Gauge mediation to effective supersymmetry through U(1)s with a dynamical supersymmetry breaking, and string compactification

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We investigate the possibility of $U(1)'$ mediation, leading to an effective SUSY where the first two family sfermions are above 100 TeV but the third family sfermions and the Higgs doublets are in the TeV region (or the light stop ($\tilde{t}_l$) case). The $U(1)'$ gaugino, *Zprimino* (Z'-ino), needs not to be at a TeV scale, but needs to be somewhat lighter than the messenger scale. We consider two cases, one the mediation is only through $U(1)'$ and the other through $U(1)'$ and the electroweak hypercharge $U(1)_Y$. In the SUSY field theory framework, we calculate the superpartner mass spectra for these two cases. We also point out that the particle species needed for these mechanisms are already obtained from a $\mathbf{Z}_{12-I}$ orbifold compactification.

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I. INTRODUCTION

Supersymmetry (SUSY) and its breaking mechanism have been the most active particle theory research in the last three decades. In particular, the SUSY flavor problem has led to the gauge mediated SUSY breaking (GMSB) [1,2]. The attractive gravity mediation scenario for transmitting SUSY breaking down to the visible sector probably violates the flavor independence of interactions, but there are ways in the gravity mediation also to suppress the flavor-changing neutral couplings (FCNC) of the standard model (SM) fermions in the effective SUSY (effSUSY) framework [3]. In the effSUSY, the first two family sfermions are sufficiently heavy above 5–20 TeV while the third family sfermion masses are in the 100 GeV–1 TeV region. The SUSY flavor solution by the GMSB relies on the family independence of the sfermion interaction, for which the gauge interactions do not distinguish family members. The family independence of sfermion masses needs the dominant SUSY breaking source with color $SU(3)_c$ charge, the weak $SU(2)_W$ charge and the weak hypercharge Y. There exists the *SUSY breaking source at some hidden sector scale* Λ_h *below* 10^{12} GeV *for the GMSB* is useful [4] and the messengers, carrying the visible sector gauge charges, acquire SUSY breaking F (or D) terms. The visible sector sfermions obtain masses via these messenger F terms and sometimes the grand unification (GUT) messenger multiplets have been considered for this transmitting purpose [2].

Even though the original GMSB seems to be attractive, similar related ideas in terms of U(1)s have been suggested by Langacker, Pas, Wang and Yavin [5], Mohapatra and Nandi[6], and Kikuchi and Kubo [7]. The Langacker *et al.* mechanism employs an extra Z' gauge interaction instead of the whole $SU(3)_c \times SU(2)_W \times U(1)_Y$ interactions of the SUSY breaking source. The messengers and the SM fields carry the Z' charges, and the $\tilde{Z}'$ gaugino mass is triggering the superpartner masses of the SM fields through the messengers. In addition, they assume a TeV scale Z', but their low energy scale is not needed in general just for a mediation mechanism alone. On the other hand, the Mohapatra-Nandi mechanism uses $U(1)_{Y_1}$ and $U(1)_{B\text{-}L}$ and both of these U(1)s participate in the breaking of SUSY and $U(1)_{Y_1}$ to obtain the $U(1)_Y$ of the SM and also the transfer of SUSY breaking to the superpartners of the SM. The SUSY breaking source can be of dynamical origin as suggested by the well-known dynamical SUSY breaking (DSB) models in $SO(10)'$ with $\mathbf{16}'$ or $\mathbf{16}' + 10'$ [8], or in an $SU(5)'$ model with $\mathbf{10}' + \bar{\mathbf{5}}'$ [9]. We understand that the effective Polonyi form for SUSY breaking [10] is parametrizing the DSB models. Therefore, for a full description of GMSB or *mixed mediation*, we should rely on the string origin of SUSY breaking endowing one hidden family of $SU(5)'$ or $SO(10)'$. There already exist models from the superstring orbifold compactification implementing the SUSY breaking source $SU(5)'$ with the visible sector $SU(3)_c \times SU(2)_W \times U(1)_Y$ [11] or with the flipped SU(5) [12]. In particular, the one hidden sector family models of $SO(10)'$ and $SU(5)'$ cannot carry $SU(3)_c$ color and $SU(2)_W$ charges, or the hidden sector does not satisfy the one family condition. Then, the gauge mediation is better through U(1) s, and it is not expected that the SM families carry the same $U(1)'$ charges, which does not satisfy the chief merit of the GMSB family independence of the mediation. The best we can anticipate for low energy SUSY is an effSUSY [3] in which the superpartners of two light family members are much heavier than the TeV scale.

The recent Large Hadron Collider (LHC) reports exclude squarks in the TeV region [13] even though these analyses are based on the R-parity conserving constrained MSSM. Therefore, in the SUSY framework, the effSUSY is the next serious candidate to be analyzed thoroughly [14]. If the R-parity conserved, the axino [15] or gravitino LSP [16] models are not free from the LHC problem. For example, for $F_a = 10^{11}$ GeV and 1 TeV squark mass, the squark decay line to the axino vertex is estimated to be on the order of a few mm, which is swamped by decays to NLSPs.

In this regard, we note that a simpler model building exists in SUSY field theory framework via the Intrilligator, Seiberg and Shih (ISS) mechanism where the vacuum is unstable but have a sufficiently long lifetime [17–19]. As noted from the string compactification, the total number of 4-dimensional (4D) chiral fields are somewhere between 100 and 200 and it is very difficult for the SUSY breaking source to carry all the $SU(3)_c \times SU(2)_W \times U(1)_Y$ charges. The ISS mechanism is not free of this problem, if not impossible, since, for example, $SU(N_h)$, where N_f flavors need $N_h + 1 \le N_f < \frac{3}{2} N_h$ chiral fields for an unstable minimum. The simplest case $SU(5)'$ needs 6 or 7 vectorlike flavors, which has been realized in string compactifications [4], where, however, the $SU(2)_W$ is broken at the hidden sector scale and the messengers do not carry the color charges. So, it is likely that the original idea of GMSB in the ISS form needing a baroque representation may not be realizable from string compactification.

On the other hand, the Langacker *et al.*-type or the Mohapatra-Nandi-type mediation, employing only U(1)s for mediation, can be easily realizable in SUSY breaking models of one family $SU(5)'$ or of ISS.

We note that there result some phenomenologically acceptable string vacua, where *light stop Z' mediation* (LSTZPM or $\tilde{t}_l Z'$M) and *light stop mixed mediation* (LSTMM or simply MM)[1] to an effSUSY, from 10D string to a 4D minimal supersymmetric SM (MSSM) [20–23]:

(i) Model $\tilde{t}_l Z'$M

(1) Many U(1)s may contribute in the mediation. Here, we choose the simplest possibility that only one $U(1)'$ with the superpartner Zprimino (Z'-ino), $\tilde{Z}'$, is effective in the mediation.

(2) The SUSY breaking source at Λ_h does not carry the weak hypercharge Y, or the low-energy SM does not result. The messenger sector at $M_{\rm mess}$ carries the Z' charge Y' but does not carry the weak hypercharge Y.

(3) The superpartners of the third family fermions, (t, b, τ, ν_τ) do not carry the Z' charge Y'. This item realizes the effSUSY.

(4) The Higgs doublets do not carry the Z' charge Y'. The $SU(2)_W \times U(1)_Y$ breaking is naturally achieved by a running of Higgs boson masses.

(ii) Model MM

(1) Many U(1)s may contribute in the mediation. Here, we choose the simplest possibility that only one $U(1)'$ is effective in addition to $U(1)_Y$ of the SM. These gauge bosons are Z' and B, and their superpartners are called Zprimino $\tilde{Z}'$ and Bino.

