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Flavor and CP conserving moduli mediated SUSY breaking in flux compactification

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Abstract: In certain class of flux compactification, moduli mediated supersymmetry (SUSY) breaking preserves flavor and CP at leading order in the perturbative expansion controlled by the vacuum expectation value of the messenger modulus. Nevertheless there still might be dangerous flavor or CP violation induced by higher order Kähler potential. We examine the constraints on such SUSY breaking scheme imposed by low energy flavor and/or CP violating observables. It is found that all phenomenological constraints can be satisfied even for generic form of higher order Kähler potential and sparticle spectra in the sub-TeV range, under plausible assumptions on the size of higher order correction and flavor mixing angles. This implies for instance that mirage mediation scheme of SUSY breaking, which involves such modulus mediation together with an anomaly mediation of comparable size, and also the modulus-dominated mediation realized in flux compactification can be free from the SUSY flavor and CP problems, while giving gaugino and sfermion masses in the sub-TeV range.

Keywords: Flux compactifications, Supersymmetry Breaking, Supergravity Models, Supersymmetry Phenomenology.

Contents

1. Introduction

Weak scale supersymmetry (SUSY) is one of the prime candidates for physics beyond the standard model at the TeV scale [1] which will be probed soon at the LHC. One key question on weak scale SUSY is the origin of the soft SUSY breaking terms of visible gauge and matter superfields in low energy effective lagrangian [2]. Those soft terms are required to preserve flavor and CP with high accuracy, which severely constrains the possible mediation mechanism of SUSY breaking. There are certain schemes such as gauge mediation [3] and anomaly mediation [4] in which the standard model gauge interactions play a dominant role for the mediation, thereby automatically yield flavor and CP conserving soft terms. On the other hand, it is commonly thought that gravity mediation [5] generically leads to dangerous flavor and/or CP violation, and therefore needs an additional ingredient in order to be consistent with low energy observations.

The messenger scale of gravity mediation is near the Planck scale $M_{\rm Pl} = 2.4\times 10^{18}$ GeV which might be identified as the scale of quantum gravity. As string theory is the only known candidate for a theory of quantum gravity, it is natural to ask if string theory can provide a framework for flavor and CP conserving gravity mediation. In compactified string theory, moduli (including the dilaton) which determine the 4-dimensional (4D) gauge and Yukawa couplings are the most plausible candidate for a messenger of SUSY breaking, giving a gravity mediated contribution to gaugino and sfermion masses [6]. Then, constraints from low energy flavor and/or CP violations imply that the dominant messenger modulus should have flavor universal and CP conserving coupling to the minimal supersymmetric standard model (MSSM) chiral matter fields. As the mechanism of moduli stabilization determines which modulus is the dominant messenger, this in turn leads to a nontrivial constraint on the possible moduli stabilization scheme.

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Moduli mediated SUSY breaking and its phenomenological consequences have been studied before while regarding the moduli F components as a generic background without specifying the underlying stabilization scheme [6]. It has been noticed that a particular form of mediation dominated by the heterotic string dilaton gives universal and CP conserving soft terms at string tree level. If such dilaton domination can be realized while keeping the quantum correction to the Kähler potential small enough, the resulting soft terms would satisfy the constraints from flavor and CP violation with sparticle masses in sub-TeV range [7].

Recent progress of flux compactification suggests that string flux might play key roles to achieve a phenomenologically viable string vacuum state [8]. Flux can stabilize moduli while producing a huge landscape of vacua which contains a de Sitter vacuum with nearly vanishing cosmological constant. Flux compactification can also provide a SUSY breaking scheme in which soft terms preserve flavor and CP at leading order in the string coupling $g_{\rm st}$ or the slope parameter α' [9 – 11]. In this SUSY breaking scheme, only a particular modulus which is unfixed by flux and whose vacuum expectation value controls the $g_{\rm st}$ or α' expansion of 4D effective action can be an important messenger of SUSY breaking. The couplings of such messenger modulus to chiral matter fields are naturally flavor universal and CP conserving at leading order since the perturbative expansion is controlled by the messenger modulus itself.

While providing a good starting point, this scheme does not assure yet the absence of dangerous flavor or CP violation even when all other mediations in the model preserve flavor and CP. There can be higher order correction to the messenger modulus-matter couplings in the Kähler potential, which is expected to be flavor non-universal in general [7]. Then the modulus mediation itself associated with such higher order term might lead to a flavor or CP violation exceeding the current experimental bound. In this paper, we first discuss some features of flux compactification leading to a SUSY breaking scheme which preserves flavor and CP at leading order in the perturbative expansion controlled by the messenger modulus, and then examine the constraints on the scheme coming from flavor and/or CP violation induced by higher order Kähler potential. It is found that all phenomenological constraints can be satisfied even for generic form of higher order Kähler potential and sparticle spectra in the sub-TeV range, under plausible assumptions on the size of higher order correction and flavor mixing angles. This implies for instance that mirage mediation [10, 12 – 14] involving such modulus mediation together with an anomaly mediation of comparable size and also the modulus-dominated mediation [9, 15, 16] realized in flux compactification can be free from the SUSY flavor and CP problems, while giving gaugino and sfermion masses in the sub-TeV range. Same statement applies also to the axionic or deflected mirage mediation [17] in which gauge mediation of comparable size is added to mirage mediation.

The organization of this paper is as follows. In section 2, we discuss the relevant features of flux compactification leading to the SUSY breaking scheme under consideration. In section 3, we examine the structure of soft terms induced by higher order Kähler potential together with the constraints from various flavor and/or CP violating observables. Section 4 is the conclusion.

2. Relevant features of flux compactification

2.1 Moduli mass hierarchy

In this paper, we will be focusing on flux compactification which can realize the weak scale SUSY together with the high unification scale[1] $M_{\rm GUT} \sim 2\times 10^{16}$ GeV. In such compactification, both the string scale $M_{\rm st}$ and the compactification scale $M_{\rm KK}$ are comparable to the 4D Planck scale $M_{\rm Pl} \approx 2.4\times 10^{18}$ GeV or $M_{\rm GUT}$. This results in a big mass hierarchy between the heavy moduli U stabilized by flux and the light moduli T unfixed by flux. In this subsection, we briefly discuss this moduli mass hierarchy, while ignoring the little mass hierarchies of $\mathcal{O}(10-10^2)$ between $M_{\rm Pl}$, $M_{\rm st}$, and $M_{\rm KK}$, i.e. while regarding

$$M_{\rm st} \sim M_{\rm KK} \sim M_{\rm Pl}. \tag{2.1}$$

If one introduces nonzero flux over a cycle $\mathcal{C}$ in compact internal space, the modulus parameterizing the size of $\mathcal{C}$ is stabilized generically with a SUSY preserving mass m_U comparable to $M_{\rm st}$ [8]. In the language of 4D effective theory, one finds

$$m_U \sim \left\langle \frac{\partial^2 W_{\rm flux}}{\partial U^2}\right\rangle \sim M_{\rm st}, \tag{2.2}$$

where $W_{\rm flux}$ is the flux-induced superpotential. (More precisely, m_U is given by $m_U \sim M_{\rm Pl} e^{K/2}\partial_U^2 W_{\rm flux}/\partial_U\partial_{\bar{U}}K \sim g_{\rm st}M_{\rm st}/\sqrt{M_{\rm st}^6 V}$, where $g_{\rm st}$ and V denote the string coupling and the compactification volume, respectively.) Most string compactifications allow the NS or RR 3-form fluxes over the 3-cycles of internal space, which would stabilize all complex structure moduli. Depending upon the model, string dilaton or Kähler moduli might be stabilized also by flux. For instance, in type IIB compactification, the dilaton can be stabilized by RR 3-form flux [18]. It has been noticed that Kähler moduli in heterotic compactification might be stabilized by intrinsic torsion flux [19], suggesting the possibility that all complex structure and Kähler moduli in heterotic compactifications are stabilized by nonzero NS and torsion fluxes.

In many flux compactifications, there remains a modulus T which can *not* be fixed by flux. One example of such modulus is the volume modulus in type IIB flux compactification. The dilaton in heterotic string compactification can be another example. Eventually, this modulus should be stabilized by other means, e.g. nonperturbative dynamics [20]. It is expected that the resulting modulus mass m_T is tied to the scale of SUSY breaking, and thus

$$m_T \sim m_{3/2} \tag{2.3}$$

up to a little hierarchy of $\mathcal{O}(10-10^2)$.

