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Unification of SUSY breaking and GUT breaking

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Abstract: We build explicit supersymmetric unification models where grand unified gauge symmetry breaking and supersymmetry (SUSY) breaking are caused by the same sector. Besides, the SM-charged particles are also predicted by the symmetry breaking sector, and they give the soft SUSY breaking terms through the so-called gauge mediation. We investigate the mass spectrums in an explicit model with SU(5) and additional gauge groups, and discuss its phenomenological aspects. Especially, nonzero A-term and B-term are generated at one-loop level according to the mediation via the vector superfields, so that the electro-weak symmetry breaking and 125 GeV Higgs mass may be achieved by the large B-term and A-term even if the stop mass is around 1 TeV.

Keywords: Supersymmetry Breaking, Beyond Standard Model, Extended Supersymmetry, Gauge Symmetry

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Contents

1 Introduction

As well-known, the Standard Model (SM) is very successful in describing our nature, and it is firmly established by the Higgs discovery at the LHC [1, 2]. There are still some ambiguities in not only the signal strength of the Higgs particle but also the other observations such as flavor physics, but it would be getting more difficult to consider new-physics effects in any signals.

On the other hand, we are sure that the SM remains several mysteries about our nature: the origin of the fermion generations, the hyper-charge assignment, the Higgs mass, and so on. Many Beyond Standard Models (BSM) were proposed so far motivated by those mysteries, and some of them are expected to be found near future. One of the candidates is the supersymmetric grand unified theory (GUT), which reveals the origin of the Higgs mass and the fermion charges. There are some issues in Yukawa couplings, for instance, how to generate realistic Yukawa couplings and heavy colored Higgs, but it succeeds in the charge quantization ($|Q_e + Q_p| < 10^{-21}$ [3]) and naturally deriving the electro-weak (EW) scale, if the supersymmetry (SUSY) scale ($\Lambda_{\rm SUSY}$) is close to the EW scale. The supersymmetric GUT scenario is constrained by the observation of the proton decay, the direct search of SUSY particles, and the SM measurements. Especially, the Higgs discovery around 125 GeV may require high-scale SUSY ($\Lambda_{\rm SUSY} \gg O(1)$TeV) [4, 5], which may discard the strong motivation of SUSY, that is, the natural explanation of the EW scale. Furthermore, the gauge coupling unification of supersymmetric SU(5) GUT might be lost in high-scale SUSY, depending on the mass spectrum of the SUSY particles. The supersymmetric models could have so many parameters in the bottom-up approach, so that we could have some solutions for the Higgs mass and the gauge coupling unification. However, it is very important to find how to derive such a specific SUSY mass spectrum.

In this paper, we propose an explicit supersymmetric GUT with $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ gauge groups. We discard the miracle of the gauge coupling unification in the Minimal Supersymmetric SM (MSSM), but SUSY breaking and GUT breaking sectors are unified.[1] The SM-charged particles also appear after the symmetry breaking, so the messenger fields for the gauge mediation is also introduced by the breaking sector in our model.[2] The SM fields are only charged under the $\mathrm{SU}(5)_F$ gauge group, so that the charge quantization is realized.

The breaking sector consists of one $\mathrm{SU}(5)_F$ adjoint plus singlet filed (Φ) and $\mathrm{SU}(5)_F$ fundamental and anti-fundamental fields (ϕ, $\widetilde{\phi}$). The vector-like pairs (ϕ, $\widetilde{\phi}$) are also charged under $\mathrm{SU}(2) \times \mathrm{U}(1)_\phi$. As discussed in ref. [23], this type of gauge theory causes SUSY breaking along with the gauge symmetry breaking. In our model, $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ symmetry breaks down to the SM gauge groups, $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$, where $\mathrm{SU}(3)_c$ is from the subgroup of $\mathrm{SU}(5)_F$, and $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ are the linear combinations of the subgroup of $\mathrm{SU}(5)_F$ and $\mathrm{SU}(2) \times \mathrm{U}(1)_\phi$. SUSY is broken by the F-component of the part of Φ. After the symmetry breaking, SM-charged particles are generated by the fluctuation of Φ and (ϕ, $\widetilde{\phi}$) around the vacuum expectation values (VEVs). One interesting point is that the massive gauge boson of $\mathrm{SU}(5)_F$ and the fermionic partners could mediate the SUSY breaking effect through the gauge coupling with Φ, and play a crucial role in generating the non-zero A-term and B-term as discussed in refs. [24, 25]. It is well-known that SUSY-scale A-term could shift the upper bound on the lightest Higgs mass in the MSSM, even if squark is light, and the SUSY-scale B-term is required to realize the EW symmetry breaking. Our A-term and B-term are given at one-loop level, so that they are the same order as the squark masses and gaugino masses. In fact, we will see that Higgs mass could be around 125 GeV, even if $\Lambda_{\rm SUSY}$ is less than $O(1)$ TeV, and the B-term could be consistent with the EW symmetry breaking.

In section 2, we introduce the SUSY and GUT breaking sector in generic $\mathrm{SU}(N_F)_F \times \mathrm{SU}(N) \times \mathrm{U}(1)_\phi$ gauge theory. There, we discuss not only the symmetry breaking, but also the behavior of the gauge couplings and soft SUSY breaking terms according to the gauge mediation with the mediators of the chiral superfields and the vector superfields. In section 3, we apply the breaking sector to the $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ gauge theory. As we mentioned above, an interesting aspect of this model is the improvement of the consistency with the EW symmetry breaking and Higgs mass in the case with low-scale

[1] This type of scenario has been proposed in refs. [6–11].

[2] The messenger sector and SUSY breaking sector are unified, for instance, in refs. [12–22].

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	ϕ	$\widetilde{\phi}$	Φ
$\mathrm{SU}(N_F)_F$	$\mathbf{N_F}$	$\overline{\mathbf{N_F}}$	$\mathbf{adj}_{N_F}\mathbf{+1}$
$\mathrm{SU}(N)$	$\overline{\mathbf{N}}$	$\mathbf{N}$	$\mathbf{1}$
$\mathrm{U}(1)_\phi$	Q_ϕ	$-Q_\phi$	0

Table 1. Chiral superfields in $\mathrm{SU}(N_F)_F \times \mathrm{SU}(N) \times \mathrm{U}(1)_\phi$ gauge theory.

SUSY. We investigate the soft SUSY breaking terms, and discuss how well it is achieved in our scenario. In section 4, we give a comment on the possibility that the breaking sector is applied to other GUT models. Section 5 is devoted to the summary. In appendix A, we give the mass spectrum in the SUSY breaking sector. In appendix B, we show examples of mass spectrums in the MSSM sector.

2 $\mathrm{SU}(N) \times \mathrm{SU}(N_F)_F \times \mathrm{U}(1)_\phi$ gauge theory

In this section, we introduce the model which causes SUSY breaking together with gauge symmetry breaking, based on ref. [23].

We consider $\mathrm{SU}(N_F)_F \times \mathrm{SU}(N) \times \mathrm{U}(1)_\phi$ gauge theory with $N_F > N$. The matter content is shown in table 1: Φ is the $\mathrm{SU}(N_F)_F$ adjoint plus singlet field and $(\phi, \widetilde{\phi})$ pair is the vector-like under $\mathrm{SU}(N_F)_F \times \mathrm{SU}(N) \times \mathrm{U}(1)_\phi$ gauge group.

The superpotential is given by

$$W_R = -hTr_N(\widetilde{\phi}\Phi\phi) + h\Lambda_G \mathrm{Tr}_{N_F}(\Phi), \tag{2.1}$$

assigning $\mathrm{U}(1)_R$ symmetry: the R-charge of Φ is 2 and the R-charge of $(\phi, \widetilde{\phi})$ is vanishing. However, there would be an issue about how to break R-symmetry and how to avoid the massless particle according the $\mathrm{U}(1)_R$ symmetry breaking. Let us introduce explicit $\mathrm{U}(1)_R$ breaking terms,

$$W_{\not R} = m_\phi Tr_N(\widetilde{\phi}\phi) + c, \tag{2.2}$$

and discuss the superpotential as $W_{\mathrm{SB}} = W_R + W_{\not R}$. In ref. [23], W_R is generated, considering the dual side of $\mathrm{SU}(N_F)_F \times \mathrm{SU}(N + N_F)$ gauge theory with the N_F vector-like pairs $(q_d, \widetilde{q}_d)$ of $\mathrm{SU}(N + N_F)$ gauge group. Φ is interpreted as the composite operator as $\Phi \equiv \widetilde{q}_d q_d$, and $h\Lambda_G \mathrm{Tr}_{N_F}(\Phi)$ in W_R corresponds to the mass term of the $(q_d, \widetilde{q}_d)$.

Some ideas to induce $W_{\not R}$ have been proposed in ref. [26], where the small wave-function factor of Φ suppresses Φ^2 and Φ^3 terms according to the strong dynamics or the profile in the extra dimension. In ref. [27], the effect of the explicit R-symmetry breaking terms is well studied. Here, we simply start the discussion from the superpotential W_{SB} assuming that such a mechanism, as discussed in ref. [26], works in underlying theories above the GUT, and study the symmetry breaking. In the global SUSY with canonical Kähler potential, the scalar potential is given by $V = |\partial_\Phi W_{\mathrm{SB}}|^2 + |\partial_\phi W_{\mathrm{SB}}|^2 + |\partial_{\widetilde{\phi}} W_{\mathrm{SB}}|^2$, and SUSY vacua satisfy $\partial_\Phi W_{\mathrm{SB}} = \partial_\phi W_{\mathrm{SB}} = \partial_{\widetilde{\phi}} W_{\mathrm{SB}} = 0$. In this model, $\partial_\Phi W_{\mathrm{SB}}$ is given by

$$\partial_{\Phi_{ji}} W_{\mathrm{SB}} = -h(\phi\widetilde{\phi})_{ij} + h\Lambda_G \delta_{ij}, \tag{2.3}$$

and all elements cannot be vanishing, because $N_F \times N_F$ matrix $(\phi\widetilde{\phi})$ has the rank N $(< N_F)$. This means that SUSY is broken by the F-components of $(N_F - N)$ elements in Φ and $\mathrm{SU}(N_F)_F$ would be also broken.