(2) The SUSY breaking source does not carry the weak hypercharge Y, or the low energy SM does not result. The messenger sector carries both the weak hypercharge Y and the Z' charge Y'.

(3) The superpartners of the third family fermions do not carry the Z' charge Y'. This item realizes the effective SUSY [3].

(4) Higgs doublets do not carry the Z' charge Y'.

(5) The $SU(2)_W \times U(1)_Y$ breaking is done by a fine-tuning between parameters of the Higgs boson mass matrix [5].

These two cases are the effSUSY generalization of the $U(1)'$ mediation [5] and the mixed U(1)s mediation [6]. In fact, both of these cases are explicitly found in $\mathbf{Z}_{12-I}$ orbifold compactification [11].

In Sec. II, we discuss the general features of $\tilde{t}_l Z'$M and MM on the spectra of superpartners of the SM, and in Sec. III we present such realizations from a published string compactification model [11]. Here, the gauge symmetry breaking to the MSSM is achieved by the vacuum expectation values (VEVs) of some scalar fields obtained from the orbifold compactification. In Appendix A, we present the renormalization group (RG) running inputs and the relevant formulae. In Appendix B, we present two tables on charged and neutral singlets used in Sec. III. Section IV is a conclusion.

II. SUSY BREAKING MEDIATION BY $U(1)'$

We argue that $U(1)'$ mediation of SUSY breaking is of general nature in string compactification. A prototype example has been given in Ref. [11], where the SUSY breaking source is provided by the confining hidden sector $SU(5)'$ with one family $\mathbf{10}' + \bar{\mathbf{5}}'$. In Ref. [11], the original GMSB idea has been commented by assuming the Planck scale singlet VEVs, but it is probable that some needed singlets do not have that large VEVs. Then, the unremovable SUSY-breaking mediation is through U(1)s. This explicit string model will be commented after we present phenomenological aspects of $\tilde{t}_l Z'$M and MM schemes.

We consider two U(1) gauge bosons, B_μ corresponding to Y of the SM and Z'_μ corresponding to an additional hypercharge Y'. The messenger matter fields f and $\bar{f}$ have the quantum numbers of $f(Y, Y')$ and $\bar{f}(-Y, -Y')$, respectively. Both of these possibilities are possible with the

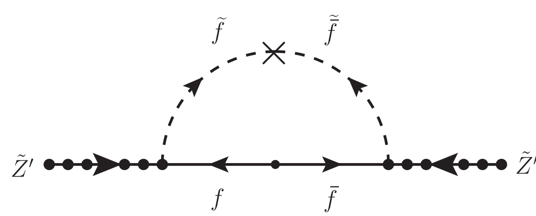

FIG. 1. The mass diagram of Zprimino. The SUSY breaking insertion from DSB is $\times$. The bulleted line is $\tilde{Z}'$. This soft mass is added to the SUSY mass.

[1] We pick up the light stop among the third family members because the RG evolution is dominated by the top-Yukawa coupling.

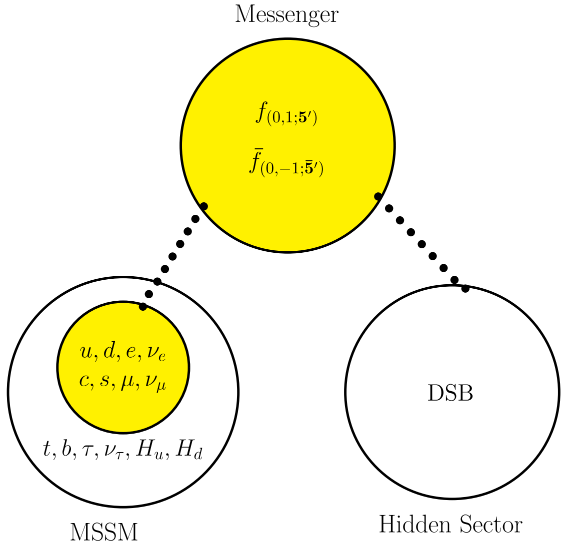

FIG. 2 (color online). An effSUSY through $\tilde{t}_l Z' M$. The Z' line is the bulletted one. The particles in gray disks (yellow online) are neutral under U(1)Y.

model of Ref. [11]. In these models, the Zprimino mass diagram appears as in Fig. 1. We emphasize the Z' mediation by representing Z' as a sequence of bullets and Zprimino as bullets connected with a line.

A. Light stop Z' mediation

The messenger fields, carrying the hidden sector color, such as the SU(5)$'$ charge, have the following $(Y, Y'; \mathrm{SU}(5)')$

$$f(0, 1; \mathbf{5}'), \qquad \bar{f}(0, -1; \bar{\mathbf{5}}'), \tag{1}$$

and the third family members do not carry the Y' charge. In addition, Higgs doublets also do not carry the Y' charge. In this case, the mediation mechanism is shown pictorially in Fig. 2, which is called $\tilde{t}_l Z'$M. In this case, a light Higgs boson and the light third family members are obtained naturally. In Fig. 2, the U(1)$'$ charged sectors are gray (yellow on line).

As a field theory example, we consider an anomaly-free U(1)$'$ charge assignment as $Y' = B - L$ for the first two families, and $Y' = 0$ for the third family members as listed in Table I. Certainly, it may be difficult for this model to produce a successful flavor structure if the U(1)$'$breaking scale is below 10^{12} GeV and the messenger scale is at the GUT scale. So, we assume the messenger scale is low, i.e. only a factor of 100 larger than the DSB scale. Here, our main concern is obtaining the superparticle spectrum.

The Zprimino $\tilde{Z}'$ obtains mass through the diagram shown in Fig. 1, where the SUSY-breaking insertion is shown as $\times$. Below the messenger mass scale $M_{\rm mess}$, the Zprimino soft mass is estimated as[2]

[2]This soft mass is added to the supersymmetric mass.

TABLE I. The Y' charges of the SM fermions, Higgs doublets and heavy neutrinos.

Light families	Y	Y'	Third family and $H_{d,u}$	Y	Y'
$q_{1,2}$	$\frac{1}{6}$	$\frac{1}{3}$	(t, b)	$\frac{1}{6}$	0
$u^c_{1,2}$	$\frac{-2}{3}$	$\frac{-1}{3}$	t^c	$\frac{-2}{3}$	0
$d^c_{1,2}$	$\frac{1}{3}$	$\frac{-1}{3}$	b^c	$\frac{1}{3}$	0
$l_{1,2}$	$\frac{-1}{2}$	-1	(ν_τ, τ)	$\frac{-1}{2}$	0
$e^c_{1,2}$	1	1	τ^c	1	0
$N^c_{1,2}$	0	1	N^c_3	0	0
			H_d	$\frac{-1}{2}$	0
			H_u	$\frac{1}{2}$	0

$$\frac{M_{\tilde{Z}'}(\mu)}{g^2_{Y'}(\mu)} = -\frac{1}{8\pi^2}\frac{F_{\rm mess}}{M_{\rm mess}}, \tag{2}$$

where $F_{\rm mess}$ is the relevant F-term of the messenger sector.

