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Diluting SUSY flavour problem on the Landscape

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ABSTRACT

We consider an explicit effective field theory example based on the Bousso-Polchinski framework with a large number N of hidden sectors contributing to supersymmetry breaking. Each contribution comes from four form quantized fluxes, multiplied by random couplings. The soft terms in the observable sector in this case become random variables, with mean values and standard deviations which are computable. We show that this setup naturally leads to a solution of the flavor problem in low-energy supersymmetry if N is sufficiently large. We investigate the consequences for flavor violating processes at low-energy and for dark matter.

1. Introduction and motivation

Supersymmetry breaking in MSSM (Minimal Supersymmetric Standard Model) is introduced in terms of explicit soft breaking terms. These are large in number $\sim$ 105, most of which violate flavor and CP symmetries. Phenomenologically there are strong constraints on the flavor off-diagonal entries, requiring them to be suppressed (compared to the flavour diagonal ones) by one to several orders of magnitude. The flavor violation constraints on the first two generations are significantly stronger compared to the ones involving the third generation. These bounds are well documented in the literature [1] (for reviews, see [2–4]).

To solve the problem with flavour violating soft terms, several solutions have been proposed. If supersymmetry breaking is mediated purely in terms of gauge interactions, the resulting soft terms would not contain any flavor violation [5–8]. However, the discovery of the Higgs boson puts constraints on such models. If the Higgs boson is of supersymmetric origin, one would expect the mass of the lightest CP even Higgs boson in Minimal Supersymmetric Standard Model (MSSM) to be rather light. This already puts severe constraints on the supersymmetric parameter space, in particular that of the third generation up-type squarks (the stops): they are either required to mix almost maximally or to be heavy, between 3 to 4 TeV (See for example, [9]). Several supersymmetry breaking models like minimal gauge mediation and its variations are disfavoured in the light of the Higgs discovery [10] or require a rather heavy spectra in the range of multi-TeV [9].

In string or supergravity based models, it has been long known that it in general scenarios it is hard to escape flavour violation, unless some specific conditions are chosen [11,12]. For example, if the Kähler potential of the matter fields is canonical and independent of the moduli/hidden sector fields, one could expect an universal, flavor-independent form for the soft terms as in minimal Supergravity. On the other hand, the problem can also be avoided if supersymmetry breaking is dominantly dilaton mediated [13]. Other solutions include decoupling of the first two generations [14] or imposing flavour symmetries (See for example, [15–18] and references there in).

In the present letter, we would like to address these issues from a different point of view, inspired by the landscape of string theory vacua. We will consider a large number N of sectors contributing to supersymmetry breaking. Large number of sequestered hidden sectors have also been considered recently in [19–21], in models with multiple (pseudo)goldstini. Other works which have addressed supersymmetric soft spectrum phenomenology from the landscape following [22] include [23–26]. In particular a solution to flavour and CP problems in the landscape through heavy first two generations was proposed in [27].

In the present work, we consider quantized four form fluxes à la Bousso-Polchinski [28]. Each sector contributes in a quantized way, with a quantum that will be taken to be below the electroweak scale. Due to the large number of contributions, the observable soft terms become random variables with Normal-type distributions around an average value. The setup has also the virtue of minimizing fine-tuning of the electroweak scale, due to the small contributions of each sector. We find that our setup can address at the same time the flavor problem of low-energy SUSY by generating FCNC effects proportional to standard deviation of

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soft terms from their mean value, which are parametrically suppressed as $1/\sqrt{N}$. By performing a RG analysis from high to low-energy, the setup makes also concrete predictions for low-energy flavor observables.

The letter is organized as follows. In Section 2, we review the Bousso-Polchinski setup and follow it in Section 3 with a review of four-forms fluxes. In the same section, we derive the soft terms and also show the impact on the flavour violating soft terms in the limit of large N. We also set up the boundary conditions for the scanning. In Section 4 we discuss the numerical results and present constraints from $K^0 - \bar{K}^0$ oscillations and $\mu \to e + \gamma$. The framework also has interesting implications for dark matter which we discuss briefly at the end of the section. We end with a small section of conclusions and outlook.