Following ref. [23], we decompose Φ and $(\phi, \tilde{\phi})$ as

$$\Phi = \begin{pmatrix} (v_Y)\,\mathbf{1}_N + \hat{Y} & \widetilde{Z} \\ Z & (v_X)\,\mathbf{1}_{\widetilde{N}} + \hat{X} \end{pmatrix}, \tag{2.4}$$

$$\phi = \begin{pmatrix} (v_\chi)\,\mathbf{1}_N + \hat{\chi} \\ \rho \end{pmatrix}, \qquad \widetilde{\phi}^T = \begin{pmatrix} (v_\chi)\,\mathbf{1}_N + \hat{\tilde{\chi}}^T \\ \widetilde{\rho}^T \end{pmatrix}, \tag{2.5}$$

where $\hat{Y}$, $\hat{\chi}$ and $\hat{\tilde{\chi}}$ are $N \times N$ matrices, $\hat{X}$ is an $\widetilde{N} \times \widetilde{N}$ matrix ($\widetilde{N} = N_F - N$), Z and ρ ($\widetilde{Z}$ and $\widetilde{\rho}$) are $N \times \widetilde{N}$ matrices ($\widetilde{N} \times N$ matrices). The VEVs, v_Y and v_χ, are fixed by the stationary conditions

$$v_Y = \frac{m_\phi}{h}, \tag{2.6}$$

$$v_\chi = \Lambda_G. \tag{2.7}$$

This solution also satisfies the D-flat conditions. v_X is a flat direction in global SUSY. If we consider gravity and one-loop corrections, it would be stabilized at the nonzero value [23, 28].

The nonzero VEVs break $\mathrm{SU}(N) \times \mathrm{SU}(N_F)_F \times \mathrm{U}(1)_\phi$ gauge symmetry to $\mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_D \times \mathrm{U}(1)_Y$. $\mathrm{SU}(N)_D$ and $\mathrm{U}(1)_Y$ are the linear combinations of the subgroups of $\mathrm{SU}(N_F)_F$ and $\mathrm{SU}(N) \times \mathrm{U}(1)_\phi$.

2.1 Gauge bosons

After the symmetry breaking, massive gauge bosons appear according to the Higgs mechanism. Let us decompose the vector field (V_F^μ) for $\mathrm{SU}(N_F)_F$ as

$$V_F^\mu = \begin{pmatrix} W_F^\mu - aB'^\mu & \frac{1}{\sqrt{2}}(X^\mu)^\dagger \\ \frac{1}{\sqrt{2}}X^\mu & G^\mu + \frac{N}{\widetilde{N}}aB'^\mu \end{pmatrix}, \tag{2.8}$$

where $a = \frac{\sqrt{\widetilde{N}}}{\sqrt{2N(N+\widetilde{N})}}$ is defined. W_F^μ and G^μ are the adjoint representations of the subgroups of $\mathrm{SU}(N_F)_F$: $\mathrm{SU}(N)_F$ and $\mathrm{SU}(\widetilde{N})$. X_μ is the anti-fundamental and fundamental representations of $\mathrm{SU}(N)_F \times \mathrm{SU}(\widetilde{N})$, and B'_μ is the $\mathrm{U}(1)_F$ vector field, where $\mathrm{U}(1)_F$ is from $\mathrm{SU}(N_F)_F$.

The nonzero VEVs generate the following mass terms,

$$\mathcal{L}_g = M_X^2 X_\mu^\dagger X^\mu + \frac{1}{2}M_{W'}^2 W_\mu'^A W'^{A\mu} + \frac{1}{2}M_{Z'}^2 Z'_\mu Z'^\mu, \tag{2.9}$$

$$M_X^2 = g_F^2(v_\chi^2 + \Delta v^2), \tag{2.10}$$

$$M_{W'}^2 = 2(g_F^2 + g_N'^2)v_\chi^2, \tag{2.11}$$

$$M_{Z'}^2 = 4N(Q_\phi^2 g_\phi^2 + a^2 g_F^2)v_\chi^2, \tag{2.12}$$

	Z	$\widetilde{Z}$	ρ	$\widetilde{\rho}$	Y	χ	$\widetilde{\chi}$	X
$\mathrm{SU}(\widetilde{N})$	$\widetilde{\mathbf{N}}$	$\overline{\widetilde{\mathbf{N}}}$	$\widetilde{\mathbf{N}}$	$\overline{\widetilde{\mathbf{N}}}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{adj}_{\widetilde{N}}$
$\mathrm{SU}(N)_D$	$\overline{\mathbf{N}}$	$\mathbf{N}$	$\overline{\mathbf{N}}$	$\mathbf{N}$	$\mathbf{adj}_N$	$\mathbf{adj}_N$	$\mathbf{adj}_N$	$\mathbf{1}$
$\mathrm{U}(1)_Y$	$\frac{N+\widetilde{N}}{N\widetilde{N}}$	$\frac{-N-\widetilde{N}}{N\widetilde{N}}$	$\frac{N+\widetilde{N}}{N\widetilde{N}}$	$\frac{-N-\widetilde{N}}{N\widetilde{N}}$	0	0	0	0

Table 2. Extra Chiral superfields charged under the $\mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_D \times \mathrm{U}(1)_Y$.

where $\Delta v = v_X - v_Y$ is defined. $W'^{A\mu}$ and Z'^{μ} are given by the linear combinations of $W_F^{A\mu}$ and $\mathrm{SU}(N)$ gauge boson ($W_N^{A\mu}$), and B'^{μ} and $\mathrm{U}(1)_\phi$ gauge boson (A_ϕ^μ) respectively:

$$\begin{pmatrix} B'^{\mu} \\ A_\phi^\mu \end{pmatrix} = \begin{pmatrix} \cos\theta_Y & \sin\theta_Y \\ -\sin\theta_Y & \cos\theta_Y \end{pmatrix} \begin{pmatrix} B^\mu \\ Z'^{\mu} \end{pmatrix}, \tag{2.13}$$

$$\begin{pmatrix} W_F^{A\mu} \\ W_N^{A\mu} \end{pmatrix} = \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} W^{A\mu} \\ W'^{A\mu} \end{pmatrix}, \tag{2.14}$$

where $\cos\theta_Y$ and $\cos\theta$ are defined as

$$\cos\theta_Y = \frac{Q_\phi g_\phi}{\sqrt{Q_\phi^2 g_\phi^2 + a^2 g_F^2}}, \quad \cos\theta = \frac{g'_N}{\sqrt{g_N'^2 + g_F^2}}. \tag{2.15}$$

G_μ, W_μ, and B_μ are the gauge bosons for $\mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_D \times \mathrm{U}(1)_Y$ gauge symmetry, and their gauge couplings are given by

$$g_N = g_F \cos\theta, \ g_{\widetilde{N}} = g_F, \ g'_1 = aNg_1 = aNg_F \cos\theta_Y. \tag{2.16}$$

2.2 SM-charged fields from symmetry breaking sector

According to the decomposition in eqs. (2.4) and (2.5), we introduce the charge assignment of $(Z, \widetilde{Z})$, $(\rho, \widetilde{\rho})$, Y, $(\chi, \widetilde{\chi})$, and X in table 2. Y, $(\chi, \widetilde{\chi})$, and X are the adjoint parts of $\hat{Y}$, $(\hat{\chi}, \hat{\widetilde{\chi}})$, and $\hat{X}$. The singlet parts are not charged under the SM, and they are not so relevant to our analysis. The mass matrices are studied in appendix A.

These fields obtain masses according to the nonzero VEVs, v_χ, v_Y and v_X as we see in the appendix A. They decouple at some scales above the EW scale. In the next subsection, we investigate the RG flows of the gauge couplings including the threshold corrections and discuss the soft SUSY breaking terms mediated by the heavy fields.

2.3 RG flows of the gauge couplings

In this model, two kinds of symmetry breaking actually happen: one is $\mathrm{SU}(N_F)_F$ breaking, $\mathrm{SU}(N_F)_F \to \mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_F \times \mathrm{U}(1)_F$, and the other is $\mathrm{SU}(N)_F \times \mathrm{U}(1)_F \times \mathrm{SU}(N) \times \mathrm{U}(1)_\phi$ breaking: $\mathrm{SU}(N)_F \times \mathrm{SU}(N) \to \mathrm{SU}(N)_D$ and $\mathrm{U}(1)_F \times \mathrm{U}(1)_\phi \to \mathrm{U}(1)_Y$. The former is caused by Δv, and the later is by v_χ. We consider a simple scenario assuming $\Delta v \gg v_\chi$.