Since the messengers are not charged under the SM gauge group, the MSSM gaugino masses are induced only through RG running from the loops shown in Fig. 3, where the Zprimino mass is shown as $\times$. Assuming that U(1)$'$ is broken at a scale much larger than $M_{\tilde{Z}'}$ but below $M_{\rm mess}$, one can obtain

$$\frac{M_a(\mu)}{g^2_a(\mu)} = -\frac{c_a g^2_{Y'}(M')}{(8\pi^2)^2} M_{\tilde{Z}'}(M') \ln\left(\frac{M_{\rm mess}}{M'}\right), \tag{3}$$

for $\mu < M'$, with M' being the U(1)$'$-breaking scale. For the U(1)$'$ charge assignment given in Table I, c_a are given by

$$c_Y = \frac{92}{27}, \qquad c_2 = \frac{8}{3}, \qquad c_3 = \frac{8}{9}, \tag{4}$$

and thus the MSSM gauginos have a compressed mass spectra compared to the ordinary gauge mediation.

On the other hand, the first two family sfermions directly couple to U(1)$'$ and obtain masses as

$$m^2_{\tilde{q}_{1,2},\tilde{l}_{1,2}} = Y'^2_{q_{1,2},l_{1,2}} M^2_{\tilde{Z}'}, \tag{5}$$

at the messenger scale. The dominant effect on the RG running of their masses comes from the loops involving the Zprimino as shown in Fig. 4.

Because of the desired gauge coupling hierarchy, the MSSM gauginos are lighter than the first two family sfermions. The third family sfermions obtain mass through the diagram shown in Fig. 5, where the SUSY-breaking mass of the MSSM gauginos are shown as $\times$. The soft scalar masses for the third family sfermions and Higgs bosons can be obtained from the RG running equations, which are the same as those in the MSSM at the leading order.

The electroweak symmetry breaking is achieved radiatively [24] by the RG running equations. For a successful electroweak symmetry breaking in SUSY models, we need a TeV scale μ term. Because it is a superpotential term, it is

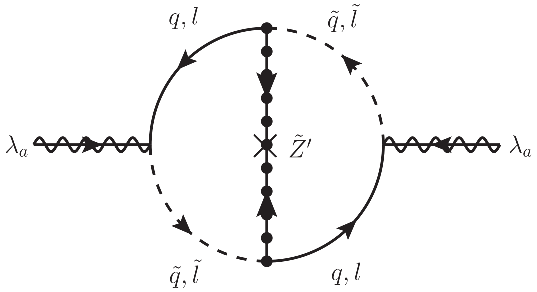

FIG. 3. The mass diagram of the SM gauginos. The SUSY breaking from Zprimino sector is shown as $\times$. The $\tilde{Z}'$ line is a bulleted line.

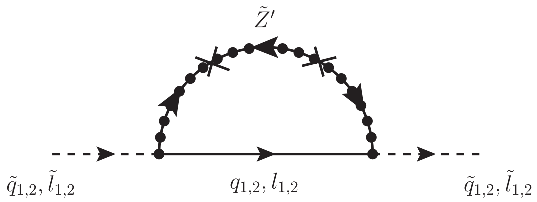

FIG. 4. The first two family sfermion($\tilde{q}_{1,2}$, $\tilde{l}_{1,2}$) mass diagrams. The SUSY breaking from Zprimino sector is shown as $\times$.

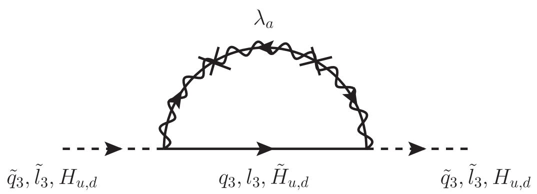

FIG. 5. The mass diagrams for the third family sfermion (q_3, l_3) and Higgs bosons. The SUSY breaking from the SM gauginos are shown as $\times$.

not generated radiatively. In the GMSBs and in our $\tilde{t}_l Z'$M and MM, we need to introduce it independently. The gravitational (or more explicitly in string models, the moduli) interactions introduce a nonrenormalizable term of the form for the Higgsino doublet pair [25],

$$\sim \frac{S_1 S_2}{M_P} H_u H_d, \tag{6}$$

where $M_P \simeq 2.44 \times 10^{18}$ GeV, and $S_{1,2}$ are the SM singlet (s). With VEVs of $S_{1,2}$ in the 10^{10-11} GeV region, we obtain the needed magnitude of μ. Without this additional gravity effect, it may be difficult if not impossible to obtain a successful electroweak symmetry breaking. The Giudice-Masiero mechanism [26] does not introduce the right order of μ from the Kähler potential since the gravitino mass $m_{3/2}$ is required to be much smaller than the electroweak scale.

Below, we introduce the needed μ independently from the $\tilde{t}_l Z'$M and MM.

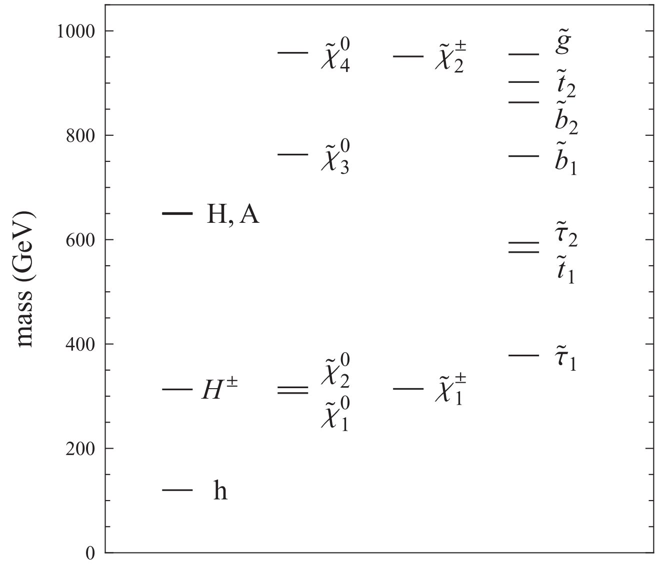

FIG. 6. The sparticle and Higgs boson mass spectra in the $\tilde{t}_l Z'$M. We have taken $M_{\rm mess} = 10^{14}$ GeV, $M_{Z'} = 10^8$ GeV and $M_{\tilde{Z}'}(M_{\rm mess}) = 1.8 \times 10^6$ GeV, for which the squark masses of the first two families are above 10^6 GeV.

We present the spectra based on Table I in Fig. 6.[3] From Fig. 6, we note that the lightest Higgs boson mass is around 120 GeV. We also note that the stop mass is near 580 GeV which is lower than the recent CMS bound of 1.2 TeV [13]. However, the latter is not a serious problem, for we can achieve this CMS bound by enlarging the Zprimino mass,[4] which will subsequently raise the MSSM gaugino masses. Note that the third family sfermions and Higgs bosons acquire soft masses through RG running from the loops involving MSSM gauginos. This explains why they are lighter than the MSSM gauginos.

In the example, $\tilde{\chi}_4^0$ and $\tilde{\chi}_2^\pm$ come mostly from the neutral and charged wino, respectively, while $\tilde{\chi}_3^0$ is bino-like. The lightest ordinary sparticle is $\tilde{\chi}_1^0$, which is higgsino-like and has mass around 306 GeV. Since the gluino mass is comparable to the wino/bino mass, the third generation squarks are not so heavy at high energy scales. As a result, $m_{H_u}^2$ is slowly driven to negative as the energy scale goes down compared to the ordinary gauge mediation, and a small μ-term is required for the electroweak symmetry breaking. In the example, $\mu = 313$ GeV and $B = 133$ GeV at the weak scale, and $\tan\beta = 10$.

Meanwhile, the lightest SUSY particle(LSP) is given by the gravitino having mass $\sim\Lambda_h^3/M_{Pl}^2$.

[3]But in the string example discussed in Sec. III, since the first and second generation $SU(2)_L$ doublet quarks and leptons are not charged under $U(1)'$, the spectra may be distinct from Fig. 6.