2. The Bousso-Polchinski setup

The logic and setup we put forward is mainly originating from Bousso and Polchinski approach to the cosmological constant [28]. We are regarding the implications of the string theory landscape for observable sector soft SUSY breaking terms. We will assume a large $N \gg 1$ number of SUSY-breaking sectors communicating through gravitational couplings to the Supersymmetric Standard Model (SSM). Such models could naturally appear in string theory, where there may be several independent sources of supersymmetry breaking.

Higher-dimensional operators and gravitational interactions lead to interactions between the moduli and the visible sector. One writes an effective action for the visible sector fields at a high scale, treating the hidden sector fields as non-dynamical background fields. This is justified if they are very heavy compared to the observable fields. One then writes a set of renormalization group equations for the higher-dimensional operators and evolves them into the infrared, ignoring the hidden sector dynamics. Supersymmetry breaking F and D components of the background hidden sector fields then give rise to visible sector soft supersymmetry breaking parameters.

Usually SUSY breaking is parametrized in terms of a single hidden sector field. This gives to a spectra at high scale (up to O(1) parameters) in terms of the scale of SUSY breaking F_α. In our case, we consider instead MSSM interacting with N hidden sectors at the Planck scale. In the models we are considering, auxiliary fields of the hidden sector fields contain quantized four-form fluxes, with discrete charges contributing to supersymmetry breaking, as in the Bousso-Polchinski solution to the cosmological constant problem.

The main features of our framework are:

- Integer quanta parameterising the soft supersymmetry breaking contribution from each hid- den sector
- Minimal number of parameters representing coupling between hidden sector fields and MSSM. Since these couplings depend on moduli vev's and interactions, we parametrize them by random continuous parameters taking values inside a compact interval around zero.
- Assume gravity mediation for simplicity. A similar scan can be done for gauge mediation, although the details will be quantitatively different.
- We consider a flat probability distribution of flux in each hidden sector. Since each flux is a random variable, due to the central limit theorem, this will lead to Normal-type distributions for the soft terms.

While we will impose the cancellation of the cosmological constant à la Bousso-Polchinski, we do not necessarily use our framework to address the cosmological constant problem. Instead we use the framework as a network of hidden sectors each contributing individually to soft supersymmetry breaking.

3. Four-forms and fluxes

Three form gauge potentials with (non-dynamical) four-form field strengths were considered longtime ago for addressing the cosmological constant problem [29–36], the gauge hierarchy problem [37–40] (see also [41]), the strong CP problem [42,43], inflation [44–47] and supersymmetry breaking [48]. On the other hand, it turned out to play an important role in the landscape of string theory compactifications [49–51] (for a recent review see e.g. [52]). Here we briefly review the main points of a theory containing three-forms with quantized form-forms field strengths.

Let us start from a lagrangian containing some scalar fields φ_i and three-form fields C^α_{mnp}, with the action

$$\mathcal{S}_0 = \int d^4x \{ -\frac{1}{2}(\partial\varphi_i)^2 - \Lambda_0 - \frac{1}{2\times 4!} F^{\alpha,2}_{mnpq} + \frac{1}{24} f_\alpha(\varphi_i)\epsilon^{mnpq} F^\alpha_{mnpq} \} , \tag{1}$$

where

$$F^\alpha_{mnpq} = \partial_m C^\alpha_{npq} + 3 \text{ perms.} . \tag{2}$$

For future convenience we define

$$F^\alpha = \frac{1}{4!}\epsilon^{mnpq} F^\alpha_{mnpq} \ , \ F^\alpha_{mnpq} = -\epsilon_{mnpq} F^\alpha \ . \tag{3}$$

