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Sparticle masses in deflected mirage mediation

Kiwoon Choi,^a Kwang Sik Jeong,^b Shuntaro Nakamura,^c Ken-Ichi Okumura^c and Masahiro Yamaguchi^c

^a*Department of Physics, KAIST,
335 Gwahangno, Yuseong-gu, Daejeon 305-701, Korea*

^b*School of Physics, Korea Institute for Advanced Study,
87 Hoegiro, Dongdaemun-gu, Seoul 130-722, Korea*

^c*Department of Physics, Tohoku University,
6-3 Aza-Aoba, Aramaki, Aoba-ku, Sendai 980-8578, Japan*

E-mail: kchoi@muon.kaist.ac.kr, ksjeong@kias.re.kr,
shuntaro@tuhep.phys.tohoku.ac.jp, okumura@tuhep.phys.tohoku.ac.jp,
yama@tuhep.phys.tohoku.ac.jp

ABSTRACT: We discuss the sparticle mass patterns that can be realized in deflected mirage mediation scenario of supersymmetry breaking, in which the moduli, anomaly, and gauge mediations all contribute to the MSSM soft parameters. Analytic expression of low energy soft parameters and also the sfermion mass sum rules are derived, which can be used to interpret the experimentally measured sparticle masses within the framework of the most general mixed moduli-gauge-anomaly mediation. Phenomenological aspects of some specific examples are also discussed.

KEYWORDS: Supersymmetry Breaking, Supergravity Models

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1 Introduction

Weak scale supersymmetry (SUSY) is one of the prime candidates for physics beyond the standard model at the TeV scale [1]. Low energy phenomenology of weak scale SUSY is determined mostly by the soft SUSY breaking terms of the visible gauge and matter superfields. Those soft terms are required to preserve flavor and CP with a good accuracy, which severely constrains the possible mediation mechanism of SUSY breaking. Presently, there are three known mediation schemes to yield flavor and CP conserving soft terms:¹ gauge mediation [2, 3], anomaly mediation [4], and string dilaton or volume-moduli mediation [5].² In gauge and anomaly mediations, the radiative corrections due to the standard model (SM) gauge interaction play dominant role for the mediation, and thereby the resulting soft terms automatically preserve flavor and CP. For dilaton/moduli mediation, soft terms induced by the dilaton/moduli F -components preserve flavor and CP by different reasons. The couplings between the messenger dilaton/moduli and the MSSM matter fields preserve flavor

¹Even in these schemes, there can be dangerous CP violation from the Higgs μ and B parameters, which should be considered separately.

²Dilaton/moduli mediation can be considered as a particular version of more general gravity mediation [6], which is one of the most plausible source of soft terms in string theory and naturally preserves flavor and CP at least at leading order in the string coupling or α' expansion.

as they are determined by family-universal rational numbers called the modular weight [5], and preserve CP as a consequence of the associated axionic shift symmetries [7].

So far, most studies of SUSY phenomenology have been focused on the cases that SUSY breaking is dominated by one of the dilaton/moduli, gauge and anomaly mediations. However, recent progress in moduli stabilization suggests that it is a rather plausible possibility that moduli mediation and anomaly mediation are comparable to each other [8–10], which can be naturally realized in KKLT-type moduli stabilization scenario [11]. The resulting soft terms show a distinct feature that sparticle masses are unified at a mirage messenger scale hierarchically lower than the scale of gauge coupling unification [10]. Also a mixed scheme of anomaly and gauge mediations has been proposed before as a solution to the tachyonic slepton problem of anomaly mediation [12]. Recently, it has been pointed out that these schemes can be generalized to include the three known flavor and CP conserving mediations altogether [13, 14]. Such a most general mixed mediation has been dubbed ‘deflected mirage mediation’ as sfermion masses are deflected from the mirage unification trajectory due to the presence of gauge mediation.

In this paper, we wish to examine in more detail the sparticle mass pattern in deflected mirage mediation scenario, together with some phenomenological aspects of the scheme. The organization of this paper is as follows. In section II, we discuss a class of string-motivated effective supergravity models that realize deflected mirage mediation scenario. In section III, we analyze the renormalization group running of soft parameters to derive the (approximate) analytic expression of low energy sparticle masses in deflected mirage mediation, which can be used to interpret the experimentally measured sparticle masses within the framework of the most general flavor and CP conserving mediation scheme. We discuss in section IV the phenomenological feature of two specific examples, one with an accidental little hierarchy between $m_{H_u}^2$ and other soft mass-squares and another with gluino NLSP, that can be obtained within deflected mirage mediation scenario. Section V is the conclusion.

2 Effective supergravity for deflected mirage mediation

In this section, we discuss a class of 4-dimensional (4D) N=1 supergravity (SUGRA) models that realize the deflected mirage mediation scenario. The models discussed here may arise as a low energy effective theory of KKLT-type flux compactification or its variants in string theory. The model contains first of all the MSSM gauge and matter superfields, V^a and Q_i , and also vector-like MSSM-charged exotic matter superfields, $\Phi + \Phi^c$, which live on the visible sector brane. There are light moduli T_I , e.g. the Kähler moduli, stabilized by non-perturbative effects encoded in the superpotential, and also heavy moduli U_p stabilized by flux, e.g. the complex structure moduli. Typically $\text{Im}(T_I)$ corresponds to an axion-like field, and thus the couplings of T_I are invariant under the axionic shift symmetry:

$$\text{U}(1)_{T_I} : T_I \rightarrow T_I + \text{imaginary constant}, \quad (2.1)$$

upon ignoring exponentially small non-perturbative effects.

In KKLT-type compactification, moduli stabilization dynamics itself does not break SUSY since the flux and non-perturbative effects stabilize moduli at a supersymmetric AdS vacuum. Thus, to break SUSY and lift the vacuum to dS state, one needs to introduce a SUSY breaking brane separately. This SUSY breaking brane might be an anti-brane that exists in the underlying string theory, or a brane carrying a 4D dynamics that breaks SUSY spontaneously. An important feature of KKLT-type compactification is that it involves a highly warped throat produced by flux. In the presence of such a warped throat, SUSY breaking brane is stabilized at the tip of throat where the potential of the position modulus is minimized. On the other hand, to implement the high scale gauge coupling unification in the MSSM, the visible sector brane should be stabilized within the internal space at the UV end of throat. This results in a warped separation between the visible brane and the SUSY breaking brane, making the visible sector and the SUSY breaking sector to be sequestered from each other [15]. To be specific, here we will consider a SUSY breaking sector described by a Polony-like superfield Z having a linear superpotential. However, it should be stressed that the visible sector soft terms which are of our major concern are independent of how SUSY is broken at the tip of throat, and therefore our subsequent discussion is valid in cases that SUSY is broken by other means, e.g. by an anti-brane [16].

2.1 Effective supergravity action

With the above features of KKLT-type compactification, the 4D effective action can be written as

$$\mathcal{L}_{4D} = \int d^4\theta CC^* \left[\Omega_{\text{mod}} + \Omega_{\text{matter}} + \Omega_{\text{polony}} \right] + \left[\int d^2\theta \left(\frac{1}{4} f_a W^{a\alpha} W_\alpha^a + C^3 \{ W_{\text{mod}} + W_{\text{matter}} + W_{\text{polony}} \} \right) + \text{h.c.} \right], \quad (2.2)$$

where

$$\begin{aligned} \Omega_{\text{mod}} &= \Omega_{\text{mod}}(U_p, U_p^*, T_I + T_I^*), \\ W_{\text{mod}} &= W_{\text{flux}}(U_p) + \sum_I A_I(U_p) e^{-8\pi^2 a_I T_I}, \end{aligned}$$

for a flux-induced superpotential $W_{\text{flux}}(U_p)$ and the nonperturbative term $e^{-8\pi^2 a_I T_I}$ with real parameter a_I of order unity,

$$\begin{aligned} \Omega_{\text{matter}} &= \sum_A \mathcal{Y}_A(U_p, U_p^*, T_I + T_I^*) \Phi^{A*} \Phi^A \quad (\Phi^A = Q_i, \Phi, \Phi^c, X), \\ W_{\text{matter}} &= \lambda_\Phi(U_p) X \Phi \Phi^c + \frac{\kappa(U_p) X^n}{M_{Pl}^{n-3}} + \frac{1}{6} \lambda_{ijk}(U_p) Q_i Q_j Q_k, \end{aligned} \quad (2.3)$$

where Q_i are the MSSM matter superfields, $\Phi + \Phi^c$ are exotic vector-like matter superfields, and X is a singlet superfield giving a mass to $\Phi + \Phi^c$, and

$$\begin{aligned} \Omega_{\text{polony}} &= ZZ^* - \frac{(ZZ^*)^2}{4M_*^2}, \\ W_{\text{polony}} &= M_{\text{SUSY}}^2 Z, \end{aligned} \quad (2.4)$$

for a Polony-like field Z which breaks SUSY at the tip of throat. Here C is the chiral compensator superfield, f_a are the gauge kinetic functions of the MSSM gauge fields, and we are using the SUGRA unit with $M_{Pl} = 1$, where $M_{Pl} = \sqrt{G_N/8\pi} \simeq 2 \times 10^{18}$ GeV.

As Z is localized at the tip of throat, and thus is sequestered from the visible sector [15], there are no contact interactions between Z and the visible sector fields in the superspace action, which means \mathcal{Y}_A , f_a , λ_{ijk} , λ_Φ , and κ are all independent of Z . Also the axionic shift symmetry (2.1) requires that \mathcal{Y}_A is a function of the invariant combination $T_I + T_I^*$, the holomorphic couplings λ_{ijk} , λ_Φ and κ are independent of T_I , and $\partial f_a/\partial T_I$ are real constants. To incorporate the anomaly mediated SUSY breaking, one needs to include the logarithmic C -dependence of f_a and \mathcal{Y}_A , which is associated with the renormalization group (RG) running of the gauge and Yukawa couplings. Then, under the constraints from the axionic shift symmetry, f_a and \mathcal{Y}_A can be written as

$$\begin{aligned}
 f_a &= \tilde{f}_a(T_I, U_p) + \frac{b_a}{16\pi^2} \ln C = \sum_I k_I T_I + \epsilon(U_p) + \frac{b_a}{16\pi^2} \ln C, \\
 \ln \mathcal{Y}_A &= \ln \tilde{\mathcal{Y}}_A(T_I + T_I^*, U_p, U_p^*) + \frac{1}{8\pi^2} \int_{M_{\text{GUT}}}^{\mu/\sqrt{CC^*}} \frac{d\mu'}{\mu'} \gamma_A,
 \end{aligned}
 \tag{2.5}$$

where b_a and γ_A are the one-loop beta function coefficient and the anomalous dimension of Φ^A , respectively, and k_I are real parameters of order unity. Here we assume the gauge coupling unification around the scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV, which requires k_I to be universal for the MSSM gauge kinetic functions f_a . In fact, as f_a corresponds to an Wilsonian gauge coupling, the one-loop coefficient of $\ln C$ in f_a depends on the corresponding regularization scheme. On the other, the 1PI gauge coupling does not have such scheme dependence. Here we have chosen a specific scheme that the one loop C -dependence of the 1PI gauge coupling is fully encoded in the Wilsonian coupling, for which b_a is given by the one-loop beta function coefficient.

As for the stabilization of X , one can consider two scenarios. The first scenario is that X is stabilized by the combined effects of the SUSY breaking by F^C and the non-renormalizable operator $\kappa X^n/M_{Pl}^{n-3}$ ($n > 3$). Another possibility is that $\kappa = 0$, and X is stabilized by the radiative correction to its Kähler potential. In fact, both scenarios give a similar size of F^X/X , while the resulting mass of X is quite different. In the first scenario, all components of X get a mass comparable to the gravitino mass which is of $\mathcal{O}(10)$ TeV [12]. On the other hand, in the second scenario dubbed ‘axionic mirage mediation’, the pseudo-scalar component of X can be identified as the nearly massless QCD axion solving the strong CP problem, and its fermionic partner, the axino, gets a two-loop suppressed small mass relative to the gravitino mass [13].

As for the SUSY breaking sector, we have taken a simple example given by

$$\begin{aligned}
 \Omega_{\text{polony}} &= ZZ^* - \frac{(ZZ^*)^2}{4M_*^2}, \\
 W_{\text{polony}} &= M_{\text{SUSY}}^2 Z,
 \end{aligned}
 \tag{2.6}$$

where M_{SUSY} and M_* are the two mass parameters for SUSY-breaking dynamics. Generically, some moduli may have a non-negligible wavefunction value at the tip of throat. Then

those moduli can have a sizable contact interaction with Z , which means that M_{SUSY} and M_* , as well as the coefficient of ZZ^* in Ω_{polony} , become a nontrivial function of moduli. However, such an additional complexity does not affect our subsequent discussion, and thus here we consider the simple case that M_{SUSY} and M_* are moduli-independent constants. At any rate, with the above form of Ω_{polony} and W_{polony} , the vacuum value of $Z = z + \theta\tilde{z} + F^Z\theta^2$ is determined as

$$\langle Z \rangle = -\frac{C_0^{r*2}}{C_0} M_{\text{SUSY}}^2 \theta^2, \quad (2.7)$$

where C_0 is the scalar component of the compensator superfield C . The scalar component z gets a mass

$$m_z \sim M_{\text{SUSY}}^2 / M_*, \quad (2.8)$$

while the fermion component \tilde{z} corresponds to the Goldstino. Due to the warping, both M_{SUSY} and M_* are red-shifted by an exponentially small warp factor at the tip of throat:

$$M_{\text{SUSY}} \sim M_* \sim e^{-A} M_{Pl}, \quad (2.9)$$

where

$$g_{\mu\nu}|_{\text{tip}} = e^{-2A} \eta_{\mu\nu}. \quad (2.10)$$

2.2 Integrating out heavy moduli and Polony-like field

The flux-induced superpotential of U_p can be expanded around its stationary point:

$$\begin{aligned} W_{\text{flux}}(U_p) &= W_{\text{flux}}(\tilde{U}_p) + \frac{1}{2} \frac{\partial^2 W_{\text{flux}}(\tilde{U})}{\partial U_p \partial U_q} (U_p - \tilde{U}_p)(U_q - \tilde{U}_q) + \dots \\ &\equiv w_0 + \frac{1}{2} (M_U)_{pq} (U_p - \tilde{U}_p)(U_q - \tilde{U}_q) + \dots, \end{aligned} \quad (2.11)$$

where \tilde{U}_p denotes the stationary point of W_{flux} :

$$\left. \frac{\partial W_{\text{flux}}}{\partial U_p} \right|_{U_q=\tilde{U}_q} = 0. \quad (2.12)$$

Due to the quantization of flux, for generic flux configuration, both w_0 and M_U would be of order unity in the unit with $M_{Pl} = 1$. However, if SUSY breaking is initiated at the tip of throat with a red-shifted $M_{\text{SUSY}} \sim e^{-A} M_{Pl}$, we are required to consider a special type of flux configuration giving an exponentially small

$$w_0 = W_{\text{flux}}|_{U_p=\tilde{U}_p} \sim e^{-2A}, \quad (2.13)$$

in order to get a nearly vanishing vacuum energy density schematically given by

$$V_{\text{vac}} = |M_{\text{SUSY}}|^4 - |w_0|^2. \quad (2.14)$$

Still the flux-induced moduli mass matrix M_U generically has the eigenvalues of order unity. Therefore, in flux compactification scenario with SUSY breaking initiated at the tip of throat, one has the mass hierarchy:

$$m_{3/2} \sim e^{-A} m_z \sim e^{-2A} M_U, \quad (2.15)$$

where $m_{3/2}$ is the gravitino mass, M_U denotes the supersymmetric mass of the flux-stabilized moduli U_p , and $m_z \sim M_{\text{SUSY}}^2/M_*$ is the non-supersymmetric mass of the scalar component of the Polony-like superfield localized at the tip of throat.

With the mass hierarchy (2.15), we can integrate out U_p and z to construct the effective theory of light fields including the visible sector fields $\Phi^A = (Q_i, \Phi, \Phi^c, X)$, chiral compensator C , light moduli T_I , and the Goldstino \tilde{z} . This can be done by solving the following superfield equations of motion within the expansion in powers of the warp factor e^{-A} :

$$\begin{aligned} \frac{1}{4} \bar{\mathcal{D}}^2 \left(C C^* \frac{\partial \Omega_{\text{mod}}}{\partial U_p} \right) + C^3 \frac{\partial W_{\text{mod}}}{\partial U_p} &= 0, \\ \frac{1}{4} \bar{\mathcal{D}}^2 \left(C C^* \frac{\partial \Omega_{\text{polony}}}{\partial Z} \right) + C^3 \frac{\partial W_{\text{polony}}}{\partial Z} &= 0, \end{aligned} \quad (2.16)$$

where $\bar{\mathcal{D}}^2 = \bar{\mathcal{D}}^{\dot{\alpha}} \bar{\mathcal{D}}_{\dot{\alpha}}$ is the superspace covariant derivative. It is straightforward to find that the solutions are given by

$$\begin{aligned} U_p^{\text{sol}} &= \tilde{U}_p + \mathcal{O} \left(\frac{\bar{\mathcal{D}}^2}{M_U}, \frac{m_{3/2}}{M_U} \right), \\ Z^{\text{sol}} &= -\frac{C^{*2}}{C} M_{\text{SUSY}}^2 \Lambda^2 + \mathcal{O} \left(\frac{\bar{\mathcal{D}}^2}{m_z}, \frac{m_{3/2}}{m_z} \right), \end{aligned} \quad (2.17)$$

where Λ^α is the Goldstino superfield defined as

$$\Lambda^\alpha = \theta^\alpha + \frac{1}{M_{\text{SUSY}}^2} \tilde{z}^\alpha + \dots, \quad (2.18)$$

with the ellipsis denoting the Goldstino-dependent higher order terms. Note that $\bar{\mathcal{D}}^2$ acting on light field eventually gives rise to an F -component which is of $\mathcal{O}(m_{3/2})$ up to a factor of $\mathcal{O}(8\pi^2)$.