In order to realize the weak scale SUSY, the gravitino mass $m_{3/2}$ is required to be smaller than $M_{\rm Pl}$ by many orders of magnitudes. For $M_{\rm st}\sim M_{\rm Pl}$, this is nontrivial to be achieved in flux compactification as generic flux configuration yields $\langle W_{\rm flux}\rangle = \mathcal{O}(1)$ (in the unit with $M_{\rm Pl}=1$) due to the quantization of flux. On the other hand, if SUSY is broken

[1] In fact, our analysis of flavor and CP constraints in section 3 applies also to the intermediate string scale scenario proposed in [15].

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by nonperturbative dynamics such as gaugino condensation [21], or warped dynamics [22], the resulting SUSY breaking scale is hierarchically lower than $M_{\rm Pl}$:

$$M_{\rm SUSY} \sim e^{-A} M_{\rm Pl}, \tag{2.4}$$

where e^{-A} is an exponentially small nonperturbative or warp factor. In 4D effective theory, the vacuum energy density at leading order is given by

$$V_{\rm vac} = M_{\rm SUSY}^4 - 3m_{3/2}^2 M_{\rm Pl}^2, \tag{2.5}$$

where $m_{3/2}/M_{\rm Pl} \sim \langle W_{\rm flux}\rangle$. As a result, in nonperturbative or warped SUSY breaking scenario, only a particular class of flux vacua with an exponentially small vacuum value of the flux-induced superpotential, i.e.

$$\langle W_{\rm flux}\rangle \sim e^{-2A}, \tag{2.6}$$

can have a (nearly) vanishing cosmological constant.

With the above observation, one can make the following assumptions to achieve a phenomenologically viable vacuum state with weak scale SUSY: (i) the underlying compactification involves a large number $N \gg 1$ of cycles each of which can carry a quantized flux in the range $[-L, L]$ for $L \gg 1$, which would allow a huge number of different flux configurations of $\mathcal{O}(L^N)$, (ii) such flux configurations provide a fine discretum of $\langle W_{\rm flux}\rangle$ varying from $\mathcal{O}(1)$ to a nearly vanishing value, (iii) SUSY is broken by nonperturbative or warped dynamics, yielding an exponentially small $M_{\rm SUSY}/M_{\rm Pl} \sim e^{-A} \sim 10^{-6} - 10^{-7}$. To be able to tune the vacuum energy density to the observed value $\sim (3\times 10^{-12}{\rm GeV})^4$, the spacing between different values of $\langle W_{\rm flux}\rangle$ should be as small as

$$\delta\langle W_{\rm flux}\rangle \lesssim \left(\frac{M_{\rm Pl}}{m_{3/2}}\right)\left(\frac{3\times 10^{-12}{\rm GeV}}{M_{\rm Pl}}\right)^4 \sim 10^{-104}. \tag{2.7}$$

Such extremely fine spacing might be achieved in flux compactification with $N \sim L = \mathcal{O}(100)$ as in the case of flux energy density discussed in [23].

Under the assumptions specified above, the fine tuning for vanishing cosmological constant selects a particular class of flux vacua with $\langle W_{\rm flux}\rangle \sim e^{-2A}$. For such vacua, still the moduli mass $m_U \sim \langle \partial^2 W_{\rm flux}/\partial U^2\rangle$ is generically of order unity due to the flux quantization. This results in a big moduli mass hierarchy:

$$\frac{m_T}{m_U} \sim \frac{\langle W_{\rm flux}\rangle}{\langle \partial^2 W_{\rm flux}/\partial U^2\rangle} \sim \frac{m_{3/2}}{M_{\rm Pl}} \sim e^{-2A}, \tag{2.8}$$

where again the little hierarchy factors of $\mathcal{O}(10-10^2)$ are ignored. It should be stressed that this moduli mass hierarchy is an outcome of the fine tuning of the cosmological constant and the assumed hierarchy (2.4) between the SUSY breaking scale and the Planck scale.

Generically, both the heavy moduli U and the light modulus T couple to SUSY breaking sector, therefore developing nonzero F-components. However, regardless of the details of SUSY breaking, the F-component of the flux stabilized U is given by

$$F^U \sim \frac{m_{3/2}^2}{m_U}, \tag{2.9}$$

which is negligibly small for m_U comparable to the string or GUT scale. (Note that moduli are normalized to be dimensionless, so their F components have a mass dimension one.) On the other hand, the light modulus T can develop a sizable F^T, e.g.

$$F^T \sim m_{3/2} \quad \text{or} \quad \frac{m_{3/2}}{\ln(M_{\rm Pl}/m_{3/2})}, \tag{2.10}$$

and therefore can be an important messenger of SUSY breaking [10, 15, 16].

2.2 4D effective action expanded in the inverse powers of the messenger modulus

Quite often, the messenger modulus T which is unfixed by flux has the following features: (a) $1/\mathrm{Re}(T)$ is proportional to certain powers of the string coupling $g_{\rm st}$ or the inverse of the compactification radius (in the unit with $\alpha' = 1$), thus its vacuum expectation value controls the $g_{\rm st}$ or α' expansion of the 4D action, (b) $\mathrm{Im}(T)$ is an axion whose non-linear PQ symmetry

$$U(1)_T: \ \mathrm{Im}(T) \to \mathrm{Im}(T) + \text{constant} \tag{2.11}$$

is respected at any finite order in the $g_{\rm st}$ and α' expansion. As a concrete example of such messenger modulus, one might consider the volume modulus and its RR axion partner in type IIB flux compactification or the dilaton-axion in heterotic flux compactification.

For such messenger modulus T, the couplings of moduli to the visible gauge and matter fields are given by

$$\int d^4\theta \; Y_{I\bar{J}}(T+T^*, U, U^*) Q^I Q^{J*} \\ + \left(\int d^2\theta \left[\frac{1}{4} f_a(T,U) W^{a\alpha} W^a_\alpha + \frac{1}{6}\lambda_{IJK}(U) Q^I Q^J Q^K \right] + \text{h.c.} \right), \tag{2.12}$$

where $W^{a\alpha}$ and Q^I denote the visible gauge and matter superfields, respectively. Here the matter kinetic function $Y_{I\bar{J}}$ is given by

$$Y_{I\bar{J}} = e^{-K_0/3} Z_{I\bar{J}} \tag{2.13}$$

for the Kähler potential

$$K = K_0 + Z_{I\bar{J}} Q^I Q^{J*}, \tag{2.14}$$

where K_0 is the moduli Kähler potential and $Z_{I\bar{J}}$ are the matter Kähler metrics. Expanding the 4D action in powers of $g_{\rm st}$ or α' while preserving the non-linear PQ symmetry $\mathrm{U}(1)_T$, the matter and gauge kinetic functions can be written as

$$\begin{aligned} Y_{I\bar{J}} &= (T+T^*)^{n_{I\bar{J}}} \Gamma_{I\bar{J}}(U, U^*) \left(1 - \frac{\Delta_{I\bar{J}}(U,U^*)}{[8\pi^2(T+T^*)]^{k_{I\bar{J}}}} + \cdots \right), \\ f_a &= k_a T + \frac{1}{8\pi^2} \Delta_a(U), \end{aligned}$$

where $n_{I\bar{J}}, k_{I\bar{J}}$, and k_a are all rational numbers. The successful unification of the MSSM gauge couplings at $M_{\rm GUT} \sim 2 \times 10^{16}$ GeV suggests that k_a are universal for the MSSM gauge group. In the following, we take the normalization of T for which $k_a = 1$, and thus

$$\langle {\rm Re}(T) \rangle \simeq \langle {\rm Re}(f_a) \rangle \simeq \frac{1}{g^2_{\rm GUT}}. \tag{2.15}$$

As was noticed before [10, 24, 25], if the MSSM chiral matter fields with same gauge charge originate from branes with same world volume dimension, the matter modular weights $n_{I\bar{J}}$ are automatically flavor universal (see appendix A for a more discussion of matter modular weights):

$$n_{I\bar{J}} = \text{flavor universal } n_I. \tag{2.16}$$

Also, in view of that T determines the 4D gauge coupling, it is expected that the messenger modulus expansion of 4D action is controlled by

$$\frac{1}{8\pi^2(T+T^*)} \sim \frac{\alpha_{\rm GUT}}{4\pi}, \tag{2.17}$$

and thus

$$k_{I\bar{J}} = 1, \quad \Delta_{I\bar{J}} = \mathcal{O}(1), \quad \Delta_a = \mathcal{O}(1). \tag{2.18}$$

In the following, we will assume this feature of the messenger modulus expansion, and examine its phenomenological consequences. Note that the non-linear PQ symmetry ${\rm U}(1)_T$ and the holomorphicity assure that λ_{IJK} are independent of T.