[4]To get $m_{\tilde{t}_1} \gtrsim 1.2$ TeV in the example, one can take $M_{\tilde{Z}'} \gtrsim 3.8 \times 10^6$ GeV at the messenger scale. The electroweak symmetry breaking would then require a rather large Higgs μ term: $\mu \gtrsim 670$ GeV.

B. Mixed mediation

The only difference of MM from the $\tilde{t}_l Z'$M is that the messenger fields carry the Y charge,

$$f(Y_f, 1; \mathbf{5}'), \bar{f}(-Y_f, -1; \bar{\mathbf{5}}'). \quad (7)$$

The particles in the MM mechanism are shown in Fig. 7. If the bino is much heavier than other MSSM gauginos, the top-Yukawa interaction would drive not only $m^2_{H_u}$ but also the left-handed stop mass squared to negative at around the weak scale.[5] To avoid such a problem in the MM scenario, we need $Y_f^2 g_Y^2 \lesssim 10^{-3} g^2_{Y'}$ so that the bino mediation induces soft scalar masses at most of the order of the wino/gluino mass. In this case, the effective SUSY can be obtained since the bino mediation is much weaker than the Z' mediation. Nonetheless, the bino mediation can still change the mass spectra of light sparticles and Higgs bosons.

For the case that the bino mediation is as important as the Z' mediation, a tachyonic stop can be avoided if the third family sfermions are charged under $U(1)'$. Only the wino and gluino then remain light while all the scalars acquire quite large soft masses. Hence, we need to fine-tune the Higgs mass parameters to achieve the correct electroweak symmetry breaking.

III. STRING EXAMPLE

In this section, we discuss the $\tilde{t}_l Z'$M and MM based on the $\mathbf{Z}_{12-I}$ orbifold model of Ref. [11], where the 4D gauge group is

$$SU(3)_c \times SU(3)_W \times U(1)^3 \times SU(2)_n \times SU(5)' \times SU(3)' \times U(1)'^2.$$

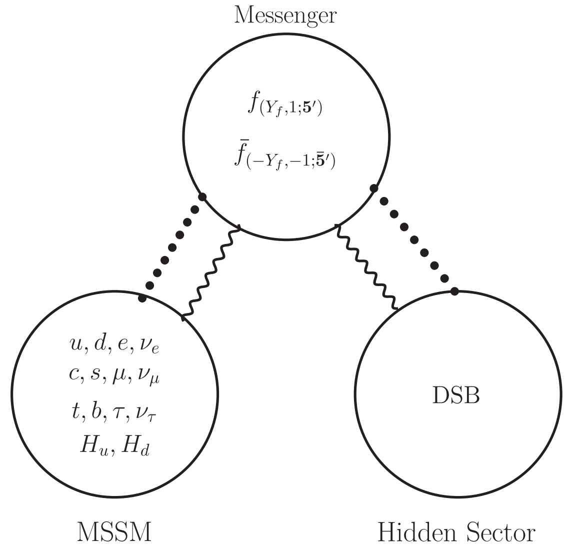

FIG. 7. An effSUSY through MM.

The gauge groups $SU(2)_n$ and $SU(3)'$ are completely broken by the Higgs mechanism. The gauge group $SU(2)_n$ is neutral, i.e. it does not contribute to the SM hypercharge Y. But the hidden sector gauge group $SU(3)'$ contributes to Y: $\mathbf{3}' \to (\frac{-1}{3}\,\frac{-1}{3}\,\frac{2}{3})$ and $\bar{\mathbf{3}}' \to (\frac{1}{3}\,\frac{1}{3}\,\frac{-2}{3})$. To break $SU(3)'$ completely with these hypercharge contributions, we assign VEVs to two independent even Γ fields in the lower box of Table V. Thus at the GUT scale, $SU(3)_c \times SU(3)_W \times U(1)$ is broken to $SU(3)_c \times SU(2)_W \times U(1)_Y$. Then, at the electroweak scale, we have the following gauge group

$$SU(3)_c \times SU(2)_W \times U(1)_Y \times SU(5)' \times U(1)_1 \times U(1)_2 \times U(1)_3 \times U(1)_4 \times U(1)_5, \quad (9)$$

where the five additional U(1) charges are

$$\begin{aligned}
Q_1 &= (6\ \ 6\ \ -6\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)(0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)'\\
Q_2 &= (0\ \ 0\ \ 0\ \ 6\ \ 6\ \ 6\ \ 0\ \ 0)(0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)'\\
Q_3 &= (0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 2\ \ 2)(0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)'\\
Q_4 &= (0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)(4\ \ 4\ \ 4\ \ 4\ \ 4\ \ 0\ \ 0\ \ 0)'\\
Q_5 &= (0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 0)(0\ \ 0\ \ 0\ \ 0\ \ 0\ \ 4\ \ 4\ \ 4)'.
\end{aligned} \quad (10)$$

[5] In the situation under consideration, the left-handed stop and up-type Higgs boson would acquire soft masses as

$$\begin{aligned}
m^2_{H_u}(\mu) &\simeq \frac{1}{4} g_Y^2 P_1 + \frac{1}{2} g_Y^2 P_2 - 3 y_t^2 P_3,\\
m^2_{\tilde{t}_L}(\mu) &\simeq \frac{1}{36} g_Y^2 P_1 + \frac{1}{18} g_Y^2 P_2 - y_t^2 P_3,
\end{aligned} \quad (8)$$

at $\mu < M_{Z'}$. Here $P_{1,2,3}$ are positive numbers of $\mathcal{O}(M^2_{\tilde{B}})$. Also note that the first term corresponds to the bino-mediated contribution at $M_{\rm mess}$, while the latter two are from the RG effects associated with $U(1)_Y$ gauge and top-Yukawa coupling, respectively. It is obvious that $3m^2_{\tilde{t}_L} < m^2_{H_u}$ at a low energy scale.

TABLE II. Hidden sector $SU(5)'$ representations under $SU(2)_n \times SU(5)' \times SU(3)'$. After removing vectorlike representations by $\Gamma =$ even integer singlets, the starred representations remain.

$P + n[V \pm a]$	Γ	(Reports)$Y[Q_1, Q_2, Q_3, Q_4, Q_5]$
$(\frac{1}{6}^2\,\frac{-1}{6}\,\frac{1}{6}^3\,\frac{1}{4}^2)(\underline{\frac{-3}{4}\,\frac{1}{4}^4}\,\frac{-1}{4}^3)'_{T1_-}$	2	$(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[3,3,1;1,-1]}$
$(\frac{1}{6}^2\,\frac{-1}{6}^4\,0^2)(\underline{\frac{1}{2}\,\frac{1}{2}\,\frac{-1}{2}^3}\,\frac{-1}{6}^3)'_{T2_+}$	-1	$\star(\mathbf{1};\mathbf{10}',\mathbf{1})^L_{0[3,-3,0;-2,-2]}$
$(0^6\,\underline{\frac{1}{4}\,\frac{-3}{4}})(\underline{\frac{3}{4}\,\frac{-1}{4}^4}\,\frac{1}{4}^3)'_{T3}$	-1	$(\mathbf{2}_n;\mathbf{5}',\mathbf{1})^L_{0[0,0,-1;-1,3]}$
$(0^6\,\underline{\frac{3}{4}\,\frac{-1}{4}})(\underline{\frac{-3}{4}\,\frac{1}{4}^4}\,\frac{-1}{4}^3)'_{T9}$	1	$(\mathbf{2}_n;\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,0,1;1,-3]}$
$(0^3\,\frac{-1}{3}^3\,\frac{1}{4}\,\frac{1}{4})(\underline{\frac{-3}{4}\,\frac{1}{4}^4}\,\frac{1}{12}^3)'_{T7_0}$	-1	$\star(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,-6,1;1,1]}$
$(\frac{1}{6}^2\,\frac{-1}{6}\,\frac{1}{6}^3\,\frac{-1}{4}^2)(\underline{\frac{3}{4}\,\frac{-1}{4}^4}\,\frac{1}{4}^3)'_{T7_-}$	0	$(\mathbf{1};\mathbf{5}',\mathbf{1})^L_{0[3,3,-1;-1,3]}$
$(0^6\,\frac{-1}{2}\,\frac{-1}{2})(\underline{-10^4}\,0^3)'_{T6}$	-2	$3\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{0[0,0,-2;-4,0]}$
$(0^6\,\frac{-1}{2}\,\frac{-1}{2})(\underline{10^4}\,0^3)'_{T6}$	-2	$2\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})^L_{1[0,0,-2;4,0]}$
$(0^6\,\frac{1}{2}\,\frac{1}{2})(\underline{-10^4}\,0^3)'_{T6}$	2	$2\cdot(\mathbf{1};\bar{\mathbf{5}}',\mathbf{1})^L_{-1[0,0,2;-4,0]}$
$(0^6\,\frac{1}{2}\,\frac{1}{2})(\underline{10^4}\,0^3)'_{T6}$	2	$3\cdot(\mathbf{1};\mathbf{5}',\mathbf{1})^L_{0[0,0,2;4,0]}$