As we see in appendix A, there will be several intermediate scales, where heavy particles in the symmetry breaking sector are decoupled and the RG flow of gauge couplings is modify. According to the one-loop RG equations, the gauge couplings at the EW scale (M_Z) are evaluated as follows: $\mathrm{SU}(N)_F$, $\mathrm{SU}(N)$ and $\mathrm{SU}(N)_D$ gauge couplings ($\alpha_{F_N}, \alpha'_N, \alpha_N$) are

$$4\pi\alpha_N^{-1}(M_Z) = 4\pi\alpha_{F_N}^{-1}(T_{\chi_N}) + 4\pi\alpha_N'^{-1}(T_{\chi_N}) + b_N \ln\left(\frac{M_Z^2}{T_{\chi_N}^2}\right) + \Delta b_{\rm ex}^N \left(\frac{T_{\rm ex}^2}{\Lambda^2}\right), \tag{2.17}$$

$$4\pi\alpha_{F_N}^{-1}(T_{\chi_N}) = 4\pi\alpha_G^{-1}(\Lambda) + \Delta b_N \ln\left(\frac{T_N^2}{\Lambda^2}\right) + (b_N - 2N)\ln\left(\frac{T_{\chi_N}^2}{\Lambda^2}\right), \tag{2.18}$$

$$4\pi\alpha_N'^{-1}(T_{\chi_N}) = 4\pi\alpha_N'^{-1}(\Lambda) + \Delta b_{\rho_N} \ln\left(\frac{T_{\rho_N}^2}{\Lambda^2}\right) - N\ln\left(\frac{T_{\chi_N}^2}{\Lambda^2}\right). \tag{2.19}$$

$\mathrm{SU}(\widetilde{N})$ gauge coupling ($\alpha_{\widetilde{N}}$) is

$$\begin{aligned}4\pi\alpha_{\widetilde{N}}^{-1}(M_Z) &= 4\pi\alpha_G^{-1}(\Lambda) + b_{\widetilde{N}} \ln\left(\frac{M_Z^2}{\Lambda^2}\right) + \Delta b_{\widetilde{N}} \ln\left(\frac{T_{\widetilde{N}}^2}{\Lambda^2}\right)\\ &\quad + \Delta b_{\rm ex}^{\widetilde{N}} \ln\left(\frac{T_{\rm ex}^2}{\Lambda^2}\right) + \Delta b_{\rho_{\widetilde{N}}} \ln\left(\frac{T_{\rho_{\widetilde{N}}}^2}{\Lambda^2}\right) + \Delta b_X \ln\left(\frac{T_{\chi_{\widetilde{N}}}^2}{\Lambda^2}\right).\end{aligned} \tag{2.20}$$

$\mathrm{U}(1)_F$, $\mathrm{U}(1)_\phi$, and $\mathrm{U}(1)_Y$ gauge couplings ($\alpha_{F_1}, \alpha_\phi, \alpha_1$) are

$$4\pi\alpha_1^{-1}(M_Z) = 4\pi\alpha_{F_1}^{-1}(T_{\chi_1}) + \frac{4\pi a^2}{Q_\phi^2}\alpha_\phi^{-1}(T_{\chi_1}) + b_1 \ln\left(\frac{M_Z^2}{T_{\chi_1}^2}\right) + \Delta b_{\rm ex}^1 \ln\left(\frac{T_{\rm ex}^2}{\Lambda^2}\right), \tag{2.21}$$

$$4\pi\alpha_{F_1}^{-1}(T_{\chi_1}) = 4\pi\alpha_G^{-1}(\Lambda) + \Delta b_1 \ln\left(\frac{T_1^2}{\Lambda^2}\right)! + \Delta b_{\rho_1} \ln\left(\frac{T_{\rho_1}^2}{\Lambda^2}\right) + (b_1 + \Delta b_{\chi_1})\ln\left(\frac{T_{\chi_1}^2}{\Lambda^2}\right), \tag{2.22}$$

$$4\pi\alpha_\phi^{-1}(T_{\chi_1}) = 4\pi\alpha_\phi^{-1}(\Lambda) + \Delta b_{\rho_\phi} \ln\left(\frac{T_{\rho_1}^2}{\Lambda^2}\right) + \Delta b_{\chi_\phi} \ln\left(\frac{T_{\chi_1}^2}{\Lambda^2}\right). \tag{2.23}$$

Λ is the cut-off scale and T_i, T_{χ_i} and T_{ρ_i} ($i = N, \widetilde{N}, 1$) are the intermediate scales where X_μ, $\chi_i(\chi_{\widetilde{N}} \equiv X)$, and ρ_i decouple respecitvely. According to the mass spectrums at each scale in appendix A, T_i, T_{χ_i} and T_{ρ_i} ($i = N, \widetilde{N}, 1$) are estimated as

$$(T_N, T_{\rho_N}, T_{\chi_N}) = (M_X, h\Delta v, \sqrt{h}M_{G'}), \tag{2.24}$$

$$(T_{\widetilde{N}}, T_{\rho_{\widetilde{N}}}, T_{\chi_{\widetilde{N}}}) = (M_X, h\Delta v, m_X), \tag{2.25}$$

$$(T_1, T_{\rho_1}, T_{\chi_1}) = (M_X, h\Delta v, \sqrt{h}M_{Z'}). \tag{2.26}$$

The factor in front of each intermediate scale describes the freedom of the particles decoupling at the scale:

$$(\Delta b_N, \Delta b_{\rho_N}) = (2(N_F - N), -\widetilde{N}), \tag{2.27}$$

$$(\Delta b_{\widetilde{N}}, \Delta b_{\rho_{\widetilde{N}}}, \Delta b_X) = (2(N_F - \widetilde{N}), -N, -\widetilde{N}), \tag{2.28}$$

$$(\Delta b_1, \Delta b_{\rho_1}, \Delta b_{\chi_1}) = \left(2N_F, -a^2N^2\frac{2N}{\widetilde{N}}, -2a^2N^2\right), \tag{2.29}$$

$$(\Delta b_{\rho_\phi}, \Delta b_{\chi_\phi}) = (-2N\widetilde{N}Q_\phi^2, -2N^2Q_\phi^2). \tag{2.30}$$

We may also have to introduce additional particles charged under the gauge symmetry, in order to achieve realistic mass spectrums. For instance, colored Higgs would be necessary to derive the MSSM Higgs doublet at the low scale in section 3, and it is charged under $\mathrm{SU}(\widetilde{N}) \times \mathrm{U}(1)_F$ in our explicit model. Such an extra intermediate scale and the coefficient is defined as T_{ex} and Δb^J_{ex} $(J = \widetilde{N}, 1)$.

We also study the soft SUSY breaking terms of sfermions in the next subsection. Let us also introduce the wave function renormalization factor (Z_q) for $\mathrm{SU}(N)_F$-charged field (q). The one-loop renormalization group for Z_q can be integrated analytically, if the Yukawa coupling is negligible,

$$\begin{aligned}\ln Z_q(M_Z) = \ln Z_q(\Lambda) &+ \frac{2c^q_G}{b_G} \ln\left(\frac{\alpha_G(\Lambda)}{\alpha_G(T_i)}\right) + \frac{2c^q_i}{b_G - \Delta b_i} \ln\left(\frac{\alpha_{F_i}(T_i)}{\alpha_{F_i}(T_{\mathrm{ex}})}\right) \\ &+ \frac{2c^q_i}{b_G - \Delta b_i - \Delta b^i_{\mathrm{ex}}} \ln\left(\frac{\alpha_{F_i}(T_{\mathrm{ex}})}{\alpha_{F_i}(T_{\rho_i})}\right) \\ &+ \frac{2c^q_i}{b_G - \Delta b_i - \widetilde{\Delta b_{\rho_i}} - \Delta b^i_{\mathrm{ex}}} \ln\left(\frac{\alpha_{F_i}(T_{\rho_i})}{\alpha_{F_i}(T_{\chi i})}\right) + \frac{2c^q_i}{b_i} \ln\left(\frac{\alpha_i(T_{\chi i})}{\alpha_i(M_Z)}\right), \end{aligned} \tag{2.31}$$

where $(\widetilde{\Delta b_{\rho_N}}, \widetilde{\Delta b_{\rho_{\widetilde{N}}}}, \widetilde{\Delta b_{\rho_1}}) = (0, \Delta b_{\rho_{\widetilde{N}}}, \Delta b_{\rho_1})$ is defined and $T_i \geq T_{\mathrm{ex}} \geq T_{\rho_i}$ is assumed. c^q_G and c^q_i are the second Casimir of the field q, corresponding to the gauge groups. The masses squared of sfermions can be derived by the v_X-dependence in Z_q. v_X appears in the gauge couplings, so that v_X-dependence on the gauge couplings is only relevant to the sfermion masses [29].

2.4 Soft SUSY breaking terms

Based on the above results, we investigate soft SUSY breaking terms which relate to particles charged under the gauge symmetry. Soft SUSY breaking terms in $\mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_D \times \mathrm{U}(1)$ are calculated by substituting $v_X + \theta^2 F_X$ for v_X in the gauge couplings [29]. Compared to typical gauge mediation, where messengers are only chiral superfields, massive gauge bosons and the fermionic partners also work as the mediators to generate the soft SUSY breaking terms, in our models [24, 25, 30, 31].

In eqs. (2.17) , (2.20), and (2.21), the only intermediate scales, T_i, T_{ρ_i}, and T_{ex} depend on v_X. This leads the masses $(M_{\widetilde{N}}, M_N, M_1)$ of the gauginos, which are the superpartner of $\mathrm{SU}(N)_D \times \mathrm{SU}(\widetilde{N}) \times \mathrm{U}(1)$ gauge bosons, as follows:

$$M_N(\mu) = -(\Delta b_N + \Delta b_{\rho_N} + \Delta b^N_{\mathrm{ex}} \xi_N) \frac{\alpha_N(\mu)}{4\pi} \frac{F_X}{|\Delta v|}, \tag{2.32}$$

$$M_{\widetilde{N}}(\mu) = -(\Delta b_{\widetilde{N}} + \Delta b_{\rho_{\widetilde{N}}} + \Delta b^{\widetilde{N}}_{\mathrm{ex}} \xi_{\widetilde{N}}) \frac{\alpha_{\widetilde{N}}(\mu)}{4\pi} \frac{F_X}{|\Delta v|}, \tag{2.33}$$

$$M_1(\mu) = -\left(\Delta b_1 + \Delta b_{\rho_1} + \frac{a^2}{Q^2_\phi} \Delta b_{\rho_\phi} + \Delta b^1_{\mathrm{ex}} \xi_1\right) \frac{\alpha_1(\mu)}{4\pi} \frac{F_X}{|\Delta v|}. \tag{2.34}$$

ξ_N, $\xi_{\widetilde{N}}$, and ξ_1 describe the v_X dependence on the mass scale of extra particles, T_{ex}. For example, the holomorphic mass of extra particles may be given by $m_{\mathrm{ex}} + \lambda_{\mathrm{ex}}(v_X + \theta^2 F_X)$, where m_{ex} and λ_{ex} are a supersymmetric mass term and Yukawa coupling involving the extra particles. That is, the gaugino mass contribution of $\ln(T_{\rm ex})$ would be proportional to $\lambda_{\rm ex}F_X/m_{\rm ex}$, if $m_{\rm ex}$ is larger than $\lambda_{\rm ex}v_X$. In this case, ξ_i is approximately given by $\xi_i = \lambda_{\rm ex}|\Delta v|/m_{\rm ex}$.