The lagrangian (1) has actually to be supplemented with a boundary term

$$\mathcal{S}_b = \frac{1}{6}\int d^4x \, \partial_m \left(F^{mnpq}_\alpha C^\alpha_{npq} - f_\alpha(\varphi_i)\epsilon^{mnpq} C^\alpha_{npq} \right) \ . \tag{4}$$

The total action is

$$\mathcal{S} = \mathcal{S}_0 + \mathcal{S}_b = \int d^4x \{ -\frac{1}{2}(\partial\varphi_i)^2 - \Lambda_0 - \frac{1}{2\times 4!} F^{\alpha,2}_{mnpq} - \frac{1}{6}\epsilon^{mnpq}\partial_m f_\alpha(\varphi_i) \, C^\alpha_{npq} \} + \frac{1}{6}\int d^4x \, \partial_m \left(F^{mnpq}_\alpha C^\alpha_{npq} \right) \ . \tag{5}$$

A massless three-form gauge field in four spacetime dimensions has no on-shell degrees of freedom. As such, it can be integrated out via its field eqs.

$$\partial^m F^\alpha_{mnpq} \ = \ \epsilon_{mnpq}\partial^m f_\alpha(\varphi_i) \ , \tag{6}$$

whose solution is given by

$$F_\alpha = -f_\alpha(\varphi_i) + c_\alpha \ , \tag{7}$$

where c_α is a constant, which is to be interpreted as a flux. It was argued in [28] that c_α are quantized in units of the fundamental membrane coupling $c_\alpha = m_\alpha e$, fact that has important consequences for the landscape of string theory. After doing so, the final lagrangian takes the form

$$\mathcal{S} = \int d^4x \{ -\frac{1}{2}(\partial\varphi_i)^2 - \Lambda_0 - \frac{1}{2}\sum_\alpha (f_\alpha(\varphi_i) - c_\alpha)^2 \} \ . \tag{8}$$

The final resulting cosmological constant is therefore scanned by the flux

$$\Lambda = \Lambda_0 + \frac{1}{2}\sum_\alpha (f_\alpha(\varphi_i) - c_\alpha)^2 \ . \tag{9}$$

Notice that the boundary term $\mathcal{S}_b$ is crucial in obtaining the correct action. Ignoring it leads to the wrong sign of the last term in (8), fact that created confusion in the past.

3.1. Supersymmetric formulation

The embedding of four-form fluxes in supersymmetry and supergravity proceeds by introducing three-form multiplets, defined as the real superfields [53–59]

$$
\begin{aligned}
U_\alpha = \bar{U}_\alpha = & B_\alpha + i(\theta\chi_\alpha - \bar{\theta}\bar{\chi}_\alpha) + \theta^2 \bar{M}_\alpha + \bar{\theta}^2 M_\alpha \\
& + \frac{1}{3}\theta\sigma^m\bar{\theta}\epsilon_{mnpq}C_\alpha^{npq} + \theta^2\bar{\theta}(\sqrt{2}\bar{\lambda}_\alpha \\
& + \frac{1}{2}\bar{\sigma}^m\partial_m\chi_\alpha) + \bar{\theta}^2\theta(\sqrt{2}\lambda_\alpha - \frac{1}{2}\sigma^m\partial_m\bar{\chi}_\alpha) \\
& + \theta^2\bar{\theta}^2(D_\alpha - \frac{1}{4}\Box B_\alpha) \, .
\end{aligned} \tag{10}
$$

The difference between U_α and a regular vector superfield V is the replacement of the vector potential V_m by a three-form C_α^{npq}. In order to find correct kinetic terms, the analog of the chiral field strength superfield W_α for a vector multiplet is replaced by the chiral superfield [53]

$$
T_\alpha = -\frac{1}{4}\bar{D}^2 U_\alpha \quad , \quad T_\alpha(y^m,\theta) = M_\alpha + \sqrt{2}\theta\lambda_\alpha + \theta^2(D_\alpha + iF_\alpha) \, , \tag{11}
$$

with F_α defined as in (3). The definition (11) is invariant under the gauge transformation $U_\alpha \to U_\alpha - L_\alpha$, where L_α are linear multiplets. Correspondingly, lagrangians expressed as a function of T_α will have this gauge freedom. One can therefore choose a gauge in which $B_\alpha = \chi_\alpha = 0$ in (10) and the physical fields are complex scalars M_α and Weyl fermions λ_α.

