

Hide and Seek with Massive Fields in Modulus Mediation

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We study the modulus mediation of supersymmetry breaking motivated by superstring theory. We show that the renormalization group running of the soft supersymmetry breaking parameters due to the interactions of massive fields is canceled by the threshold corrections at one-loop order, if their mass is given by nonperturbative dynamics controlled by the same modulus that mediates supersymmetry breaking and a sum rule of the modular weights holds for the Yukawa couplings. As an example, we discuss order reduction of lepton flavor violation in the supersymmetric seesaw mechanism, which revives the parameter space already excluded.

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Introduction.—Supersymmetry (SUSY) is an attractive candidate for physics beyond the standard model (SM). It stabilizes the big hierarchy between the Planck scale and the electroweak (EW) scale, which is otherwise vulnerable to corrections from the heavy scale. Furthermore, the minimal SUSY SM (MSSM) predicts gauge coupling unification and the existence of dark matter protected by R parity. The observed Higgs boson mass is perfectly consistent with the prediction of MSSM and nonobservation of the SUSY particles at the Large Hadron Collider, although the heavy superpartners generally introduce a little fine-tuning in the EW symmetry breaking.

The phenomenology of low energy SUSY mostly depends on the parameters that softly break SUSY and determine the spectrum of the superpartners. The supertrace constraint requires these parameters to be generated via the nonrenormalizable Kähler potential that couples the visible MSSM sector to the hidden sector where SUSY is spontaneously broken [1,2]. Phenomenologically, the flavor structure of the soft SUSY breaking parameters is tightly constrained by the flavor and CP violating processes providing hints for the origin of the Kähler potential. The gauge mediation [3] and the anomaly mediation [4,5] are widely studied bottom-up models that readily satisfy the constraint.

The anomaly mediation particularly has a special property dubbed ultraviolet insensitivity, which means the soft SUSY breaking parameters in the low energy effective theory are completely determined by the SUSY preserving couplings at the same renormalization scale irrespective of

the ultraviolet theory. This stems from the fact that the SUSY breaking effect is encapsulated in the SUSY regulators or, in other words, the renormalization scale dependence and its analytic continuation to the superspace [6,7]. In view of the renormalization group (RG) running of the soft parameters, running due to interactions with massive particles is canceled by the threshold corrections and disappears in low energy theory. This is useful to evade the flavor- CP constraints because the soft SUSY breaking parameters in the MSSM are functions of the gauge and Yukawa couplings at the EW scale. Unfortunately, the minimal model predicts a tachyonic slepton mass and it must be circumvented in some way, while it is reasonable to think that this property will be lost if other sources of SUSY breaking are introduced simultaneously, since it relies on the fact that the mediator only couples to the Weyl anomaly.

Another well-motivated semi-top-down mechanism is the modulus mediation [8–10]. It is a string realization of the gravity mediation based on supergravity [11,12]. The modern view of supergravity is the low energy effective theory of superstring. The 10D spacetime is compactified on the Calabi-Yau threefold and the low-energy description of the massless modes is given by the 4D supergravity. The visible and hidden sectors can be geometrically separated and direct interaction between the two sectors could be strongly suppressed [4,13–19]. On the other hand, the Calabi-Yau manifold is parametrized by the moduli which characterize the size of internal 3 and 4 cycles. They are 4D massless fields originating from the 6D metric. Combined with the higher form fields depending on the type of string, they are organized into the chiral multiplets U_i , T_i called complex structure moduli and Kähler moduli, respectively. The 10-dimensional dilaton also produces a massless chiral multiplet S called a complex dilaton. Because of their gravitational origin their interactions with the matter fields are intrinsically nonrenormalizable and often preserve flavor at least at the leading order of g_s and α' expansion

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[8,9,20,21]. Thus, the modulus and dilaton fields are recognized as good candidates of the SUSY-breaking mediator and their phenomenology has been studied for almost three decades. In such a study, including more general gravity mediation, the SUSY spectrum at the EW scale is calculated using the RG equations with the given boundary condition at the string (or Planck) scale. All the relevant interactions between the mediation scale and the EW scale are believed to leave their traces in the low energy parameters via RG running. This machinery sometimes leads to important phenomenological predictions [22–36].