Let us consider the soft SUSY breaking terms corresponding to the trilinear (A-term) and bilinear couplings (B-term) of the scalar components of the SU$(N_F)_F$-charged fields (q_I). They are relevant to the v_X-dependence of the wave renormalization factor. For instance, the A-terms corresponding to the Yukawa couplings $y_{IJK}q_Iq_Jq_K$ in the superpotential are given by $A_{IJK} = A_I + A_J + A_K$, where $A_I = \frac{\partial \ln Z_I}{\partial \ln v_X}$ is defined and the trilinear coupling is described as $y_{IJK}A_{IJK}q_Iq_Jq_K$.

Eventually, A_I is obtained from eq. (2.31),

$$
\begin{aligned}
A_I = \Bigg\{ & 2c_G^I \frac{\alpha_G(T_i)}{4\pi} - \frac{2b_G c_i^I}{b_G - \Delta b_i}\frac{\alpha_{F_i}(T_i)}{4\pi} \\
& + \left(\frac{2c_i^I}{b_G-\Delta b_i} - \frac{2c_i^I}{b_G - \Delta b_i - \Delta b^i_{\rm ex}}\right)(b_G\xi + \Delta b_i(1-\xi))\frac{\alpha_{F_i}(T_{\rm ex})}{4\pi} \\
& + \left(\frac{2c_i^I}{b_G-\Delta b_i-\Delta b^i_{\rm ex}} - \frac{2c_i^I}{b_G - \Delta b_i - \Delta b^i_{\rm ex} - \widetilde{\Delta b}_{\rho_i}}\right)(b_G - \Delta b^i_{\rm ex}(1-\xi))\frac{\alpha_{F_i}(T_{\rho_i})}{4\pi} \\
& + \left(\frac{2(\Delta b_i + \Delta b^i_{\rm ex}\xi + \widetilde{\Delta b}_{\rho_i})c_i^I}{b_G - \Delta b_i - \Delta b^i_{\rm ex} - \widetilde{\Delta b}_{\rho_i}}\right)\frac{\alpha_{F_i}(T_{\chi i})}{4\pi} - \left(\frac{2(\Delta b_i + \Delta b^i_{\rm ex}\xi + \Delta b'_{\rho_i})c_i^I}{b_i}\right)\frac{\alpha_i(T_{\chi i})}{4\pi} \\
& + \left(\frac{2(\Delta b_i + \Delta b^i_{\rm ex}\xi + \Delta b'_{\rho_i})c_i^I}{b_i}\right)\frac{\alpha_i(\mu)}{4\pi}\Bigg\}\frac{F_X}{|\Delta v|},
\end{aligned}
\tag{2.35}
$$

assuming $\xi_N = \xi_{\widetilde{N}} = \xi_1 = \xi$. $\alpha_{F_{\widetilde{N}}} \equiv \alpha_{\widetilde{N}}$ is defined.

The masses squared (m_q^2) of q could be also estimated by the eq. (2.31), seeing the $|v_X|^2$-dependence of Z_q [29]. As discussed in ref. [30], the gauge mediation with gauge messengers may contribute to the masses squared at the one-loop level, if the gauge symmetry breaking and SUSY breaking are caused by the VEVs and F-components of several fields. In our case, we simply assume $v_\chi \ll \Delta v$, so that the gauge symmetry breaking and SUSY breaking are caused by only Δv and the F-component of Δv.[3] The one-loop correction is strongly suppressed by $(v_\chi/\Delta v)^2$ according to ref. [30], so that we have to investigate the two-loop corrections, as discussed in refs. [24, 29].

Following refs. [24, 29], m_q^2 could be written as

$$
\begin{aligned}
m_q^2(\mu) = \Bigg\{ & 2b_G c_G^q \frac{\alpha_G^2(T_i)}{(4\pi)^2} - \frac{2b_G^2 c_i^q}{b_G - \Delta b_i}\frac{\alpha_{F_i}^2(T_i)}{(4\pi)^2} \\
& + \left(\frac{2c_i^q}{b_G-\Delta b_i} - \frac{2c_i^q}{b_G - \Delta b_i - \Delta b^i_{\rm ex}}\right)(b_G\xi + \Delta b_i(1-\xi))^2\frac{\alpha_{F_i}^2(T_{\rm ex})}{(4\pi)^2} \\
& + \left(\frac{2c_i^q}{b_G-\Delta b_i-\Delta b^i_{\rm ex}} - \frac{2c_i^q}{b_G - \Delta b_i - \Delta b^i_{\rm ex} - \widetilde{\Delta b}_{\rho_i}}\right)(b_G - \Delta b^i_{\rm ex}(1-\xi))^2\frac{\alpha_{F_i}^2(T_{\rho_i})}{(4\pi)^2}
\end{aligned}
$$

[3] Δv corresponds to the VEV of one adjoint field.

$$
\begin{aligned}
&+\left(\frac{2(\Delta b_i+\Delta b^i_{\rm ex}\xi+\widetilde{\Delta b_{\rho_i}})^2 c_i^q}{b_G-\Delta b_i-\Delta b^i_{\rm ex}-\widetilde{\Delta b_{\rho_i}}}\right)\frac{\alpha^2_{F_i}(T_{\chi i})}{(4\pi)^2}\\
&-\left(\frac{2(\Delta b_i+\Delta b^i_{\rm ex}\xi+\Delta b'_{\rho_i})^2 c_i^q}{b_i}\right)\frac{\alpha^2_{i}(T_{\chi i})}{(4\pi)^2}\\
&+\left(\frac{2(\Delta b_i+\Delta b^i_{\rm ex}\xi+\Delta b'_{\rho_i})^2 c_i^q}{b_i}\right)\frac{\alpha^2_{i}(\mu)}{(4\pi)^2}\Bigg\}\frac{F_X^2}{|\Delta v|^2},
\end{aligned}
\tag{2.36}
$$

where $\Delta b'_{\rho_i}$ is $(\Delta b'_{\rho_3}, \Delta b'_{\rho_2}, \Delta b'_{\rho_1}) = (\Delta b_{\rho_3}, \Delta b_{\rho_2}, \Delta b_{\rho_1} + a^2 \Delta b_{\rho_\phi}/Q_\phi^2)$.

In the next section, we discuss one explicit model, where $\mathrm{SU}(\widetilde{N}) \times \mathrm{SU}(N)_D \times \mathrm{U}(1)_Y$ is the SM gauge groups corresponding to $(N_F, N) = (5, 2)$. In the explicit model, we see that a few parameters control all soft SUSY breaking terms according to this analysis. Then, $\Lambda_{\rm SUSY}$ is roughly given by $(\alpha_G/(4\pi)) \times (F_X/|\Delta v|)$, and A-term and B-term are of $O(\Lambda_{\rm SUSY})$, which we could expect that are consistent with the condition for the EW symmetry breaking. We study the compatibility with the EW condition and the Higgs mass, in section 3.4.

3 $\mathbf{SU(5)_F \times SU(2) \times U(1)_\phi}$ gauge theory: $\boldsymbol{(N_F, N) = (5, 2)}$

In this section, we consider a $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ gauge symmetric model, which correspond to the $(N_F, N) = (5, 2)$ case. We expect that the MSSM fields are embedded in to **10** and $\overline{\mathbf{5}}$ representation as in the Georgi-Glashow SU(5) GUT. Involving **5**-representation Higgs $(H, \overline{H})$, the superpotential for the Yukawa couplings in the visible sector is

$$
W_{vis} = \hat{y}^u_{kl} H \mathbf{10}^k \mathbf{10}^l + \hat{y}^d_{kl} \overline{H}\,\overline{\mathbf{5}}^k \mathbf{10}^l, \tag{3.1}
$$

where $\overline{\mathbf{5}}^k$ and $\mathbf{10}^l$ are defined as the matter fields. As well-known, $\hat{y}^u_{kl}$ and $\hat{y}^d_{kl}$ may require Φ and $(\phi, \widetilde{\phi})$ dependences in order to generate realistic mass matrices at the EW scale according to the higher-dimensional operators. Here, we simply assume that the contributions to the soft SUSY breaking terms are enough small.

One serious problem in the SU(5) GUT is how to generate the mass splitting between the colored Higgs and the MSSM Higgs doublet. The mass of colored Higgs should be around the GUT scale to avoid the too short life time of proton: $m_{H_c} \gtrsim 10^{16}\mathrm{GeV} \times (1\mathrm{TeV}/\Lambda_{\rm SUSY})$ [32, 33]. In our $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ model, the relevant terms to the Higgs masses is written as

$$
W_H = \mu \overline{H} H + \lambda_H \overline{H} \Phi H. \tag{3.2}
$$

After the symmetry breaking, the colored Higgs mass and MSSM Higgs mass are given by $\mu + \lambda v_X$ and $\mu + \lambda_H v_Y$. If $v_Y = m_\phi/h$ is the GUT scale, μ should be also around the GUT scale and then the fine-tuning between μ and λv_Y is required: $\mu + \lambda_H v_Y \approx O(M_Z)$. On the other hand, we could expect that the colored Higgs is enough heavy because of μ, if there is no cancellation between μ and $\lambda_H v_X$. Let us also consider the case that v_X is the GUT scale. In this case, the MSSM Higgs mass could be light if μ and v_Y are around the weak scale, and the colored Higgs is heavy: $m_{H_c} \approx \lambda_H v_X$.