At leading order in the messenger modulus expansion, the non-linear PQ symmetry ${\rm U}(1)_T$ and the flavor universality of matter modular weights n_I assure that

$$\begin{aligned}
\frac{\partial}{\partial T}\ln(Y_{I\bar{J}}) &= \frac{n_I}{T+T^*} = \text{real and flavor universal,} \\
\frac{\partial}{\partial T}\ln(\lambda_{IJK}) &= 0, \\
\frac{\partial}{\partial T}\ln({\rm Re}(f_a)) &= \frac{k_a g_a^2}{2} = \text{real,}
\end{aligned} \tag{2.19}$$

with which the T-mediated SUSY breaking preserves flavor and CP [10, 15, 26]. On the other hand, $\frac{\partial}{\partial U}\ln(Y_{I\bar{J}})$ and $\frac{\partial}{\partial U}\ln(\lambda_{IJK})$ are flavor non-universal and complex, so the U-mediated SUSY breaking violates flavor and CP in general. However, as we will see shortly,

$$F^U \sim \frac{m^2_{3/2}}{m_U} \sim \frac{m^2_{3/2}}{M_{\rm Pl}} \tag{2.20}$$

regardless of the details of SUSY breaking, and thus the moduli mass hierarchy (2.8) assures that the U-mediated SUSY breaking is absolutely negligible. Still there might be a dangerous CP violation associated with the phase of Higgs μ and B parameters. Even for this, the non-linear PQ symmetry ${\rm U}(1)_T$ is useful as it allows the relative phase between F^T and $m_{3/2}$ to be rotated away. With real $F^T/m_{3/2}$, if μ is generated dominantly by

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the Chun-Kim-Nilles mechanism [27], or by the Giudice-Masiero mechanism [28], or by a singlet vacuum value as in the next to minimal supersymmetric standard model, the resulting Higgs mass parameters preserve CP [13].

One can now integrate out the heavy moduli U to derive the effective action of the visible fields and the light messenger modulus T. Let us start with the full 4D action which is generically given by

$$\int d^4\theta\, CC^*\, \Omega(U, U^*, \Phi, \Phi^*) + \left[\int d^2\theta C^3 \Big(W_{\text{flux}}(U) + \tilde{W}(U, \Phi)\Big) + \text{h.c.}\right], \tag{2.21}$$

where C is the chiral compensator superfield and Φ stands for all light superfields including the visible gauge and matter fields as well as the light modulus T. Here W_{flux} is the flux-induced superpotential depending only on U, and $\tilde{W}$ denotes the other part of superpotential which might include a $\mathrm{U}(1)_T$ breaking non-perturbative term, e.g.

$$\tilde{W} = A(U)e^{-aT} + \frac{1}{6}\lambda_{IJK}(U)Q^I Q^J Q^K. \tag{2.22}$$

To integrate out U, we note that flux quantization implies

$$M_U \equiv \frac{\partial^2 W_{\text{flux}}(U = U_0)}{\partial U^2} \sim M_{\text{st}}, \tag{2.23}$$

and the fine tuning of the cosmological constant in the presence of non-perturbative or warped SUSY breaking requires

$$W_{\text{flux}}(U = U_0) \sim e^{-2A}, \tag{2.24}$$

where $e^{-A} = M_{\text{SUSY}}/M_{\text{Pl}}$ is an exponentially small non-perturbative or warp factor and U_0 is the globally supersymmetric stationary point of the flux-induced superpotential:

$$\frac{\partial W_{\text{flux}}(U = U_0)}{\partial U} = 0. \tag{2.25}$$

Apparently the physical moduli mass m_U is dominated by the globally supersymmetric mass M_U in the limit when $M_U \gg m_{3/2}$, and U_0 and M_U are independent of light superfields.

The heavy moduli U can be integrated out by replacing U in the action (2.21) with the solution of the following superfield equation of motion:

$$\frac{1}{4}\bar{\mathcal{D}}^2 \left(CC^* \frac{\partial \Omega}{\partial U}\right) + C^3 \frac{\partial W}{\partial U} = 0, \tag{2.26}$$

where $\bar{\mathcal{D}}^2 = \bar{\mathcal{D}}^{\dot{\alpha}}\bar{\mathcal{D}}_{\dot{\alpha}}$ denotes the supercovariant derivative, $W = W_{\text{flux}} + \tilde{W}$ and all light fields Φ and also the compensator C are considered to be generic background superfields. In the limit with $m_{3/2}/M_U \sim e^{-2A} \ll 1$, the solution can be expanded in powers of $\bar{\mathcal{D}}^2/M_U$ and $\tilde{W}/M_U$ both of which are of the order of $m_{3/2}/M_U$. Note that $\partial^n \tilde{W}/\partial U^n \sim m_{3/2}$ for arbitrary $n \geq 0$ if the mass scale of the visible sector, e.g. the weak scale, is determined by SUSY breaking. In the perturbative expansion in powers of $\bar{\mathcal{D}}^2/M_U$ and $\tilde{W}/M_U$, the solution is given by

$$U = U_0 - \frac{1}{M_U}\left[\frac{1}{4}\bar{\mathcal{D}}^2\left(\frac{C^*}{C^2}\frac{\partial\Omega(U_0, U_0^*, \Phi, \Phi^*)}{\partial U}\right) + \frac{\partial\tilde{W}(U_0, \Phi)}{\partial U}\right] + \cdots, \tag{2.27}$$

where M_U is given in (2.23), and the ellipsis denotes higher order terms. One immediate consequence of this superfield solution is

$$F^U \sim \frac{m_{3/2}}{M_U}F^\Phi \sim \frac{m_{3/2}^2}{M_U}, \tag{2.28}$$

which assures that F^U is negligibly small compared to $F^\Phi \sim m_{3/2}$ when $M_U \sim M_{\rm st}$.

It is now obvious that, upon ignoring the small corrections suppressed by $m_{3/2}/M_U$, the low energy effective lagrangian can be obtained by replacing U in (2.21) with U_0. After this, one can make a proper redefinition of Q^I under which

$$\Gamma_{I\bar{J}}(U_0, U_0^*) \to \delta_{I\bar{J}}, \quad \Delta_{I\bar{J}}(U_0, U_0^*) \to \Delta_I\delta_{I\bar{J}}. \tag{2.29}$$

After such field redefinition, the effective couplings of the messenger modulus T to the visible gauge and matter fields are given by

$$\int d^4\theta\, Y_I Q^I Q^{I*} + \left(\int d^2\theta\left[\frac{1}{4}f_a W^a W^a + \frac{1}{6}\lambda_{IJK}Q^I Q^J Q^K\right] + \text{h.c.}\right) + \mathcal{O}\left(\frac{m_{3/2}}{M_U}\right), \tag{2.30}$$

where

$$\begin{aligned} Y_I &= (T+T^*)^{n_I}\left(1 - \frac{\Delta_I}{8\pi^2(T+T^*)}\right), \\ f_a &= k_a T + \frac{1}{8\pi^2}\Delta_a, \end{aligned} \tag{2.31}$$

where Δ_I and Δ_a are constants of order unity, and λ_{IJK} are constants with which the canonically normalized Yukawa couplings are determined as

$$y_{IJK} = \frac{\lambda_{IJK}}{\sqrt{Y_I Y_J Y_K}} \simeq \frac{\lambda_{IJK}}{(T+T^*)^{(n_I+n_J+n_K)/2}}. \tag{2.32}$$

The soft SUSY-breaking terms of canonically normalized sfermion fields $\tilde{Q}^I$ can be written as

$$\mathcal{L}_{\rm soft} = -\frac{1}{2}m_I^2|\tilde{Q}^I|^2 - \frac{1}{6}A_{IJK}y_{IJK}\tilde{Q}^I\tilde{Q}^J\tilde{Q}^K + \text{h.c.}, \tag{2.33}$$

which include the modulus mediated contribution [6] at $M_{\rm GUT}$ as

$$\begin{aligned} m_I^2 &= -F^T F^{\bar{T}}\partial_T\partial_{\bar{T}}\ln(Y_I) + \cdots \\ &= \left(n_I + \frac{g_{\rm GUT}^2}{8\pi^2}\Delta_I\right)M_0^2 + \cdots, \end{aligned}$$

$$
\begin{aligned}
A_{IJK} &= -F^T \partial_T \ln\left(\frac{\lambda_{IJK}}{Y_I Y_J Y_K}\right) + \cdots \\
&= \left[(n_I + n_J + n_K) + \frac{g^2_{\rm GUT}}{16\pi^2}\left(\Delta_I + \Delta_J + \Delta_K\right)\right] M_0 + \cdots\,,
\end{aligned} \tag{2.34}
$$

where

$$
M_0 = \frac{F^T}{T+T^*} \tag{2.35}
$$

corresponds to the modulus mediated contribution to the gaugino mass at $M_{\rm GUT}$, and the ellipses stand for the contribution from other mediation in the model.