One can now derive the effective action of light fields by replacing U_p and Z with U_p^{sol} and Z^{sol} . At leading order in e^{-A} , the effective action is obtained by simply replacing U_p with \tilde{U}_p , and Z with $-M_{\text{SUSY}}^2 \Lambda^2 C^{*2}/C$:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{4D}|_{U_p=\tilde{U}_p, Z=-M_{\text{SUSY}}^2 \Lambda^2 C^{*2}/C}. \quad (2.19)$$

The resulting effective action is given by

$$\begin{aligned} \mathcal{L}_{\text{eff}} &= \int d^4\theta \left[-3CC^* e^{-K_{\text{eff}}/3} - C^2 C^{*2} M_{\text{SUSY}}^4 \Lambda^2 \bar{\Lambda}^2 \right] \\ &+ \left[\int d^2\theta \left(\frac{1}{4} f_a^{\text{eff}} W^{a\alpha} W_\alpha^a + C^3 W_{\text{eff}} \right) + \text{h.c.} \right], \end{aligned} \quad (2.20)$$

where

$$\begin{aligned}
 K_{\text{eff}} &= K_0(T_I + T_I^*) + \mathcal{Z}_A(T_I + T_I^*)\Phi^A\Phi^{A*}, \\
 f_a^{\text{eff}} &= \tilde{f}_a(T_I) + \frac{b_a}{16\pi^2} \ln C = \sum_I k_I T_I + \epsilon + \frac{b_a}{16\pi^2} \ln C, \\
 W_{\text{eff}} &= w_0 + \sum_I A_I e^{-8\pi^2 a_I T_I} + \frac{\kappa X^n}{M_{Pl}^{n-3}} + \lambda_\Phi X \Phi \Phi^c + \frac{1}{6} \lambda_{ijk} Q_i Q_j Q_k, \quad (2.21)
 \end{aligned}$$

for the moduli Kähler potential K_0 and the matter Kähler metric \mathcal{Z}_A determined as

$$\begin{aligned}
 -3e^{-K_0/3} &= \Omega_{\text{mod}}(U_p, U_p^*, T_I + T_I^*)|_{U_p=\tilde{U}_p}, \\
 e^{-K_0/3} \mathcal{Z}_A &= \mathcal{Y}_A(U_p, U_p^*, \ln(CC^*), T_I + T_I^*)|_{U_p=\tilde{U}_p}. \quad (2.22)
 \end{aligned}$$

The above effective action is defined at a renormalization point μ below the compactification scale, but above the mass of the exotic matter field $\Phi + \Phi^c$. Note that this renormalization point can be higher than the masses of the integrated heavy moduli U_p and Polony scalar z , while it should be lower than the mass scale characterizing the non-renormalizable interactions between the integrated fields and the remained light fields, which is of order the compactification scale or the GUT scale in our case. In the procedure to integrate out the Polony scalar to obtain the Akulov-Volkov action of the Goldstino superfield, we have used the following identity for the Goldstino superfield in the flat spacetime limit [17]:

$$\frac{1}{4} \bar{\mathcal{D}}^2 (\Lambda^2 \bar{\Lambda}^2) = -\Lambda^2 (1 - 2i\partial_\mu \Lambda \sigma^\mu \bar{\Lambda} - 4\bar{\Lambda}^2 \partial_\mu \Lambda \sigma^{\mu\nu} \partial_\nu \Lambda), \quad (2.23)$$

and ignored the higher order Goldstino operators as well as the higher derivative operators.

2.3 Supersymmetry breaking

In the class of models discussed here, the SUSY breaking field Z is sequestered from the visible sector, and then the MSSM soft terms are determined by F^C , F^{T_I} and F^X , which characterize the anomaly, moduli, and gauge mediation, respectively. As we have noticed, in order to get a nearly vanishing cosmological constant with $M_{\text{SUSY}}/M_{Pl} \sim e^{-A}$, one needs a special type of flux configuration yielding $m_{3/2}/M_{Pl} \sim w_0/M_{Pl}^3 \sim e^{-2A}$. On the other hand, nonperturbative dynamics generating the superpotential term $A_I e^{-8\pi^2 a_I T_I}$ originates from the UV end of throat, and thus there is no significant red-shift for A_I . This suggests that A_I are generically of order unity in the unit with $M_{Pl} = 1$, and then

$$\ln(A_I/w_0) \simeq \ln(M_{Pl}/m_{3/2}) \sim 4\pi^2 \quad (2.24)$$

for $m_{3/2} = \mathcal{O}(10)$ TeV. In the presence of such a big hierarchy between w_0 and A_I , much of the physical properties of T_I and X can be determined without knowing the explicit form of their Kähler potential. For instance, T_I are stabilized near the supersymmetric solution of $\partial_I W + (\partial_I K)W = 0$ with a mass

$$m_{T_I} \sim m_{3/2} \ln(M_{Pl}/m_{3/2}). \quad (2.25)$$

If $\kappa \neq 0$ for some $n > 3$, X is stabilized at an intermediate scale with

$$m_X \sim m_{3/2}. \quad (2.26)$$

The resulting vacuum expectation values of the scalar and F components of T_I and X (in the Einstein frame) are given by (see refs. [14, 19] for explicit derivations)

$$\begin{aligned} a_I T_I &\simeq \frac{\ln(M_{Pl}/m_{3/2})}{8\pi^2}, \\ X &\sim \left(\frac{M_{Pl}^{n-3} m_{3/2}}{\kappa} \right)^{1/(n-2)}, \\ \frac{F^{T_I}}{T_I + T_I^*} &\simeq \frac{1}{\ln(M_{Pl}/m_{3/2})} \frac{F^C}{C}, \\ \frac{F^X}{X} &\simeq -\frac{2}{n-1} \frac{F^C}{C}, \end{aligned} \quad (2.27)$$

where

$$\begin{aligned} \frac{F^C}{C} &= m_{3/2}^* + \frac{1}{3} F^P \partial_P K, \\ F^P &= -e^{K/2} K^{P\bar{Q}} (\partial_{\bar{Q}} W + (\partial_{\bar{Q}} K) W)^* \end{aligned} \quad (2.28)$$

for $\Phi^P = (T_I, X)$, and we have used $\ln(A_I/w_0) \simeq \ln(M_{Pl}/m_{3/2})$.

As is well known, F^C generates the anomaly-mediated soft parameters of $\mathcal{O}(F^C/8\pi^2)$ [4] at a high messenger scale around the compactification scale, while F^{T_I} generates moduli-mediated soft parameters of $\mathcal{O}(F^{T_I})$ [5] at a similar high messenger scale. In addition to these, the exotic vector-like matter fields $\Phi + \Phi^c$ give rise to a gauge-mediated contribution of $\mathcal{O}(F^{M_\Phi}/8\pi^2 M_\Phi)$ [2, 3] at the messenger scale M_Φ , where M_Φ and F^{M_Φ} denote the scalar component and the F component, respectively, of the messenger mass given by

$$\int d^2\theta C^3 (M_\Phi + \theta^2 F^{M_\Phi}) \Phi \Phi^c. \quad (2.29)$$

In the SUGRA models of the form (2.21), $\Phi + \Phi^c$ get a mass through the superpotential coupling $\lambda_\Phi X \Phi \Phi^c$, and then

$$\frac{F^{M_\Phi}}{M_\Phi} = \frac{F^X}{X} = -\frac{2}{n-1} \frac{F^C}{C} \quad \text{with } M_\Phi = \lambda_\Phi \langle X \rangle. \quad (2.30)$$

The most interesting feature of these SUGRA models is that

$$\frac{F^{T_I}}{T_I + T_I^*} \sim \frac{1}{8\pi^2} \frac{F^C}{C} \sim \frac{1}{8\pi^2} \frac{F^X}{X}, \quad (2.31)$$

independently of the Kähler potential. As a result, the MSSM soft parameters receive a similar size of contribution from all of the moduli, anomaly, and gauge mediations.

Another interesting feature is that the phases of F components are dynamically aligned to each other as

$$\arg \left[\frac{F^C}{C} \left(\frac{F^{T_I}}{T_I + T_I^*} \right)^* \right] = \arg \left[\frac{F^C}{C} \left(\frac{F^X}{X} \right)^* \right] = 0. \quad (2.32)$$

With this feature, soft terms preserve CP, although they receive a comparable contribution from three different origins. For this dynamical alignment, the axionic shift symmetry (2.1) plays an essential role [7–9]. To see this, let us note that one can always make w_0 in the superpotential to be real by an appropriate $U(1)_R$ transformation of the Grassmann variables, make A_I real by an axionic shift of T_I , and finally make κ real by a phase rotation of X , under which the Kähler potential is invariant. In this field basis, it is straightforward to see that $\text{Im}(T_I)$ and $\arg(X)$ are stabilized at a CP conserving value, and therefore W , $e^{-8\pi^2 a_I T_I}$, and X have real vacuum values. As the Kähler potential is invariant under the axionic shift symmetry (2.1) and the phase rotation of X , the resulting vacuum values of $\frac{F^C}{C}$, $\frac{F^X}{X}$, and F^{T_I} are all real.

In deflected mirage mediation, soft parameters can preserve flavor in a natural way. To satisfy the FCNC constraints, the following moduli-mediated sfermion masses and A -parameters are required to be (approximately) family-independent:

$$\begin{aligned} \tilde{m}_i^2 &= -F^{T_I} F^{T_J^*} \partial_{T_I} \partial_{T_J^*} \ln \left(e^{-K_0/3} \mathcal{Z}_i \right), \\ \tilde{A}_{ijk} &= F^{T_I} \partial_{T_I} \ln \left(e^{-K_0} \mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k \right), \end{aligned} \quad (2.33)$$

where \mathcal{Z}_i is the Kähler metric of the MSSM matter field Q_i . At leading order in the string coupling g_{st} or the string slope parameter α' , the T_I -dependence of \mathcal{Z}_i is typically given by [5]

$$\mathcal{Z}_i = \prod_I (T_I + T_I^*)^{n_I^i}, \quad (2.34)$$

where n_I^i is the modular weight of Q_i . If different families with the same gauge charges originate from the same type of branes or brane intersections, which is indeed the case in most of semi-realistic string models, the matter modular weights are *family-independent* rational numbers [8, 18–20], for which the resulting \tilde{m}_i^2 and \tilde{A}_{ijk} are family-independent.

In the above, we have considered the models of deflected mirage mediation, in which the gauge messengers get a mass through the superpotential coupling $\lambda_\Phi X \Phi \Phi^c$ with X stabilized at an intermediate scale, either by radiative effects or by the higher dimensional operator $\kappa X^n / M_{Pl}^{n-3}$ ($n > 3$). In fact, one can consider a different way to generate the gauge messenger mass, which would still give $\frac{F^{M_\Phi}}{M_\Phi} \sim \frac{F^C}{C}$. For instance, the gauge messengers may get a mass through the Kähler potential operator [21]

$$\int d^4\theta C C^* \left(c_\Phi \Phi \Phi^c + \text{h.c.} \right), \quad (2.35)$$

where c_Φ is a generic function of moduli. In this case, we have

$$\frac{F^{M_\Phi}}{M_\Phi} \simeq -2 \frac{F^C}{C} \quad \text{with} \quad M_\Phi = \mathcal{O}(m_{3/2}). \quad (2.36)$$

One may consider a more involved model [22], in which the gauge messengers get a mass from

$$\int d^4\theta CC^* \left(\frac{\mathcal{Y}_X}{2} X^* X + c_\Phi \Phi \Phi^c + \frac{c_X}{2} X^2 \right) + \int d^2\theta C^3 \left(\frac{\kappa}{6} X^3 + \lambda_\Phi X \Phi \Phi^c \right) + \text{h.c.}, \quad (2.37)$$

where $\mathcal{Y}_X, c_\Phi, c_X, \kappa$, and λ_Φ are again generic functions of moduli.³ One then finds [22]

$$\frac{F^{M_\Phi}}{M_\Phi} \simeq -\frac{8 - x_1 x_2}{4(1 - x_1)} \frac{F^C}{C} \quad \text{with} \quad M_\Phi = \mathcal{O}(m_{3/2}), \quad (2.38)$$

where

$$\begin{aligned} x_1 &= \frac{\lambda_\Phi (3c_X + \sqrt{c_X(c_X - 8\mathcal{Y}_X)})}{2\kappa c_\Phi}, \\ x_2 &= \frac{c_X + 4\mathcal{Y}_X - \sqrt{c_X(c_X - 8\mathcal{Y}_X)}}{\mathcal{Y}_X}. \end{aligned} \quad (2.39)$$

It is also possible to have a model in which the ratio $\frac{F^{M_\Phi}}{M_\Phi} / \frac{F^C}{C}$ takes a positive value of order unity, while the messenger scale is at an arbitrary intermediate scale [23]. One such an example would be the model with a composite X having an Affleck-Dine-Seiberg superpotential:

$$\int d^2\theta C^3 \left(\frac{\Lambda_X^{3-l}}{X^l} + \lambda_\Phi X \Phi \Phi^c \right), \quad (2.40)$$

where Λ_X is a dynamical scale hierarchically lower than M_{GUT} and l is a positive rational number. One then finds

$$\frac{F^{M_\Phi}}{M_\Phi} = \frac{2}{l+1} \frac{F^C}{C} \quad \text{with} \quad M_\Phi \sim \left(\frac{\Lambda_X^{3+l}}{m_{3/2}} \right)^{1/(l+2)}. \quad (2.41)$$

3 Soft parameters

In this section, we examine the renormalization group (RG) running of soft parameters in deflected mirage mediation. In particular, we derive (approximate) analytic expressions of low energy soft parameters, expressed in terms of the SUGRA model parameters defined in the previous section. Our results can be used to interpret the TeV scale sparticle masses measured in future collider experiments within the framework of the most general mixed mediation scheme preserving flavor and CP.

3.1 Soft parameters at scales above the gauge threshold scale

We first examine the soft parameters at scales above the gauge threshold scale set by the gauge messenger mass M_Φ . Our starting point is the effective SUGRA action (2.21) which has been obtained after integrating out the flux-stabilized heavy moduli and the sequestered

³Here we assume that all of these coefficients have real vacuum values. Unless, the model generically suffers from the SUSY CP problem.

SUSY breaking sector. At high scales above M_Φ , but below the gauge coupling unification scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV, the running gauge coupling and the running Kähler metric of the MSSM matter superfield Q_i are given by

$$\begin{aligned} \frac{1}{g_a^2(\mu/\sqrt{CC^*})} &= \text{Re}(\tilde{f}_a) - \frac{b_a^H}{16\pi^2} \ln\left(\frac{\mu^2}{CC^*M_{\text{GUT}}^2}\right), \\ \ln \mathcal{Z}_i(\mu/\sqrt{CC^*}) &= \ln \tilde{\mathcal{Z}}_i + \frac{1}{8\pi^2} \int_{M_{\text{GUT}}}^{\mu/\sqrt{CC^*}} \frac{d\mu'}{\mu'} \gamma_i(\mu'), \end{aligned} \quad (3.1)$$

where

$$\begin{aligned} \tilde{f}_a &= \sum_I k_I T_I + \epsilon, \\ \tilde{\mathcal{Z}}_i &= \prod_I (T_I + T_I^*)^{n_i^I}, \end{aligned} \quad (3.2)$$

and b_a^H and γ_i are the one loop beta function coefficients and the anomalous dimensions at scales between M_Φ and M_{GUT} :

$$\begin{aligned} b_a^H &= -3T_a(\text{Adj}) + \sum_i T_a(Q_i) + \sum_\Phi (T_a(\Phi) + T_a(\Phi^c)), \\ \gamma_i &= 2 \sum_a C_2^a(Q_i) g_a^2 - \frac{1}{2} \sum_{jk} |y_{ijk}|^2, \end{aligned} \quad (3.3)$$

where y_{ijk} are the canonical Yukawa couplings given by

$$y_{ijk}(\mu) = \frac{\lambda_{ijk}}{\sqrt{e^{-K_0} \mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k}}. \quad (3.4)$$

Here we have ignored the T_I -dependent Kähler and Konishi anomaly contributions to the running gauge coupling constants [25], which are determined by K_0 and \mathcal{Z}_A , and also the UV sensitive string and KK threshold corrections. Those T_I -dependent loop corrections give a contribution of $\mathcal{O}\left(\frac{FT_I}{8\pi^2}\right)$ to soft parameters, which are subleading compared to the contributions which will be discussed below. We also put the superscript H for the high scale beta function coefficients b_a^H in order to distinguish them from the low scale MSSM beta function coefficients. Note that the vacuum value of $\text{Re}(\tilde{f}_a)$ corresponds to the unified gauge coupling constant at M_{GUT} :

$$\text{Re}(\tilde{f}_a) = \sum_I k_I \text{Re}(T_I) + \text{Re}(\epsilon) = \frac{1}{g_{\text{GUT}}^2}. \quad (3.5)$$

