tilde{g}_u'^2 + \tilde{g}_d \tilde{g}'_d \tilde{g}'_u + \tilde{g}_u(T + c_{S_1}S_1 + c_{S_2}S_2) \tag{A.22}$$

$$16\pi^2 \frac{d}{dt}\tilde{g}_d = -3\tilde{g}_d c_i^d g_i^2 + \frac{5}{4}\tilde{g}_d^3 - \frac{1}{2}\tilde{g}_d \tilde{g}_u^2$$
$$+\frac{1}{4}\tilde{g}_d \tilde{g}_d'^2 + \tilde{g}_u \tilde{g}'_u \tilde{g}'_d + \tilde{g}_d(T + c_{S_1}S_1 + c_{S_2}S_2) \tag{A.23}$$

  with the coefficient

$$c_i^{u,d} = \left(\frac{3}{20}, \frac{11}{4}, 0\right), \quad c_{S_1} = c_{S_2} = 1. \tag{A.24}$$

  Below $\max(M_2, M_{\tilde{H}})$, the coupling are switched off.

JHEP03(2015)050

- Between $[\max(M_2, M_{\tilde{H}}), M_S]$, the RGE for $\tilde{g}_{u,d}$ are given by

$$16\pi^2 \frac{d}{dt}\tilde{g}'_u = -3\tilde{g}'_u \tilde{c}^u_i g_i^2 + \frac{3}{4}\tilde{g}'^3_u + \frac{3}{2}\tilde{g}'_u \tilde{g}'^2_d + \frac{3}{4}\tilde{g}'_u \tilde{g}_u^2 + 3\tilde{g}'_d \tilde{g}_d \tilde{g}_u + \tilde{g}'_u (T + c_{S_1} S_1 + c_{S_2} S_2) \tag{A.25}$$

$$16\pi^2 \frac{d}{dt}\tilde{g}'_d = -3\tilde{g}'_d \tilde{c}^d_i g_i^2 + \frac{3}{4}\tilde{g}'^3_d + \frac{3}{2}\tilde{g}'_d \tilde{g}'^2_u + \frac{3}{4}\tilde{g}'_d \tilde{g}_d^2 + 3\tilde{g}'_u \tilde{g}_u \tilde{g}_d + \tilde{g}'_d (T + c_{S_1} S_1 + c_{S_2} S_2) \tag{A.26}$$

with the coefficient

$$\tilde{c}^{u,d}_i = \left(\frac{3}{20}, \frac{3}{4}, 0\right), \quad c_{S_1} = c_{S_2} = 1. \tag{A.27}$$

Between $[\max(M_1, M_{\tilde{H}}), \max(M_2, M_{\tilde{H}})]$, the RGE reads

$$16\pi^2 \frac{d}{dt}\tilde{g}'_u = -3\tilde{g}'_u \tilde{c}^u_i g_i^2 + \frac{3}{4}\tilde{g}'^3_u + \frac{3}{2}\tilde{g}'_u \tilde{g}'^2_d + \tilde{g}'_u (T + c_{S_1} S_1 + c_{S_2} S_2), \tag{A.28}$$

$$16\pi^2 \frac{d}{dt}\tilde{g}'_d = -3\tilde{g}'_d \tilde{c}^d_i g_i^2 + \frac{3}{4}\tilde{g}'^3_d + \frac{3}{2}\tilde{g}'_d \tilde{g}'^2_u + \tilde{g}'_d (T + c_{S_1} S_1 + c_{S_2} S_2), \tag{A.29}$$

with the coefficient

$$\tilde{c}^{u,d}_i = \left(\frac{3}{20}, \frac{3}{4}, 0\right), \quad c_{S_1} = 1, \quad c_{S_2} = 0. \tag{A.30}$$

Below $\max(M_1, M_{\tilde{H}})$, the coupling are switched off.

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JHEP03(2015)050

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JHEP03(2015)050