In this Letter, we argue that the modulus mediation (including dilaton) could have similar properties to the anomaly mediation in some circumstances. The RG running of the soft SUSY breaking parameters due to interactions of massive fields is canceled by their threshold corrections at one-loop order, if their mass is given by nonperturbative dynamics controlled by the same modulus that mediates the SUSY breaking and a sum rule of the modular weights holds for the Yukawa couplings. Considering intermediate scales in string theory often emerge from such nonperturbative dynamics, this fact could have significant phenomenological implications. As an example, we discuss order reduction of lepton flavor violation in the SUSY seesaw mechanism [29–33].

Modulus mediation.—We first review the modulus mediation of SUSY breaking [37]. We employ the superspace notation in Ref. [40] unless otherwise specified. In the superconformal formulation, the scalar part of $N = 1$ 4D supergravity action is given by

$$\mathcal{S} = \int d^4x \sqrt{-g_C} \left[\int d^4\theta \left\{ -3|C|^2 \exp\left(-\frac{\mathcal{K}}{3}\right) \right\} + \left\{ \int d^2\theta \left(\frac{f_a}{4} W^a W^a + C^3 \mathcal{W} \right) + \text{H.c.} \right\} + \dots \right], \quad (1)$$

where C denotes the conformal compensator field. Suppressed terms depend on the curvature or explicit SUSY breaking required for the uplifting of the AdS vacuum as, e.g., in the KKLT construction [41]. In the formula, \mathcal{K} , f_a , and \mathcal{W} represent the Kähler potential, gauge kinetic function, and superpotential, respectively. The real part of the gauge kinetic function gives the gauge coupling constant

$$g_a^2 = 1/\text{Re}(f_a)|_{\theta^2=\bar{\theta}^2=0}. \quad (2)$$

Fixing the conformal gauge at $C_0 = \exp(\mathcal{K}/6)$ and integrating out F terms using the equation of motion,

$$\frac{F^C}{C_0} = \frac{1}{3} F^I \mathcal{K}_I + m_{3/2}, \quad F^I = -e^{\mathcal{K}/2} K^{I\bar{J}} \overline{D_{\bar{J}} \mathcal{W}}. \quad (3)$$

we can obtain the component supergravity action in the Einstein frame. In these formulas the subscript means

derivative by the corresponding field and $K^{I\bar{J}} = (K_{I\bar{J}})^{-1}$, $D_I \mathcal{W} = \mathcal{W}_I + \mathcal{K}_I \mathcal{W}$ are understood. The 4D metric in the conformal frame $g_{\mu\nu}^C$ is related to that of the Einstein frame via $g_{\mu\nu}^E = |C|^{-2} \exp(\mathcal{K}/3) g_{\mu\nu}^C$.

In the following, we assume that the Kähler potential and the superpotential are expanded in terms of the chiral matter superfields Φ^i in the visible sector as follows:

$$\mathcal{K} = \mathcal{K}_0(T + T^\dagger) + Z_i(T + T^\dagger) \Phi^i \Phi^{\dagger i}, \quad (4)$$

$$\mathcal{W} = W_0 + \frac{1}{2} \mu_{ij} \Phi^i \Phi^j + \frac{1}{3!} \lambda_{ijk} \Phi^i \Phi^j \Phi^k, \quad (5)$$

where T denotes the modulus fields collectively. Here the hidden sector is sequestered from the visible sector. We also explicitly consider the moduli having the shift symmetry $T \rightarrow T + ic$ with a real constant c in the Lagrangian, up to total derivatives. Thus, the Kähler potential has modulus dependence via $T + T^\dagger$ and no perturbative T dependence in superpotential due to holomorphy. For instance, the complex dilaton in the heterotic string and the Kähler (complex structure) moduli in type II B (A) orientifolds possess such a property [8,9,20].

The soft SUSY breaking terms of the canonically normalized visible fields can be parametrized as

$$-\mathcal{L}_{\text{Soft}} = m_i^2 |\phi^i|^2 + \left[\frac{1}{2} M_a \bar{\lambda}_L^a \lambda_L^a + \frac{1}{2} B_{ij} M_{ij} \phi^i \phi^j + \frac{1}{3!} A_{ijk} Y_{ijk} \phi^i \phi^j \phi^k + \text{H.c.} \right], \quad (6)$$

where canonical mass parameters and the Yukawa couplings are given by

$$M_{ij} = \mu_{ij} / \sqrt{Y_i Y_j}, \quad Y_{ijk} = \lambda_{ijk} / \sqrt{Y_i Y_j Y_k}, \quad (7)$$