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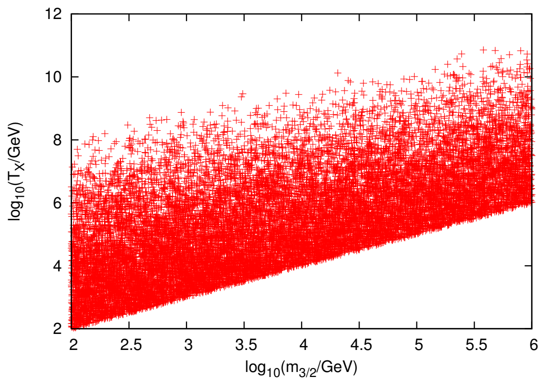

Figure 1. Gravitino mass ($m_{3/2}$) and the scale T_X with $10^{16}\,\text{GeV} \leq T_G \leq M_p$. T_X should be small to rase the GUT scale above 10^{16} GeV. T_χ is $T_\chi \approx \sqrt{M_p m_{3/2}}$. The constraints, $T_\rho > T_\chi$ and $T_X > m_{3/2}$, are also assigned. All gauge couplings and Yukawa couplings satisfy the perturbative bounds as $\alpha_{F_i} < 4\pi$.

In both cases, the colored Higgs couples with $v_X + F_X\theta^2$, so it mediates the SUSY breaking effect to the soft SUSY breaking terms. The supersymmetric mass for $\text{SU}(2)_L$ Higgs doublet is $\mu_2 = \mu + \lambda_H v_Y \approx O(M_Z)$. On the other hand, the colored-Higgs mass is $m_{H_c} = \mu + \lambda_H v_X \approx \lambda_H(v_Y - v_X)$, so that ξ for the colored Higgs in soft SUSY breaking terms is approximately estimated as $\xi \approx sign(\lambda_H)$. The one-loop correction of H_c to m_q^2 would be suppressed, because the m_{H_c}-dependence appears in Z_q as $\ln(|m_{H_c} + \lambda_H F_X \theta^2|^2)$ according to the study in ref. [29]. We could apply our analysis in section 2.4 to this scenario.

3.1 Gauge couplings

In this model, $\text{SU}(5)_F \times \text{SU}(2) \times \text{U}(1)_\phi$ breaks down to the SM gauge group, $\text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y$. $\text{SU}(3)_c$ is the subgroup of $\text{SU}(5)_F$ and $\text{SU}(2)_L \times \text{U}(1)_Y$ are the linear combination of $\text{SU}(2)_F \times \text{U}(1)_F$ and $\text{SU}(2) \times \text{U}(1)_\phi$.

On the other hand, there are several intermediate scales: $(T_G, T_\rho, T_\chi, T_X)$.[4] T_G is the GUT scale, where X_μ decouples, and T_ρ is the messenger scale fixed by the parameter h and the GUT scale. T_χ is interpreted as the SUSY breaking scale, because $T_\chi \approx \sqrt{F_X} = \sqrt{M_p m_{3/2}}$, so that it is almost fixed around $O(10^{10})$ GeV when $m_{3/2} = O(100)$ GeV. T_X is fixed by the mass scale of X (m_X), which is massless at the tree-level. X could be expected to be $O(\Lambda_{\text{SUSY}})$, because the one-loop corrections shift the mass, but it may be difficult to clearly fix the masses of bosonic and fermonic X in our model. Let us simply treat m_X as the free parameter, and figure 1 shows the allowed region for T_X, which may not be far from $O(\Lambda_{\text{SUSY}})$. Figure 2 shows the gauge couplings, $(\alpha_{F2}, \alpha_2, \alpha_{F1}, \alpha_\phi)$ at the SUSY breaking scale. Figure 3 describes RG flows of the gauge couplings $(\alpha_3, \alpha_2(\alpha_{F_2}), \alpha_1(\alpha_{F_1}))$, when $T_X = 10^7$ GeV, $T_\chi = 3.8 \times 10^{10}$ GeV, $T_\rho = 7.9 \times 10^{11}$ GeV, and $T_{\text{GUT}} = 2 \times 10^{16}$ GeV.

[4] $T_{\chi 1} = T_{\chi 2}$, $T_{\rho 1} = T_{\rho 3}$, and $T_G = T_1 = T_2 = T_3$ are assumed.

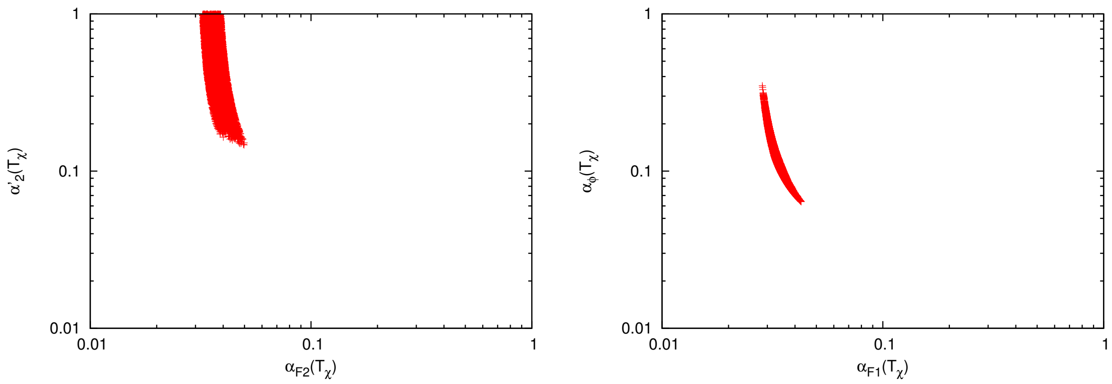

Figure 2. α_{F_2} vs. α'_2 and α_{F_1} vs. α_ϕ at the symmetry breaking scale, T_χ.

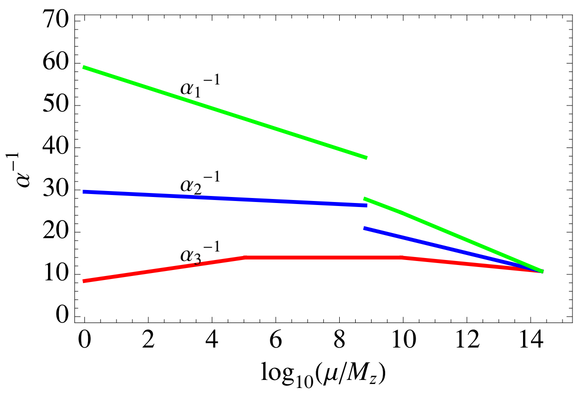

Figure 3. RG flows of the subgroups of $SU(5)_F$, with $T_X = 10^7$ GeV, $T_\chi = 3.8 \times 10^{10}$ GeV, $T_\rho = 7.9 \times 10^{11}$ GeV, and $T_{\rm GUT} = 2 \times 10^{16}$ GeV. The green, blue, and red lines correspond to the gauge couplings of $U(1)_Y \times SU(2)_L \times SU(3)_c$ below T_χ and $U(1)_F \times SU(2)_F \times SU(3)_c$ above T_χ respectively. The input parameters for the couplings are in eq. (3.6).

3.2 Soft SUSY breaking terms

We qualitatively evaluate the soft SUSY breaking terms in this scenario. According to the analysis in section 2.4, the gaugino masses at $\mu < T_\chi$ are written as

$$M_2(\mu) = -3\frac{\alpha_2(\mu)}{4\pi}\frac{F_X}{|\Delta v|}, \tag{3.3}$$

$$M_3(\mu) = -(2-\xi)\frac{\alpha_3(\mu)}{4\pi}\frac{F_X}{|\Delta v|}, \tag{3.4}$$

$$M_1(\mu) = -\left(\frac{37-2\xi}{5}\right)\frac{\alpha_1(\mu)}{4\pi}\frac{F_X}{|\Delta v|}. \tag{3.5}$$

Let us consider the case with $\xi = 1$ and the gaugino masses at the EW scale. The gauge couplings at the EW scale are [3]

$$\alpha_1(M_Z) \approx 0.01695,\ \ \alpha_2(M_Z) \approx 0.03382,\ \ \alpha_3(M_Z) \approx 0.1185, \tag{3.6}$$

so that we could derive the following mass relation:

$$\frac{M_1(M_Z)}{M_3(M_Z)} \approx 1.001, \ \frac{M_2(M_Z)}{M_3(M_Z)} \approx 0.856. \tag{3.7}$$

The masses are almost degenerate, and this may be a specific feature of the gauge messenger model [24, 34].[5] If all intermediate scales are close to the GUT scale, the fine-tuning of μ term may be drastically reduced, as discussed in ref. [36]. Figure 2 tells us that the extra SU(3)-adjoint field reside in the low-scale, so that the condition for the small μ-term would be modified. The one-loop running correction of $m^2_{H_u}$ with $T_X = 10^7$ GeV from T_χ to M_Z is estimated as

$$\Delta m^2_{H_u} \approx -0.276 M_3(M_Z)^2 - 0.047 M_2(M_Z) M_3(M_Z) + 0.221 M_2(M_Z)^2 + \dots, \tag{3.8}$$

where the ellipsis denotes the terms including A-term and scalar masses and those are not important when they are comparable to the gluino mass. This leads that the condition to cancel the large contribution of gluino is $M_2/M_3(M_Z) \approx 1.23$, which suggests the almost degenerate mass spectrum. However, we have a large A-term contribution to $\Delta m^2_{H_u}$ in our model, so that it may be difficult to avoid a certain fine-tuning even if the gaugino masses are degenerate.