The matter modular weights n_I are typically flavor universal, however there is no a priori reason for the higher order coefficients Δ_I to be flavor universal also. Even when Δ_I are flavor non-universal, there would not be any dangerous flavor or CP violation if sfermion masses are much heavier than 1 TeV, which actually happens for instance in the scheme proposed in [29]. In this paper, we are concerned with the possibility that modulus mediation including higher order effects satisfies the flavor and CP constraints with sparticle spectra in the sub-TeV range. To see this, we will examine in the next section the constraints on Δ_I imposed by low energy flavor and/or CP violating observables under the assumption that they are the dominant origin of non-minimal flavor or CP violation.

3. Constraints from flavor and/or CP violation

Let us first set up the notation. We start with the field basis for which the matter kinetic functions are diagonal as in the effective action (2.30). The MSSM matters and their $N=1$ superspace kinetic functions are denoted as

$$
\begin{aligned}
Q^I &= \{q_i, u_i^c, d_i^c, l_i, e_i^c\}, \\
Y_I &= \{Y_i^q, Y_i^u, Y_i^d, Y_i^l, Y_i^e\},
\end{aligned} \tag{3.1}
$$

where q_i ($i=1,2,3$) are the ${\rm SU(2)}_W$ doublet quarks, u_i^c and d_i^c are the ${\rm SU(2)}_W$ singlet anti-quarks, l_i are the ${\rm SU(2)}_W$ doublet leptons, e_i^c are the ${\rm SU(2)}_W$ singlet leptons, and the matter kinetic functions include higher order correction as

$$
Y_i^\phi = (T+T^*)^{n_\phi}\left(1 - \frac{\Delta_i^\phi}{8\pi^2(T+T^*)}\right) \qquad (\phi = q,u,d,l,e). \tag{3.2}
$$

Here we are interested in the flavor or CP violations associated with $\Delta_i^\phi - \Delta_j^\phi \neq 0$ for $i \neq j$ as the higher order correction to the gauge kinetic function, i.e. Δ_a of (2.31), obviously preserves flavor, and also does not give any CP violation.

Yukawa couplings and soft SUSY breaking terms of the canonically normalized MSSM matters at the *weak scale* are parameterized as

$$
\begin{aligned}
\mathcal{L}_{\rm Yukawa} &= y^u_{ij} H_u q_i u^c_j + y^d_{ij} H_d q_i d^c_j + y^e_{ij} H_d l_i e^c_j + \kappa^\nu_{ij} H_u l_i H_u l_j + {\rm h.c.}, \\
\mathcal{L}_{\rm soft} &= -\left(A^u_{ij} y^u_{ij} H_u \tilde{q}_i \tilde{u}^c_j + A^d_{ij} y^d_{ij} \tilde{q}_i H_d \tilde{d}^c_j + A^e_{ij} y^e_{ij} H_d \tilde{l}_i \tilde{e}^c_j + {\rm h.c.}\right) \\
&\quad - \left(m^{2(\tilde{q})}_{ij} \tilde{q}^*_i \tilde{q}_j + m^{2(\tilde{u})}_{ij} \tilde{u}^c_i \tilde{u}^{c*}_j + m^{2(\tilde{d})}_{ij} \tilde{d}^c_i \tilde{d}^{c*}_j + m^{2(\tilde{l})}_{ij} \tilde{l}^*_i \tilde{l}_j + m^{2(\tilde{e})}_{ij} \tilde{e}^c_i \tilde{e}^{c*}_j \right),
\end{aligned} \tag{3.3}
$$

where we include the $D = 5$ operator for neutrino masses in $\mathcal{L}_{\rm Yukawa}$. Soft parameters can be decomposed as

$$\begin{aligned} m_{ij}^{2(\tilde{\phi})} &= m_0^{2(\tilde{\phi})}\delta_{ij} + \Delta m_{ij}^{2(\tilde{\phi})} \qquad && (\tilde{\phi} = \tilde{q}, \tilde{u}, \tilde{d}, \tilde{l}, \tilde{e}), \\ A_{ij}^{\psi} &= A_0^{\psi} + \Delta A_{ij}^{\psi} && (\psi = u, d, e), \end{aligned} \tag{3.4}$$

where $m_0^{2(\tilde{\phi})}$ and A_0^{ψ} stand for flavor universal sfermion masses and A-parameters, respectively, while $\Delta m_{ij}^{2(\tilde{\phi})}$ and ΔA_{ij}^{ψ} represent flavor non-universal part. Depending upon the underlying SUSY breaking scheme, $m_0^{2(\tilde{\phi})}$ and A_0^{ψ} might receive contributions from various sources, e.g. modulus mediation, gauge mediation, anomaly mediation, renormalization group effect, e.t.c., whose relative importance will depend on the details of the model. Here we do not specify the full origin of the flavor universal $m_0^{2(\tilde{\phi})}$ and A_0^{ψ}, however the flavor non-universal part is assumed to be dominated by the modulus mediated contribution associated with non-universal Δ_i^{ϕ}:

$$\begin{aligned} \Delta m_{ij}^{2(\tilde{\phi})} &\simeq -\left[F^T F^{\bar{T}} \partial_T \partial_{\bar{T}} \ln\left(1 - \frac{\Delta_i^{\phi}}{8\pi^2(T+T^*)}\right)\right]\delta_{ij} \\ &\simeq \frac{g_{\rm GUT}^2}{8\pi^2}\Delta_i^{\phi} M_0^2 \delta_{ij} \qquad (\phi = q, u, d, l, e), \\ \Delta A_{ij}^{u} &\simeq F^T \partial_T \ln\left(1 - \frac{\Delta_i^{q}}{8\pi^2(T+T^*)}\right)\left(1 - \frac{\Delta_j^{u}}{8\pi^2(T+T^*)}\right) \\ &\simeq \frac{g_{\rm GUT}^2}{16\pi^2}(\Delta_i^q + \Delta_j^u) M_0, \\ \Delta A_{ij}^{d} &\simeq F^T \partial_T \ln\left(1 - \frac{\Delta_i^{q}}{8\pi^2(T+T^*)}\right)\left(1 - \frac{\Delta_j^{d}}{8\pi^2(T+T^*)}\right) \\ &\simeq \frac{g_{\rm GUT}^2}{16\pi^2}(\Delta_i^q + \Delta_j^d) M_0, \\ \Delta A_{ij}^{e} &\simeq F^T \partial_T \ln\left(1 - \frac{\Delta_i^{l}}{8\pi^2(T+T^*)}\right)\left(1 - \frac{\Delta_j^{e}}{8\pi^2(T+T^*)}\right) \\ &\simeq \frac{g_{\rm GUT}^2}{16\pi^2}(\Delta_i^l + \Delta_j^e) M_0, \end{aligned} \tag{3.5}$$

where

$$M_0 = \frac{F^T}{T + T^*} \tag{3.6}$$

corresponds to the modulus mediated contribution to the gaugino mass at $M_{\rm GUT}$. In fact, there are renormalization group (RG) corrections to the above non-universal part of soft parameters at the weak scale, which are mostly due to the 3rd generation Yukawa couplings. However such RG corrections can be safely ignored here as all the meaningful flavor and CP constraints on the modulus mediated SUSY breaking under consideration come from the first two generations for which the Yukawa induced RG corrections are negligibly small.