with $Y_i = e^{-\mathcal{K}_0/3} Z_i$. These terms can be read off by integrating out the F terms of the visible fields Φ^i in the supergravity action using their equation of motion as

$$M_a = F^I \partial_I \ln(\text{Re} f_a), \quad m_i^2 = -F^I F^{\bar{J}} \partial_I \partial_{\bar{J}} \ln Y_i, \quad (8)$$

$$A_{ijk} = A_i + A_j + A_k - F^I \partial_I \ln(\lambda_{ijk}), \quad (9)$$

$$B_{ij} = -\frac{F^C}{C_0} + A_i + A_j - F^I \partial_I \ln(\mu_{ij}), \quad (10)$$

where $A_i = F^I \partial_I \ln Y_i$, F^I , and ∂_I denote the F terms and the derivative for the moduli (or hidden sector fields in general). In the formulas, we leave the field dependence in the holomorphic couplings for later purposes. We also omit the effects of the flavor mixing to avoid unnecessary complications. Their inclusion by the Kähler connection is straightforward and not essential in the following

discussion. We will consider the gauge modulus having the following tree-level property at the string scale,

$$f_a = T, \quad Y_i = (T + T^\dagger)^{c_i}. \quad (11)$$

The complex dilaton in heterotic string and the overall Kähler modulus in type II B toroidal orientifolds are among the examples. c_i is related to the modular weight of the matter field n_i as $c_i = 1 - n_i$ [38,39] and typically given by a ratio of small integers depending on the origin of the field. Putting them into Eqs. (8)–(10), we obtain

$$M_a = M_0, \quad A_{ijk} = (c_i + c_j + c_k)M_0, \quad m_i^2 = c_i |M_0|^2, \quad (12)$$

with $M_0 \equiv F^T/(T + T^\dagger)$ at the string scale. For the phenomenological purpose, we need the values evolved by the renormalization group equations down to the EW scale. There is also contribution from the anomaly mediation radiatively mediated by C . We suppress the anomaly mediation as a subleading contribution; however, it is straightforward to check that the following discussions also held under the simultaneous presence of anomaly mediation (the mirage mediation) [38,39,42–44].

Method of similarity.—We discuss evolution of the soft SUSY breaking parameters under the 1-loop RG equations in terms of the method of similarity [45–47]. The 1-loop RG equations for gauge coupling and the matter Kähler metric is given by

$$\frac{dg_a^{-2}}{d \ln \mu} = -\frac{b_a}{8\pi^2}, \quad b_a = -3T_a^G + \sum_i T_a^i, \quad (13)$$

$$\frac{d \ln Y_i}{d \ln \mu} = \frac{\gamma_i}{8\pi^2}, \quad \gamma_i = 2 \sum_a C_2^{ai} g_a^2 - \frac{1}{2} \sum_{jk} |Y_{ijk}|^2, \quad (14)$$

where T_a^G (T_a^i) is the Dynkin index of the adjoint (ϕ^i) representation for the corresponding gauge group. C_2^{ai} denotes the quadratic Casimir of ϕ^i . This results in the RG equation for the canonical Yukawa coupling,

$$\frac{dY_{ijk}}{d \ln \mu} = -\frac{1}{16\pi^2} (\gamma_i + \gamma_j + \gamma_k) Y_{ijk}, \quad (15)$$

since the holomorphic couplings are not renormalized.

It is known that these first order differential equations are invariant under the following scaling transformation with a real constant Δ ,

$$\ln \mu \rightarrow \Delta \cdot \ln \mu, \quad g_a^{-2} \rightarrow \Delta \cdot g_a^{-2}, \quad Y_{ijk} \rightarrow \Delta^{-1/2} \cdot Y_{ijk}. \quad (16)$$

Then, if $g_a^{-2}[g_{aS}^{-2}; (\mu/M_S)]$ represents a solution for the boundary condition, $g_a^{-2} = g_{aS}^{-2}$ at $\mu = M_S$, a function,

$\Delta^{-1} \cdot g_a^{-2}[\Delta \cdot g_{aS}^{-2}; (\mu/M_S)^\Delta]$ also satisfies the equation. The boundary condition of this solution is g_{aS}^{-2} ; therefore, the uniqueness of the solution leads to

$$g_a^{-2}[g_{aS}^{-2}; (\mu/M_S)] = \Delta^{-1} \cdot g_a^{-2}[\Delta \cdot g_{aS}^{-2}; (\mu/M_S)^\Delta]. \quad (17)$$