The soft SUSY breaking terms are parameterized as

$$-\mathcal{L}_{\text{soft}} = m_i^2 |\phi_i|^2 + \left[\frac{1}{2} M_a \lambda^a \lambda^a + \frac{1}{6} A_{ijk} y_{ijk} \tilde{Q}_i \tilde{Q}_j \tilde{Q}_k + \text{h.c.} \right], \quad (3.6)$$

where λ^a and \tilde{Q}_i are canonically normalized gauginos and sfermions, respectively. Then at scales between M_Φ and M_{GUT} , the running soft parameters are given by

$$\begin{aligned}
M_a(\mu) &= -\left(F^{T_I} \partial_{T_I} + F^C \partial_C\right) \ln(g_a^2) \\
&= -F^{T_I} \partial_{T_I} \ln(g_a^2) + \frac{b_a^H}{16\pi^2} g_a^2(\mu) \frac{F^C}{C}, \\
A_{ijk}(\mu) &= -\left(F^{T_I} \partial_{T_I} + F^C \partial_C\right) \ln\left(\frac{\lambda_{ijk}}{e^{-K_0} \mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k}\right) \\
&= F^{T_I} \partial_{T_I} \ln(e^{-K_0} \mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k) - \frac{1}{16\pi^2} (\gamma_i + \gamma_j + \gamma_k) \frac{F^C}{C}, \\
m_i^2(\mu) &= -\left(F^{T_I} \partial_{T_I} + F^C \partial_C\right) \left(F^{T_J} \partial_{T_J} + F^C \partial_C\right)^* \ln(e^{-K_0/3} \mathcal{Z}_i) \\
&= -F^{T_I} F^{T_J*} \partial_{T_I} \partial_{T_J} \ln(e^{-K_0/3} \mathcal{Z}_i) + \left[\frac{\tilde{\gamma}_i^*}{16\pi^2} \frac{F^C}{C} + \text{h.c.}\right] - \frac{\dot{\gamma}_i}{32\pi^2} \left|\frac{F^C}{C}\right|^2, \quad (3.7)
\end{aligned}$$

where

$$\begin{aligned}
\gamma_i &= 8\pi^2 \frac{d \ln \mathcal{Z}_i}{d \ln \mu} = 2 \sum_a C_2^a(Q_i) g_a^2 - \frac{1}{2} \sum_{jk} |y_{ijk}|^2, \\
\tilde{\gamma}_i &= F^{T_I} \partial_{T_I} \gamma_i = 2 \sum_a C_2^a(Q_i) F^{T_I} \partial_{T_I} g_a^2 + \frac{1}{2} \sum_{jk} |y_{ijk}|^2 F^{T_I} \partial_{T_I} \ln(e^{-K_0} \mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k), \\
\dot{\gamma}_i &= \frac{d\gamma_i}{d \ln \mu} = \frac{1}{4\pi^2} \sum_a C_2^a(Q_i) b_a^H g_a^4 + \frac{1}{16\pi^2} \sum_{jk} |y_{ijk}|^2 (\gamma_i + \gamma_j + \gamma_k). \quad (3.8)
\end{aligned}$$

The above running soft parameters correspond to the solution of the following RG equations [26]:⁴

$$\begin{aligned}
\frac{dM_a}{d \ln \mu} &= \frac{b_a^H}{8\pi^2} g_a^2 M_a, \\
\frac{dA_{ijk}}{d \ln \mu} &= -\frac{1}{4\pi^2} \sum_a \left[C_2^a(Q_i) + C_2^a(Q_j) + C_2^a(Q_k)\right] g_a^2 M_a \\
&\quad + \frac{1}{16\pi^2} \sum_{lm} \left(A_{ilm} |y_{ilm}|^2 + A_{jlm} |y_{jlm}|^2 + A_{klm} |y_{klm}|^2\right), \\
\frac{dm_i^2}{d \ln \mu} &= \frac{1}{16\pi^2} \left[-8 \sum_a C_2^a(Q_i) g_a^2 |M_a|^2 + \sum_{jk} \left(m_i^2 + m_j^2 + m_k^2 + |A_{ijk}|^2\right) |y_{ijk}|^2\right] \\
&\quad + \frac{1}{8\pi^2} g_Y^2 Y_i \left[\sum_j Y_j m_j^2 + \sum_\Phi \left(Y_\Phi m_\Phi^2 + Y_{\Phi^c} m_{\Phi^c}^2\right)\right], \quad (3.9)
\end{aligned}$$

⁴ It is noted that the gaugino masses and A -parameters are a linear superposition of the solutions for two mediations and $\tilde{\gamma}_i(\mu)$ is determined by the solution for moduli mediation at that scale. Thus, once we obtain the soft parameters for moduli mediation at an arbitrary scale, we can reconstruct those of mirage mediation without solving the RG equation again.

with the boundary condition at the scale just below M_{GUT} :

$$\begin{aligned}
 M_a(M_{\text{GUT}}) &= M_0 + \frac{b_a^H}{16\pi^2} g_{\text{GUT}}^2 \frac{F^C}{C}, \\
 A_{ijk}(M_{\text{GUT}}) &= \tilde{A}_{ijk} - \frac{1}{16\pi^2} \left(\gamma_i(M_{\text{GUT}}) + \gamma_j(M_{\text{GUT}}) + \gamma_k(M_{\text{GUT}}) \right) \frac{F^C}{C}, \\
 m_i^2(M_{\text{GUT}}) &= \tilde{m}_i^2 + \left[\frac{\tilde{\gamma}_i^*(M_{\text{GUT}})}{16\pi^2} \frac{F^C}{C} + \text{h.c.} \right] - \frac{\dot{\gamma}_i(M_{\text{GUT}})}{32\pi^2} \left| \frac{F^C}{C} \right|^2,
 \end{aligned} \tag{3.10}$$

where Y_i , Y_Φ and Y_{Φ^c} denote the $U(1)_Y$ charges of Q_i , Φ and Φ^c , respectively, and

$$\begin{aligned}
 M_0 &\equiv F^{T_I} \partial_{T_I} \ln(\text{Re}(\tilde{f}_a)) = \frac{g_{\text{GUT}}^2}{2} \sum_I k_I F^{T_I}, \\
 \tilde{A}_{ijk} &\equiv F^{T_I} \partial_{T_I} \ln(e^{-K_0} \tilde{Z}_i \tilde{Z}_j \tilde{Z}_k), \\
 \tilde{m}_i^2 &\equiv -F^{T_I} F^{T_J^*} \partial_{T_I} \partial_{T_J^*} \ln(e^{-K_0/3} \tilde{Z}_i).
 \end{aligned} \tag{3.11}$$

In view of (3.1), \tilde{f}_a and \tilde{Z}_i correspond to the gauge kinetic function and the matter Kähler metric at M_{GUT} , and thus M_0 , \tilde{A}_{ijk} and \tilde{m}_i^2 correspond to the moduli-mediated soft parameters at M_{GUT} . As most of our discussion will be independent of their explicit form, in the following, we will not use any specific form of the moduli Kähler potential K_0 and the matter Kähler metric \tilde{Z}_i , but instead treat \tilde{A}_{ijk} and \tilde{m}_i^2 as family-independent free parameters constrained by the $SU(5)$ unification relations.

As was noticed in [10], the RG equations (3.9) with the boundary conditions (3.10) have a useful form of analytic solution. For the gaugino masses at $\mu > M_\Phi$, one easily finds

$$M_a(\mu) = M_0 \left[1 + \frac{b_a^H}{8\pi^2} g_a^2(\mu) \ln \left(\frac{\mu}{M_{\text{mir}}} \right) \right], \tag{3.12}$$

with the running gauge coupling constants:

$$\frac{1}{g_a^2(\mu)} = \frac{1}{g_{\text{GUT}}^2} - \frac{b_a^H}{8\pi^2} \ln \left(\frac{\mu}{M_{\text{GUT}}} \right), \tag{3.13}$$

and the mirage scale M_{mir} given by

$$M_{\text{mir}} = M_{\text{GUT}} \left(\frac{m_{3/2}}{M_{\text{Pl}}} \right)^{\alpha/2}, \tag{3.14}$$

where α parameterizes the anomaly to moduli mediation ratio:

$$\alpha = \frac{F^C/C}{M_0 \ln(M_{\text{Pl}}/m_{3/2})} \simeq \frac{m_{3/2}}{M_0 \ln(M_{\text{Pl}}/m_{3/2})}. \tag{3.15}$$

For the A -parameters and sfermion masses, similar analytic expressions are available *if*

- (i) the involved Yukawa couplings are negligible, or
- (ii)

$$\frac{\tilde{A}_{ijk}}{M_0} = \frac{\tilde{m}_i^2 + \tilde{m}_j^2 + \tilde{m}_k^2}{M_0^2} = 1 \text{ for non-negligible Yukawa coupling } y_{ijk}. \tag{3.16}$$

In such cases, one finds [10]

$$\begin{aligned}
 A_{ijk}(\mu) &= \tilde{A}_{ijk} - \frac{1}{8\pi^2} \left(\gamma_i(\mu) + \gamma_j(\mu) + \gamma_k(\mu) \right) M_0 \ln \left(\frac{\mu}{M_{\text{mir}}} \right), \\
 m_i^2(\mu) &= \tilde{m}_i^2 - \frac{1}{4\pi^2} \gamma_i(\mu) M_0^2 \ln \left(\frac{\mu}{M_{\text{mir}}} \right) - \frac{1}{8\pi^2} \dot{\gamma}_i(\mu) M_0^2 \left[\ln \left(\frac{\mu}{M_{\text{mir}}} \right) \right]^2 \\
 &\quad + \frac{1}{8\pi^2} Y_i \text{Tr}(Y \tilde{m}^2) g_Y^2(\mu) \ln \left(\frac{\mu}{M_{\text{GUT}}} \right),
 \end{aligned} \tag{3.17}$$

where

$$\text{Tr}(Y \tilde{m}^2) = \sum_i Y_i \tilde{m}_i^2 + \sum_{\Phi} \left(Y_{\Phi} \tilde{m}_{\Phi}^2 + Y_{\Phi^c} \tilde{m}_{\Phi^c}^2 \right) \tag{3.18}$$

for the $U(1)_Y$ charge operator Y .

The analytic solutions of (3.12) and (3.17) show that

$$M_a(M_{\text{mir}}) = M_0, \quad A_{ijk}(M_{\text{mir}}) = \tilde{A}_{ijk}, \quad m_i^2(M_{\text{mir}}) = \tilde{m}_i^2, \tag{3.19}$$

if $\text{Tr}(Y \tilde{m}^2) = 0$, which is satisfied for instance when the moduli-mediated sfermion masses at M_{GUT} satisfy the $SU(5)$ unification condition and $\tilde{m}_{H_u}^2 = \tilde{m}_{H_d}^2$. In other words, the soft parameters renormalized at $\mu = M_{\text{mir}}$ become identical to the pure moduli-mediated parameters renormalized at M_{GUT} , obeying the unification condition. With this observation, M_{mir} has been dubbed the mirage messenger scale as it does not correspond to any physical threshold scale [10]. Note that still the gauge couplings are unified at the conventional GUT scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV.

If there exist non-negligible Yukawa couplings for which the mirage condition (3.16) is not satisfied, the above analytic solutions of A_{ijk} and m_i^2 are not valid anymore. However, if there is only one such Yukawa coupling, one can still find a useful analytic expression for the running A_{ijk} and m_i^2 . The results are presented in the appendix for m_i^2 ($i = H_u, q_3, u_3$) and $A_{H_u q_3 u_3}$ in the MSSM, including only the effects of the top quark Yukawa coupling.

3.2 Soft parameters below the gauge threshold scale

So far, we have discussed the soft parameters at scales above the gauge threshold scale M_{Φ} . Those high scale soft parameters are determined by anomaly and moduli mediations, and the gaugino and light-family sfermion masses follow the mirage unification trajectory given by (3.12) and (3.17). If there were no exotic matter fields $\Phi + \Phi^c$, soft parameters would follow these analytic solutions down to the TeV scale. However, in deflected mirage mediation, soft parameters at scales below M_{Φ} are deflected from the mirage unification trajectory due to the gauge mediation by $\Phi + \Phi^c$.

Let us examine how the low energy soft parameters are affected by $\Phi + \Phi^c$ which have a mass-superfield:

$$\int d^2\theta C^3 \left(M_{\Phi} + \theta^2 F^{M_{\Phi}} \right) \Phi \Phi^c + \text{h.c.} \tag{3.20}$$

To compute the low energy soft parameters, one can add the gauge threshold contribution at $\mu = M_{\Phi}^-$ to the soft parameters of (3.12) and (3.17) evaluated at $\mu = M_{\Phi}^+$, and then apply the RG equation at lower scales. (Here M_{Φ}^+ and M_{Φ}^- denote the mass scale just above M_{Φ} and the scale just below M_{Φ} , respectively.) The gauge threshold contributions at M_{Φ}^- are given by

$$\begin{aligned}
 \Delta M_a(M_{\Phi}) &= -F \cdot \partial \ln(g_a^2(\mu)) \Big|_{\mu=M_{\Phi}^-} + F \cdot \partial \ln(g_a^2(\mu)) \Big|_{\mu=M_{\Phi}^+} \\
 &= -\frac{N_{\Phi}}{16\pi^2} g_a^2(M_{\Phi}) \left(\frac{F^{M_{\Phi}}}{M_{\Phi}} + \frac{F^C}{C} \right), \\
 \Delta A_{ijk}(M_{\Phi}) &= F \cdot \partial \ln(\mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k) \Big|_{\mu=M_{\Phi}^-} - F \cdot \partial \ln(\mathcal{Z}_i \mathcal{Z}_j \mathcal{Z}_k) \Big|_{\mu=M_{\Phi}^+} \\
 &= -\frac{1}{16\pi^2} \left(\gamma_i(M_{\Phi}^-) - \gamma_i(M_{\Phi}^+) \right) \left(\frac{F^{M_{\Phi}}}{M_{\Phi}} + \frac{F^C}{C} \right) = 0, \\
 \Delta m_i^2(M_{\Phi}) &= - (F \cdot \partial)(\bar{F} \cdot \bar{\partial}) \ln \mathcal{Z}_i(\mu) \Big|_{\mu=M_{\Phi}^-} + (F \cdot \partial)(\bar{F} \cdot \bar{\partial}) \ln \mathcal{Z}_i(\mu) \Big|_{\mu=M_{\Phi}^+} \\
 &= -\frac{1}{32\pi^2} \left(\dot{\gamma}_i(M_{\Phi}^-) - \dot{\gamma}_i(M_{\Phi}^+) \right) \left| \frac{F^{M_{\Phi}}}{M_{\Phi}} + \frac{F^C}{C} \right|^2 \\
 &= \frac{N_{\Phi}}{(16\pi^2)^2} 2C_2^a(\phi^i) g_a^4(M_{\Phi}) \left| \frac{F^{M_{\Phi}}}{M_{\Phi}} + \frac{F^C}{C} \right|^2, \tag{3.21}
 \end{aligned}$$

where $F \cdot \partial = F^{T_I} \partial_{T_I} + F^C \partial_C + F^{M_{\Phi}} \partial_{M_{\Phi}}$, and we have assumed that there are N_{Φ} flavors of $\Phi + \Phi^c$ which form the $5 + \bar{5}$ representation of $SU(5)$.

We can now obtain the soft parameters at $\mu < M_{\Phi}$ by solving the RG equation with the boundary conditions:

$$\begin{aligned}
 M_a(M_{\Phi}^-) &= M_a(M_{\Phi}^+) + \Delta M_a(M_{\Phi}), \\
 A_{ijk}(M_{\Phi}^-) &= A_{ijk}(M_{\Phi}^+) + \Delta A_{ijk}(M_{\Phi}), \\
 m_i^2(M_{\Phi}^-) &= m_i^2(M_{\Phi}^+) + \Delta m_i^2(M_{\Phi}), \tag{3.22}
 \end{aligned}$$

where the soft parameters at M_{Φ}^+ can be obtained from the high scale solutions, (3.12) and (3.17), by replacing μ with M_{Φ}^+ . Like the case of high scale solutions, it turns out that the resulting low energy solutions allow analytic expression which can be used to interpret the TeV scale sparticle masses. For instance, gaugino masses are given by [27]

$$M_a(\mu) = M_0^{\text{eff}} \left[1 + \frac{1}{8\pi^2} b_a g_a^2(\mu) \ln \left(\frac{\mu}{M_{\text{mir}}^{\text{eff}}} \right) \right], \tag{3.23}$$

where

$$\begin{aligned}
 M_0^{\text{eff}} &= R M_0, \\
 M_{\text{mir}}^{\text{eff}} &= M_{\text{GUT}} \left(\frac{m_{3/2}}{M_{\text{Pl}}} \right)^{\alpha/2R}, \\
 b_a &= -3T_a(\text{Adj}) + \sum_i T_a(Q_i) \tag{3.24}
 \end{aligned}$$

for

$$\begin{aligned}
M_0 &= F^{T_I} \partial_{T_I} \ln(\text{Re}(\tilde{f}_a)), \\
R &= 1 + \frac{N_\Phi g_0^2}{8\pi^2} \left[\frac{\alpha}{2\beta} \ln\left(\frac{M_{Pl}}{m_{3/2}}\right) - \ln\left(\frac{M_{GUT}}{M_\Phi}\right) \right], \\
\alpha &= \frac{F^C/C}{M_0 \ln(M_{Pl}/m_{3/2})}, \\
\beta &= -\frac{F^C/C}{F M_\Phi/M_\Phi}.
\end{aligned} \tag{3.25}$$

Here α and β parameterize the anomaly to moduli mediation ratio and the anomaly to gauge mediation ratio, respectively, M_0 is the moduli-mediated gaugino mass at M_{GUT} , and $g_0^2 \simeq 1/2$ corresponds to the unified gauge coupling constant in the absence of $\Phi + \Phi^c$, i.e.

$$\frac{1}{g_0^2} = \frac{1}{g_{GUT}^2} + \frac{N_\Phi}{8\pi^2} \ln\left(\frac{M_{GUT}}{M_\Phi}\right).$$

The running A -parameters and sfermion masses at $\mu < M_\Phi$ take a more involved form. Neglecting the effects of Yukawa couplings, we find

$$\begin{aligned}
A_{ijk}(\mu) &= \tilde{A}_{ijk}^{\text{eff}} - \frac{1}{8\pi^2} (\gamma_i(\mu) + \gamma_j(\mu) + \gamma_k(\mu)) M_0^{\text{eff}} \ln\left(\frac{\mu}{M_{\text{mir}}^{\text{eff}}}\right), \\
m_i^2(\mu) &= (\tilde{m}_i^{\text{eff}})^2 - \frac{1}{4\pi^2} \gamma_i(\mu) (M_0^{\text{eff}})^2 \ln\left(\frac{\mu}{M_{\text{mir}}^{\text{eff}}}\right) - \frac{1}{8\pi^2} \dot{\gamma}_i(\mu) (M_0^{\text{eff}})^2 \left[\ln\left(\frac{\mu}{M_{\text{mir}}^{\text{eff}}}\right) \right]^2 \\
&\quad + \frac{1}{8\pi^2} Y_i \left(\sum_j Y_j m_j^2(M_\Phi^-) \right) g_Y^2(\mu) \ln\left(\frac{\mu}{M_\Phi}\right),
\end{aligned} \tag{3.26}$$

where

$$\begin{aligned}
\tilde{A}_{ijk}^{\text{eff}} &= \tilde{A}_{ijk} + \frac{1}{8\pi^2} (1-R) M_0 (\gamma_i(M_\Phi) + \gamma_j(M_\Phi) + \gamma_k(M_\Phi)) \ln\left(\frac{M_{GUT}}{M_\Phi}\right), \\
(\tilde{m}_i^{\text{eff}})^2 &= \tilde{m}_i^2 + 2(1-R)^2 M_0^2 \sum_a C_2^a(Q_i) \left[\frac{1}{N_\Phi} \frac{g_a^4(M_\Phi)}{g_0^4} \right. \\
&\quad \left. + \frac{g_a^2(M_\Phi)}{8\pi^2} \left(\frac{1+R}{1-R} - \frac{g_a^2(M_\Phi)}{g_0^2} \right) \ln\left(\frac{M_{GUT}}{M_\Phi}\right) \right].
\end{aligned} \tag{3.27}$$