To examine the flavor and/or CP violating observables induced by $\Delta m_{ij}^{2(\tilde{\phi})}$ and ΔA_{ij}^{ψ}, it is convenient to use the super-CKM basis in which the quark and lepton mass matrices are diagonal [30]. Starting from the Yukawa coupling matrices y^{ψ}_{ij} ($\psi = u, d, e$) defined in the field basis for which the matter kinetic functions are diagonal, the super-CKM basis can be achieved by the unitary rotations of the matter superfields under which the Yukawa matrices become real and diagonal:

$$\begin{aligned} (V^{\psi}_L)^T y^{\psi} V^{\psi}_R &= \mathrm{Diag}(\hat{y}^{\psi}_1, \hat{y}^{\psi}_2, \hat{y}^{\psi}_3), \\ (V^{\nu}_L)^T \kappa^{\nu} V^{\nu}_L &= \mathrm{Diag}(\hat{\kappa}^{\nu}_1, \hat{\kappa}^{\nu}_2, \hat{\kappa}^{\nu}_3), \end{aligned} \tag{3.7}$$

where $V^{\psi}_{L,R}$ and V^{ν}_L are unitary matrices.

In supersymmetric limit, flavor and/or CP violations are all described by the CKM and PMNS mixing matrices given by

$$V_{\rm CKM} = V^{u\dagger}_L V^d_L, \quad V_{\rm PMNS} = V^{e\dagger}_L V^{\nu}_L. \tag{3.8}$$

However, in the presence of soft SUSY breaking terms, there can be further flavor and/or CP violations induced by non-universal $\Delta m^{2(\tilde{\phi})}_{ij}$ and ΔA^{ψ}_{ij}. Most of those non-minimal flavor violations can be described by the following mass-insertion parameters with $i \neq j$ [31, 32]:

$$\begin{aligned} (\delta^d_{LL})_{ij} &= \frac{(V^{d\dagger}_L \Delta m^{2(\tilde{q})} V^d_L)_{ij}}{m^2_{\tilde{q}}} \simeq \frac{g^2_{GUT}}{8\pi^2}\frac{M_0^2}{m^2_{\tilde{q}}}(\Delta^d_{LL})_{ij}, \\ (\delta^d_{RR})_{ij} &= \frac{(V^{dT}_R \Delta m^{2(\tilde{d})} V^{d*}_R)_{ij}}{m^2_{\tilde{q}}} \simeq \frac{g^2_{GUT}}{8\pi^2}\frac{M_0^2}{m^2_{\tilde{q}}}(\Delta^d_{RR})_{ij}, \\ (\delta^d_{LR})_{ij} &= \frac{(V^{dT}_L \Delta \mathcal{A}^d V^d_R)_{ij}\langle H_d\rangle}{m^2_{\tilde{q}}} \simeq \frac{g^2_{GUT}}{16\pi^2}\frac{M_0^2}{m^2_{\tilde{q}}}\left((\Delta^d_{LL})_{ij}\frac{m^d_j}{M_0} + \frac{m^d_i}{M_0}(\Delta^d_{RR})_{ij}\right), \\ (\delta^e_{LL})_{ij} &= \frac{(V^{e\dagger}_L \Delta m^{2(\tilde{l})} V^e_L)_{ij}}{m^2_{\tilde{l}}} \simeq \frac{g^2_{GUT}}{8\pi^2}\frac{M_0^2}{m^2_{\tilde{l}}}(\Delta^e_{LL})_{ij}, \\ (\delta^e_{RR})_{ij} &= \frac{(V^{eT}_R \Delta m^{2(\tilde{e})} V^{e*}_R)_{ij}}{m^2_{\tilde{l}}} \simeq \frac{g^2_{GUT}}{8\pi^2}\frac{M_0^2}{m^2_{\tilde{l}}}(\Delta^e_{RR})_{ij}, \\ (\delta^e_{LR})_{ij} &= \frac{(V^{eT}_L \Delta \mathcal{A}^e V^e_R)_{ij}\langle H_d\rangle}{m^2_{\tilde{l}}} \simeq \frac{g^2_{GUT}}{16\pi^2}\frac{M_0^2}{m^2_{\tilde{l}}}\left((\Delta^e_{LL})_{ij}\frac{m^e_j}{M_0} + \frac{m^e_i}{M_0}(\Delta^e_{RR})_{ij}\right), \end{aligned} \tag{3.9}$$

where $m_{\tilde{q}}$ and $m_{\tilde{l}}$ denote the average squark and slepton masses, m^d_i and m^e_i ($i = 1, 2, 3$) are the down-type quark and charged lepton mass eigenvalues, and

$$\begin{aligned} (\Delta^{d,e}_{LL})_{ij} &= \sum_k (V^{d,e}_L)^*_{ki}(V^{d,e}_L)_{kj}\Delta^{q,l}_k, \\ (\Delta^{d,e}_{RR})_{ij} &= \sum_k (V^{d,e}_R)_{ki}(V^{d,e}_R)^*_{kj}\Delta^{d,e}_k, \\ \Delta\mathcal{A}^{d,e}_{ij} &= y^{d,e}_{ij}\Delta A^{d,e}_{ij}. \end{aligned} \tag{3.10}$$

According to our assumption that the messenger modulus expansion is controlled by $1/8\pi^2\mathrm{Re}(T)$, all of the above mass-insertion parameters are suppressed by a factor of

$\mathcal{O}(g_{\rm GUT}^2/8\pi^2)$. In fact, the flavor changing mass-insertion parameters with $i \neq j$ can be further suppressed by small mixing angle in the unitary matrices $V^\psi_{L,R}$ ($\psi = u, d, e$). To see this, we note that the observed quark and charged lepton masses and the CKM mixing angles suggest that the Yukawa couplings take the form

$$y^u_{ij} \sim \epsilon^q_i \epsilon^u_i, \quad y^d_{ij} \sim \epsilon^q_i \epsilon^d_j, \quad y^e_{ij} \sim \epsilon^l_i \epsilon^e_j. \tag{3.11}$$

This form of Yukawa couplings can be naturally obtained by assuming either the localization of matter fields in extra dimension [33−35] or a spontaneously broken flavor symmetry [36]. In the scheme utilizing localization, different flavors with the same gauge charge are assumed to be localized at different positions in extra dimension, and then the flavor parameters ϵ^ϕ_i ($\phi = q, u, d, l, e$) determined by the wavefunction of matter fields show hierarchical pattern. Similar result can be obtained also in the scheme which assumes a broken flavor symmetry under which different flavors have different charges. In both schemes, the above form of Yukawa couplings is maintained even after the kinetic terms of matter fields are diagonalized. Note that neither localization nor flavor symmetry does provide a further suppression of Δ^ϕ_i in the matter kinetic functions.

The Yukawa couplings of (3.11) give rise to the mass hierarchy:

$$m^u_i/m^u_j \sim |\epsilon^q_i\epsilon^u_i|/|\epsilon^q_j\epsilon^u_j|, \quad m^d_i/m^d_j \sim |\epsilon^q_i\epsilon^d_i|/|\epsilon^q_j\epsilon^d_j|, \quad m^e_i/m^e_j \sim |\epsilon^l_i\epsilon^e_i|/|\epsilon^l_j\epsilon^e_j|, \tag{3.12}$$

and also the mixing angle pattern for $i \leq j$:

$$\begin{aligned} \left(V^{u,d}_L\right)_{ij} \sim \left(V^{u,d}_L\right)_{ji} \sim \epsilon^q_i/\epsilon^q_j, &\qquad \left(V^e_L\right)_{ij} \sim \left(V^e_L\right)_{ji} \sim \epsilon^l_i/\epsilon^l_j, \\ \left(V^\psi_R\right)_{ij} \sim \left(V^\psi_R\right)_{ji} \sim \epsilon^\psi_i/\epsilon^\psi_j, &\qquad (\psi = u, d, e), \end{aligned} \tag{3.13}$$

where we have assumed the normal hierarchy structure:

$$|\epsilon^\phi_1| \lesssim |\epsilon^\phi_2| \lesssim |\epsilon^\phi_3|. \tag{3.14}$$