A similar discussion for the Yukawa coupling follows:

$$Y_{ijk}[\{g_{aS}, Y_S\}; (\mu/M_S)] = \Delta^{1/2} \cdot Y_{ijk}[\{\Delta^{-1/2} \cdot g_{aS}, \Delta^{-1/2} \cdot Y_S\}; (\mu/M_S)^\Delta], \quad (18)$$

where $\{g_{aS}, Y_S\}$ represents the boundary condition of the relevant gauge and Yukawa couplings for the set of RG equations. In the following, we take M_S as the string scale where the boundary condition, Eq. (11) is satisfied. If the sum rule, $a_{ijk} = c_i + c_j + c_k = 1$ holds for all the Yukawa couplings, which is often the case for the large Yukawa couplings [18,21,48–51], we set Δ as

$$\Delta = \langle T + T^\dagger \rangle / (T + T^\dagger), \quad (19)$$

where the rectangle parentheses denote the vacuum expectation value. Then we obtain the final expressions,

$$g_a^{-2}[g_{aS}^{-2}; (\mu/M_S)] = \Delta^{-1} \cdot g_a^{-2}[\langle g_{aS}^{-2} \rangle; (\mu/M_S)^\Delta], \quad (20)$$

$$Y_{ijk}[\{g_{aS}, Y_S\}; (\mu/M_S)] = \Delta^{1/2} \cdot Y_{ijk}[\{\langle g_{aS} \rangle, \langle Y_S \rangle\}; (\mu/M_S)^\Delta]. \quad (21)$$

In these formulas, the modulus dependence originally residing in the boundary conditions are squeezed into the overall factors and the renormalization scale dependence. Then the analytic continuation into superspace [6,7] and Eqs. (8)–(10) with Z_i obtained by formal integration of γ_i yields a closed form for the soft SUSY breaking parameters at the scale μ using the SUSY parameters at the same scale [39],

$$\frac{M_a(\mu)}{M_0} = 1 + \frac{b_a}{8\pi^2} g_a^2(\mu) L, \quad (22)$$

$$\frac{A_{ijk}(\mu)}{M_0} = a_{ijk} - \frac{1}{8\pi^2} (\gamma_i + \gamma_j + \gamma_k) L, \quad (23)$$

$$\frac{m_i^2(\mu)}{|M_0|^2} = c_i - \frac{1}{4\pi^2} \gamma_i L - \frac{1}{8\pi^2} \frac{d\gamma_i}{d \ln \mu} L^2, \quad (24)$$

where $L = \ln(\mu/M_S)$. This property is analogous to the anomaly mediation where C dependence is encapsulated in the renormalization scale dependence as, $\mu/(CM_S)$ after the redefinition of the visible fields, $C\Phi_i \rightarrow \Phi_i$, however, only valid within the one-loop approximation.

Trace of massive fields.—In this Letter, we will examine the corrections to the soft SUSY breaking parameters due

to massive fields integrated out at some scale. As an example, we consider chiral superfields Ψ^α ($\alpha = 1, 2$) having a mass X and the Yukawa couplings λ' in the superpotential as

$$\delta\mathcal{W} = X\Psi^1\Psi^2 + \frac{1}{2}\lambda'_{ija}\Phi^i\Phi^j\Psi^\alpha + \frac{1}{2}\lambda'_{iab}\Phi^i\Psi^\alpha\Psi^\beta, \quad (25)$$

where Φ^i represents fields remain in the low energy effective theory. At the threshold $\mu = X$, we connect the RG solutions for their gauge and Yukawa couplings continuously, while switching the gauge beta function using a step function [52]. Applying the formula in the previous paragraph recursively, we obtain,

$$\begin{aligned} g_a^{-2}[g_{aX}^{-2}; (\mu/|X|)] \\ &= \Delta^{-1} \cdot g_a^{-2}[g_{aX}^{-2}[\Delta \cdot g_{aS}^{-2}; (|X|/M_S)^\Delta]; (\mu/|X|)^\Delta] \\ &= \Delta^{-1} \cdot g_a^{-2}[\Delta \cdot g_{aS}^{-2}; (|X|/M_S)^\Delta, (\mu/M_S)^\Delta] \\ &= \Delta^{-1} \cdot g_a^{-2}[\langle g_{aS}^{-2} \rangle; (|X|/M_S)^\Delta, (\mu/M_S)^\Delta], \end{aligned} \quad (26)$$