Here

$$\frac{1}{g_a^2(M_\Phi)} = \frac{1}{g_{GUT}^2} + \frac{b_a + N_\Phi}{8\pi^2} \ln\left(\frac{M_{GUT}}{M_\Phi}\right) = \frac{1}{g_0^2} + \frac{b_a}{8\pi^2} \ln\left(\frac{M_{GUT}}{M_\Phi}\right), \tag{3.28}$$

where b_a are the MSSM beta function coefficients. For the RG contribution associated with $\text{Tr}(Y m^2)$, using the gauge invariance of Yukawa interactions and the anomaly cancellation conditions, we find

$$\sum_i Y_i m_i^2(M_\Phi^-) = \frac{5}{3} \frac{g_Y^2(M_\Phi)}{g_0^2} \sum_i Y_i \tilde{m}_i^2 + \left(\frac{5}{3} \frac{g_Y^2(M_\Phi)}{g_0^2} - 1 \right) \sum_\Phi (Y_\Phi \tilde{m}_\Phi^2 + Y_{\Phi^c} \tilde{m}_{\Phi^c}^2), \tag{3.29}$$

which vanishes when the moduli-mediated sfermion masses at M_{GUT} are SU(5)-invariant and $\tilde{m}_{H_u}^2 = \tilde{m}_{H_d}^2$.⁵

Like the case that soft masses are dominated by one particular mediation, one can consider the sum rules of sfermion masses in deflected mirage mediation, which may be useful for identifying the structure of moduli mediation at M_{GUT} . For instance, from (3.26) and (3.27), we find the following relations amongst the light-family squark and slepton masses:⁶

$$\begin{aligned}
 & m_{\tilde{q}_L}^2(\mu) - 2m_{\tilde{u}_R}^2(\mu) + m_{\tilde{d}_R}^2(\mu) - m_{\tilde{l}_L}^2(\mu) + m_{\tilde{e}_R}^2(\mu) \\
 &= \tilde{m}_{\tilde{q}_L}^2 - 2\tilde{m}_{\tilde{u}_R}^2 + \tilde{m}_{\tilde{d}_R}^2 - \tilde{m}_{\tilde{l}_L}^2 + \tilde{m}_{\tilde{e}_R}^2 + \frac{5}{3} \frac{g_Y^2(\mu)}{4\pi^2} \ln\left(\frac{\mu}{M_\Phi}\right) \sum_i Y_i m_i^2(M_\Phi^-), \\
 & 2m_{\tilde{q}_L}^2(\mu) - m_{\tilde{u}_R}^2(\mu) - m_{\tilde{d}_R}^2(\mu) - 2m_{\tilde{l}_L}^2(\mu) + m_{\tilde{e}_R}^2(\mu) \\
 &= 2\tilde{m}_{\tilde{q}_L}^2 - \tilde{m}_{\tilde{u}_R}^2 - \tilde{m}_{\tilde{d}_R}^2 - 2\tilde{m}_{\tilde{l}_L}^2 + \tilde{m}_{\tilde{e}_R}^2 + \frac{4}{3} \frac{g_Y^2(\mu)}{4\pi^2} \ln\left(\frac{\mu}{M_\Phi}\right) \sum_i Y_i m_i^2(M_\Phi^-), \quad (3.30)
 \end{aligned}$$

where $\tilde{q}_L, \tilde{q}_R = (\tilde{u}_R, \tilde{d}_R)$, \tilde{l}_L , and \tilde{e}_R denote the squark-doublet, squark-singlet, slepton-doublet, and slepton-singlet, respectively. If the moduli-mediated sfermion masses at M_{GUT} are SU(5)-invariant, i.e.

$$\begin{aligned}
 \tilde{m}_{\tilde{q}_L}^2 &= \tilde{m}_{\tilde{u}_R}^2 = \tilde{m}_{\tilde{e}_R}^2 = \tilde{m}_{10}^2, \\
 \tilde{m}_{\tilde{d}_R}^2 &= \tilde{m}_{\tilde{l}_L}^2 = \tilde{m}_5^2, \quad (3.31)
 \end{aligned}$$

and also $\tilde{m}_{H_u}^2 = \tilde{m}_{H_d}^2$, the above sum rules give

$$\begin{aligned}
 m_{\tilde{q}_L}^2(\mu) - 2m_{\tilde{u}_R}^2(\mu) + m_{\tilde{d}_R}^2(\mu) - m_{\tilde{l}_L}^2(\mu) + m_{\tilde{e}_R}^2(\mu) &= 0, \\
 2m_{\tilde{q}_L}^2(\mu) - m_{\tilde{u}_R}^2(\mu) - m_{\tilde{d}_R}^2(\mu) - 2m_{\tilde{l}_L}^2(\mu) + m_{\tilde{e}_R}^2(\mu) &= 2\tilde{m}_{10}^2 - 3\tilde{m}_5^2, \quad (3.32)
 \end{aligned}$$

indicating that these sum rules can be used to ascertain the existence of nonzero moduli-mediation as well as the GUT relations of the moduli-mediated sfermion masses [29].

For the effective SUGRA model (2.21), which is a representative class of model for deflected mirage mediation, it is straightforward to compute α and β , which gives

$$\begin{aligned}
 \alpha &\simeq 1 + \frac{\text{Re}(\epsilon)}{\sum_I k_I \text{Re}(T_I)}, \\
 \beta &= \frac{n-1}{2} \quad (n \geq 3), \quad (3.33)
 \end{aligned}$$

⁵ We note that $\tilde{A}_{ijk}^{\text{eff}}$ and $(\tilde{m}_i^{\text{eff}})^2$ correspond to the soft parameters at $M_{\text{mir}}^{\text{eff}}$ after removing the trace term proportional to Y_i . They are obtained by extrapolating the weak scale soft terms (subtracted anomaly mediation and the trace term) to M_{GUT} neglecting the gauge threshold scale. It is obvious that the low energy soft parameters are summarized in the mirage mediation pattern using these effective parameters because the superposition of anomaly mediation and other mediations closes at each scale as in (3.7) and anomaly mediation can not distinguish the origin of other contributions at higher scale [28].

⁶After the electroweak symmetry breaking, these sum rules are affected by the D -term contribution, which is of order M_Z^2 .

where we have used $\ln(A_I/w_0) \simeq \ln(M_{Pl}/m_{3/2})$. Here, the value of β for $n > 3$ applies to the model in which X is stabilized by the non-renormalizable superpotential term $\kappa X^n/M_{Pl}^{n-3}$, while the value of β for $n = 3$ applies to the model with $\kappa = 0$, in which X is stabilized by the radiative correction to its Kähler potential. In view of underlying string theory, $\text{Re}(\epsilon)$ corresponds to a higher order correction to the gauge kinetic function in the g_{st} or α' expansion. This suggests that $\text{Re}(\epsilon)$ is significantly smaller than $\sum_I k_I \text{Re}(T_I)$, and thus α has a value close to the unity.

With (3.23) and (3.26), providing the analytic expression of low energy soft parameters in deflected mirage mediation, one can take an appropriate limit to obtain the soft parameters in more familiar case dominated by a single mediation. Specifically, each single mediation corresponds to the limit:

$$\begin{aligned}
 * \text{ Anomaly mediation: } & \quad R \rightarrow 1, \quad \frac{1}{\alpha} \rightarrow 0, \quad \alpha M_0 = \text{finite}, \quad \tilde{A}_{ijk} = \tilde{m}_i^2 = 0, \\
 * \text{ Gauge mediation: } & \quad \frac{1}{R} \rightarrow 0, \quad RM_0 = \text{finite}, \quad \tilde{A}_{ijk} = \tilde{m}_i^2 = 0, \\
 * \text{ Moduli mediation: } & \quad R \rightarrow 1, \quad \alpha \rightarrow 0,
 \end{aligned}$$

while the mixed gauge-anomaly mediation (= deflected anomaly mediation) and the mixed moduli-anomaly mediation (= mirage mediation) can be obtained as

$$\begin{aligned}
 * \text{ Deflected anomaly mediation: } & \quad \frac{1}{R} \rightarrow 0, \quad \frac{\alpha}{R} = \text{finite}, \\
 & \quad RM_0 = \text{finite}, \quad \tilde{A}_{ijk} = \tilde{m}_i^2 = 0, \\
 * \text{ Mirage mediation: } & \quad R \rightarrow 1.
 \end{aligned}$$

Figure 1 summarizes these different limits of deflected mirage mediation in the parameter space spanned by α and R .

Soft parameters in case of multi-step gauge thresholds can be obtained by applying our results recursively. For instance, the gaugino masses after the n -step of thresholds are given by (3.23) with

$$\begin{aligned}
 M_0^{\text{eff}} &= R_n R_{n-1} \cdots R_1 M_0, \\
 \alpha^{\text{eff}} &= \frac{\alpha}{R_n R_{n-1} \cdots R_1},
 \end{aligned} \tag{3.34}$$

where

$$R_n = 1 + \frac{N_{\Phi_n} g_0^2}{8\pi^2} \left[\frac{1}{R_{n-1} \cdots R_1} \frac{\alpha}{2\beta_n} \ln \left(\frac{M_{Pl}}{m_{3/2}} \right) - \ln \left(\frac{M_{\text{GUT}}}{M_{\Phi_n}} \right) \right]$$

for N_{Φ_n} denoting the number of the gauge messenger pairs at the n -th threshold scale M_{Φ_n} , and β_n is the anomaly to gauge mediation ratio for the n -th gauge threshold. Light-family sfermion soft parameters also can be written as (3.26) with appropriately defined $\tilde{A}_{ijk}^{\text{eff}}$ and \tilde{m}_i^{eff} . As we will see below, such parametrization provides a useful set-up to interpret the TeV scale sparticle masses within the framework of the most general mixed moduli-anomaly-gauge mediation.

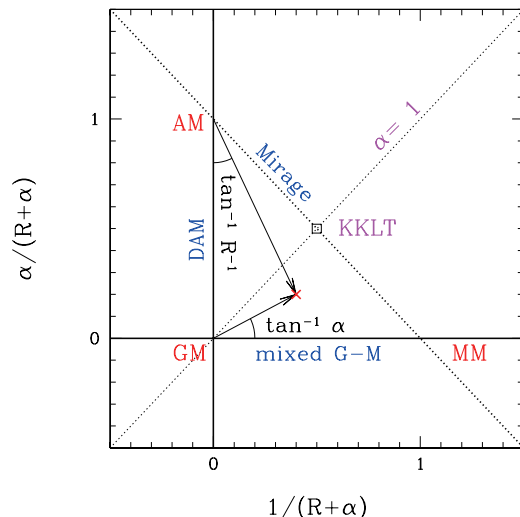


Figure 1. Parameter space of deflected mirage mediation spanned by α and R . Here AM, GM and MM denote anomaly mediation, gauge mediation and moduli mediation, respectively.

3.3 Sparticle masses at the TeV scale

From (3.23) and (3.26), one can obtain the low energy sparticle masses at the TeV scale. If one assumes that the moduli-mediated sfermion masses at M_{GUT} satisfy the SU(5) unification condition and also $\tilde{m}_{H_u}^2 = \tilde{m}_{H_d}^2$, the gaugino and light-family sfermion masses in generic deflected mirage mediation at the renormalization point $\mu = 500$ GeV are given by⁷

$$\begin{aligned}
 M_1 &= M_0^{\text{eff}} [0.43 + 0.29\alpha_{\text{eff}}], \\
 M_2 &= M_0^{\text{eff}} [0.83 + 0.084\alpha_{\text{eff}}], \\
 M_3 &= M_0^{\text{eff}} [2.5 - 0.74\alpha_{\text{eff}}], \\
 m_{\tilde{q}_L}^2 &= \tilde{m}_{10}^2 + (M_0^{\text{eff}})^2 [5.0 - 3.48\alpha_{\text{eff}} + 0.48\alpha_{\text{eff}}^2 + \delta_{\tilde{q}_L}], \\
 m_{\tilde{u}_R}^2 &= \tilde{m}_{10}^2 + (M_0^{\text{eff}})^2 [4.6 - 3.29\alpha_{\text{eff}} + 0.49\alpha_{\text{eff}}^2 + \delta_{\tilde{u}_R}], \\
 m_{\tilde{e}_R}^2 &= \tilde{m}_{10}^2 + (M_0^{\text{eff}})^2 [0.15 - 0.045\alpha_{\text{eff}} - 0.015\alpha_{\text{eff}}^2 + \delta_{\tilde{e}_R}], \\
 m_{\tilde{d}_R}^2 &= \tilde{m}_5^2 + (M_0^{\text{eff}})^2 [4.5 - 3.27\alpha_{\text{eff}} + 0.49\alpha_{\text{eff}}^2 + \delta_{\tilde{d}_R}], \\
 m_{\tilde{l}_L}^2 &= \tilde{m}_5^2 + (M_0^{\text{eff}})^2 [0.5 - 0.22\alpha_{\text{eff}} - 0.014\alpha_{\text{eff}}^2 + \delta_{\tilde{l}_L}],
 \end{aligned} \tag{3.35}$$

where

$$\begin{aligned}
 M_0^{\text{eff}} &= RM_0, \quad \alpha_{\text{eff}} = \alpha/R, \\
 \delta_i &\equiv \frac{(\tilde{m}_i^{\text{eff}})^2 - \tilde{m}_i^2}{(M_0^{\text{eff}})^2} = \sum_a C_2^a(Q_i)\delta_a,
 \end{aligned} \tag{3.36}$$

⁷ For colored sparticles, there can be a sizable difference between this running mass at $\mu = 500$ GeV and the physical mass [30].

for

$$\delta_a = \frac{2(1-R)^2}{R^2} \left[\frac{1}{N_\Phi} \frac{g_a^4(M_\Phi)}{g_0^4} + \frac{g_a^2(M_\Phi)}{8\pi^2} \left(\frac{1+R}{1-R} - \frac{g_a^2(M_\Phi)}{g_0^2} \right) \ln \left(\frac{M_{\text{GUT}}}{M_\Phi} \right) \right]. \quad (3.37)$$

One interesting limit of deflected mirage mediation is the pure mirage mediation in which there is no gauge-mediated contribution. In this limit, $R = 1$, and therefore

$$M_0^{\text{eff}} = M_0, \quad \alpha_{\text{eff}} = \alpha, \quad \delta_i = 0. \quad (3.38)$$

Since the deflected mirage mediation provides a framework that involves all three prominent flavor and CP conserving mediation mechanisms, it is important to understand how does each mediation reveal its existence in low energy sparticle masses. From (3.35), one easily notices that anomaly mediation reveals itself through a nonzero value of α_{eff} , which can be read off from the gaugino mass pattern [27]. Once M_0^{eff} and α_{eff} could be determined from the gaugino masses, one may examine $m_{\tilde{q}_L}^2 - m_{\tilde{e}_R}^2$ and $m_{\tilde{d}_R}^2 - m_{\tilde{l}_L}^2$ to see the existence of gauge mediation, from which $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ can be determined.

It is obvious that δ_i , particularly $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$, are crucial for identifying the underlying mediation mechanism from the sparticle masses at TeV. Let us thus examine the possible values of δ_i in various models of deflected mirage mediation. The overall size of δ_i is determined by

$$R - 1 = \frac{N_\Phi g_0^2}{8\pi^2} \left[\frac{\alpha}{2\beta} \ln \left(\frac{M_{Pl}}{m_{3/2}} \right) - \ln \left(\frac{M_{\text{GUT}}}{M_\Phi} \right) \right], \quad (3.39)$$

where α/β represents the gauge to moduli mediation ratio. If $\alpha/\beta > 0$, there is a cancellation between gauge and moduli mediations, reducing the size of $R - 1$, and thus of δ_i . In particular, if the gauge messenger mass M_Φ is close to the scale

$$M_{\text{GUT}} \left(\frac{m_{3/2}}{M_{Pl}} \right)^{\alpha/2\beta} \simeq 10^{16.3-6.9\alpha/\beta} \text{ GeV}, \quad (3.40)$$

the cancellation is most efficient. For many of the representative SUGRA models discussed in the previous section, although the strength of gauge mediation is comparable to those of anomaly and moduli mediations, the resulting δ_i are small because of this cancellation. In such models, the predicted pattern of sparticle masses is quite similar to that of pure mirage mediation.