This pattern of mixing angles implies for instance

$$\left|\left(V^d_L\right)_{12}\left(V^d_R\right)_{12}\right| \sim m_d/m_s, \quad \left|\left(V^e_L\right)_{12}\left(V^e_R\right)_{12}\right| \sim m_e/m_\mu. \tag{3.15}$$

Using the mass hierarchy (3.12) and the mixing angle pattern (3.13) together with

$$\begin{aligned} \sum_k V^*_{ki}V_{kj}\Delta_k &= \delta_{ij}\Delta_1 + V^*_{2i}V_{2j}(\Delta_2 - \Delta_1) + V^*_{3i}V_{3j}(\Delta_3 - \Delta_1) \\ &= \delta_{ij}\Delta_2 + V^*_{1i}V_{1j}(\Delta_1 - \Delta_2) + V^*_{3i}V_{3j}(\Delta_3 - \Delta_2), \end{aligned} \tag{3.16}$$

it is straightforward to find (for $i \neq j$)

$$\begin{aligned} (\delta^d_{LL})_{ij} &\sim \frac{g^2_{GUT}}{8\pi^2}\left(\frac{M_0^2}{m^2_{\tilde q}}\right)(\Delta^q_j - \Delta^q_i)(V^d_L)_{ij}, \\ (\delta^d_{RR})_{ij} &\sim \frac{g^2_{GUT}}{8\pi^2}\left(\frac{M_0^2}{m^2_{\tilde q}}\right)(\Delta^d_j - \Delta^d_i)(V^d_R)_{ij}, \end{aligned}$$

$$
\begin{aligned}
(\delta^e_{LL})_{ij} &\sim \frac{g^2_{GUT}}{8\pi^2}\left(\frac{M_0^2}{m^2_{\tilde{l}}}\right)(\Delta^l_j - \Delta^l_i)(V^e_L)_{ij}, \\
(\delta^e_{RR})_{ij} &\sim \frac{g^2_{GUT}}{8\pi^2}\left(\frac{M_0^2}{m^2_{\tilde{l}}}\right)(\Delta^e_j - \Delta^e_i)(V^e_R)_{ij}, \\
(\delta^{d,e}_{LR})_{ij} &\sim \frac{m^{d,e}_i}{2M_0}(\delta^{d,e}_{RR})_{ij} + \frac{m^{d,e}_j}{2M_0}(\delta^{d,e}_{LL})_{ij}.
\end{aligned}
\tag{3.17}
$$

Let us now consider the phenomenological constraints on the mass-insertion parameters. For the quark sector, the most stringent constraint comes from the CP violating K-$\bar{K}$ mixing parameter ϵ_K. Requiring that the SUSY contribution to ϵ_K should be less than the standard model value,[2] while assuming the gluino mass $m_{\tilde{g}} \sim m_{\tilde{q}}$, one finds [37]

$$
\begin{aligned}
\sqrt{\left|\,\mathrm{Im}\left[(\delta^d_{LL})_{12}(\delta^d_{RR})_{12}\right]\right|} &\lesssim 4\times 10^{-4}\left(\frac{m_{\tilde{q}}}{1\,\mathrm{TeV}}\right), \\
\sqrt{\left|\,\mathrm{Im}\left[(\delta^d_{LR,RL})^2_{12}\right]\right|} &\lesssim 8\times 10^{-4}\left(\frac{m_{\tilde{q}}}{1\,\mathrm{TeV}}\right).
\end{aligned}
\tag{3.18}
$$

For the mass-insertion parameters of (3.17), the second bound is easily satisfied, while the first bound leads to

$$
\frac{M_0^2}{m^2_{\tilde{q}}}\left(\frac{m_d}{m_s}\right)^{1/2}\sqrt{\left|(\Delta^q_2-\Delta^q_1)(\Delta^d_2-\Delta^d_1)\sin\eta_d\right|} \lesssim 7\times 10^{-2}\left(\frac{m_{\tilde{q}}}{1\,\mathrm{TeV}}\right), \tag{3.19}
$$

where η_d is a CP violating phase coming from the unitary rotation matrices, and we have used the relation $|(V^d_L)_{12}(V^d_R)_{12}| \sim m_d/m_s$. Due to the renormalization group evolution, the squark mass $m_{\tilde{q}}$ at the weak scale is typically bigger than the modulus mediated gaugino mass M_0 at M_{GUT}. Then, with the help from the small mixing angle $|(V^d_L)_{12}(V^d_R)_{12}| \sim m_d/m_s \sim 1/20$ and also an additional minor suppression by $M_0^2/m^2_{\tilde{q}} \sim 1/3$, the above bound can be satisfied even when $|\Delta^\phi_1 - \Delta^\phi_2| \sim 1$, $|\sin\eta_d| \sim 1$, and $m_{\tilde{q}} \sim 1\,\mathrm{TeV}$.

One might consider the $b \to s\gamma$ process to see if the higher order matter kinetic functions give rise to a contribution exceeding the current experimental bound. Requiring that the SUSY contribution to the branching ratio of $b \to s\gamma$ should be less than 10^{-5}, again with $m_{\tilde{g}} \sim m_{\tilde{q}}$, one finds [38]

$$
\left|(\delta^d_{LR})_{23}\right| \lesssim 5\times 10^{-3}\left(\frac{m_{\tilde{q}}}{1\,\mathrm{TeV}}\right), \tag{3.20}
$$

which is well satisfied by the mass-insertion parameters estimated in (3.17).

One might consider also the atomic and neutron electric dipole moments (EDMs) induced by the imaginary part of the diagonal LR mass-insertion parameter [39]. However,

[2] In view of that the CKM phase explains rather accurately all the observed CP violating phenomena including those of the B meson system, one might require a stronger condition that the SUSY contribution to ϵ_K should be less than about 10% of the standard model prediction. This would result in a factor of few stronger bound than (3.18), but does not change our conclusion.

in our case those LR parameters are given by

$$(\delta^d_{LR})_{ii} = \frac{(V_L^{dT}\Delta\mathcal{A}^d V_R^d)_{ii}\langle H_d\rangle}{m_{\tilde{q}}^2} \simeq \frac{g_{GUT}^2}{16\pi^2}\frac{M_0^2}{m_{\tilde{q}}^2}\frac{m_i^d}{M_0}\left[(\Delta^d_{LL})_{ii}+(\Delta^d_{RR})_{ii}\right],$$
$$(\delta^e_{LR})_{ii} = \frac{(V_L^{eT}\Delta\mathcal{A}^e V_R^e)_{ii}\langle H_d\rangle}{m_{\tilde{l}}^2} \simeq \frac{g_{GUT}^2}{16\pi^2}\frac{M_0^2}{m_{\tilde{l}}^2}\frac{m_i^e}{M_0}\left[(\Delta^e_{LL})_{ii}+(\Delta^e_{RR})_{ii}\right], \tag{3.21}$$

which are manifestly real. As a result, the atomic and neutron EDMs induced by higher order matter kinetic function are far below the current experimental limits.

In fact, the most stringent constraints on the modulus mediated SUSY breaking scheme come from the $\mu \to e\gamma$ process. Requiring that $\mathrm{Br}(\mu \to e\gamma) \le 1.2 \times 10^{-11}$, while assuming the Wino mass $m_{\tilde{W}} \sim m_{\tilde{l}}$ and the Higgsino mass $\mu \sim 2m_{\tilde{l}}$, one finds [40, 38]

$$\left|(\delta^e_{LL})_{12}\right| \lesssim \frac{7\times 10^{-3}}{\tan\beta}\left(\frac{m_{\tilde{l}}}{300\,\mathrm{GeV}}\right)^2,$$
$$\left|(\delta^e_{RR})_{12}\right| \lesssim \frac{2\times 10^{-2}}{\tan\beta}\left(\frac{m_{\tilde{l}}}{300\,\mathrm{GeV}}\right)^2,$$
$$\left|(\delta^e_{LR,RL})_{12}\right| \lesssim 6\times 10^{-6}\left(\frac{m_{\tilde{l}}}{300\,\mathrm{GeV}}\right). \tag{3.22}$$