where the subscript X stands for the value at $\mu = X$. A similar expression holds for the Yukawa coupling,

$$\begin{aligned} Y_{ijk}[\{g_{aX}, Y_X\}; (\mu/|X|)] \\ &= \Delta^{1/2} \cdot Y_{ijk}[\{\langle g_{aS} \rangle, \Delta^{(a-1)/2} \cdot \langle Y_S \rangle, \Delta^{(a'-1)/2} \cdot \langle Y'_S \rangle\}; \\ &(|X|/M_S)^\Delta, (\mu/M_S)^\Delta], \end{aligned} \quad (27)$$

where Y' represents the canonical Yukawa coupling corresponds to λ' .

If the sum rule of c_i holds for all the Yukawa couplings ($a_{ijk} = a'_{ija} = a'_{iab} = 1$), the modulus dependence can be squeezed into the overall factor and the dependence on the threshold mass and the renormalization scale. Furthermore, if the threshold is given by the nonperturbative effect like

$$X = \Lambda \exp(-bT), \quad (28)$$

and Λ is given by the cutoff $\Lambda = M_S$, the threshold dependence disappears. Note that such a nonperturbative mass due to field theoretic or string instantons is a typical origin of the intermediate scale in string theory [53–59] (see also Ref. [60] and the references therein). In such a case, the low energy soft SUSY breaking parameters again given by the closed form at μ as in Eqs. (22)–(24). This kind of ultraviolet insensitivity means that the threshold correction integrating out Ψ^α cancels the effect of RG running due to Y' as in the anomaly mediation [61]. The leftover is suppressed by $\ln(\Lambda/M_S)/8\pi^2$. If some Yukawa coupling Y does not satisfies the sum rule, the leading effect of Y' in the soft SUSY breaking appears through the interference terms like $|Y_S|^2|Y'_S|^2[\ln(|X|/M_S)/8\pi^2]^2$ in the Kähler metric, which is loop suppressed while $\log(|X|/M_S)$ enhanced. Inclusion of multithresholds is straightforward

and we will not repeat the discussion here. The massive field can be a gauge multiplet spontaneously broken by the scalar with a nonperturbative mass.

Example.—As a phenomenological example, we discuss lepton flavor violation in the SUSY seesaw mechanism [29–33]. We introduce three generations of right-hand neutrinos \bar{N}_i in the MSSM with the superpotential

$$\delta\mathcal{W} = \frac{1}{2}M_R\bar{N}_i\bar{N}_i + \lambda_{ij}^{ij}H_uL_i\bar{N}_j, \quad (29)$$

where H_u, L_i are up-type Higgs and lepton doublets in the MSSM (MSSMRN). We choose a flavor universal Majorana mass for simplicity. After integrating out \bar{N}_i , small Majorana neutrino masses are generated through dimension 5 operators [69,70]. In this scenario, it is well known that the RG running due to the neutrino Yukawa coupling induces flavor mixing in the left-hand slepton mass matrix, which results in the lepton flavor violating processes like $\mu \rightarrow e\gamma$ forbidden in the SM.

In Fig. 1, we calculate $\text{BR}(\mu \rightarrow e\gamma)$ for the constant M_R (dotted line) and M_R given by Eq. (28) (solid curve) as a function of $\log_{10}(\Lambda/M_S)$. We set the boundary condition (12) at the unification scale, $M_G = 2 \times 10^{16}$ GeV identified with the string scale and evolve the soft parameters down to the EW scale. We choose the modular weights of the MSSM fields as $c_i = 1/3$ so that all the Yukawa couplings satisfy the sum rule. The universal gaugino mass is set to 1 TeV, which corresponds to 2.3 TeV gluino with the one-loop RG equation. The threshold correction at M_R is calculated following the method in Refs. [6,7]. For the