Let us first examine $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ in the effective SUGRA model (2.21). In this model with $\kappa \neq 0$, the gauge messenger mass is induced by the superpotential coupling $\lambda_\Phi X \Phi \Phi^c$ with X stabilized by $\kappa X^n / M_{Pl}^{n-3}$ ($n > 3$), which results in

$$\beta = \frac{n-1}{2}, \quad M_\Phi = x M_{Pl} \left(\frac{m_{3/2}}{M_{Pl}} \right)^{1/(n-2)}, \quad (3.41)$$

where $x = \lambda_\Phi \kappa^{1/(2-n)}$. As $\alpha \simeq 1$ at leading order in the g_{st} or α' expansion in underlying string theory, we first focus on the case with $\alpha = 1$. We then find

$$|\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}| < 0.02 N_\Phi, \quad |\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}| < 0.01 N_\Phi, \quad (3.42)$$

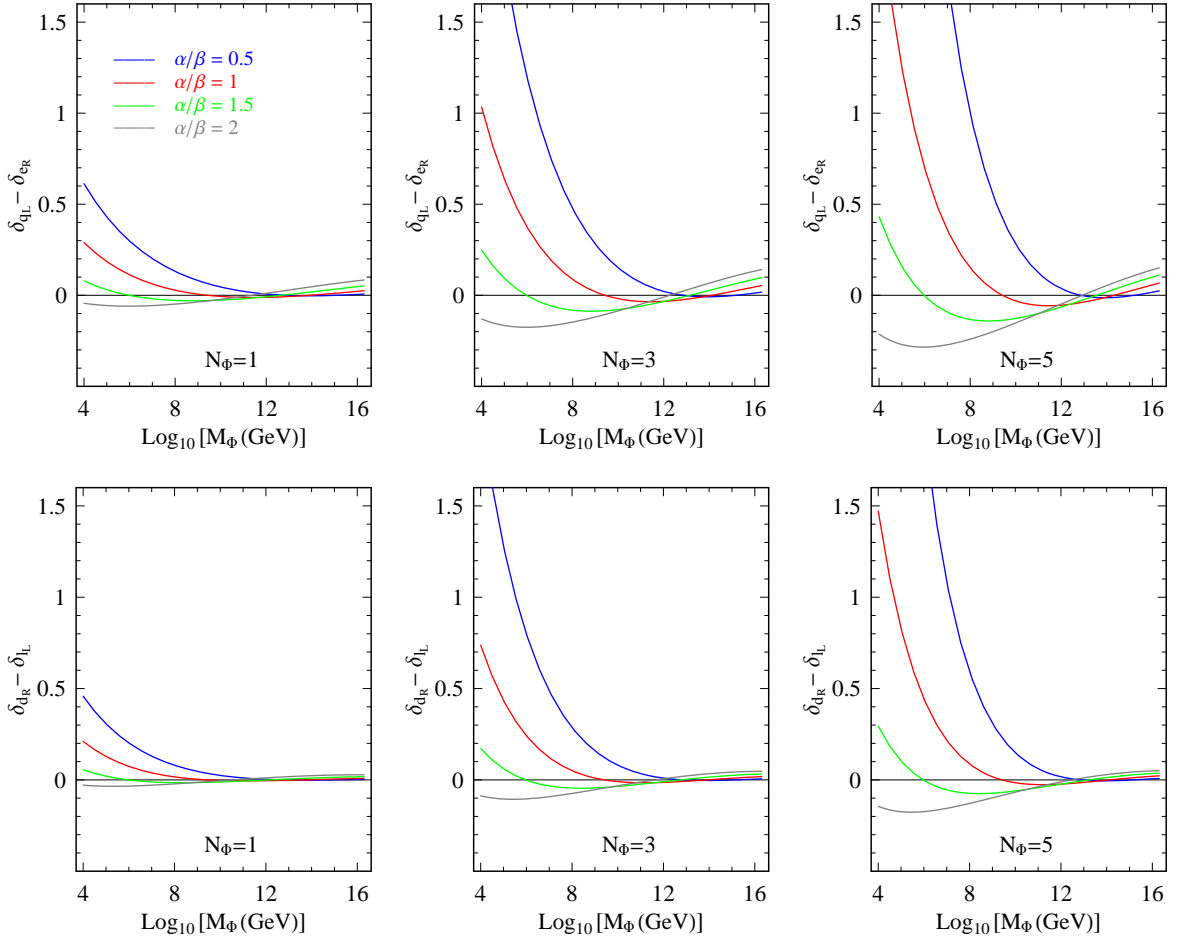


Figure 2. Difference of δ_i in deflected mirage mediation with a generic value of M_Φ . The upper panels show $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ while the lower ones show $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ for $N_\Phi = 1, 3, 5$.

for the parameter range: $\alpha = 1, 4 \leq n \leq 6$ and $10^{-(n-3)} \leq x$. For different value of α , they can have a bigger value, but still bounded as

$$|\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}| < 0.04N_\Phi, \quad |\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}| < 0.02N_\Phi, \quad (3.43)$$

for the parameter range: $0.5 \leq \alpha \leq 2, 4 \leq n \leq 6, 10^{-(n-4)} \leq x$, and $N_\Phi \leq 8$. A less stringent bound is obtained for the axionic mirage mediation model [13], in which $\kappa = 0$ and X is stabilized by the radiative correction to its Kähler potential, yielding $\beta = 1$. Provided that $\langle X \rangle$ is fixed at a scale between 10^9 GeV and 10^{12} GeV as required for $\text{Im}(X)$ to be the QCD axion, it is found that

$$|\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}| < 0.16N_\Phi, \quad |\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}| < 0.08N_\Phi, \quad (3.44)$$

for $0.5 \leq \alpha \leq 2$ and $N_\Phi \leq 8$. The above results for the models of (2.21) show that $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ are small over a reasonable range of model parameters, and therefore the

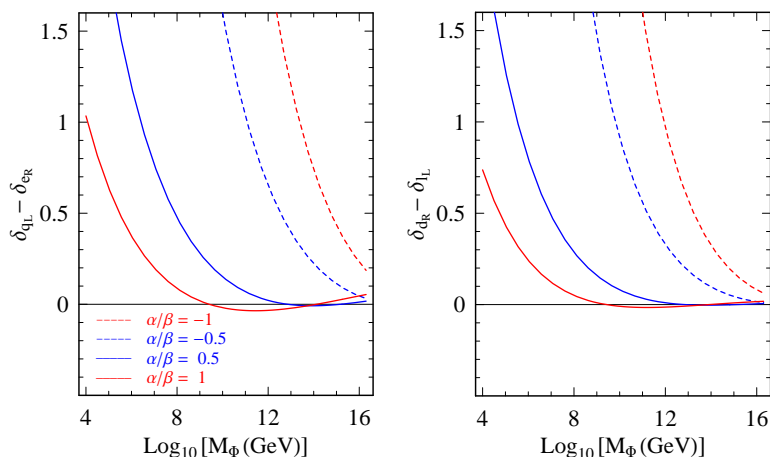


Figure 3. Difference of δ_i in deflected mirage mediation. The solid lines show the values for $0.5 \leq \alpha/\beta \leq 1$ while the dashed ones for $-1 \leq \alpha/\beta \leq -0.5$, with $N_\Phi = 3$.

predicted sparticle mass pattern is close to the pure mirage pattern obtained from (3.38). In figure 2, we depict the values of $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ in deflected mirage mediation scenario with a generic value of M_Φ , where α, β and N_Φ are assumed as $0.5 \leq \alpha/\beta \leq 2$ and $1 \leq N_\Phi \leq 5$. The models of (2.21) typically give $M_\Phi \geq 10^9$ GeV and $0.5 \leq \alpha/\beta \leq 2$, for which $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ have a small value as long as N_Φ is not unreasonably large.

There are in fact some models which can give a sizable value of $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$. One such example is a model with a negative value of α/β . For an illustration, we depict in figure 3 the values of $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ for $-1 \leq \alpha/\beta \leq -0.5$, and compare them with the values for $0.5 \leq \alpha/\beta \leq 1$.

Another scheme which can give a sizable value of $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ would be the deflected anomaly mediation [12, 21–24], in which there is no moduli mediation. Soft parameters in deflected anomaly mediation can be obtained by taking the limit: $1/R \rightarrow 0$ and $\tilde{A}_{ijk} = \tilde{m}_i^2 = 0$, while keeping RM_0 and α/R to have a nonzero finite value. The resulting sparticle masses at $\mu = 500$ GeV are given by (3.35) with

$$\begin{aligned}
 M_0^{\text{eff}} &= \frac{N_\Phi g_0^2 m_{3/2}}{16\pi^2 \beta} \simeq 3 \times 10^{-3} \frac{N_\Phi m_{3/2}}{\beta}, \\
 \alpha^{\text{eff}} &= \frac{16\pi^2}{g_0^2 \ln(M_{Pl}/m_{3/2})} \frac{\beta}{N_\Phi} \simeq \frac{10\beta}{N_\Phi}, \\
 (\tilde{m}_i^{\text{eff}})^2 &= (M_0^{\text{eff}})^2 \delta_i,
 \end{aligned} \tag{3.45}$$

where $\delta_i = \sum_a C_2^a(Q_i) \delta_a$ with

$$\delta_a = 2 \left[\frac{1}{N_\Phi} \frac{g_a^4(M_\Phi)}{g_0^4} - \frac{g_a^2(M_\Phi)}{8\pi^2} \left(1 + \frac{g_a^2(M_\Phi)}{g_0^2} \right) \ln \left(\frac{M_{\text{GUT}}}{M_\Phi} \right) \right]. \tag{3.46}$$

Thus, in deflected anomaly mediation limit, δ_i are determined by just M_Φ and N_Φ .

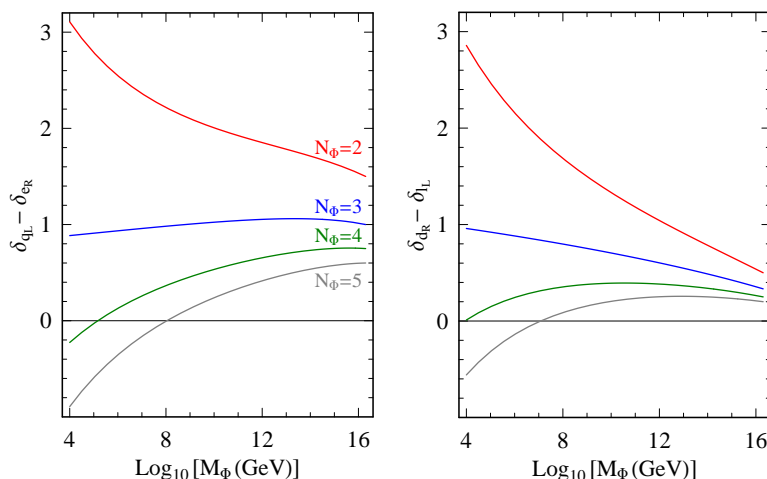


Figure 4. Difference of δ_i in deflected anomaly mediation for $2 \leq N_\Phi \leq 5$ with a generic value of M_Φ .

Let us examine the values of $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ in some specific models of deflected anomaly mediation. For the model (2.37), we have $M_\Phi = \mathcal{O}(m_{3/2})$, while β can take any value of order unity. Keeping the perturbative gauge coupling unification requires $N_\Phi \leq 5$, and tachyonic slepton can be avoided for

$$-0.25 - 0.35N_\Phi \lesssim \beta \lesssim 0.25 + 0.05N_\Phi. \quad (3.47)$$

We then find

$$\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R} \simeq -3.2 + \frac{12.4}{N_\Phi}, \quad \delta_{\tilde{d}_R} - \delta_{\tilde{l}_L} \simeq -2.5 + \frac{10.4}{N_\Phi}, \quad (3.48)$$

for $M_\Phi = \mathcal{O}(10)$ TeV. In figure 4, we consider more general situation with arbitrary value of M_Φ , and depict $\delta_{\tilde{q}_L} - \delta_{\tilde{e}_R}$ and $\delta_{\tilde{d}_R} - \delta_{\tilde{l}_L}$ for $2 \leq N_\Phi \leq 5$.

In fact, some models of deflected anomaly mediation are severely constrained by the condition to avoid tachyonic slepton, which typically requires a large value of N_Φ . An example would be the model (2.21) without the moduli T_I , which gives $\beta = (n-1)/2$ and $M_\Phi \sim M_{Pl}(m_{3/2}/M_{Pl})^{1/(n-2)}$. For the case of $n=4$, we need $N_\Phi \geq 10$ to avoid tachyonic slepton. On the other hand, the corresponding δ_i are given by

$$\begin{aligned} \delta_{\tilde{q}_L} &\simeq -0.74 + \frac{5.6}{N_\Phi}, & \delta_{\tilde{u}_R} &\simeq -0.59 + \frac{4.6}{N_\Phi}, & \delta_{\tilde{e}_R} &\simeq -0.10 + \frac{0.6}{N_\Phi}, \\ \delta_{\tilde{d}_R} &\simeq -0.56 + \frac{4.4}{N_\Phi}, & \delta_{\tilde{l}_L} &\simeq -0.22 + \frac{1.44}{N_\Phi}. \end{aligned} \quad (3.49)$$

For $N_\Phi = 10$, which is the minimal value avoiding tachyonic slepton, δ_i are all small, and then the model is difficult to be distinguished from the mirage mediation with $\alpha = 3/2$ and $\tilde{m}_5^2 = \tilde{m}_{10}^2 = 0$. Similar situation occurs for the case that X is stabilized by the radiative correction to its Kähler potential, e.g. with $N_\Phi = 10$, $M_\Phi \sim 10^{12}$ GeV and $\alpha = 1$.

4 Phenomenology of some examples

In the previous section, we have examined generic feature of mass spectrum of deflected mirage mediation, assuming the SU(5) unification of matter multiplets. The effective supergravity models only containing the mass scales $F^C/C \approx m_{3/2}$ and $M_{Pl/GUT}$, e.g. those of (2.21) and (2.35), predict the relation (3.40) with $\alpha \approx 1$. In such models, the low energy mass spectrum of mirage mediation ($\alpha \approx 1$) is robust against the gauge threshold corrections. On the other hand, once we arrange the special form of Kähler and super potentials as in the model of (2.37), or introduce (explicitly or dynamically) a new mass scale other than $M_{Pl/GUT}$ as in the model of (2.40), the mass spectrum can dramatically change as expected from figure 3, although a realization of such models is rather obscure in the string framework. In the following, we discuss two phenomenological applications of deflected mirage mediation representing these two cases.

4.1 Accidental little SUSY hierarchy

One of the virtues of the MSSM is the radiative electroweak symmetry breaking [31]. On the other hand, the radiative electroweak symmetry breaking depends on how the Higgs μ and B parameters are generated. In SUSY breaking scenarios with large gravitino mass $m_{3/2} \gg 1$ TeV, which includes the deflected mirage mediation, one can not simply implement the conventional supergravity mechanism [32] to generate μ and B as it gives a too large $B \sim m_{3/2}$. Still one can consider various ways to obtain a weak scale size of B within the framework of deflected mirage mediation, for instance the NMSSM extension of the Higgs sector [10] or the mechanism to use a radiatively stabilized flat direction which couples to $H_u H_d$ in a specific manner [13]. Here we assume that proper values of μ and B are generated by one of those mechanisms, and simply replace μ and B by M_Z and $\tan \beta$ after the electroweak symmetry breaking. We also note that although a large $\tan \beta$ generically requires a significant fine-tuning of parameters in pure dilaton mediation scenario [33], the degree of fine-tuning can be reduced in deflected mirage mediation depending on the relative importance of each of the involved gauge, anomaly, and dilaton/moduli mediations.

The Higgs mass parameter, $m_{H_u}^2$ is automatically driven to negative due to the renormalization group running by the top Yukawa coupling, even though it is given a positive value at some high energy scale, Λ . This radiative correction is controlled by the average stop mass, $m_{\tilde{t}}$,

$$\delta m_{H_u}^2 \sim -\frac{3}{4\pi^2} y_t^2 m_{\tilde{t}}^2 \ln \left(\frac{\Lambda}{m_{\tilde{t}}} \right), \quad (4.1)$$

therefore, barring fine-tuning of the initial condition, we anticipate $|m_{H_u}^2| \sim m_{\tilde{t}}^2$ for Λ hierarchically larger than $m_{\tilde{t}}$. On the other hand, the lightest Higgs boson mass in the MSSM is approximated by

$$m_{h^0}^2 \approx M_Z^2 \cos^2 2\beta + \frac{3y_t^2}{4\pi^2} m_{\tilde{t}}^2 \ln \left(\frac{m_{\tilde{t}}}{m_t} \right), \quad (4.2)$$

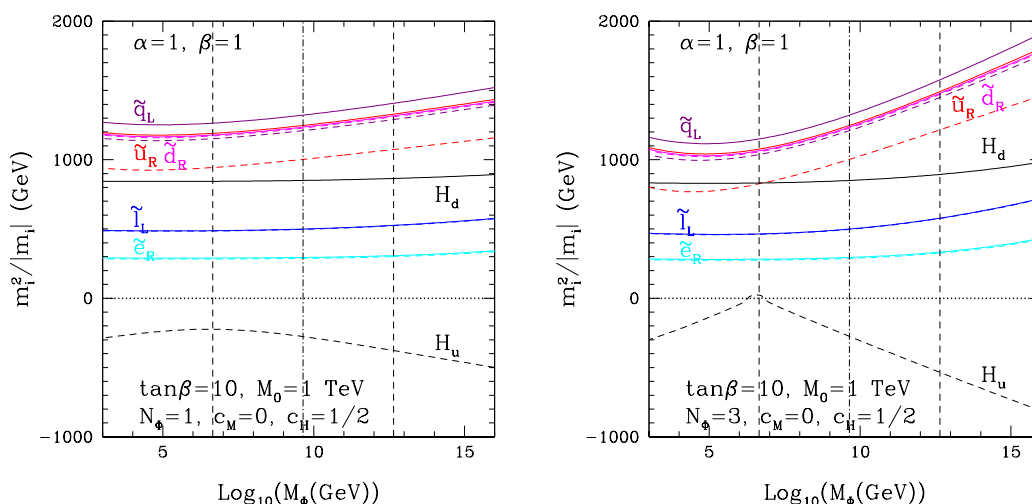


Figure 5. Accidental little SUSY hierarchy in deflected mirage mediation ($\alpha = 1$) for the case that X is stabilized by the radiative effects in Kähler potential. The left panel shows the case for $N_\Phi = 1$ while the right panel for $N_\Phi = 3$. In all of them, the modular weights are chosen as $c_M \equiv \tilde{m}_{\tilde{q}_L, \tilde{u}_R, \tilde{d}_R, \tilde{l}_L, \tilde{e}_R}^2/M_0^2 = 0$ and $c_H \equiv \tilde{m}_{H_u, H_d}^2/M_0^2 = 1/2$. Other SUSY parameters are set to $\tan\beta = 10$ and $M_0 = 1$ TeV. The vertical dashed lines indicate the predicted range of M_Φ . The vertical dot-dashed line represents the gauge threshold scale leading to $R = 1$.

where $\tan\beta = \langle H_u \rangle / \langle H_d \rangle$. (Note that in the previous section β has been used to parameterize the anomaly to gauge mediation ratio.) In order to fulfill the lower bound of the SM Higgs boson mass obtained at LEP II, $m_{h_0} > 114$ GeV, we need the stop mass as heavy as $m_{\tilde{t}} \gtrsim 600$ GeV. Thus the Higgs mass parameter is generally expected to be $|m_{H_u}| \gtrsim 600$ GeV. While one of the conditions of the electroweak symmetry breaking tells us

$$\frac{M_Z^2}{2} \approx -m_{H_u}^2 - |\mu|^2, \quad (4.3)$$

for $\tan\beta$ not too close to 1. Here μ is the higgsino mass parameter which does not break SUSY. (Note that in the previous section μ has been used to parameterize the renormalization point of running soft parameters.) This means that we are forced to fine-tune the parameters, $m_{H_u}^2$ and $|\mu|^2$ at less than 1% level to obtain the observed size of M_Z , despite these two parameters are expected to be not correlated. This not fatal but uncomfortable fine-tuning in the MSSM arising from the hierarchy between the electroweak scale and the SUSY breaking mass scale is called ‘the little SUSY hierarchy problem’ [34].