This model leads to one $SU(5)'$ family, breaking SUSY at an intermediate scale Λ_h, and the SM gauge group with three families and one pair of Higgs doublets. It contains the ingredients of the MSSM and the DSB source.

Supersymmetry breaking by one $\mathbf{10}'$ and one $\bar{\mathbf{5}}'$ is achieved by the starred fields of Table II. A possible combination with one $\mathbf{10}'$ and one $\bar{\mathbf{5}}'$ is possible with the hidden sector gauginos [27,28]

$$X' \propto \epsilon_{acfgh}\tilde{G}'^a{}_b\tilde{G}'^c{}_d\mathbf{10}'^{eb}\bar{\mathbf{5}}'_e\mathbf{10}'^{fd}\mathbf{10}'^{gh}, \tag{11}$$

which carries $Y' = Q_4 + \frac{1}{5}Q_5 = 0$, but breaks its orthogonal combination $U(1)_{Y'_\perp}$. Also, the Y value of X' is also zero. Thus, the model realizes the scenarios discussed in Sec. II. The messenger sector charges determine whether it is $\tilde{t}_l Z'$M or MM.

A. Hidden sector $SU(5)'$, gauge mediation, messengers, and R-parity

The hidden sector $SU(5)'$ representations are shown in Table II. Removing vectorlike pairs, we obtain one family $SU(5)'$ model below the GUT scale. The light $\mathbf{10}'$ and $\bar{\mathbf{5}}'$ are marked with a star. These chiral fields carry the vanishing Y charge, and SUSY breaking at the scale Λ_h does not break $U(1)_Y$ of the SM. As pointed out in Ref. [9], one family $\mathbf{10}' + \bar{\mathbf{5}}'$ of a confining $SU(5)'$ breaks SUSY. For the $\tilde{t}_l Z'$M and MM, we require the confining scale Λ_h below 10^{12} GeV [2,4].

Even though the singlet combination $\mathbf{10}'\mathbf{10}'\mathbf{10}'\bar{\mathbf{5}}'$ is not possible with one $\mathbf{10}'$ and one $\bar{\mathbf{5}}'$, SUSY breaking can be parameterized by the hidden sector gauge field strength $\mathcal{W}'^\alpha\mathcal{W}'_\alpha$ [11,29], for the messenger f and $\bar{f}$,

$$\mathcal{L} = \int d^2\theta[\xi(\cdots)f\bar{f}\mathcal{W}'^\alpha\mathcal{W}'_\alpha + \eta(\cdots)f\bar{f}] + \text{H.c.}, \tag{12}$$

where we have in general the holomorphic functions ξ and η of singlet chiral fields, $S_1, S_2, \cdots$. Assuming the singlet VEVs at the string scale, Ref. [11] discussed the GMSB. On the other hand, it is generally expected that some of singlet VEVs are smaller than the string scale. Then, f and $\bar{f}$ carrying $SU(3)_c$ and $SU(2)_W$ charges have negligible couplings to $\mathcal{W}'^\alpha\mathcal{W}'_\alpha$, and the original GMSB [1] is probably not realized in the model of Ref. [11]. The main reason is that $SU(5)'$-colored f and $\bar{f}$ are multiplied to $\mathcal{W}'^\alpha\mathcal{W}'_\alpha$ together with many small VEV singlet fields. We argue that this may be of general nature. This leads us to consider $\tilde{t}_l Z'$M and MM discussed in Sec. II. For simplicity, we choose only one relatively light (at the 10^{13-15} GeV) pair of f and $\bar{f}$ from Table II, and assume

TABLE III. Three families of quarks and leptons and a pair of Higgs doublets of [11]. The quark singlets d^c and b^c are interchanged from those of Ref. [11] to have an effSUSY. The lepton singlets are taken from Table IV in Appendix B.

Sector	(Reports)$Y[Q_1, Q_2, Q_3, Q_4, Q_5]$	Γ	Label
T_{4_-}	$3\cdot(\mathbf{3},\mathbf{2})^L_{1/6[0,0,0;0,0]}$	1	q_1, q_2, q_3
T_{4_-}	$2\cdot(\bar{\mathbf{3}},\mathbf{1})^L_{-2/3[-3,3,2;0,0]}$	3	u^c, c^c
T_{7_+}	$(\bar{\mathbf{3}},\mathbf{1})^L_{-2/3[0,6,-1;5,1]}$	1	t^c
T_{2_0}	$(\bar{\mathbf{3}},\mathbf{1})^L_{1/3[3,-3,0;0,-4]}$	-1	b^c
T_{4_-}	$2\cdot(\bar{\mathbf{3}},\mathbf{1})^L_{1/3[-3,3,-2;0,0]}$	1	d^c, s^c
T_{4_-}	$3\cdot(\mathbf{1},\mathbf{2})^L_{-1/2[-6,6,0;0,0]}$	1	l_1, l_2, l_3
U_1	$\mathbf{1}^L_{1[9,3,-2;0,0]}$	1	e^c
T_3	$\mathbf{1}_{1[0,-6,1;5,-3]}$	1	μ^c
T_{1_0}	$\mathbf{1}^L_{1[0,6,-1;5,1]}$	3	τ^c
T_{1_0}	$(\mathbf{1},\mathbf{2})^L_{1/2[0,6,-1;5,1]}$	0	H_u
T_{7_+}	$(\mathbf{1},\mathbf{2})^L_{-1/2[-6,0,-1;5,1]}$	-2	H_d

TABLE IV. The charged singlets. The fields in the lower box get Y contribution from the $SU(3)'$ generators, and the underline means permutations.