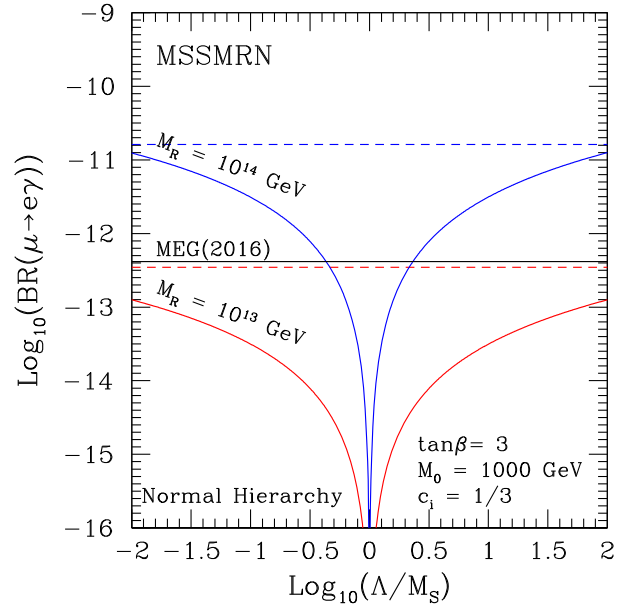


FIG. 1. $\text{BR}(\mu \rightarrow e\gamma)$ in the MSSM with right-handed neutrinos. The solid (dashed) curve represents the result for the non-perturbative (constant) mass term. The gluino mass is 2.3 TeV at the one-loop level.

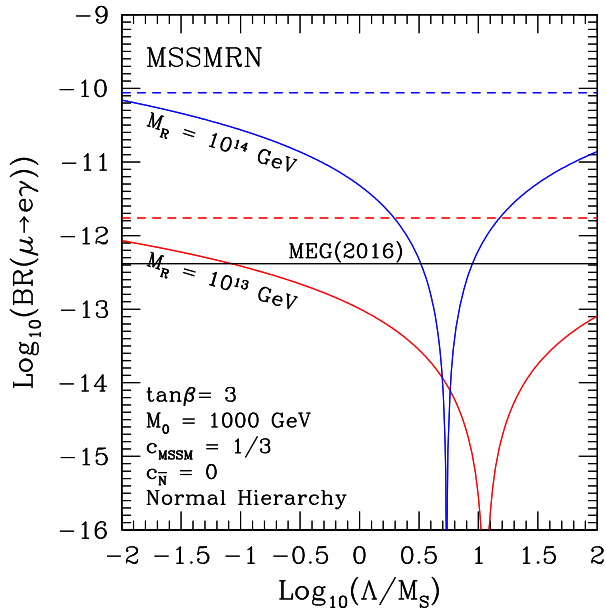


FIG. 2. The same as in Fig. 1 except the sum rule is violated in the neutrino Yukawa coupling ($c_{\tilde{N}=0}$).

estimation of the neutrino mass matrix, we adopt the normal hierarchy and the result of the global fit in Ref. [71] with the assumption of vanishing lightest neutrino mass. The horizontal line in the figure indicates the current experimental bound [72]. We show two cases, $M_R = 10^{13}$, 10^{14} GeV. In both cases, the branching ratio for the nonperturbative threshold disappears at $\Lambda = M_S$ due to cancellation. Even an order difference in Λ/M_S can lead to an order reduction in $\text{BR}(\mu \rightarrow e\gamma)$. This revives the parameter space already excluded by the experimental bound. In the calculation of $\text{BR}(\mu \rightarrow e\gamma)$, we confirm that the sum rule for the top Yukawa coupling is numerically irrelevant, while breaking it in the neutrino Yukawa coupling shifts the point of cancellation from $\Lambda = M_S$ as shown in Fig. 2, where we take $c_{\tilde{N}} = 0$ with the other modular weights and mass parameters are intact.

Conclusion.—We argued that the RG running of the soft SUSY breaking parameters due to interaction of massive fields is canceled by their threshold correction at one-loop order in the modulus mediation if their mass is given by nonperturbative dynamics controlled by the same modulus that mediates SUSY breaking and the modular-weight sum rule is satisfied in the case of the Yukawa couplings. This could have significant implications for phenomenology. As an example, we showed that lepton flavor violation in the SUSY seesaw mechanism could have order reduction. The cancellation also works with the anomaly mediation and it is interesting to introduce the seesaw mechanism or vector-like fields in the TeV scale mirage mediation models [43,47,73–79] without spoiling its little SUSY hierarchy. There might be a plethora of other phenomenological applications including massive gauge bosons as in the

grand unified theories. We also note that the mechanism itself is not limited to low energy SUSY models. The scale of the soft SUSY breaking is arbitrary. On the other hand, exploration of string model space realizing the cancellation is also an interesting issue, although it is beyond the scope of this Letter. We leave them to future work.

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