Notice that for the purpose of finding the correct on-shell lagrangian and scalar potential, there is a simpler formulation in which T_α are treated as standard chiral superfields with $D_\alpha + iF_\alpha$ as auxiliary fields, no boundary terms are included, but the superpotential of the theory is changed according to [57–59]

$$
W(\phi_i, T_\alpha) \to W'(\phi_i, T_\alpha) \;=\; W(\phi_i, S_\alpha) + ic_\alpha T_\alpha \, , \tag{12}
$$

where c_α are the quantized fluxes. The linear terms in the superpotential shift linearly the auxiliary fields. In supergravity, the (F-term) scalar potential can be written as ($M_P = 1$ in what follows)

$$
V = K_{\alpha\bar{\beta}}F^\alpha F^{\bar{\beta}} - 3m_{3/2}^2 \quad , \quad \text{where} \quad F^\alpha = e^{\frac{K}{2}}K^{\alpha\bar{\beta}}\overline{D_\beta W} \, . \tag{13}
$$

The linear flux terms shift therefore the auxiliary fields according to

$$
F^\alpha = e^{\frac{K}{2}}K^{\alpha\bar{\beta}}\overline{D_\beta W'} = e^{\frac{K}{2}}K^{\alpha\bar{\beta}}\overline{D_\beta W} - ie^{\frac{K}{2}}K^{\alpha\bar{\gamma}}(\bar{c}_{\bar{\gamma}} + K_{\bar{\gamma}}\bar{T}_{\bar{\beta}}\bar{c}_{\bar{\beta}}) \, , \tag{14}
$$

leading to a scanning of the cosmological constant.

3.2. Soft terms in supergravity

We start from a supergravity lagrangian containing hidden sector (moduli) fields T_α, whose auxiliary fields contain the four-form fluxes we introduced previously, coupled to matter fields called Q_i in what follows. The Kahler potential and superpotential are defined by

$$
\begin{aligned}
K \;= & \hat{K}(T_\alpha,\bar{T}_\alpha) + K_{i\bar{j}}(T_\alpha,\bar{T}_\alpha)Q^i\bar{Q}^{\bar{j}} \\
& + \frac{1}{2}\left(Z_{ij}(T_\alpha,\bar{T}_\alpha)Q^iQ^j + \text{h.c.}\right) \, ,
\end{aligned}
$$

$$
\begin{aligned}
W = & \hat{W}(T_\alpha) + \frac{1}{2}\tilde{\mu}_{ij}(T_\alpha)Q^iQ^j \\
& + \frac{1}{3}\tilde{Y}_{ijk}(T_\alpha)Q^iQ^jQ^k + \cdots \, .
\end{aligned} \tag{15}
$$

The low-energy softly broken supersymmetric lagrangian is defined by the superpotential and soft scalar potential

$$
\begin{aligned}
W_{\text{eff}} = & \frac{1}{2}\mu_{ij}Q^iQ^j + \frac{1}{3}Y_{ijk}Q^iQ^jQ^k \, , \\
\mathcal{L}_{\text{soft}} = & -m_{i\bar{j}}^2 q^i q^{\bar{j}} - \left(\frac{1}{2}B_{ij}q^iq^j + \frac{1}{3}A_{ijk}q^iq^jq^k \right. \\
& \left. + \frac{1}{2}M_a\lambda_a\lambda_a + \text{h.c.}\right) \, ,
\end{aligned} \tag{16}
$$

where $Y_{ijk} = e^{K/2}\tilde{Y}_{ijk}$. After imposing the cancellation of the cosmological constant, the various soft terms and the supersymmetric masses are given by [11–13,60,61]