According to eqs. (2.36) and (2.35), the masses squared of superpartners and A-term are evaluated explicitly. Setting $T_G = T_{H_c} > T_\rho > T_\chi$ and $\xi = 1$, stop masses at T_χ are given by

$$m^2_Q(T_\chi) \approx \left(8.83 - 6.67 \frac{\alpha^2_3(T_\rho)}{\alpha^2_G} - 10.80 \frac{\alpha^2_{F_2}(T_\chi)}{\alpha^2_G} - 0.33 \frac{\alpha^2_{F_1}(T_\chi)}{\alpha^2_G}\right) \Lambda^2_{\rm SUSY}, \tag{3.9}$$

$$m^2_U(T_\chi) \approx \left(8.60 - 6.67 \frac{\alpha^2_3(T_\rho)}{\alpha^2_G} - 5.30 \frac{\alpha^2_{F_1}(T_\chi)}{\alpha^2_G} - 0.01 \frac{\alpha^2_{F_1}(T_\rho)}{\alpha^2_G}\right) \Lambda^2_{\rm SUSY}. \tag{3.10}$$

As we see, large stop masses are generated by the large second casimir ($c^t_2 = 18/5$), but they might be driven to the tachyonic if T_χ and T_ρ are close to the GUT scale. The SUSY scale ($\Lambda_{\rm SUSY}$) from the gauge mediation is defined as

$$\Lambda_{\rm SUSY} = \frac{\alpha_G}{(4\pi)} \frac{F_X}{|\Delta v|} \approx \frac{\alpha_G}{(4\pi)} \frac{M_p}{T_G} m_{3/2}. \tag{3.11}$$

α_G is of $O(0.1)$ when T_G is around 10^{16} GeV, so that $\Lambda_{\rm SUSY}$ might be compatible with $m_{3/2}$. If T_G is smaller, the situation, $\Lambda_{\rm SUSY} \gg m_{3/2}$, is achieved but suffers from the constraint from proton decay. The correction from the gravity mediation is naively estimated as $O(m_{3/2})$. It is almost the same order as the one from the gauge mediation in our model, and it may make it difficult to control flavors. In fact, the gauge-mediation contributions are typically at least 5 times as large as the gravitino mass in our model, as we see in table 4. In this case, we could expect the gravity-mediation effect is sub-dominant, and the SUSY scale is governed by the gauge-mediation. However, the gravity-mediation contribution

[5] The gaugino masses are degenerate in the TeV-scale mirage mediation scenario, too [35].

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should be $O(10^{-2})$ times suppressed, if it contributes to the sparticles masses squared flavor-universally [37]. In order to realize such a suppression and control flavor in the MSSM, we have to consider flavor symmetry or some dynamics above the GUT scale, as discussed in refs. [38–44].[6] Indeed, explicit contributions on soft masses through the gravity mediation depend on the UV completion of our model. In this letter, one of our main motivations is to achieve 125 GeV Higgs mass and realistic EW symmetry breaking, which may be independent of this issue about the constraint from flavor physics, so that we will discuss our SUSY mass spectrums assuming that the gauge-mediation is dominant. The underlying theory above the GUT scale will be studied in ref. [55].

A_t, which is the trilinear coupling of stops $(\widetilde{t})$ as $y_t A_t \widetilde{t}_L H_u \widetilde{t}_R$ is given by

$$A_t(T_\chi) \approx \left(22.57 - 8.00\frac{\alpha_3(T_\rho)}{\alpha_G} - 3.6\frac{\alpha_{F_2}(T_\chi)}{\alpha_G} - 0.98\frac{\alpha_{F_1}(T_\chi)}{\alpha_G} + 0.01\frac{\alpha_{F_1}(T_\rho)}{\alpha_G}\right)\Lambda_{\rm SUSY}, \tag{3.12}$$

and the B-term, which is the bilinear coupling of two Higgs $\mu B H_u H_d$, is estimated as

$$B(T_\chi) \approx \left(10.27 - 3.60\frac{\alpha_{F_2}(T_\chi)}{\alpha_G} - 0.68\frac{\alpha_{F_1}(T_\chi)}{\alpha_G} + 0.01\frac{\alpha_{F_1}(T_\rho)}{\alpha_G}\right)\Lambda_{\rm SUSY}. \tag{3.13}$$

As we see, the A-term and B-term might be large as $O(10)\Lambda_{\rm SUSY}$. This may be good to achieve the EW symmetry breaking, but too large A-term makes the stop masses tachyonic because of the running correction such as

$$\Delta m_U^2(M_Z) \approx -0.08 A_t(T_\chi)^2 + 1.54 M_3(T_\chi)^2 - 0.15 A_t(T_\chi) M_3(T_\chi). \tag{3.14}$$

In our model, the gluino mass M_3 is relatively small as wee see in eq. (3.4), so $\Delta m_U^2(M_Z)$ becomes easily negative and stop mass becomes tachyonic even if the positive m_U^2 is generated at the SUSY breaking scale T_χ. In order to avoid the tachyonic stop masses, we add an extra contribution to the gluino mass, as we see below.

3.3 Shift of the gluino mass

We consider an extra term, which contributes to the gluino mass,

$$W = \frac{1}{\Lambda_0} Tr_5(\Phi W_5 W_5). \tag{3.15}$$

There are several ways to introduce this term, such as gravity effect. Here, we simply assume that $N_{\rm extra}$ extra heavy SU(5) vector-like pairs $(\psi, \overline{\psi})$ with the masses $\overline{\psi}(\Lambda_0 + \lambda_X \Phi)\psi$ induce this term, integrating out them at the scale Λ_0. After the SU(5) breaking, the gauge coupling would have the extra v_X dependence as

$$\alpha_3^{-1} \to \alpha_3^{-1} - \frac{N_{\rm extra}}{4\pi} \ln\left(\frac{(|\Lambda_0 + \lambda_X(v_X + F_X\theta^2)|^2)}{\Lambda^2}\right). \tag{3.16}$$

This additional coupling could shift the gluino mass as

$$M_3 \to M_3 - \frac{\alpha_3 N_{\rm eff}}{4\pi}\frac{F_X}{|\Delta v|}, \tag{3.17}$$

[6] In fact, such strong dynamics has been proposed not only to suppress flavor changing currents but also to realize the superpotential $W_{\rm SB}$ in section 2 [26].

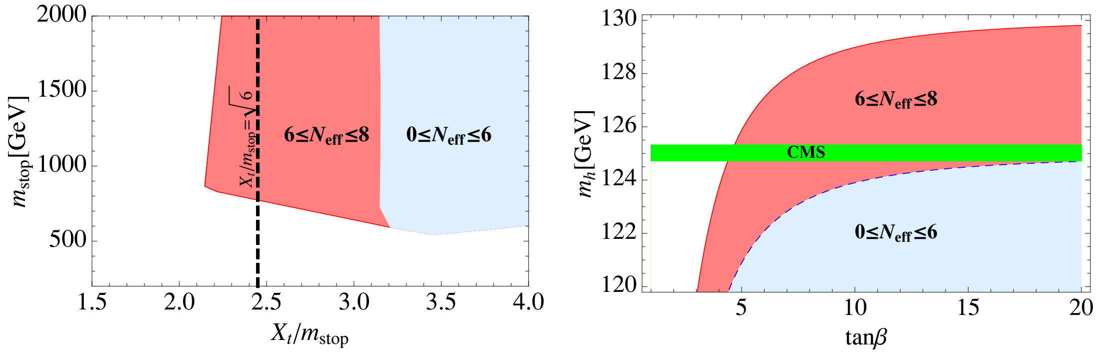

Figure 4. $X_t/m_{\rm stop}$ vs. $m_{\rm stop}$ and $\tan\beta$ vs. the lightest Higgs mass in the case with $(T_{\rm GUT}, T_X) = (2\times 10^{16}\text{GeV},\ 10^7\text{GeV})$ and $0 \leq N_{\rm eff} \leq 6$ (light blue), $6 \leq N_{\rm eff} \leq 8$ (light red). The dashed line corresponds to $X_t/m_{\rm stop} = \sqrt{6}$. In the right figure, m_h is calculated at the two-loop level using $m_t = 172.9\,\text{GeV}$, and $m_{\rm stop}$ is lighter than $2\,\text{TeV}$. The green band is the CMS result on Higgs mass from $h \to \gamma\gamma$, ZZ channels [47].

where $N_{\rm eff}$ may not be $N_{\rm extra}$ because of the scale difference between Λ_0 and the GUT scale. Including $N_{\rm eff}$, the gluino mass becomes

$$M_3(\mu) = -(1 - N_{\rm eff})\frac{\alpha_3(\mu)}{4\pi}\frac{F_X}{|\Delta v|}, \tag{3.18}$$

so $N_{\rm eff}$ should be bigger than 2 in order to shift M_3. In fact, we discuss large $N_{\rm eff}$ cases and find that $N_{\rm eff}$ enables us to evade the negative squared masses and achieve the large SM Higgs mass.

3.4 Consistency with the Higgs mass and the EW symmetry breaking

One issue in supersymmetric models is how to realize the μ and B terms which are consistent with the EW scale. Especially, μ relates to the lightest Higgs mass, because of the upper bound in MSSM, so that the recent Higgs discovery with the mass 125 GeV may impose unnatural SUSY scenarios on us. In fact, 125 GeV Higgs mass may require $\Lambda_{\rm SUSY} \gtrsim O(10)$ TeV in the simple scenarios as discussed in refs. [4, 5]. $O(10)$-TeV SUSY scale would require 0.01% fine-tuning against μ without any cancellation in $m^2_{H_u}$. As pointed out in refs. [45, 46], it is known that a special relation between A_t and squark mass relaxes the fine-tuning, maximizing the loop corrections in the Higgs mass in the MSSM. This relation is so-called "maximal mixing" and described as $X_t/m_{\rm stop} = \sqrt{6}$, where $X_t = A_t - \mu/\tan\beta$ and $m^2_{\rm stop} = \sqrt{m^2_Q m^2_U}$ are defined. If this relation is satisfied, the 125 GeV Higgs mass could be achieved even if the stop is light. We can see our prediction on X_t and the upper bound on the Higgs mass in the case with $0 \leq N_{\rm eff} \leq 6$ (light blue), $6 \leq N_{\rm eff} \leq 8$ (light red) in figure 4. On the all regions, all masses squared of the superpartners are positive and (T_G, T_X) are fixed at $(2\times 10^{16}\text{GeV},\ 10^7\text{GeV})$. We find that our A-term is too large to realize $X_t/m_{\rm stop} = \sqrt{6}$, but the maximal mixing could be achieved, if we allow large $N_{\rm eff}$, and enhance the Higgs mass, even if $m_{\rm stop}$ is around 1 TeV.