One obvious solution is having $m_{H_u}^2 \sim M_Z^2 \ll m_{\tilde{t}}^2$ by accident due to a choice of the boundary condition at Λ . However, in mirage mediation, it is not apparent whether one can achieve such a pattern or not, because the choice of the modular weights is discrete. The moderate modification of the spectrum by the deflection may help to obtain the desired mass pattern.⁸ In figure 5 and 6, we show such an accidental little SUSY hierarchy

⁸For a different approach to this problem in mirage mediation, see [35]. It has also been argued that a negative stop mass-square at high renormalization point can reduce the fine-tuning [36].

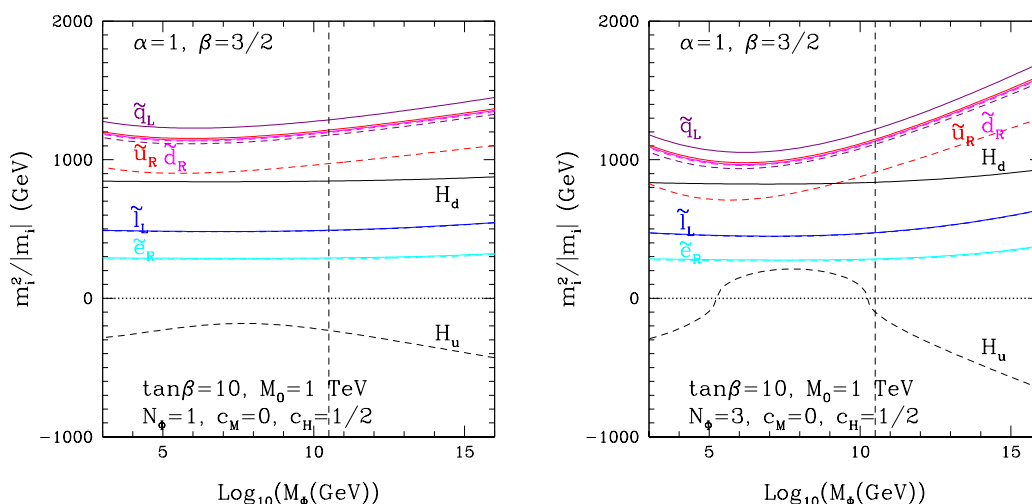


Figure 6. Accidental little SUSY hierarchy in deflected mirage mediation ($\alpha = 1$) for the case with X stabilized by the higher dimensional operator in superpotential. SUSY parameters and modular weights are same as in figure 5. The vertical dashed line indicates the predicted value of M_Φ .

achieved by the deflected mirage mediation at $\alpha = 1$, where we chose $c_{\tilde{q}_L, \tilde{u}_R, \tilde{d}_R, \tilde{l}_L, \tilde{e}_R} \equiv \tilde{m}_{\tilde{q}_L, \tilde{u}_R, \tilde{d}_R, \tilde{l}_L, \tilde{e}_R}^2 / M_0^2 = 0$, $c_{H_u, H_d} \equiv \tilde{m}_{H_u, H_d}^2 / M_0^2 = 1/2$ and $M_0 = 1$ TeV. In figure 5, we present the case that X is stabilized by the radiative effects in Kähler potential, while in figure 6, the case for stabilization by the higher dimensional operator in superpotential. The left panels adopt $N_\Phi = 1$ and the right ones $N_\Phi = 3$. In all of them, the dashed curves denote the 3rd generations. In figure 5, the dot-dashed line indicates the gauge threshold scale at which the cancellation between gauge and moduli mediations leads to $R = 1$. For $R = 1$, neglecting the effects of Yukawa couplings, the sfermion masses simply reduce to the values in pure mirage mediation. The deviation associated with Yukawa couplings is also vanishing for $R = 1$ if the mirage condition (3.16) is satisfied for the moduli-mediated soft terms.

For $N_\Phi = 1$, the effect of deflection is limited. However, for the case that X is stabilized by the radiative effects, we can obtain an improved hierarchy for $N_\Phi = 3$ and $M_\Phi \sim 10^{6-7}$ GeV, which is within the plausible range of this stabilization mechanism indicated by the vertical dashed lines [13]. In this case we need to break the PQ symmetry slightly to make the axion heavy so that it will not be produced in astrophysical processes and evade the bound, 10^9 GeV $\lesssim M_\Phi$ coming from the burst duration of supernova SN1987A and the cooling of globular-cluster stars and white dwarf [37]. In figure 6, which is for the case with X stabilized by the non-renormalizable superpotential term, the vertical dashed line shows the gauge threshold scale predicted for $n = 4$ with $x = 1$ in (3.41). Again, we have an improved hierarchy for $N_\Phi = 3$. In both cases, we need a mechanism to generate μ and $B\mu$ terms of appropriate size, which does not disturb the mass spectrum. For instance, in the first case we can employ the mechanism described in [12, 13] in weak coupling limit, which employs a term $(X^\dagger/X) H_d H_u + \text{h.c.}$ in the Kähler potential to generate the desired values of μ and $B\mu$. In the second case, we can use the same mechanism by

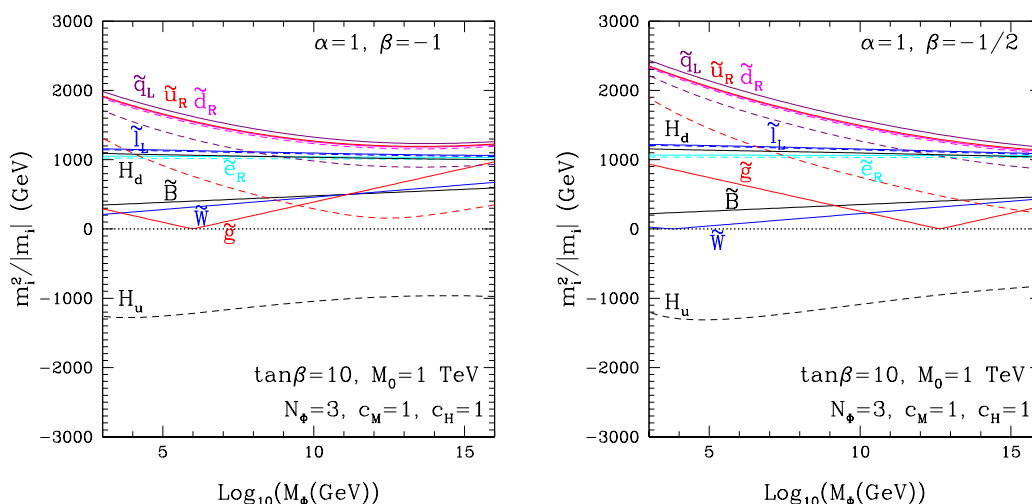


Figure 7. Gluino and wino LSP in deflected mirage mediation ($\alpha = 1$). The dashed curves indicate the 3rd generation. The left panel shows the case for $\beta = -1$ and the right for $\beta = -1/2$. All the modular weights are chosen so that $c_i \equiv \tilde{m}_i^2/M_0^2 = 1$. Other parameters are chosen as $\tan\beta = 10$, $M_0 = 1$ TeV and $N_\Phi = 3$.

introducing another singlet X' stabilized at M_{mir} by the Kähler potential, which minimizes the deflection (see (3.40)).

4.2 Gluino lightest supersymmetric particle in deflected mirage mediation

As we have seen in the previous section, we can considerably reduce R (or increase α_{eff}) from 1 if we chose the special form of Kähler and superpotential (2.37) [22] or use the negative power superpotential for X (2.40) [14, 23]. In such a case, we have possibilities that the lightest supersymmetric particle (LSP) becomes gluino ($\alpha_{\text{eff}} \approx 3$) or wino ($\alpha_{\text{eff}} \gg 1$) as shown in figure 1 of [10]. On the other hand, in contrast to mirage mediation, $(\tilde{m}_i^{\text{eff}})^2$ does not vanish in the limit $M_0^{\text{eff}}(R) \rightarrow 0$ as seen in (3.27). This considerably dilutes the renormalization scale dependent part in (3.26) relative to the constant term, which leads to the ‘quasi-infrared-fixed-point’ behavior of the scalar mass as first observed in [14]. This effect makes squarks and sleptons somewhat heavier than the gauginos for small R (M_Φ) as shown in figure 7. Therefore a direct production of the squarks or sleptons at hadron collider is suppressed. Their on-shell states also can not appear in the cascade decay of gluino in the wino LSP case. As the physics of the wino LSP has been extensively examined in association with the anomaly mediation [38, 39], here we will focus on the more exotic case: the gluino LSP.

Phenomenology of the gluino LSP or meta-stable gluino has been investigated from various motivations [40–44], particularly in the context of split supersymmetry recently [45–50]. Most of the analyses can be directly applied to our case. We summarize some of the results below.

Since squarks are somewhat heavier than gluino in the present scenario, gluino is

mainly produced via the processes, $q\bar{q} \rightarrow g^* \rightarrow \tilde{g}\tilde{g}$ or $gg \rightarrow g^* \rightarrow \tilde{g}\tilde{g}$ at the hadron collider such as Tevatron and LHC. While at the lepton collider it is always produced in association with quarks, $e\bar{e} \rightarrow q\bar{q} \rightarrow q\bar{q}\tilde{g}\tilde{g}$ since the leptons are color singlets. The produced gluino hadronizes into a color-singlet composite state called R-hadron [51]. The bound states of gluino and color-octet hadron, $\tilde{g}g$, $\tilde{g}q\bar{q}$ and $\tilde{g}qqq$ are known as R-gluon, R-meson and R-baryon, respectively. Phenomenologically, the most relevant question is the stability of the charged particles which leave tracks in the detector. It depends on the identity of the lightest R-hadron and their mass differences. The mass spectrum of R-hadron is estimated by the MIT bag model [52] and the quenched lattice simulation [53] which predict isotriplet vector R-meson and R-gluon ($J^{PC} = 1^{+-}$) as the lightest R-hadron, respectively. In both cases their mass difference is smaller than the pion mass. Thus the vector R-meson is stable against hadronic decay. The mass difference of the lightest R-baryon and the ground state of R-gluon or R-meson is also estimated to be smaller than the nucleon mass and the R-baryon is stable [42, 52, 53]. Therefore the probability for gluino to hadronize into the charged states, P is expected to be non-negligible, although the reliability of such a conclusion is limited by those of the calculation method and adopted assumptions.

Inside the detector the R-hadron deposits energy via hadronic interaction with nucleus and, if it is charged, ionization of the detector material. The gluino is typically produced with momentum similar order of its mass ($m_R \sim 100$ GeV) and thus the R-hadron is relativistic but slow, $\beta < 1$. The energy involved in the hadronic interaction is approximated by $Q = \sqrt{s} - m_R - m_N \approx (\gamma - 1)m_N$ with $\gamma = 1/\sqrt{1 - \beta^2}$ [42]. Therefore the interaction is soft nevertheless the energy carried by the R-hadron is huge (~ 100 GeV). Then the neutral R-hadron traverses the detector repeatedly kicking off the nucleon inside the nucleus and the soft secondary particles dissipating a small fraction of its kinetic energy at each collision. Eventually it penetrates the detector carrying away significant amount of missing energy. This behavior is in contrast to the ordinary hadron which develops shower and exponentially dumps its energy in the calorimeter. The ionization energy loss of the charged R-hadron is calculated by the standard Bethe-Bloch formula [37]. Since the energy loss, $-dE/dx$ is proportional to the inverse of β^2 while the average hadronic energy loss per collision behaves like $\langle \Delta E \rangle \propto \gamma$, the ionization plays a minor role for large β ($\gtrsim 0.9$), however, it quickly dominates over the hadronic interaction as the R-hadron is slowed down [47]. The heavily ionizing track provides a characteristic signal of the slowly moving charged R-hadron. It is noted that even if the R-hadron is neutral it can change its identity in the hadronic interaction. If the neutral and charged states are both stable, it transforms from neutral to charged and vice versa along with its path (flipper scenario).

Based on the careful inference about the nature of R-hadron as explained above, the pioneering work by [40] excluded the gluino mass range, $3 \text{ GeV} \leq m_{\tilde{g}} \leq 130 \text{ GeV}$ at 95% CL almost independent of P using jet+ \cancel{E} channel in LEP and Tevatron CDF RUN I, while $50 \text{ GeV} \leq m_{\tilde{g}} \leq 200 \text{ GeV}$ for $P \geq 1/2$ by a heavily-ionizing track search in Tevatron CDF. On the other hand, the analysis in [47] emphasized a model independent role of the high p_T monojet which is produced in association with the gluino pair. They obtained a conservative bound, $m_{\tilde{g}} \geq 170 \text{ GeV}$, independent of P using the Tevatron Run I data and projected it to $m_{\tilde{g}} \geq 210 \text{ GeV}$ for Run II and $m_{\tilde{g}} \geq 1.1 \text{ TeV}$ for LHC. They also estimated a

reach of the charged track search ($P = 1$) as $m_{\tilde{g}} = 270$ GeV for CDF Run I with integrated luminosity 100 ps^{-1} , $m_{\tilde{g}} = 430$ GeV for Run II with 2 fb^{-1} and $m_{\tilde{g}} = 2.4$ TeV for LHC with 100 fb^{-1} . More serious study on the R-hadron discovery potential in LHC ATLAS detector has been performed using the ATLAS fast simulation framework [42]. They have taken into account the flipper behavior of the R-hadron and perform Monte Carlo simulation based on GEANT3. Using event selection with global event variables such as the missing transverse energy \cancel{E}_T , the total sum of the transverse energy E_T^{tot} and the transverse momentum (p_T) measured in the muon chamber, they concluded that the R-hadron is discovered for the mass up to 1.4 TeV with the integrated luminosity 30 fb^{-1} . While the reach extends to 1.7 TeV if the time-of-flight information for slow-moving R-hadron between the muon chambers is used in the event selection in stead of \cancel{E}_T and E_T^{tot} . From these studies we can conclude that still the gluino LSP scenario of our interest has a considerable portion of parameter space, however, it will be fully examined in the relatively early stage of LHC.

The stable R-hadron in cosmological time scale having electric or hadronic interaction will conflict with various phenomenological constraints such as heavy isotope searches [56–58] unless the relic abundance is sufficiently small [40, 59, 60]. An alternative interesting possibility is that the gluino is unstable but decays outside of the detector as in the split supersymmetry, which circumvents the cosmological difficulty but will not change the collider signature we have discussed above. A concrete example is given by the axionic extension of deflected mirage mediation as discussed in [13]. Because of the recursive nature of the mass spectrum of deflected mirage mediation with multi-thresholds, we can introduce the axion superfield S stabilized by the Kähler potential in addition to X without disturbing the low energy mass spectrum as long as $\langle X \rangle < \langle S \rangle \approx M_{\text{mir}}$. The LSP is now axino, \tilde{a} whose mass is one-loop suppressed relative to the gauginos since there is no tree-level contribution from the superpotential. The next lightest supersymmetric particle (NLSP) is gluino, which decays to the axino via $\langle S \rangle$ -dependence in the gauge coupling constant introduced by the threshold correction. The corresponding interaction Lagrangian is given by [63]

$$\mathcal{L}_{\tilde{g}\tilde{a}g} = \frac{\alpha_s N_\Phi}{16\sqrt{2}\pi} \frac{1}{\langle S \rangle} \tilde{a} \gamma_5 \sigma^{\mu\nu} \tilde{g} g_{\mu\nu} + \text{h.c.}, \quad (4.4)$$

which yields the decay width

$$\Gamma(\tilde{g} \rightarrow \tilde{a}g) \simeq \frac{\alpha_s^2 N_\Phi^2}{32\pi^3} \frac{m_{\tilde{g}}^3}{\langle S \rangle^2}, \quad (4.5)$$

and the life time

$$\tau_{\tilde{g}} = \Gamma(\tilde{g} \rightarrow \tilde{a}g)^{-1} \simeq 5.7 \times 10^{-7} \text{ sec} \left(\frac{m_{\tilde{g}}}{200 \text{ GeV}} \right)^{-3} \left(\frac{\langle S \rangle}{10^{10} \text{ GeV}} \right)^2 N_\Phi^{-2}, \quad (4.6)$$

where N_Φ is the number of messengers coupling with the axion. For $m_{\tilde{g}} \lesssim 1$ TeV and $N_\Phi \simeq 1$, most of the decay occurs outside of the detector and the discovery prospects discussed above is applicable. On the other hand, for $m_{\tilde{g}} \gtrsim \text{TeV}$ or $N_\Phi \gg 1$, gluinos decay inside the detector, which results in a displaced jet vertex with missing energy. It would provide a clear signature and a similar experimental reach as in the case of the

heavily ionizing track, although a detailed study in a realistic circumstance is mandatory for deriving any definitive conclusion.