$P + n[V \pm a]$	Γ	No. $\times$ (Reports)$Y[Q_1, Q_2, Q_3, Q_4, Q_5]$
$(\frac{1}{2}\frac{1}{2}\frac{-1}{2}\frac{1}{2}\frac{1}{2}\frac{-1}{2}\frac{-1}{2}\frac{-1}{2})(0^8)'_{U_1}$	1	$\mathbf{1}^L_{1[9,3,-2;0,0]} \to e^c$
$(00000-1\underline{-10})(0^8)'_{U_3}$	-2	$2\cdot\mathbf{1}^L_{1[0,-6,-2;0,0]}$
$(000\frac{2}{3}\frac{2}{3}\frac{-1}{3}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{12}\frac{1}{12}\frac{1}{12})'_{T_{1_0}}$	3	$\mathbf{1}^L_{1[0,6,-1;5,1]} \to \tau^c$
$(000\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{12}\frac{1}{12}\frac{1}{12})'_{T_{1_0}}$	-3	$2\cdot\mathbf{1}^L_{1[0,-6,-1;5,1]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{4}\frac{1}{4})(\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4})'_{T_{1_-}}$	1	$2\cdot\mathbf{1}^L_{-1[3,3,1;-5,3]}$
$(00000-1\frac{1}{4}\frac{1}{4})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4})'_{T_3}$	1	$(2_L+1_R)\mathbf{1}_{1[0,-6,1;5,-3]} \to \mu^c$
$(000000\underline{\frac{1}{4}\frac{-3}{4}})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4})'_{T_3}$	-1	$(6_L+6_R)\mathbf{1}_{1[0,0,-1;5,-3]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{-5}{6}\frac{-1}{2}\frac{-1}{2})(0^8)'_{T_{4_-}}$	0	$2\cdot\mathbf{1}^L_{1[3,-3,-2;0,0]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\underline{\frac{-1}{2}\frac{-1}{2}})(0^8)'_{T_{4_-}}$	0	$12_L\cdot\mathbf{1}_{1[3,3,0;0,0]}$
$(\frac{-1}{3}\frac{-1}{3}\frac{1}{3}\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}00)(0^8)'_{T_{4_-}}$	-2	$7_R\cdot\mathbf{1}_{1[-6,-6,0;0,0]}$
$(\frac{-1}{3}\frac{-1}{3}\frac{1}{3}\frac{2}{3}\frac{2}{3}\frac{-1}{3}00)(0^8)'_{T_{4_-}}$	4	$3_R\cdot\mathbf{1}_{1[-6,6,0;0,0]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{-5}{6}\frac{-1}{2}\frac{-1}{2})(0^8)'_{T_{4_-}}$	-4	$2_R\cdot\mathbf{1}_{1[3,-3,-2;0,0]}$
$(\frac{1}{3}\frac{1}{3}\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}\frac{2}{3}\frac{1}{4}\frac{1}{4})(\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{12}\frac{-1}{12}\frac{-1}{12})'_{T_{7_+}}$	-1	$\mathbf{1}^L_{-1[6,0,1;-5,-1]}$
$(\frac{-1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\underline{\frac{3}{4}\frac{-1}{4}})(\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{12}\frac{-1}{12}\frac{-1}{12})'_{T_{7_+}}$	2	$2\cdot\mathbf{1}^L_{-1[-3,3,1;-5,-1]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{-1}{4}\frac{-1}{4})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4})'_{T_{7_-}}$	0	$2\cdot\mathbf{1}^L_{1[3,3,-1;5,-3]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{4}\frac{1}{4})(\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{1}{4}\frac{3}{4}\frac{-1}{4}\frac{-1}{4})'_{T_{1_-}}$	2	$\mathbf{1}^L_{1[3,3,1;5,1]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{-1}{6}\frac{-1}{6}\frac{-1}{6}00)(0^5\frac{1}{3}\frac{-2}{3}\frac{-2}{3})'_{T_{2_+}}$	-1	$\mathbf{1}^L_{1[3,-3,0;0,-4]}$
$(000\frac{-1}{3}\frac{-1}{3}\frac{2}{3}00)(00000\frac{-2}{3}\frac{1}{3}\frac{1}{3})'_{T_{4_0}}$	-2	$3\cdot\mathbf{1}^L_{-1[0,0,0;0,0]}$
$(\frac{1}{3}\frac{1}{3}\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}00)(0^5\frac{2}{3}\frac{-1}{3}\frac{-1}{3})'_{T_{4_+}}$	-2	$3\cdot\mathbf{1}^L_{1[6,-6,0;0,0]}$
$(\frac{-1}{3}\frac{-1}{3}\frac{1}{3}\frac{1}{3}\frac{1}{3}\frac{1}{3}\frac{-1}{2}\frac{-1}{2})(0^5\frac{2}{3}\frac{-1}{3}\frac{-1}{3})'_{T_{4_+}}$	0	$2\cdot\mathbf{1}^L_{1[-6,6,-1;0,0]}$
$(\frac{-1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{2}\frac{1}{2})(0^5\frac{-1}{3}\underline{\frac{2}{3}\frac{-1}{3}})'_{T_{4_+}}$	3	$4\cdot\mathbf{1}^L_{-1[-3,3,2;0,0]}$
$(000\frac{-1}{3}\frac{-1}{3}\frac{-1}{3}\frac{1}{4}\frac{1}{4})(\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-5}{12}\underline{\frac{-5}{12}\frac{7}{12}})'_{T_{7_0}}$	-1	$2\cdot\mathbf{1}^L_{-1[0,-6,1;-5,-1]}$
$(\frac{1}{6}\frac{1}{6}\frac{-1}{6}\frac{1}{6}\frac{1}{6}\frac{1}{6}\frac{-1}{4}\frac{-1}{4})(\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-1}{4}\frac{-3}{4}\frac{1}{4}\frac{1}{4})'_{T_{7_-}}$	0	$\mathbf{1}^L_{-1[3,3,-1;-5,-1]}$

all the other $\mathbf{5}'$ and $\bar{\mathbf{5}}'$ are sufficiently heavy such that the consideration of one pair of f and $\bar{f}$ is sufficient.

For $\tilde{t}_l Z'$M, we can choose, for example, $f = \mathbf{5}'_0$ from the T_3 sector and $\bar{f} = \bar{\mathbf{5}}'_0$ from the T_9 sector. For MM, we must choose $f = \mathbf{5}'_1$ from T_6 and $\bar{f} = \bar{\mathbf{5}}'_{-1}$ from T_6. Therefore, the model presented in [11] has the basic ingredients for $\tilde{t}_l Z'$M and MM.

As discussed in [11], there appear three SM family members of Table III. Also, only one pair of Higgs doublets results because the superpotential for three $SU(3)_W$ antitriplets must be symmetric under exchange of two superfields. But, $SU(3)_W$ invariance needs an antisymmetric $SU(3)_W$ indices, needing an antisymmetric flavor indices. This leads to one pair of massless Higgsinos naturally. So by SUSY, we have a pair of massless Higgs doublets at the $SU(3)_W$-breaking scale (the GUT scale). The TeV scale μ is generated by norenormalizable superpotential terms [25], which will be worked out explicitly in the present string model [30]. This fulfils all the requirements of $\tilde{t}_l Z'$M and MM.

The Y' quantum number is

$$Y' = Q_3 + \frac{1}{5} Q_4. \qquad (13)$$

From Table III, we find that $\tilde{t}_L$, $\tilde{t}_R$, $\tilde{b}_L$, $\tilde{b}_R$, $\tilde{\tau}_L$, $\tilde{\tau}_R$, $\tilde{\nu}_\tau$, H_u, and H_d carry the vanishing Y'. Also, X' of Eq. (11) carries the vanishing Y'. These provide the needed quantum numbers of Fig. 2 and Table I. The string model [11] is a kind of the flavor unification model, and different families need not have the same Y' quantum numbers. In GUTs descending from E_6 which is not a flavor unification model, the family distinction of Y' is not present and hence there is the problem of low-energy baryonic or leptophobic $U(1)_{Y'}$ [31]. The GUTs from F-theory construction [32] is not free from the low-energy $U(1)_{Y'}$ problems.