$$
\begin{aligned}
M_a = & \frac{1}{2}g_a^2 F^\alpha\partial_\alpha f_a \, , \\
m_{i\bar{j}}^2 = & m_{3/2}^2 K_{i\bar{j}} - F^\alpha F^{\bar{\beta}}R_{i\bar{j}\alpha\bar{\beta}} \quad , \quad \text{where} \\
R_{i\bar{j}\alpha\bar{\beta}} = & \partial_\alpha\partial_{\bar{\beta}}K_{i\bar{j}} - K^{m\bar{n}}\partial_\alpha K_{i\bar{n}}\partial_{\bar{\beta}}K_{m\bar{j}} \, , \\
A_{ijk} = & (m_{3/2} - F^\alpha\partial_\alpha\log m_{3/2})Y_{ijk} + F^\alpha\partial_\alpha Y_{ijk} - 3F^\alpha\Gamma_{\alpha(i}^l Y_{ljk)} \, ,
\end{aligned} \tag{17}
$$

where we have introduced also the Kahler connexion

$$
\Gamma_{IJ}^K = K^{K\bar{L}}\partial_I K_{J\bar{L}} \, . \tag{18}
$$

B_{ij} terms are not displayed since they will not be scanned in what follows. Similarly the μ term is determined by the radiative electroweak symmetry breaking conditions at the weak scale.

3.3. Scanning soft terms and gravitino mass

Taking into account the scanning of auxiliary fields from four-forms fluxes, in what follows we use the simplified scanning

$$
F_\alpha = m_\alpha\tilde{m}M_P \, , \tag{19}
$$

where m_α are integers and $\tilde{m}M_P$ is the quantum of scanning. Taking into account the cancellation of the cosmological constant, and setting the matter fields wavefunctions in a canonical form, the formulae for the scanning will be taken to be

$$
\begin{aligned}
m_{3/2} = & \tilde{m}(g_0 + \sum_\alpha g_\alpha m_\alpha) \quad , \\
m_{3/2}^2 = & \frac{1}{3}\sum_{\alpha=1}^N \frac{F_\alpha^2}{M_P^2} = \frac{1}{3}\tilde{m}^2\sum_\alpha m_\alpha^2 \, , \\
(m_0^2)_{i\bar{j}} = & m_{3/2}^2\delta_{i\bar{j}} + \tilde{m}^2\sum_\alpha d_{\alpha,i\bar{j}}m_\alpha^2 \, , \\
M_{1/2}^a = & \tilde{m}\sum_\alpha s_\alpha^a m_\alpha \, , \\
A_{ijk} = & m_{3/2}y_{ijk} + \tilde{m}\sum_\alpha a_{\alpha,ijk}m_\alpha \, ,
\end{aligned} \tag{20}
$$

where $\alpha = 1\cdots N$, $-M \le m_\alpha \le M$ (integers), $-d_0 \le a_\alpha, d_\alpha, s_\alpha, g_\alpha \le d_0$ (continuous). These last couplings are taking to be continuous in order to take into account the couplings of the hidden sector fields with the MSSM ones, which are dependent on the hidden sector vev's and interactions. Note that the soft terms defined above are in the so-called super CKM basis which is important for the flavor discussion below.

The scanning of the gravitino mass, combined with the cancellation of the cosmological constant (the Deser-Zumino relation) implies a constrained among the fluxes

$$g_0^2 + 2g_0 \sum_\alpha g_\alpha m_\alpha + \sum_{\alpha,\beta} g_\alpha g_\beta m_\alpha m_\beta = \frac{1}{3} \sum_\alpha m_\alpha^2 \,. \tag{21}$$