The relic axino abundance in deflected mirage mediation with the gluino NLSP can be estimated by following the discussion in [13]. As is well known, if we have a light modulus T as in (deflected) mirage mediation, its coherent oscillation dominates the energy density of the universe after the inflation [61]. Eventually it decays and produced entropy dilutes everything existed before. The saxion also oscillates, but it is harmless as long as $\langle S \rangle \lesssim 10^{11}$ GeV since it decays faster than the modulus.⁹ The reheating temperature of the modulus is estimated as [13]

$$T_R = \left(\frac{90}{\pi^2 g_*(T_R)} \right)^{1/4} \sqrt{M_{Pl} \Gamma_T} \simeq 0.15 \text{ GeV} \left(\frac{g_*(T_R)}{10} \right)^{-1/4} \left(\frac{m_T}{10^6 \text{ GeV}} \right)^{3/2} d_g, \quad (4.7)$$

where $g_*(T_R)$ denotes the effective bosonic degrees of freedom at T_R for the energy density and d_g is a model dependent parameter of order unity defined as

$$d_g \equiv 2(-3\partial_T \partial_{T^*} \ln \Omega_{\text{mod}})^{-1/2} \partial_T \ln(\text{Re}(\tilde{f}_a)). \quad (4.8)$$

It is marginally smaller than the thermal decoupling temperature of gluino [40]

$$T_F = m_{\tilde{g}}/x_F \approx 6 \text{ GeV} \left(\frac{m_{\tilde{g}}}{200 \text{ GeV}} \right). \quad (4.9)$$

Therefore the thermal relic abundance of gluino is not suitable to calculate the axino abundance. Actually, the dominant contribution to the axino relic abundance comes from the gravitino which is produced by the decay of the modulus [62]. The gravitino decay eventually produces at least one gluino (or axino) and hence one axino. The gravitino yield which is conserved in the adiabatic expansion of the universe is given by [62]

$$Y_{3/2} \equiv \frac{n_{3/2}}{s} \simeq \frac{3}{2} \frac{T_R}{m_T} B_{3/2}^T \sim 4.3 \times 10^{-9} \left(\frac{g_*(T_R)}{10} \right)^{-1/4} \left(\frac{m_T}{10^6 \text{ GeV}} \right)^{1/2} d_g. \quad (4.10)$$

The gluino decay to axino potentially competes with the gluino annihilation. The perturbative annihilation cross section in the zero relative velocity limit ($\beta = 0$) is given by $\sigma_{\text{ann}}\beta \simeq (171\pi\alpha_s^2/64m_{\tilde{g}}^2)$ [40]. Then the inverse of annihilation rate is estimated as

$$\begin{aligned} \Gamma_{\text{ann}}^{-1} &\simeq (n_{3/2} \sigma_{\text{ann}} \beta)^{-1} \simeq (Y_{3/2} s(T_{3/2}) \sigma_{\text{ann}} \beta)^{-1} \\ &\simeq 1.6 \times 10^{-5} \text{ sec} \left(\frac{m_T}{10^6 \text{ GeV}} \right)^{-1/2} \left(\frac{m_{3/2}}{10^5 \text{ GeV}} \right)^{-9/2} \left(\frac{m_{\tilde{g}}}{200 \text{ GeV}} \right)^2 d_g^{-1}, \end{aligned} \quad (4.11)$$

which is considerably larger than (4.6). Thus the annihilation is negligible and (4.10) gives a good estimation of the axino yield. Then the current axino density is given by

$$\Omega_{\tilde{a}} h^2 = \frac{Y_{3/2} s_0 h^2}{3M_P^2 H_0^2} \simeq 1.2 \left(\frac{m_{\tilde{a}}}{1 \text{ GeV}} \right) \left(\frac{g_*(T_R)}{10} \right)^{-1/4} \left(\frac{m_T}{10^6 \text{ GeV}} \right)^{1/2} d_g, \quad (4.12)$$

⁹If the saxion oscillation dominates the universe, axions produced from its decay upset the successful prediction of Big Bang nucleosynthesis. Therefore axion can not be a dominant component of the dark matter in this scenario [13].

where s_0 and H_0 are the entropy density and Hubble constant of the current universe, respectively. Thus the axino, which has one-loop suppressed mass relative to the gaugino (~ 100 MeV), accompanied with the gluino NLSP can naturally saturate the observed dark matter density as discussed in [13] for other NLSPs.

5 Conclusion

In this paper, we have discussed the sparticle mass pattern in deflected mirage mediation scenario of supersymmetry breaking, in which all of the three known flavor conserving mediations, i.e. dilaton/moduli, anomaly and gauge mediations, contribute to the MSSM soft parameters. Starting with a class of string-motivated effective supergravity models that realize deflected mirage mediation, we analyzed the renormalization group running of soft parameters to derive the (approximate) analytic expression of low energy sparticle masses at the TeV scale. We also discussed more detailed phenomenology of two specific examples, one with an accidental little hierarchy between $m_{H_u}^2$ and other soft mass-squares and another with gluino NLSP, that can be obtained within deflected mirage mediation scenario.

If some sparticles masses are in the sub-TeV range, the corresponding sparticles can be copiously produced at the CERN LHC, and then one may be able to measure their masses with the methods proposed in [64]. Our results then can be used to interpret the experimentally measured sparticle masses within the framework of the most general flavor and CP conserving mediation scheme.

Acknowledgments

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A Sfermion soft parameters

In this appendix, we present the analytic expression of sfermion soft parameters including the effects of the top quark Yukawa coupling y_t , when the moduli-mediated soft parameters of $Q_i = H_u, q_3, u_3$ at M_{GUT} do *not* satisfy the mirage condition (3.16). As the expression for generic deflected mirage mediation is too much involved, so not useful, here we present only the result in the mirage mediation limit without gauge mediation contribution.

In the low $\tan\beta$ regime of the MSSM, neglecting the Yukawa couplings other than y_t , the anomalous dimension of Q_i reads

$$\gamma_i(\mu) = 2 \sum_a C_i^a(Q_i) g_a^2(\mu) + k_i y_t^2(\mu), \tag{A.1}$$

where k_i is non-vanishing only for $Q_i = H_u, q_3, u_3$ having the top Yukawa interaction

$$k_{H_u} = -3, \quad k_{q_3} = -1, \quad k_{u_3} = -2. \quad (\text{A.2})$$

Above the gauge threshold scale M_Φ , the running top Yukawa coupling is given by

$$y_t^2(\mu) = \frac{y_t^2(M_{\text{GUT}})\dot{G}(\mu)}{1 - \frac{3}{4\pi^2}y_t^2(M_{\text{GUT}})G(\mu)}, \quad (\text{A.3})$$

where

$$G(\mu) = \int_{M_{\text{GUT}}}^{\mu} \frac{d\mu'}{\mu'} \prod_a \left(1 + \frac{b_a^H}{8\pi^2} g_{\text{GUT}}^2 \ln \left(\frac{M_{\text{GUT}}}{\mu'} \right) \right)^{2C_t^a/b_a^H}, \quad (\text{A.4})$$

with $C_t^a = C_2^a(H_u) + C_2^a(q_3) + C_2^a(u_3)$. Using this, we can find the analytic expressions for the sfermion soft parameters even when the mirage condition (3.16) is not satisfied:

$$\begin{aligned} A_{ijk}(\mu) &= \tilde{A}_{ijk} - (k_i + k_j + k_k)(\tilde{A}_{H_u q_3 u_3} - M_0)\rho(\mu) \\ &\quad - \frac{M_0}{8\pi^2} \left(\gamma_i(\mu) + \gamma_j(\mu) + \gamma_k(\mu) \right) \ln \left(\frac{\mu}{M_{\text{mir}}} \right), \\ m_i^2(\mu) &= \tilde{m}_i^2 - k_i \left[(\tilde{A}_{H_u q_3 u_3} - M_0)^2 (1 + 6\rho(\mu)) + (\tilde{m}_{H_u}^2 + \tilde{m}_{\tilde{t}_L}^2 + \tilde{m}_{\tilde{t}_R}^2 - M_0^2) \right] \rho(\mu) \\ &\quad - \frac{M_0}{4\pi^2} \left[M_0 \gamma_i(\mu) + k_i (\tilde{A}_{H_u q_3 u_3} - M_0) (1 + 6\rho(\mu)) y_t^2(\mu) \right] \ln \left(\frac{\mu}{M_{\text{mir}}} \right) \\ &\quad - \frac{M_0^2}{8\pi^2} \dot{\gamma}_i(\mu) \left[\ln \left(\frac{\mu}{M_{\text{mir}}} \right) \right]^2, \end{aligned} \quad (\text{A.5})$$

where

$$\rho(\mu) = \frac{y_t^2(\mu) G(\mu)}{8\pi^2 \dot{G}(\mu)}.$$

It is obvious that the above solutions reduce to (3.17) when the mirage condition (3.1) is satisfied.

References

- [1] H.P. Nilles, *Supersymmetry, supergravity and particle physics*, *Phys. Rept.* **110** (1984) 1 [SPIRES];
H.E. Haber and G.L. Kane, *The search for supersymmetry: probing physics beyond the standard model*, *Phys. Rept.* **117** (1985) 75 [SPIRES].
- [2] M. Dine, W. Fischler and M. Srednicki, *Supersymmetric technicolor*, *Nucl. Phys.* **B 189** (1981) 575 [SPIRES];
S. Dimopoulos and S. Raby, *Supercolor*, *Nucl. Phys.* **B 192** (1981) 353 [SPIRES];
M. Dine and W. Fischler, *A phenomenological model of particle physics based on supersymmetry*, *Phys. Lett.* **B 110** (1982) 227 [SPIRES];
C.R. Nappi and B.A. Ovrut, *Supersymmetric extension of the SU(3) × SU(2) × U(1) model*, *Phys. Lett.* **B 113** (1982) 175 [SPIRES];
L. Álvarez-Gaumé, M. Claudson and M.B. Wise, *Low-energy supersymmetry*, *Nucl. Phys.* **B 207** (1982) 96 [SPIRES];
S. Dimopoulos and S. Raby, *Geometric hierarchy*, *Nucl. Phys.* **B 219** (1983) 479 [SPIRES].

- [3] M. Dine, A.E. Nelson and Y. Shirman, *Low-energy dynamical supersymmetry breaking simplified*, *Phys. Rev. D* **51** (1995) 1362 [[hep-ph/9408384](#)] [[SPIRES](#)];
M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, *New tools for low-energy dynamical supersymmetry breaking*, *Phys. Rev. D* **53** (1996) 2658 [[hep-ph/9507378](#)] [[SPIRES](#)];
G.F. Giudice and R. Rattazzi, *Theories with gauge-mediated supersymmetry breaking*, *Phys. Rept.* **322** (1999) 419 [[hep-ph/9801271](#)] [[SPIRES](#)].
- [4] L. Randall and R. Sundrum, *Out of this world supersymmetry breaking*, *Nucl. Phys. B* **557** (1999) 79 [[hep-th/9810155](#)] [[SPIRES](#)];
G.F. Giudice, M.A. Luty, H. Murayama and R. Rattazzi, *Gaugino mass without singlets*, *JHEP* **12** (1998) 027 [[hep-ph/9810442](#)] [[SPIRES](#)];
J.A. Bagger, T. Moroi and E. Poppitz, *Anomaly mediation in supergravity theories*, *JHEP* **04** (2000) 009 [[hep-th/9911029](#)] [[SPIRES](#)].
- [5] V.S. Kaplunovsky and J. Louis, *Model independent analysis of soft terms in effective supergravity and in string theory*, *Phys. Lett. B* **306** (1993) 269 [[hep-th/9303040](#)] [[SPIRES](#)];
A. Brignole, L.E. Ibáñez and C. Muñoz, *Towards a theory of soft terms for the supersymmetric Standard Model*, *Nucl. Phys. B* **422** (1994) 125 [*Erratum ibid.* **B 436** (1995) 747] [[hep-ph/9308271](#)] [[SPIRES](#)].
- [6] A.H. Chamseddine, R.L. Arnowitt and P. Nath, *Locally supersymmetric grand unification*, *Phys. Rev. Lett.* **49** (1982) 970 [[SPIRES](#)];
R. Barbieri, S. Ferrara and C.A. Savoy, *Gauge models with spontaneously broken local supersymmetry*, *Phys. Lett. B* **119** (1982) 343 [[SPIRES](#)];
L.J. Hall, J.D. Lykken and S. Weinberg, *Supergravity as the messenger of supersymmetry breaking*, *Phys. Rev. D* **27** (1983) 2359 [[SPIRES](#)].
- [7] K. Choi, *Small SUSY phases in string inspired supergravity*, *Phys. Rev. Lett.* **72** (1994) 1592 [[hep-ph/9311352](#)] [[SPIRES](#)].
- [8] K. Choi, A. Falkowski, H.P. Nilles, M. Olechowski and S. Pokorski, *Stability of flux compactifications and the pattern of supersymmetry breaking*, *JHEP* **11** (2004) 076 [[hep-th/0411066](#)] [[SPIRES](#)];
K. Choi, A. Falkowski, H.P. Nilles and M. Olechowski, *Soft supersymmetry breaking in KKLT flux compactification*, *Nucl. Phys. B* **718** (2005) 113 [[hep-th/0503216](#)] [[SPIRES](#)].
- [9] M. Endo, M. Yamaguchi and K. Yoshioka, *A bottom-up approach to moduli dynamics in heavy gravitino scenario: superpotential, soft terms and sparticle mass spectrum*, *Phys. Rev. D* **72** (2005) 015004 [[hep-ph/0504036](#)] [[SPIRES](#)].
- [10] K. Choi, K.S. Jeong and K.-I. Okumura, *Phenomenology of mixed modulus-anomaly mediation in fluxed string compactifications and brane models*, *JHEP* **09** (2005) 039 [[hep-ph/0504037](#)] [[SPIRES](#)].
- [11] S. Kachru, R. Kallosh, A. Linde and S.P. Trivedi, *De Sitter vacua in string theory*, *Phys. Rev. D* **68** (2003) 046005 [[hep-th/0301240](#)] [[SPIRES](#)].
- [12] A. Pomarol and R. Rattazzi, *Sparticle masses from the superconformal anomaly*, *JHEP* **05** (1999) 013 [[hep-ph/9903448](#)] [[SPIRES](#)];
R. Rattazzi, A. Strumia and J.D. Wells, *Phenomenology of deflected anomaly mediation*, *Nucl. Phys. B* **576** (2000) 3 [[hep-ph/9912390](#)] [[SPIRES](#)].
- [13] S. Nakamura, K.-I. Okumura and M. Yamaguchi, *Axionic mirage mediation*, *Phys. Rev. D* **77** (2008) 115027 [[arXiv:0803.3725](#)] [[SPIRES](#)].

- [14] L.L. Everett, I.-W. Kim, P. Ouyang and K.M. Zurek, *Deflected mirage mediation: a framework for generalized supersymmetry breaking*, *Phys. Rev. Lett.* **101** (2008) 101803 [[arXiv:0804.0592](#)] [[SPIRES](#)]; *Moduli stabilization and supersymmetry breaking in deflected mirage mediation*, *JHEP* **08** (2008) 102 [[arXiv:0806.2330](#)] [[SPIRES](#)].
- [15] S. Kachru, L. McAllister and R. Sundrum, *Sequestering in string theory*, *JHEP* **10** (2007) 013 [[hep-th/0703105](#)] [[SPIRES](#)];
M. Son and R. Sundrum, *Anomaly-mediation and sequestering from a higher-dimensional viewpoint*, *JHEP* **08** (2008) 004 [[arXiv:0801.4789](#)] [[SPIRES](#)].
- [16] K. Choi and K.S. Jeong, *Supersymmetry breaking and moduli stabilization with anomalous U(1) gauge symmetry*, *JHEP* **08** (2006) 007 [[hep-th/0605108](#)] [[SPIRES](#)];
F. Brummer, A. Hebecker and M. Trapletti, *SUSY breaking mediation by throat fields*, *Nucl. Phys. B* **755** (2006) 186 [[hep-th/0605232](#)] [[SPIRES](#)].
- [17] S. Samuel and J. Wess, *A superfield formulation of the nonlinear realization of supersymmetry and its coupling to supergravity*, *Nucl. Phys. B* **221** (1983) 153 [[SPIRES](#)].
- [18] J.P. Conlon, D. Cremades and F. Quevedo, *Kähler potentials of chiral matter fields for Calabi-Yau string compactifications*, *JHEP* **01** (2007) 022 [[hep-th/0609180](#)] [[SPIRES](#)].
- [19] K. Choi and K.S. Jeong, *String theoretic QCD axion with stabilized saxion and the pattern of supersymmetry breaking*, *JHEP* **01** (2007) 103 [[hep-th/0611279](#)] [[SPIRES](#)].
- [20] K. Choi, K.S. Jeong and K.-I. Okumura, *Flavor and CP conserving moduli mediated SUSY breaking in flux compactification*, *JHEP* **07** (2008) 047 [[arXiv:0804.4283](#)] [[SPIRES](#)].
- [21] A.E. Nelson and N.T. Weiner, *Extended anomaly mediation and new physics at 10 TeV*, [hep-ph/0210288](#) [[SPIRES](#)].
- [22] K. Hsieh and M.A. Luty, *Mixed gauge and anomaly mediation from new physics at 10 TeV*, *JHEP* **06** (2007) 062 [[hep-ph/0604256](#)] [[SPIRES](#)].
- [23] N. Okada, *Positively-deflected anomaly mediation*, *Phys. Rev. D* **65** (2002) 115009 [[hep-ph/0202219](#)] [[SPIRES](#)].
- [24] M. Endo and K. Yoshioka, *Low-scale gaugino mass unification*, *Phys. Rev. D* **78** (2008) 025012 [[arXiv:0804.4192](#)] [[SPIRES](#)].
- [25] V. Kaplunovsky and J. Louis, *Field dependent gauge couplings in locally supersymmetric effective quantum field theories*, *Nucl. Phys. B* **422** (1994) 57 [[hep-th/9402005](#)] [[SPIRES](#)].
- [26] N.K. Falck, *Renormalization group equations for softly broken supersymmetry: the most general case*, *Z. Phys. C* **30** (1986) 247 [[SPIRES](#)];
D.J. Castano, E.J. Piard and P. Ramond, *Renormalization group study of the standard model and its extensions. 2. The minimal supersymmetric standard model*, *Phys. Rev. D* **49** (1994) 4882 [[hep-ph/9308335](#)] [[SPIRES](#)].
- [27] K. Choi and H.P. Nilles, *The gaugino code*, *JHEP* **04** (2007) 006 [[hep-ph/0702146](#)] [[SPIRES](#)].
- [28] K. Choi, K.S. Jeong, K.-I. Okumura and Y. Shimizu, *Deflected anomaly mediation into mirage mediation*, talk at *The physical society of Japan annual meeting, Hokkaido university Japan 24 September 2007*, unpublished.
- [29] S.P. Martin and P. Ramond, *Sparticle spectrum constraints*, *Phys. Rev. D* **48** (1993) 5365 [[hep-ph/9306314](#)] [[SPIRES](#)];