The proton longevity is the key requirement in the SUSY extension of the SM, usually achieved in terms of the R-parity. This is a parity where the SM matter superfields carry the odd parity while the Higgs superfields carry the even parity. In the orbifold compactification, one combi-

TABLE V. The same as Table IV, but for the neutral singlets. The starred even Γ quantum number fields can develop VEVs without breaking the *R*-parity.

$P+n[V\pm a]$	Γ	No. × (Reports)$Y[Q_1, Q_2, Q_3, Q_4, Q_5]$
$(00000\,-1\underline{1}0)(0^8)'_{U_1}$	$2\star$	$2\cdot\mathbf{1}^L_{0[0,-6,2;0,0]}$
$(00011000)(0^8)'_{U_2}$	$6\star$	$\mathbf{1}^L_{0[0,12,0;0,0]}\equiv S'_1$
$(000\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{4}\tfrac{-1}{4})(\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-5}{12}\tfrac{-5}{12}\tfrac{-5}{12})'_{T_{1_0}}$	-3	$\mathbf{1}^L_{0[0,-6,-1;-5,-5]}$
$(\tfrac{-1}{2}\tfrac{-1}{2}\tfrac{1}{2}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{4}\tfrac{1}{4})(\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{12}\tfrac{1}{12}\tfrac{1}{12})'_{T_{1_0}}$	$2\star$	$\mathbf{1}^L_{0[-9,3,1;5,1]}$
$(\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{4}\tfrac{1}{4})(\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\underline{\tfrac{3}{4}\tfrac{-1}{4}}\tfrac{-1}{4})'_{T_{1_-}}$	$2\star$	$2\cdot\mathbf{1}^L_{0[3,3,1;5,1]}$
$(000\tfrac{1}{3}\tfrac{1}{3}\tfrac{-2}{3}\tfrac{1}{2}\tfrac{1}{2})(0^5\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3})'_{T_{2_0}}$	$2\star$	$\mathbf{1}^L_{0[0,0,2;0,-4]}$
$(\tfrac{-1}{2}\tfrac{-1}{2}\tfrac{1}{2}\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{-1}{6}00)(0^5\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3})'_{T_{2_0}}$	-1	$2\cdot\mathbf{1}^L_{0[-9,-3,0;0,-4]}$
$(000\tfrac{1}{3}\tfrac{1}{3}\tfrac{1}{3}\underline{\tfrac{1}{2}\tfrac{-1}{2}})(0^5\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3})'_{T_{2_0}}$	$2\star$	$2\cdot\mathbf{1}^L_{0[0,6,0;0,-4]}\equiv S'_2$
$(\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{1}{3}\tfrac{1}{3}\tfrac{1}{3}\tfrac{1}{3}\tfrac{-1}{2}\tfrac{-1}{2})(0^5\tfrac{1}{3}\tfrac{1}{3}\tfrac{1}{3})'_{T_{2_+}}$	$0\star$	$\mathbf{1}^L_{0[-6,6,-2;0,4]}$
$(\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{-1}{6}00)(0^5\tfrac{1}{3}\tfrac{1}{3}\tfrac{1}{3})'_{T_{2_+}}$	-1	$\mathbf{1}^L_{0[3,-3,0;0,4]}$
$(000001\tfrac{1}{4}\tfrac{1}{4})(\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4})'_{T_3}$	1	$(1_L+2_R)\mathbf{1}^L_{0[0,6,1;5,-3]}$
$(\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{-5}{6}\tfrac{-1}{4}\tfrac{-1}{4})(\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{12}\tfrac{-1}{12}\tfrac{-1}{12})'_{T_{7_+}}$	$0\star$	$\mathbf{1}^L_{0[-3,-3,-1;-5,-1]}$
$(\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{4}\tfrac{1}{4})(\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\tfrac{1}{4}\underline{\tfrac{3}{4}\tfrac{-1}{4}}\tfrac{-1}{4})'_{T_{1_-}}$	$2\star$	$2\cdot\mathbf{1}^L_{0[3,3,1;5,1]}$
$(\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{-1}{6}00)(0^5\underline{\tfrac{1}{3}\tfrac{-2}{3}}\tfrac{-2}{3})'_{T_{2_+}}$	-1	$2\cdot\mathbf{1}^L_{0[3,-3,0;0,-4]}$
$(000\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{2}{3}00)(0^5\underline{\tfrac{1}{3}\tfrac{-2}{3}}\tfrac{1}{3})'_{T_{4_0}}$	$-2\star$	$6\cdot\mathbf{1}^L_{0[0,0,0;0,0]}\equiv S'_3$
$(\tfrac{1}{3}\tfrac{1}{3}\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3}00)(0^5\underline{\tfrac{2}{3}\tfrac{-1}{3}}\tfrac{-1}{3})'_{T_{4_+}}$	$-2\star$	$6\cdot\mathbf{1}^L_{0[6,-6,0;0,0]}$
$(\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{2}\tfrac{1}{2})(0^5\underline{\tfrac{2}{3}\tfrac{-1}{3}}\tfrac{-1}{3})'_{T_{4_+}}$	3	$4\cdot\mathbf{1}^L_{0[-3,3,2;0,0]}$
$(\tfrac{-1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{2}\tfrac{-1}{2})(0^5\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{2}{3})'_{T_{4_+}}$	-1	$2\cdot\mathbf{1}^L_{0[-3,3,-2;0,0]}$
$(000\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{-1}{3}\tfrac{1}{4}\tfrac{1}{4})(\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-5}{12}\tfrac{-5}{12}\tfrac{7}{12})'_{T_{7_0}}$	-1	$\mathbf{1}^L_{0[0,-6,1;-5,-1]}$
$(\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{1}{6}\tfrac{-1}{4}\tfrac{-1}{4})(\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\tfrac{-1}{4}\underline{\tfrac{1}{4}\tfrac{-3}{4}}\tfrac{1}{4})'_{T_{7_-}}$	$0\star$	$2\cdot\mathbf{1}^L_{0[3,3,-1;-5,-1]}$

nation, say $U(1)_\Gamma$, of U(1)s is the covering gauge symmetry of the *R*-parity. If some even Γ scalars, with the smallest $|\Gamma|$ normalized as 1, develop VEVs, then we obtain the *R*-parity [21,33]:

$$U(1)_\Gamma \to P. \tag{14}$$

If some odd Γ scalars develop VEVs also, then the *R*-parity is spontaneously broken. To have an *R*-parity, we define the following Γ,

$$\Gamma = \big(0\quad 0\quad 0\quad 3\quad 3\quad 0\quad 2\quad 2\big)(0^8). \tag{15}$$

We require that only the even Γ fields are allowed to develop VEVs. If we used some global symmetries in string models, they must be approximate [34,35]. The $\mathbf{Z}_4$ and other discrete *R* symmetries from an approximate global $U(1)_R$ have been tabulated recently, where discrete anomaly-free conditions have been imposed in addition [36].

B. VEVs leading to one Z'

There are five extra U(1)s, Eq. (10), beyond $U(1)_Y$ of the SM. To have one light Z', we need four independent singlet VEVs. Three of these are provided by the starred singlet fields of Table V:

$$S'_1\colon\ U_2(6\star)[0,12,0,0,0]\qquad S'_2\colon\ T_{2_0}(2\star)[0,6,0,0,-4]$$
$$S'_3\colon\ T_{4_+}(-2\star)[6,-6,0,0,-4], \tag{16}$$

where the sector, Γ, and five $U(1)'$ quantum numbers are shown. The fourth singlet combination is the quantum number of X' of Eq. (11)

$$X'\colon\ [0,0,1;-5,0] \tag{17}$$

(as shown in [27,28]), which carries $Q_3=1$ and $Q_4=-5$ and hence carries $Y'=0$. The other combination orthogonal to Y', say $Y'_\perp$, is broken by this dynamical composite, and we obtain one light Z' model.