For the mass-insertion parameters given by (3.17), the LR bound is easily satisfied. On the other hand, the LL and RR bounds lead to

$$\frac{M_0^2}{m_{\tilde{l}}^2}\left|\left(V_L^e\right)_{12}(\Delta_1^l-\Delta_2^l)\right| \lesssim \frac{1}{\tan\beta}\left(\frac{m_{\tilde{l}}}{300\,\mathrm{GeV}}\right)^2,$$
$$\frac{M_0^2}{m_{\tilde{l}}^2}\left|\left(V_R^e\right)_{12}(\Delta_1^e-\Delta_2^e)\right| \lesssim \frac{3}{\tan\beta}\left(\frac{m_{\tilde{l}}}{300\,\mathrm{GeV}}\right)^2. \tag{3.23}$$

It is reasonably expected that $M_0 \sim m_{\tilde{l}}$, and also the lepton mixing angles which affect $\mu \to e\gamma$ are related to the μ to e mass ratio as

$$\left|\left(V_L^e\right)_{12}\left(V_R^e\right)_{12}\right| \sim \frac{m_e}{m_\mu}. \tag{3.24}$$

If $\tan\beta \sim 1$, the above LL and RR bounds can be satisfied even when $m_{\tilde{l}} \sim 300\,\mathrm{GeV}$, $|\Delta_1^{l,e}-\Delta_2^{l,e}| \sim 1$, and $\left(V^e_{L,R}\right)_{12}$ have generic values satisfying the above relation. However, for large $\tan\beta$, the $\mu \to e\gamma$ bound requires a small $\left|\left(V_L^e\right)_{12}\right|$ unless $m_{\tilde{l}} \gg 300\,\mathrm{GeV}$ or $|\Delta_1^{l,e}-\Delta_2^{l,e}| \ll 1$. For the case with $m_{\tilde{l}} \sim 300\,\mathrm{GeV}$ and $|\Delta_1^{l,e}-\Delta_2^{l,e}| \sim 1$, which is actually the case of interest for us, the $\mu \to e\gamma$ bound can be satisfied with the following small mixing angle pattern as long as $\tan\beta \lesssim 30$:

$$\left|\left(V_L^e\right)_{12}\right| \sim \theta_C^n, \quad \left|\left(V_R^e\right)_{12}\right| \sim \theta_C^{3-n}, \quad (n=1,2), \tag{3.25}$$

where $\theta_C \sim 0.2$ is the Cabbibo angle. In this case, the large neutrino mixing angles in the PMNS matrix V_{PMNS} should originate from the unitary matrix V_L^ν diagonalizing the neutrino mass matrix as (3.7), and this might provide a nontrivial condition on the mechanism to generate the neutrino masses. It is interesting to note that this lepton mixing angle pattern allows a sizable SUSY contribution to the muon anomalous magnetic moment, [41 – 43] which is given by [44]

$$\frac{a_\mu^{\rm SUSY}}{1\times10^{-9}} \simeq \left(\frac{\tan\beta}{6}\right)\left(\frac{300{\rm GeV}}{m_{\tilde{l}}}\right)^2\left(\frac{\mu}{m_{\tilde{l}}}\right) \tag{3.26}$$

for a Wino mass $m_{\tilde{W}} \sim m_{\tilde{l}}$.

To summarize the flavor and CP constraints on moduli-mediated SUSY breaking in flux compactification, we find that most of the constraints other than those from ϵ_K and $\mu \to e\gamma$ are well satisfied even for generic form of higher order Kähler potential and sparticle spectra in the sub-TeV range, if the size of higher order Kähler potential in the messenger modulus expansion is of $\mathcal{O}(g_{\rm GUT}^2/8\pi^2)$. The constraints from ϵ_K and $\mu \to e\gamma$ can be satisfied also again for generic form of higher order Kähler potential and sparticle spectra in the sub-TeV range, if one makes a plausible assumption on flavor mixing angles motivated by the observed hierarchical structure of quark and charged lepton masses, for instance $|(V_L^d)_{12}(V_R^d)_{12}| \sim m_d/m_s$ and $|(V_L^e)_{12}(V_R^e)_{12}| \sim m_e/m_\mu$ with $|(V_L^e)_{12}| \lesssim 1/\tan\beta$.

4. Conclusion

Flux compactification can provide a SUSY breaking scheme in which soft terms preserve flavor and CP at leading order in the perturbative expansion controlled by the vacuum expectation value of the messenger modulus. In this paper, we have discussed some features of flux compactification leading to such SUSY breaking scheme, and examined the flavor and CP constraints on the higher order Kähler potential. It is found that all phenomenological constraints can be satisfied even for generic form of higher order Kähler potential and sparticle spectra in the sub-TeV range, under plausible assumptions on the size of higher order correction and flavor mixing angles. This implies that various SUSY breaking schemes involving such modulus mediation, e.g. mirage mediation and modulus-dominated mediation realized in flux compactification, can be free from the SUSY flavor and CP problems, while giving gaugino and sfermion masses in the sub-TeV range which can be probed by the LHC.

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A. Matter modular weights

In some class of compactification, the modular weights can be determined by a simple scaling argument combined with the non-linear PQ symmetry of the axion component [45,

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24, 25]. In our notation, the modular weight n_I is defined by the matter kinetic function as

$$Y_I \propto (T+T^*)^{n_I} \tag{A.1}$$

at leading order in the messenger modulus expansion. Note that

$$Y_I = e^{-K_0/3} Z_I, \tag{A.2}$$

where K_0 is the moduli Kähler potential and Z_I is the matter Kähler metric, i.e

$$K = K_0(T+T^*) + Z_I(T+T^*) Q^{I*} Q^I. \tag{A.3}$$

At leading order, e^{-K_0} and Z_I have a simple power-dependence on $\mathrm{Re}(T)$:

$$e^{-K_0} \propto (T+T^*)^{n_0}, \quad Z_I \propto (T+T^*)^{k_I}, \tag{A.4}$$

and then[3]

$$n_I = \frac{1}{3} n_0 + k_I. \tag{A.5}$$

Typically, the messenger modulus behaves (in the string unit with $\alpha' = 1$) as

$$\mathrm{Re}(T) \propto R^l / g_{\rm st}^n, \tag{A.6}$$

where $g_{\rm st}$ is the string coupling, R is the compactification radius, and l and n are (model-dependent) non-negative integers. The string coupling and compactification radius define another modulus $\propto R^{l'}/g_{\rm st}^{n'}$ which might be fixed by flux. Here, we consider two simple cases: the case (A) with $n' = 0$, in which R is stabilized by flux, while $g_{\rm st}$ remains unfixed, and another case (B) with $l' = 0$, in which $g_{\rm st}$ is stabilized by flux, while R remains unfixed. In case (A), the messenger modulus expansion can be identified as a string coupling expansion with $g_{\rm st}^n \propto 1/\mathrm{Re}(T)$. On the other hand, in case (B), the messenger modulus expansion can be identified as a radius expansion with $1/R^l \propto 1/\mathrm{Re}(T)$.

Case (A). Let us first examine the case that the messenger modulus expansion corresponds to a string coupling expansion with

$$\mathrm{Re}(T) \propto 1/g_{\rm st}^n, \tag{A.7}$$

where n is a positive integer. For this case, we assume that the kinetic terms of 4D gauge and matter fields and also the trilinear Yukawa couplings are generated at the same (leading) order in $g_{\rm st}$,and thus the $g_{\rm st}$-dependence of the 4D action is schematically given by

$$\mathcal{L} = \frac{1}{g_{\rm st}^N} \left[-\frac{1}{4} F^a_{\mu\nu} F^{a\mu\nu} + \partial_\mu \phi^{I*} \partial^\mu \phi^I + i\bar{\psi}^I \sigma^\mu \partial_\mu \psi^I + \left(\lambda_{IJK} \phi^I \psi^J \psi^K + \mathrm{h.c.} \right) \right], \tag{A.8}$$

[3] Often $-k_I$ is also called the matter modular weight.

where (ϕ^I, ψ^I) denote the chiral matter multiplets, N is a positive integer, and λ_{IJK} are independent of $g_{\rm st}$. We then have

$$g_{\rm GUT} \propto g_{\rm st}^{N/2}, \quad y_{IJK} \propto g_{\rm st}^{N/2}, \tag{A.9}$$

where $g_{\rm GUT}$ and y_{IJK} are the 4D gauge coupling and the canonically normalized Yukawa couplings, respectively.