On the other hand, we notice that there is no special cancellation in $m^2_{H_u}$ and $m^2_{H_d}$, as we see in figure 5. Large $m_{\rm stop}$ corresponds to large μ, so that 1-TeV squark mass requires

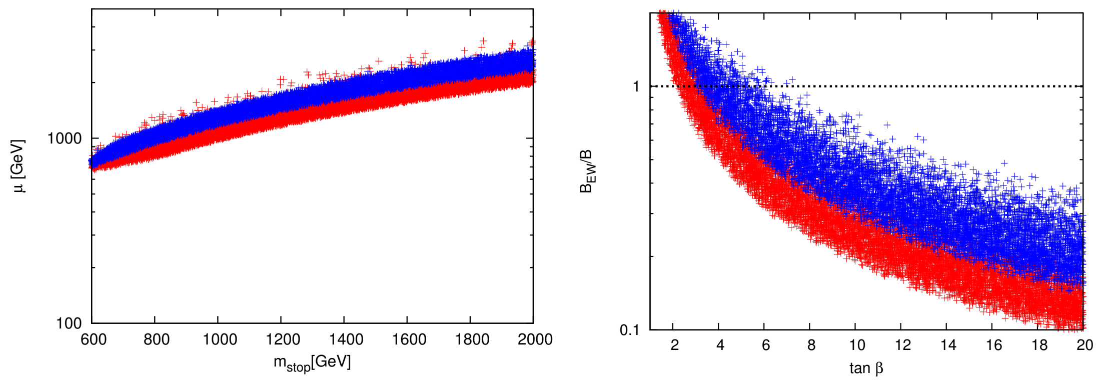

Figure 5. $m_{\rm stop}$ vs. μ and $\tan\beta$ vs. B-term in the case with $(T_{\rm GUT}, T_X) = (2\times 10^{16}{\rm GeV}, 10^7{\rm GeV})$ and $0 \le N_{\rm eff} \le$ (blue), $6 \le N_{\rm eff} \le 8$ (red). In the right figure, $m_{\rm stop}$ is lighter than 2 TeV. The dashed line is consistent with the condition for the EW symmetry breaking.

1% fine-tuning against μ. The right figure in figure 5 shows that small $\tan\beta$ is consistent with the EW symmetry breaking. $B_{\rm EW}$ is the value to realize the EW symmetry breaking,

$$B_{\rm EW} \equiv -\frac{1}{2\mu}\{(m_{H_d}^2 - m_{H_u}^2)\tan 2\beta + M_Z^2 \sin 2\beta\}, \tag{3.19}$$

and B is our prediction via the gauge mediation. It seems that $2 \lesssim \tan\beta \lesssim 6$ is necessary to achieve 125 GeV Higgs mass. The $\tan\beta$ region may be inconsistent with the one required by 125 GeV Higgs ($\tan\beta \gtrsim 4$) with $m_{\rm stop} \le 2$ TeV. Table 4 in appendix B shows the parameter sets in our model, which satisfy $m_h \approx 125$ GeV and $|B_{\rm EW}/B| \approx 1$. There, $m_{\rm stop}$ and $|\mu|$ are around 3 TeV, and $O(0.1)$ % fine-tuning is required against μ term.

4 $\mathbf{SU(5)_F \times SU(3) \times U(1)_\phi}$ gauge theory: $\mathbf{(N_F, N) = (5, 3)}$

Our symmetry breaking model could be embed into other type GUT model. One simple example would be the $\mathrm{SU}(5)_F \times \mathrm{SU}(3) \times \mathrm{U}(1)_\phi$ gauge symmetric model, and we could consider the same setup as in the $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$ gauge theory. The visible sector is given by eq. (3.1). However, the modification of the Higgs sector may be required because $\lambda\overline{H}\Phi H$ term gives the very large B-term, $\lambda F_X H_u H_d$. There may be a solution to realize the EW symmetry breaking, but the serious fine-tuning may be required. Here, we consider another solution to shift the colored Higgs mass which maybe favor high-scale SUSY.

We introduce SU(3) vector-like fields $(H_3, \overline{H}_3)$ and assign Z_3 symmetry to the fields as in table 3. Z_3 symmetry is broken by the VEV of S. The superpotential for the Higgs sector is given by

$$W_H = \lambda_S S\overline{H}H + \lambda_\phi \overline{H}\overline{H}_3\phi + \lambda_{\tilde{\phi}}\tilde{\phi}H_3 H + \frac{\lambda_3}{3}S^3. \tag{4.1}$$

After the GUT symmetry breaking, $(H, \overline{H})$ are decomposed as $((H_u, H_3'), (H_d, \overline{H}_3'))$ and the mass terms for $(H_3, \overline{H}_3')$ and $(H_3', \overline{H}_3)$ pairs appear as

$$W_H^{\rm eff} = \lambda_\phi v_\chi \overline{H}_3'\overline{H}_3 + \lambda_{\tilde{\phi}} v_\chi H_3 H_3'. \tag{4.2}$$

	$\overline{\mathbf{5}}_i$	$\mathbf{10}_i$	H	$\overline{H}$	H_3	$\overline{H}_3$	S	ϕ	$\widetilde{\phi}$	Φ
$\mathrm{SU}(5)_F$	$\overline{\mathbf{5}}$	$\mathbf{10}$	$\mathbf{5}$	$\overline{\mathbf{5}}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{5}$	$\overline{\mathbf{5}}$	$\mathbf{adj}_5+\mathbf{1}$
$\mathrm{SU}(3)$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	$\overline{\mathbf{3}}$	$\mathbf{1}$	$\mathbf{3}$	$\overline{\mathbf{3}}$	$\mathbf{1}$
$\mathrm{U}(1)_\phi$	0	0	0	0	Q_ϕ	$-Q_\phi$	0	Q_ϕ	$-Q_\phi$	0
Z_3	ω	ω	ω	ω	ω^2	ω^2	ω	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$

Table 3. Chiral superfields in $\mathrm{SU}(5)_F \times \mathrm{SU}(3) \times \mathrm{U}(1)$ gauge theory.

H_u and H_d correspond to the Higgs $\mathrm{SU}(2)_L$ doublets in MSSM, and they could get the supersymmetric mass term according to the nonzero VEV of S. In refs. [48–54], we can see not only the $\mathrm{SU}(5)_F \times \mathrm{SU}(2) \times \mathrm{U}(1)_\phi$-type but also this type of product-GUT.

In order to avoid the bound from the proton decay caused by the five dimensional operators, v_χ should be large as

$$v_\chi \gtrsim 10^{16}\mathrm{GeV} \times \left(\frac{1\mathrm{TeV}}{\Lambda_{\mathrm{SUSY}}}\right). \tag{4.3}$$

F_X is given by $-hv_\chi^2$, so that very tiny h is necessary to achieve the low-scale SUSY. When $v_\chi \approx 10^{16}$ GeV and $\Lambda_{\mathrm{SUSY}} = 1$ TeV are set, h should be around $O(10^{-10})$, because of

$$h = \frac{4\pi}{\alpha_G}\frac{\Delta v}{v_\chi^2}\Lambda_{\mathrm{SUSY}} \approx 10^{-10} \times \left(\frac{10^{16}\mathrm{GeV}}{v_\chi}\right)^2 \left(\frac{\Lambda_{\mathrm{SUSY}}}{1\mathrm{TeV}}\right) \lesssim 10^{-10} \times \left(\frac{\Lambda_{\mathrm{SUSY}}}{1\mathrm{TeV}}\right)^3. \tag{4.4}$$

We conclude that high-scale SUSY is favored to avoid such an extremely small h.

We can consider the applications of our symmetry breaking models to the other BSMs, such as

- $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L} \to \mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$,
- $\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{U}(1) \to \mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$.

We would study such patterns elsewhere [55]. In these models, all of chiral superfields appear as adjoint representations and bi-fundamental representations. Such models can be constructed in D-brane models, e.g. intersecting/magnetized D-brane models (see for a review [56, 57] and references therein). Thus, the above models are interesting from the viewpoint of superstring theory.

5 Summary

The MSSM is one of the attractive BSMs to solve the hierarchy problem in the SM and it may be expected to be found near future. One big issue in the MSSM is how to control the SUSY breaking parameters, so that many ideas and works on spontaneous SUSY breaking and mediation mechanisms of the SUSY breaking effects have been discussed so far. In this paper, we proposed an explicit and simple supersymmetric model, where the spontaneous SUSY breaking and GUT breaking are achieved by the same sector. The origin of the

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hyper-charge assignment in the MSSM is also explained by the analogy with the Georgi-Glashow SU(5) GUT. The SM-charged particles are also introduced by the breaking sector, so that we could also predict the soft SUSY breaking terms via the gauge mediation with the gauge and chiral messenger superfields. The crucial role of the gauge-messenger mediation is to induce large A-terms and B-terms at the one-loop level. We investigated the scenario with light superpartners that such a large A-term realizes the maximal mixing and shift the lightest Higgs mass. In fact, we have to introduce additional contribution to the gluino mass, but 125 GeV Higgs mass could be achieved, even if stop is light. $m_{\rm stop}$ should be as light as possible to relax the fine-tuning of μ parameter. On the other hand, the one-loop B-term could be also consistent with the EW symmetry breaking, if $\tan\beta$ is within $2 \lesssim \tan\beta \lesssim 6$. Such small $\tan\beta$ may require large stop mass, as we see in figures 4 and 5. In fact, we see that about 3 TeV $m_{\rm stop}$ can achieve 125 GeV Higgs mass and the EW symmetry breaking in table 4.