The hierarchy of VEVs is that $\langle S'_1\rangle$, $\langle S'_2\rangle$, and $\langle S'_3\rangle$ are much greater than $M_{\rm GUT}$ and the $U(1)'_\perp$ breaking scale is at the $\tilde{t}_l Z'$M scale, i.e. the hidden sector scale $\Lambda_h \lesssim 10^{12}$ GeV. Therefore, below $M_{\rm GUT}$ we may consider two light Z's: Z' and $Z'_\perp$.

Note that both the $Z'_\perp$ and $\tilde{Z}'_\perp$ mass scales are Λ_h because the SUSY breaking *F*-term carries a nonvanishing $Y'_\perp$ charge also. Compared to the other U(1)-priminos, there are two relatively light inos, $\tilde{Z}'$ and $\tilde{Z}'_\perp$. Among these, the mass splitting of the $Z'_\perp$ multiplet is greater than that of the Z' multiplet. The mass splitting of the $\tilde{Z}'_\perp$ multiplet is

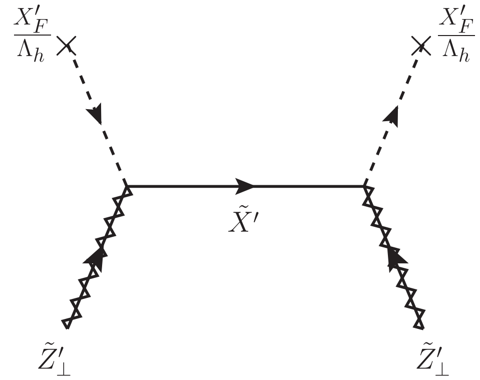

FIG. 8. The $\tilde{Z}'_\perp$ mass splitting via the SUSY breaking through X'. The $\tilde{Z}'_\perp$ line is sawed. The dimensional parameter is the confining scale Λ_h. Even though the $\tilde{Z}'_\perp$ mass is smaller than that of $\tilde{Z}'$, the mass splitting of the $\tilde{Z}'_\perp$ multiplet is larger than that of $\tilde{Z}'$.

of order Λ_h as shown in Fig. 8 while the mass splitting of the $\tilde{Z}'$ multiplet is of order $\Lambda_h^2/M_{\rm mess}$ as shown in Fig. 1. Thus, we can consider only a light $\tilde{Z}'$ mass splitting for the light superpartners as discussed in Sec. II.

The breaking scale of $U(1)_{Y'}$ or the Z' mass is required to be somewhat below the GUT scale, and it is not required for it to be less than Λ_h. The only requirement is that the SUSY breaking scale via Z' mediation, i.e. the supertrace of the Z' SUSY sector or the Zprimino mass, is of order a TeV scale.

IV. CONCLUSION

Motivated by the existing string-compactification model, we investigated the possibility of the $U(1)'$ contribution to the mediation mechanism, leading to an effective SUSY. The first two family sfermions are required to be above 100 TeV but the third family fermions and the Higgs doublets are in the TeV region. For a few parameter ranges, we calculated the spectra of superpartners in the $\tilde{t}_l Z'$M and MM. In the $\tilde{t}_l Z'$M scenario, the Higgs fields survive down to the electroweak scale by tuning the ratio of the DSB scale Λ_h and the messenger scale $M_{\rm mess}$. In the mixed-mediation scenario, it is shown that an additional fine-tuning between parameters of the Higgs boson mass matrix is required as in Ref. [5]. We noted that the Zprimino needs not be at a TeV scale. It is required that it is somewhat lighter than the messenger scale. We also discussed the needed conditions among the fields obtained in the string construction of Ref. [11] for the $\tilde{t}_l Z'$M or the MM.

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APPENDIX A: SOFT TERMS IN $U(1)'$ MEDIATION

In this Appendix, we present the RG equations for soft terms in the Zprimino mediation. We consider the case that the $U(1)'$-breaking scale is much higher than $M_{\tilde{Z}'}$, for which the $U(1)'$ vector superfield acquires a large supersymmetric mass.

At the messenger scale $M_{\rm mess}$, the Zprimino acquires soft mass at the one-loop level while the MSSM gaugino masses vanish:

$$M_{\tilde{Z}'} = -\frac{g_{Y'}^2}{8\pi^2}\frac{F_{\rm mess}}{M_{\rm mess}}, \qquad M_a = 0. \tag{A1}$$

The RG equation for gaugino masses is written as

$$\mu\frac{d}{d\mu}\left(\frac{M_{\tilde{Z}'}}{g_{Y'}^2}\right) = 0, \quad \mu\frac{d}{d\mu}\left(\frac{M_a}{g_a^2}\right) = \frac{c_a}{8\pi^2 b_{Y'}}\mu\frac{dM_{\tilde{Z}'}}{d\mu}, \tag{A2}$$

for $M' < \mu < M_{\rm mess}$. Here, $b_{Y'} = \sum_i Y_i'^2$ is the beta function coefficient for $U(1)'$, and c_a are given by

$$c_Y = \sum\left[6\left(\frac{1}{6}\right)^2 Y_Q'^2 + 3\left(\frac{1}{3}\right)^2 Y_{U^c}'^2 + 3\left(\frac{1}{3}\right)^2 Y_{D^c}'^2 + 2\left(\frac{1}{2}\right)^2 Y_L'^2 + Y_{E^c}'^2\right],$$

$$c_2 = \sum[3Y_Q'^2 + Y_L'^2], \quad c_3 = \sum[2Y_Q'^2 + Y_{U^c}'^2 + Y_{D^c}'^2], \tag{A3}$$

where the sum is over $U(1)'$-charged families. The $U(1)'$ vector multiplet decouples in a supersymmetric way at energy scales below M', which is assumed to be much higher than $M_{\tilde{Z}'}$. Hence, at $\mu < M'$, the MSSM gaugino masses are determined by

$$\mu\frac{d}{d\mu}\left(\frac{M_a}{g_a^2}\right) = 0. \tag{A4}$$

The RG equations for the soft terms associated with the third family sfermions and Higgs bosons are the same as the MSSM at energy scales below $M_{\rm mess}$. For the first two family sfermions, one finds

$$\mu\frac{dm_i^2}{d\mu} = -\frac{Y_i'^2}{2\pi^2} g_{Y'}^2 M_{\tilde{Z}'}^2, \tag{A5}$$

because they couple to the $U(1)'$ vector multiplet, and have negligible Yukawa couplings.

APPENDIX B: SINGLETS

In this Appendix, we list all charged singlets in Table IV and all neutral singlets in Table V, which were needed in Sec. III but not listed in Ref. [11]. The shift vector and the Wilson line are

$$V = \left(\frac{1}{2}\,\frac{1}{2}\,\frac{1}{2}\,\frac{1}{6}\,\frac{1}{6}\,\frac{1}{6}\,\frac{1}{4}\,\frac{1}{4}\right)\left(\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{4}\,\frac{1}{12}\,\frac{1}{12}\,\frac{1}{12}\right)' \tag{B1}$$

$$a_3 = \left(\frac{1}{3}\,\frac{1}{3}\,\frac{2}{3}00000\right)\left(00000\frac{1}{3}\,\frac{1}{3}\,\frac{-2}{3}\right)'. \tag{B2}$$

Tables IV and V are located in Sec. III.

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