In $N = 1$ superspace, this 4D action can be written as

$$\int d^4\theta\, Y_I Q^I Q^{I*} + \left[\int d^2\theta \left(\frac{1}{4} f_a W^a W^a + \frac{1}{6}\lambda_{IJK} Q^I Q^J Q^K\right) + \text{h.c.}\right], \tag{A.10}$$

where $Q^I = \phi^I + \theta\psi^I + \theta^2 F^I$. The non-linear PQ symmetry $\text{U}(1)_T$ of the axion component $\text{Im}(T)$ implies that the holomorphic Yukawa couplings λ_{IJK} are independent of T, while the gauge kinetic functions f_a are either linear in T or independent of T. Combining those constraints from $\text{U}(1)_T$ with

$$\frac{1}{g^2_{\rm GUT}} = \text{Re}(f_a), \quad y_{IJK} = \frac{\lambda_{IJK}}{\sqrt{Y_I Y_J Y_K}}, \tag{A.11}$$

one easily finds $N = n$, and

$$f_a \propto T, \quad Y_I \propto (T + T^*)^{n_I} \tag{A.12}$$

with

$$n_I + n_J + n_K = 1 \tag{A.13}$$

at leading order in the messenger modulus expansion. On the other hand, the universal $g_{\rm st}$-dependence of matter kinetic terms and Yukawa couplings suggests that the T-dependence of Y_I is universal also, so

$$n_I = 1/3. \tag{A.14}$$

To summarize, if the messenger modulus expansion corresponds to a string coupling expansion, and the gauge and matter kinetic terms and the trilinear Yukawa couplings are generated at the same (leading) order in this expansion, the matter modular weights have a universal value 1/3. One such example is the case that the messenger modulus corresponds to the heterotic dilaton, for which $n_0 = 1$ and $k_I = 0$ in (A.4), and thus $n_I = 1/3$.

Quite often, string compactification involves an anomalous $\text{U}(1)_A$ gauge symmetry [46] under which T transforms as

$$\text{U}(1)_A : T \to T - \frac{i}{2}\alpha(x)\delta_{\rm GS}, \tag{A.15}$$

where $\alpha(x)$ is the $\text{U}(1)_A$ transformation function and $\delta_{\rm GS}$ is the Green-Schwarz coefficient of $\mathcal{O}(1/8\pi^2)$. In the presence of such anomalous $\text{U}(1)_A$, the messenger modulus

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should be redefined as it mixes with the $\mathrm{U}(1)_A$ vector superfield V. This results in a shift of modular weight after the massive $\mathrm{U}(1)_A$ vector multiplet is integrated out, as will be discussed below.

Models with anomalous $\mathrm{U}(1)_A$ give a modulus-dependent Fayet-Iliopoulos (FI) D-term

$$\xi_{\mathrm{FI}} = \frac{1}{2}\delta_{\mathrm{GS}}\frac{\partial K_0}{\partial T}. \tag{A.16}$$

Then, to satisfy the D-flat condition, one needs a $\mathrm{U}(1)_A$-charged MSSM singlet X which has a large vacuum value $\langle X\rangle = \mathcal{O}(\xi_{\mathrm{FI}})$ to cancel this FI term.

Let us consider the 4D action including such field X:

$$\begin{aligned}\mathcal{L} = \int d^4\theta\ \big[&\Omega_0(T+T^*-\delta_{\mathrm{GS}}V) + Y_X(T+T^*-\delta_{\mathrm{GS}}V)X^*e^{-2V}X \\ &+Y_I(T+T^*-\delta_{\mathrm{GS}}V)Q^{I*}e^{2q_IV}Q^I\big]\,,\end{aligned} \tag{A.17}$$

where $\Omega_0 \equiv -3e^{-K_0/3}$ and the $\mathrm{U}(1)_A$ charge of X is normalized as $q_X = -1$. For $\delta_{\mathrm{GS}} = \mathcal{O}(1/8\pi^2)$, one can show [48] that the mass eigenstate vector superfield $\tilde{V}$ is given by

$$\tilde{V} \simeq V - \ln|X| \tag{A.18}$$

which has a superheavy mass $M^2_{\tilde{V}} \sim \delta_{\mathrm{GS}}M^2_{\mathrm{Pl}}$. It is straightforward to integrate out $\tilde{V}$ to obtain the effective action of the light modulus T and the visible matter fields Q^i:

$$\mathcal{L}_{\mathrm{eff}} = \int d^4\theta\ \Big[\Omega_0(T+T^*) + Y_I^{\mathrm{eff}}(T+T^*)Q^{I*}Q^I + \cdots\Big]\,, \tag{A.19}$$

where the ellipsis stands for the corrections suppressed by δ_{GS}, and the effective matter kinetic function is given by (after an appropriate redefinition of Q^I) [47, 48]

$$Y_I^{\mathrm{eff}} = \left(\frac{Y_X}{\partial_T\Omega_0}\right)^{q_I} Y_I. \tag{A.20}$$

After $\tilde{V}$ is integrated out, the effective modular weight is defined as

$$Y_I^{\mathrm{eff}} \propto (T+T^*)^{n_I^{\mathrm{eff}}} \tag{A.21}$$

at leading order in the messenger modulus expansion. In case when T corresponds to the heterotic dilaton, we have $\Omega_0, Y_X, Y_I \propto (T+T^*)^{1/3}$. The resulting effective modular weight is give by

$$n_I^{\mathrm{eff}} = \frac{1}{3} + q_I, \tag{A.22}$$

which would be flavor universal if the $\mathrm{U}(1)_A$ charges are flavor universal.

Case (B). Let us consider another case that the messenger modulus expansion corresponds to a radius expansion with

$$\mathrm{Re}(T) \propto R^l \quad (l > 0). \tag{A.23}$$

In this case, we can have more variety of possibilities.

Let us suppose that the gauge field A^a_μ propagates over l_G-dimensional internal space ($l_G > 0$), the matter field Q^I propagates over l_I-dimensional internal space, and the Yukawa coupling y_{IJK} originates from a wavefunction integral over l_{IJK}-dimensional internal space. Then, schematically, the 4D action takes the form:

$$\mathcal{L}=-\frac{1}{4}R^{l_G}F^a_{\mu\nu}F^{a\mu\nu}+R^{l_I}\left(\partial_\mu\phi^{I*}\partial^\mu\phi^I+i\bar{\psi}^I\sigma^\mu\partial_\mu\psi^I\right)+\left(R^{l_{IJK}}\lambda_{IJK}\phi^I\psi^J\psi^K+\text{h.c.}\right), \tag{A.24}$$

where

$$0 \leq l_I \leq l_G, \quad 0 \leq l_{IJK} \leq \min(l_I, l_J, l_K). \tag{A.25}$$

The resulting gauge and canonically normalized Yukawa couplings behave as

$$\begin{aligned} \frac{1}{g^2_{\rm GUT}} &= \mathrm{Re}(f_a) \,\propto\, R^{l_G}, \\ y_{IJK} &= \frac{\lambda_{IJK}}{\sqrt{Y_I Y_J Y_K}} \,\propto\, R^{l_{IJK}-\frac{l_I+l_J+l_K}{2}}. \end{aligned} \tag{A.26}$$

Again, with the non-linear PQ symmetry $\mathrm{U}(1)_T$ which requires f_a is either linear in T or independent of T, and λ_{IJK} are independent of T, these relations imply $l_G = l$, and

$$Y_I \propto (T + T^*)^{n_I}, \tag{A.27}$$

where n_I are constrained as

$$n_I + n_J + n_K = \frac{l_I + l_J + l_K - 2l_{IJK}}{l_G}. \tag{A.28}$$

For the MSSM matter fields, it is quite plausible that l_I and l_{IJK} are universal. Then, the resulting modular weights are universal and given by

$$n_I = \frac{l_I}{l_G} - \frac{2}{3}\frac{l_{IJK}}{l_G}. \tag{A.29}$$

One interesting point is that the modular weights have a universal value 1/3 as in the case (A) if all gauge and matter fields propagate over the same internal space and also the Yukawa couplings are given by wavefunction integrals over the same internal space, i.e. $l_G = l_I = l_{IJK}$. In models with an anomalous $\mathrm{U}(1)_A$, this modular weight is shifted by the $\mathrm{U}(1)_A$ charge as determined by (A.20).

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