Taking the average value of (21) this implies in particular

$$\begin{aligned} g_0^2 &= \sum_\alpha \overline{(1/3 - g_\alpha^2) m_\alpha^2} \simeq \frac{N}{9}(1-d_0^2) M^2 \quad , \\ \overline{m_{3/2}} &= \tilde{m} g_0 \sim O(\sqrt{N}) \tilde{m} \,, \end{aligned} \tag{22}$$

where we have used the large flux limit $M \gg 1$. We can also compute

$$\overline{m_{3/2}^2} = \frac{\tilde{m}^2}{3} \sum_\alpha \overline{m_\alpha^2} \sim \frac{1}{9} N M^2 \tilde{m}^2 \,. \tag{23}$$

The mean values of the soft terms are therefore computed to be

$$\overline{(m_0^2)_{i\bar{j}}} = \overline{m_{3/2}^2} \delta_{i\bar{j}}, \qquad \overline{A_{ijk}} = \overline{m_{3/2}} y_{ijk}, \qquad \overline{M_{1/2}^a} = 0 \,. \tag{24}$$

There are two type of averages: one over the flux quanta m_α and the other over the (continuous) couplings d_α. Being independent variables, one can use formulae of the type

$$\begin{aligned} \overline{f_1(d_\alpha) f_2(m_\alpha)} = \overline{f_1(d_\alpha)} \times \overline{f_2(m_\alpha)} = \frac{1}{2d_0} \int_{-d_0}^{d_0} dx f_1(x) \\ \times \frac{1}{2M+1} \sum_{m_\alpha = -M}^{M} f_2(m_\alpha) \,. \end{aligned} \tag{25}$$

By using such formulae, one finds

$$\begin{aligned} (\Delta m_0^2)^2 &= (\Delta m_{3/2}^2)^2 + \tilde{m}^4 \sum_\alpha \overline{d_\alpha^2 m_\alpha^4} \simeq \frac{N M^4}{15} \tilde{m}^4 \left(\frac{4}{27} + d_0^2 \right) , \\ \text{where} \quad (\Delta m_{3/2}^2)^2 &= \frac{\tilde{m}^4}{9} \sum_\alpha [\overline{m_\alpha^4} - (\overline{m_\alpha^2})^2] \,. \end{aligned} \tag{26}$$

Consequently, one finds

$$(\delta_{ij})_{LL/RR} \equiv \frac{\delta m_0^2}{\overline{m_0^2}} \simeq \frac{1}{\sqrt{N}} \sqrt{\frac{1}{5}\left(4 + 27 d_0^2\right)} \,. \tag{27}$$

The off-diagonal entries, which have zero average values, are governed by the standard deviation δm_0^2. One concludes then that they are suppressed compared to the diagonal entries. For a large number of hidden sector $N \geq 10^6$, the flavor problem of MSSM is therefore solved. While this discussion is considering the flavor violating entries at the supergravity scale, in practice at the weak scale, as we will see in the next section, $N \sim 100$ would be sufficient to absolve strong constraints from Δm_K. For the constraint from $\mu \to e + \gamma$, however, $N \sim 100$ is not sufficient and a larger value of N should be chosen.

It should be noted however that the above discussion is pertaining to definition of δ_{ij} at the high scale. At the weak scale, for the leptonic sector (in the absence of right handed neutrinos), there is no significant change in the mean values, where as for the hadronic (squark) sector, due to the large gluino contributions to the squark masses in RG running, the $\delta_{ij}^{q,u,d}$ would be further suppressed by a factor from 7 up to an order of magnitude.

For the gaugino masses, one finds

$$\Delta M_{1/2}^2 = \tilde{m}^2 \sum_\alpha \overline{s_\alpha^2 m_\alpha^2} \simeq N \tilde{m}^2 \frac{d_0^2 M^2}{9} \,. \tag{28}$$