Our light SUSY particles are wino, bino, and gravitino, and the mass difference is not so big. The lightest particle is bino, and wino is heavier than bino. The mass difference is $O(0.1) \times m_{3/2}$ GeV. This might be one specific feature of the gauge messenger scenario in SU(5) GUT, as discussed in ref. [34].

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A Mass spectrums of the particles in the symmetry breaking sector

We investigate the mass matrices for the remnant fields in the symmetry breaking sector. First, let us discuss $(Z, \widetilde{Z})$ and $(\rho, \widetilde{\rho})$ components. We define $Z_\pm$ and $\rho_\pm$ as

$$Z_\pm = \frac{\widetilde{Z} \pm Z^\dagger}{\sqrt{2}}, \ \rho_\pm = \frac{\rho \pm \widetilde{\rho}^\dagger}{\sqrt{2}}. \tag{A.1}$$

The fermion masses are given by

$$\mathcal{L}_f = -\left(\overline{\lambda_-}\ \overline{Z_+}\ \overline{\rho_+}\right) M^f_+ \begin{pmatrix} \lambda_- \\ Z_+ \\ \rho_+ \end{pmatrix} - \left(\overline{\lambda_+}\ \overline{Z_-}\ \overline{\rho_-}\right) M^f_- \begin{pmatrix} \lambda_+ \\ Z_- \\ \rho_- \end{pmatrix}, \tag{A.2}$$

where the mass matrices $(M^f_\pm)$ are

$$M^f_+ = \begin{pmatrix} 0 & -g\Delta v & gv_\chi \\ -g\Delta v & 0 & -hv_\chi \\ gv_\chi & -hv_\chi & -h\Delta v \end{pmatrix}, \ M^f_- = \begin{pmatrix} 0 & -g\Delta v & gv_\chi \\ -g\Delta v & 0 & hv_\chi \\ gv_\chi & hv_\chi & h\Delta v \end{pmatrix}, \tag{A.3}$$

and $\lambda_\pm$ are the linear combinations of the gauginos ($X_{(+)}$) which are the suparpartners of X_μ,

$$\lambda_\pm = \frac{X_+ \pm \overline{X}}{\sqrt{2}}. \tag{A.4}$$

The masses for the bosonic superpartners are

$$\mathcal{L}_B = -\begin{pmatrix} Z_+^\dagger & \rho_+^\dagger \end{pmatrix} M_+^2 \begin{pmatrix} Z_+ \\ \rho_+ \end{pmatrix} - \begin{pmatrix} Z_-^\dagger & \rho_-^\dagger \end{pmatrix} M_-^2 \begin{pmatrix} Z_- \\ \rho_- \end{pmatrix}, \tag{A.5}$$

where the mass matrices ($M_\pm^2$) are given by

$$M_+^2 = \begin{pmatrix} h^2 v_\chi^2 & -h^2 v_\chi \Delta v \\ -h^2 v_\chi \Delta v & h^2(v_\chi^2 + \Delta v^2) + F_X \end{pmatrix}, \tag{A.6}$$

$$M_-^2 = \begin{pmatrix} h^2 v_\chi^2 + g^2 \Delta v^2 & -(h^2+g^2)\Delta v v_\chi \\ -(h^2+g^2)\Delta v v_\chi & h^2(v_\chi^2 + \Delta v^2) + g^2 v_\chi^2 - F_X \end{pmatrix}. \tag{A.7}$$

The F-term F_X is $F_X = -h^2 v_\chi^2$, so that M_+^2 includes the Goldstone mode.

The fermion masses for the other particles are also generated by the VEVs:

$$\mathcal{L}_Y = -\frac{1}{2}\begin{pmatrix} \widetilde{W}^A & Y^A \end{pmatrix} M_Y \begin{pmatrix} \chi^A \\ \widetilde{\chi}^A \end{pmatrix} - \frac{1}{2}\begin{pmatrix} \chi^A & \widetilde{\chi}^A \end{pmatrix} M_Y^T \begin{pmatrix} \widetilde{W}^A \\ Y^A \end{pmatrix} + h.c., \tag{A.8}$$

where $\widetilde{W}$ is the superpartner of W' and M_Y are defined as

$$M_Y = \begin{pmatrix} -\frac{1}{\sqrt{2}} M_{W'} & \frac{1}{\sqrt{2}} M_{W'} \\ -h v_\chi & -h v_\chi \end{pmatrix}. \tag{A.9}$$

The eigenvalues are $M_{W'}$, $M_{W'}$, $\sqrt{2} h v_\chi$, $\sqrt{2} h v_\chi$ and the bosonic masses are given by the same mass spectrum. The imaginary part of $\chi - \widetilde{\chi}$ corresponds to the Goldstone boson, and the real part has the mass, $M_{W'}$, according to the D-term. The other masses, $\sqrt{2} h v_\chi$, correspond to the ones of $\chi + \widetilde{\chi}$ and Y.

The singlet components $(Y_0, \chi_0, \widetilde{\chi}_0)$ of $\hat{Y}$ and $(\hat{\chi}, \hat{\widetilde{\chi}})$ also get masses, according to the nonzero v_χ. The fermionic mass matrix is

$$\mathcal{L}_{Y_0} = -\frac{1}{2}\begin{pmatrix} \widetilde{Z'} & Y_0 \end{pmatrix} M_{Y_0} \begin{pmatrix} \chi_0 \\ \widetilde{\chi_0} \end{pmatrix} - \frac{1}{2}\begin{pmatrix} \chi_0 & \widetilde{\chi_0} \end{pmatrix} M_{Y_0}^T \begin{pmatrix} \widetilde{Z'} \\ Y_0 \end{pmatrix} + h.c., \tag{A.10}$$

where M_{Y_0} are defined as

$$M_{Y_0} = \begin{pmatrix} -\frac{1}{\sqrt{2}} M_{Z'} & \frac{1}{\sqrt{2}} M_{Z'} \\ -h v_\chi & -h v_\chi \end{pmatrix}. \tag{A.11}$$

The mass spectrums are given, relplacing $M_{W'}$ with $M_{Z'}$.

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	$N_{\rm eff}=6$	$N_{\rm eff}=6$	$N_{\rm eff}=6.97$	$N_{\rm eff}=7.83$
$m_{3/2}$	588.84 GeV	741.31 GeV	495.79 GeV	245.02 GeV
T_ρ	2.39×10^{11} GeV	5.76×10^{12} GeV	5.66×10^{12} GeV	2.74×10^{10} GeV
T_X	1.00×10^{7} GeV	1.00×10^{7} GeV	1.00×10^{7} GeV	1.00×10^{7} GeV
$\tan\beta$	3.69	3.93	3.43	4.04
m_h	126.20 GeV	125.89 GeV	124.65 GeV	124.03 GeV
$m_{\rm stop}$	3.05 TeV	3.61 TeV	2.93 TeV	1.90 TeV
X_t	$3.43\times\ m_{\rm stop}$	$3.41\times\ m_{\rm stop}$	$2.98\times\ m_{\rm stop}$	$2.39\times\ m_{\rm stop}$
$\lvert\mu\rvert$	3.72 TeV	4.38 TeV	3.27 TeV	1.93 TeV
$\lvert B\rvert$	4.21 TeV	4.72 TeV	3.22 TeV	1.97 TeV
$\lvert B_{\rm EW}/B\rvert$	0.92	1.00	0.91	0.5
$\lvert M_3\rvert$	$5.73\times m_{3/2}$	$5.73\times m_{3/2}$	$6.85\times m_{3/2}$	$7.83\times m_{3/2}$
$\lvert M_2\rvert$	$0.98\times m_{3/2}$	$0.98\times m_{3/2}$	$0.98\times m_{3/2}$	$0.98\times m_{3/2}$
$\lvert M_1\rvert$	$0.75\times m_{3/2}$	$0.75\times m_{3/2}$	$0.69\times m_{3/2}$	$0.64\times m_{3/2}$
$m_{\tilde t_1}$	3.52 TeV	4.12 TeV	3.31 TeV	2.15 TeV
$m_{\tilde t_2}$	2.62 TeV	3.17 TeV	2.57 TeV	1.65 TeV
$m^2_{\tilde Q_L}$	11.72 TeV^2	16.56 TeV^2	10.36 TeV^2	4.24 TeV^2
$m^2_{\tilde d_R}$	15.97 TeV^2	22.52 TeV^2	13.60 TeV^2	5.40 TeV^2
$m^2_{\tilde l_L}$	0.78 TeV^2	0.93 TeV^2	0.44 TeV^2	0.18 TeV^2
$m^2_{\tilde e_R}$	1.42 TeV^2	1.75 TeV^2	0.81 TeV^2	0.31 TeV^2

Table 4. SUSY mass spectrums and parameters with $\Lambda_{\rm GUT}=2\times10^{16}$ GeV. Higgs mass is calculated by FeynHiggs [46, 58–61].

B Concrete parameter set

The parameter sets which predict $m_h\approx 125$ GeV are in table 4. The Higgs mass is calculated by FeynHiggs [46, 58–61]. $m_{\tilde t_{1,2}}$ are the stop masses in the mass eigenstate. $m^2_{\tilde Q_L}$, $m^2_{\tilde d_R}$, $m^2_{\tilde l_L}$ and $m^2_{\tilde e_R}$ are the soft SUSY breaking terms of the squarks $(\widetilde Q_L,\tilde d_R)$ and sleptons $(\tilde l_L,\tilde e_R)$.

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