- S. Dimopoulos, S.D. Thomas and J.D. Wells, *Sparticle spectroscopy and electroweak symmetry breaking with gauge-mediated supersymmetry breaking*, *Nucl. Phys. B* **488** (1997) 39 [[hep-ph/9609434](#)] [[SPIRES](#)].
- [30] For a recent discussion of this issue see S.P. Martin, *Refined gluino and squark pole masses beyond leading order*, *Phys. Rev. D* **74** (2006) 075009 [[hep-ph/0608026](#)] [[SPIRES](#)].
- [31] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, *Aspects of grand unified models with softly broken supersymmetry*, *Prog. Theor. Phys.* **68** (1982) 927 [Erratum *ibid.* **70** (1983) 330] [[SPIRES](#)];
 L.E. Ibáñez and G.G. Ross, *SU(2)_L × U(1) symmetry breaking as a radiative effect of supersymmetry breaking in guts*, *Phys. Lett. B* **110** (1982) 215 [[SPIRES](#)];
 J.R. Ellis, D.V. Nanopoulos and K. Tamvakis, *Grand unification in simple supergravity*, *Phys. Lett. B* **121** (1983) 123 [[SPIRES](#)];
 L. Álvarez-Gaumé, J. Polchinski and M.B. Wise, *Minimal low-energy supergravity*, *Nucl. Phys. B* **221** (1983) 495 [[SPIRES](#)].
- [32] G.F. Giudice and A. Masiero, *A natural solution to the μ problem in supergravity theories*, *Phys. Lett. B* **206** (1988) 480 [[SPIRES](#)];
 E.J. Chun, J.E. Kim and H.P. Nilles, *A natural solution of the mu problem with a composite axion in the hidden sector*, *Nucl. Phys. B* **370** (1992) 105 [[SPIRES](#)].
- [33] P. Nath and T.R. Taylor, *Modular invariance, soft breaking, μ and $\tan\beta$ in superstring models*, *Phys. Lett. B* **548** (2002) 77 [[hep-ph/0209282](#)] [[SPIRES](#)].
- [34] R. Barbieri and G.F. Giudice, *Upper bounds on supersymmetric particle masses*, *Nucl. Phys. B* **306** (1988) 63 [[SPIRES](#)];
 P.H. Chankowski, J.R. Ellis and S. Pokorski, *The fine-tuning price of LEP*, *Phys. Lett. B* **423** (1998) 327 [[hep-ph/9712234](#)] [[SPIRES](#)];
 P.H. Chankowski, J.R. Ellis, M. Olechowski and S. Pokorski, *Haggling over the fine-tuning price of LEP*, *Nucl. Phys. B* **544** (1999) 39 [[hep-ph/9808275](#)] [[SPIRES](#)];
 G.L. Kane and S.F. King, *Naturalness implications of LEP results*, *Phys. Lett. B* **451** (1999) 113 [[hep-ph/9810374](#)] [[SPIRES](#)];
 M. Bastero-Gil, G.L. Kane and S.F. King, *Fine-tuning constraints on supergravity models*, *Phys. Lett. B* **474** (2000) 103 [[hep-ph/9910506](#)] [[SPIRES](#)].
- [35] K. Choi, K.S. Jeong, T. Kobayashi and K.-I. Okumura, *Little SUSY hierarchy in mixed modulus-anomaly mediation*, *Phys. Lett. B* **633** (2006) 355 [[hep-ph/0508029](#)] [[SPIRES](#)];
TeV scale mirage mediation and natural little SUSY hierarchy, *Phys. Rev. D* **75** (2007) 095012 [[hep-ph/0612258](#)] [[SPIRES](#)];
 R. Kitano and Y. Nomura, *A solution to the supersymmetric fine-tuning problem within the MSSM*, *Phys. Lett. B* **631** (2005) 58 [[hep-ph/0509039](#)] [[SPIRES](#)]; *Supersymmetry, naturalness and signatures at the LHC*, *Phys. Rev. D* **73** (2006) 095004 [[hep-ph/0602096](#)] [[SPIRES](#)];
 A. Pierce and J. Thaler, *Prospects for mirage mediation*, *JHEP* **09** (2006) 017 [[hep-ph/0604192](#)] [[SPIRES](#)].
- [36] R. Dermisek and H.D. Kim, *Radiatively generated maximal mixing scenario for the Higgs mass and the least fine tuned minimal supersymmetric standard model*, *Phys. Rev. Lett.* **96** (2006) 211803 [[hep-ph/0601036](#)] [[SPIRES](#)];
 R. Dermisek, H.D. Kim and I.-W. Kim, *Mediation of supersymmetry breaking in gauge messenger models*, *JHEP* **10** (2006) 001 [[hep-ph/0607169](#)] [[SPIRES](#)].

- [37] PARTICLE DATA GROUP collaboration, C. Amsler et al., *Review of particle physics*, *Phys. Lett. B* **667** (2008) 1 [SPIRES].
- [38] J.L. Feng, T. Moroi, L. Randall, M. Strassler and S.-F. Su, *Discovering supersymmetry at the Tevatron in Wino LSP scenarios*, *Phys. Rev. Lett.* **83** (1999) 1731 [hep-ph/9904250] [SPIRES];
T. Gherghetta, G.F. Giudice and J.D. Wells, *Phenomenological consequences of supersymmetry with anomaly-induced masses*, *Nucl. Phys. B* **559** (1999) 27 [hep-ph/9904378] [SPIRES];
G.D. Kribs, *Distinguishing anomaly-mediation from gauge-mediation with a Wino NLSP*, *Phys. Rev. D* **62** (2000) 015008 [hep-ph/9909376] [SPIRES];
U. Chattopadhyay, D.K. Ghosh and S. Roy, *Constraining anomaly mediated supersymmetry breaking framework via ongoing muon $g-2$ experiment at Brookhaven*, *Phys. Rev. D* **62** (2000) 115001 [hep-ph/0006049] [SPIRES];
H. Baer, J.K. Mizukoshi and X. Tata, *Reach of the CERN LHC for the minimal anomaly-mediated SUSY breaking model*, *Phys. Lett. B* **488** (2000) 367 [hep-ph/0007073] [SPIRES];
A.J. Barr, C.G. Lester, M.A. Parker, B.C. Allanach and P. Richardson, *Discovering anomaly-mediated supersymmetry at the LHC*, *JHEP* **03** (2003) 045 [hep-ph/0208214] [SPIRES];
M. Ibe, T. Moroi and T.T. Yanagida, *Possible signals of Wino LSP at the Large Hadron Collider*, *Phys. Lett. B* **644** (2007) 355 [hep-ph/0610277] [SPIRES];
S. Asai, T. Moroi, K. Nishihara and T.T. Yanagida, *Testing the anomaly mediation at the LHC*, *Phys. Lett. B* **653** (2007) 81 [arXiv:0705.3086] [SPIRES];
S. Asai, T. Moroi and T.T. Yanagida, *Test of anomaly mediation at the LHC*, *Phys. Lett. B* **664** (2008) 185 [arXiv:0802.3725] [SPIRES].
- [39] T. Moroi and L. Randall, *Wino cold dark matter from anomaly-mediated SUSY breaking*, *Nucl. Phys. B* **570** (2000) 455 [hep-ph/9906527] [SPIRES];
H. Baer, A. Mustafayev, E.-K. Park and S. Profumo, *Mixed Wino dark matter: consequences for direct, indirect and collider detection*, *JHEP* **07** (2005) 046 [hep-ph/0505227] [SPIRES];
H. Baer, T. Krupovnickas, S. Profumo and P. Ullio, *Model independent approach to focus point supersymmetry: from dark matter to collider searches*, *JHEP* **10** (2005) 020 [hep-ph/0507282] [SPIRES];
P. Grajek, G. Kane, D.J. Phalen, A. Pierce and S. Watson, *Neutralino dark matter from indirect detection revisited*, arXiv:0807.1508 [SPIRES].
- [40] H. Baer, K.-M. Cheung and J.F. Gunion, *A heavy gluino as the lightest supersymmetric particle*, *Phys. Rev. D* **59** (1999) 075002 [hep-ph/9806361] [SPIRES].
- [41] S. Raby and K. Tobe, *The phenomenology of SUSY models with a gluino LSP*, *Nucl. Phys. B* **539** (1999) 3 [hep-ph/9807281] [SPIRES];
A. Mafi and S. Raby, *An analysis of a heavy gluino LSP at CDF: the heavy gluino window*, *Phys. Rev. D* **62** (2000) 035003 [hep-ph/9912436] [SPIRES].
- [42] A.C. Kraan, *Interactions of heavy stable hadronizing particles*, *Eur. Phys. J. C* **37** (2004) 91 [hep-ex/0404001] [SPIRES];
A.C. Kraan, J.B. Hansen and P. Nevski, *Discovery potential of R-hadrons with the ATLAS detector*, *Eur. Phys. J. C* **49** (2007) 623 [hep-ex/0511014] [SPIRES].
- [43] P.J. Bussey, T.D. Coughlin, J.R. Forshaw and A.D. Pilkington, *Central exclusive production of longlived gluinos at the LHC*, *JHEP* **11** (2006) 027 [hep-ph/0607264] [SPIRES].

- [44] M. Fairbairn et al., *Stable massive particles at colliders*, *Phys. Rept.* **438** (2007) 1 [[hep-ph/0611040](#)] [[SPIRES](#)].
- [45] N. Arkani-Hamed and S. Dimopoulos, *Supersymmetric unification without low energy supersymmetry and signatures for fine-tuning at the LHC*, *JHEP* **06** (2005) 073 [[hep-th/0405159](#)] [[SPIRES](#)];
G.F. Giudice and A. Romanino, *Split supersymmetry*, *Nucl. Phys. B* **699** (2004) 65 [*Erratum ibid.* **706** (2005) 65] [[hep-ph/0406088](#)] [[SPIRES](#)].
- [46] W. Kilian, T. Plehn, P. Richardson and E. Schmidt, *Split supersymmetry at colliders*, *Eur. Phys. J. C* **39** (2005) 229 [[hep-ph/0408088](#)] [[SPIRES](#)].
- [47] J.L. Hewett, B. Lillie, M. Masip and T.G. Rizzo, *Signatures of long-lived gluinos in split supersymmetry*, *JHEP* **09** (2004) 070 [[hep-ph/0408248](#)] [[SPIRES](#)].
- [48] K. Cheung and W.-Y. Keung, *Split supersymmetry, stable gluino and gluinonium*, *Phys. Rev. D* **71** (2005) 015015 [[hep-ph/0408335](#)] [[SPIRES](#)].
- [49] M. Toharia and J.D. Wells, *Gluino decays with heavier scalar superpartners*, *JHEP* **02** (2006) 015 [[hep-ph/0503175](#)] [[SPIRES](#)].
- [50] A. Arvanitaki, S. Dimopoulos, A. Pierce, S. Rajendran and J.G. Wacker, *Stopping gluinos*, *Phys. Rev. D* **76** (2007) 055007 [[hep-ph/0506242](#)] [[SPIRES](#)].
- [51] G.R. Farrar and P. Fayet, *Phenomenology of the production, decay and detection of new hadronic states associated with supersymmetry*, *Phys. Lett. B* **76** (1978) 575 [[SPIRES](#)].
- [52] M.S. Chanowitz and S.R. Sharpe, *Spectrum of gluino bound states*, *Phys. Lett. B* **126** (1983) 225 [[SPIRES](#)];
F. Buccella, G.R. Farrar and A. Pugliese, *R baryon masses*, *Phys. Lett. B* **153** (1985) 311 [[SPIRES](#)].
- [53] UKQCD collaboration, M. Foster and C. Michael, *Hadrons with a heavy colour-adjoint particle*, *Phys. Rev. D* **59** (1999) 094509 [[hep-lat/9811010](#)] [[SPIRES](#)].
- [54] M. Moshe, *Recent developments in reggeon field theory*, *Phys. Rept.* **37** (1978) 255 [[SPIRES](#)].
- [55] J.F. Gunion and D.E. Soper, *Quark counting and hadron size effects for total cross-sections*, *Phys. Rev. D* **15** (1977) 2617 [[SPIRES](#)].
- [56] P.F. Smith et al., *A search for anomalous hydrogen in enriched D-2 O, using a time-of-flight spectrometer*, *Nucl. Phys. B* **206** (1982) 333 [[SPIRES](#)].
- [57] T.K. Hemmick et al., *A search for anomalously heavy isotopes of low Z nuclei*, *Phys. Rev. D* **41** (1990) 2074 [[SPIRES](#)].
- [58] D. Javorsek et al., *New experimental limits on strongly interacting massive particles at the TeV scale*, *Phys. Rev. Lett.* **87** (2001) 231804 [[SPIRES](#)].
- [59] R.N. Mohapatra and V.L. Teplitz, *Primordial nucleosynthesis constraint on massive, stable, strongly interacting particles*, *Phys. Rev. Lett.* **81** (1998) 3079 [[hep-ph/9804420](#)] [[SPIRES](#)].
- [60] D. Javorsek II, E. Fischbach and V. Teplitz, *New experimental bounds on the contributions to the cosmological density parameter Ω from strongly interacting massive particles*, *Astrophys. J.* **568** (2002) 1.
- [61] B. de Carlos, J.A. Casas, F. Quevedo and E. Roulet, *Model independent properties and cosmological implications of the dilaton and moduli sectors of 4d strings*, *Phys. Lett. B* **318** (1993) 447 [[hep-ph/9308325](#)] [[SPIRES](#)];

- T. Banks, D.B. Kaplan and A.E. Nelson, *Cosmological implications of dynamical supersymmetry breaking*, *Phys. Rev. D* **49** (1994) 779 [[hep-ph/9308292](#)] [[SPIRES](#)].
- [62] S. Nakamura and M. Yamaguchi, *Gravitino production from heavy moduli decay and cosmological moduli problem revived*, *Phys. Lett. B* **638** (2006) 389 [[hep-ph/0602081](#)] [[SPIRES](#)];
M. Endo, K. Hamaguchi and F. Takahashi, *Moduli-induced gravitino problem*, *Phys. Rev. Lett.* **96** (2006) 211301 [[hep-ph/0602061](#)] [[SPIRES](#)].
- [63] L. Covi, J.E. Kim and L. Roszkowski, *Axinos as cold dark matter*, *Phys. Rev. Lett.* **82** (1999) 4180 [[hep-ph/9905212](#)] [[SPIRES](#)];
L. Covi, H.-B. Kim, J.E. Kim and L. Roszkowski, *Axinos as dark matter*, *JHEP* **05** (2001) 033 [[hep-ph/0101009](#)] [[SPIRES](#)].
- [64] I. Hinchliffe, F.E. Paige, M.D. Shapiro, J. Soderqvist and W. Yao, *Precision SUSY measurements at CERN LHC*, *Phys. Rev. D* **55** (1997) 5520 [[hep-ph/9610544](#)] [[SPIRES](#)];
M.M. Nojiri, G. Polesello and D.R. Tovey, *Proposal for a new reconstruction technique for SUSY processes at the LHC*, [hep-ph/0312317](#) [[SPIRES](#)];
K. Kawagoe, M.M. Nojiri and G. Polesello, *A new SUSY mass reconstruction method at the CERN LHC*, *Phys. Rev. D* **71** (2005) 035008 [[hep-ph/0410160](#)] [[SPIRES](#)];
B.K. Gjelsten, D.J. Miller and P. Osland, *Measurement of SUSY masses via cascade decays for SPS 1a*, *JHEP* **12** (2004) 003 [[hep-ph/0410303](#)] [[SPIRES](#)];
W.S. Cho, K. Choi, Y.G. Kim and C.B. Park, *Gluino transverse mass*, *Phys. Rev. Lett.* **100** (2008) 171801 [[arXiv:0709.0288](#)] [[SPIRES](#)]; *Measuring superparticle masses at hadron collider using the transverse mass kink*, *JHEP* **02** (2008) 035 [[arXiv:0711.4526](#)] [[SPIRES](#)];
A.J. Barr, B. Gripaios and C.G. Lester, *Weighing wimps with kinks at colliders: invisible particle mass measurements from endpoints*, *JHEP* **02** (2008) 014 [[arXiv:0711.4008](#)] [[SPIRES](#)];
H.-C. Cheng, D. Engelhardt, J.F. Gunion, Z. Han and B. McElrath, *Accurate mass determinations in decay chains with missing energy*, *Phys. Rev. Lett.* **100** (2008) 252001 [[arXiv:0802.4290](#)] [[SPIRES](#)];
M. Burns, K. Kong, K.T. Matchev and M. Park, *Using subsystem MT_2 for complete mass determinations in decay chains with missing energy at hadron colliders*, *JHEP* **03** (2009) 143 [[arXiv:0810.5576](#)] [[SPIRES](#)].