Therefore one finds the standard deviation

$$\Delta M_{1/2} = d_0 \sqrt{\overline{m_{3/2}^2}} \,. \tag{29}$$

For A-terms, let us consider for definiteness

$$A^u = m_{3/2} y_D^u + \tilde{m} \sum_\alpha a_\alpha^u m_\alpha \,, \tag{30}$$

where in the mass basis for fermions and scalar-fermion-gaugino couplings are diagonal, y_D^u is diagonal in the flavor space. If $a_\alpha^u \sim y^u \tilde{a}_\alpha$, then one expects the flavor violation in this case to be under control. However, if this is not the case, we can use the same arguments as above. One then finds the standard deviation

$$\begin{aligned} (\Delta A^u)^2 &= \overline{m_{3/2}^2} (y_D^u)^2 + \tilde{m}^2 \sum_\alpha \overline{(a_\alpha^u)^2 m_\alpha^2} - \left(\overline{m_{3/2}} y_D^u \right)^2 \\ &\simeq \frac{N M^2}{9} \tilde{m}^2 d_0^2 \left(1 + (y_D^u)^2 \right) \,. \end{aligned} \tag{31}$$

The A-terms are such that additional flavor violation (other than from Yukawa couplings) would be from the variance of the distribution. Thus we have

$$(\delta_{ij}^u)_{LR/RL} \equiv \frac{\Delta A^u v_u}{\overline{m_0^2}} \sim \frac{3 d_0 v_u}{\sqrt{N} \tilde{m} M} \,. \tag{32}$$

From the above it is clear that, similarly to the case of scalar masses, there is a suppression $1/\sqrt{N}$ coming from the large number of hidden-sector fields.

Notice that our starting expressions for soft terms (17) and the scanning we performed above is different compared to one based on a naive spurion-type parameterization of soft terms:

$$\sum_{\alpha=1}^N \frac{s_\alpha^a}{M_P} \int d^2\theta T_\alpha W^a W^a \ \to \ M_{1/2}^a = \frac{1}{M_P} \sum_{\alpha=1}^N s_\alpha^a F_{T_\alpha} \,, \tag{33}$$

$$\sum_{\alpha=1}^N \frac{d_{\alpha,ij}}{M_P^2} \int d^4\theta T_\alpha^\dagger T_\alpha Q_i^\dagger Q_j \ \to \ m_{\tilde{f}_{ij}}^2 = \frac{1}{M_P^2} \sum_{\alpha=1}^N d_{\alpha,ij} F_{T_\alpha}^\dagger F_{T_\alpha} \,, \tag{34}$$

$$\sum_{\alpha=1}^N \frac{a_{\alpha,ijk}}{M_P} \int d^2\theta T_\alpha Q_i Q_j Q_k \ \to \ A_{ijk} = \frac{1}{M_P} \sum_{\alpha=1}^N a_{\alpha,ijk} F_{T_\alpha} \,, \tag{35}$$

$$\sum_{\alpha=1}^N \frac{b_\alpha}{M_P^2} \int d^4\theta T_\alpha T_\alpha^\dagger H_u H_d \ \to \ B_{H_u H_d} = \frac{1}{M_P^2} \sum_{\alpha=1}^N b_\alpha F_{T_\alpha} F_{T_\alpha}^\dagger \,, \tag{36}$$

$$\sum_{\alpha=1}^N \frac{q_\alpha}{M_P} \int d^4\theta T_\alpha^\dagger H_u H_d \ \to \ \mu = \frac{1}{M_P} \sum_{\alpha=1}^N q_\alpha F_{T_\alpha}^\dagger \,. \tag{37}$$

The difference is that in the SUGRA expressions (17) the flavor-blind contributions proportional to $m_{3/2}^2$ scan coherently (add up) in soft terms, whereas the other contributions, which are similar to the global SUSY expressions (37), being multiplied by random couplings scanned around zero, average to zero. A similar scan we performed above, but starting from (37) would not lead a suppression of FCNC effects, unlike our scan above.

4. Numerical analysis

Using eqs. (20) as boundary conditions at the high scale, we perform a numerical analysis of the resulting soft spectrum at the weak scale and studied the phenomenology. For the numerical analysis, we have considered N to be 100, with m_α varying

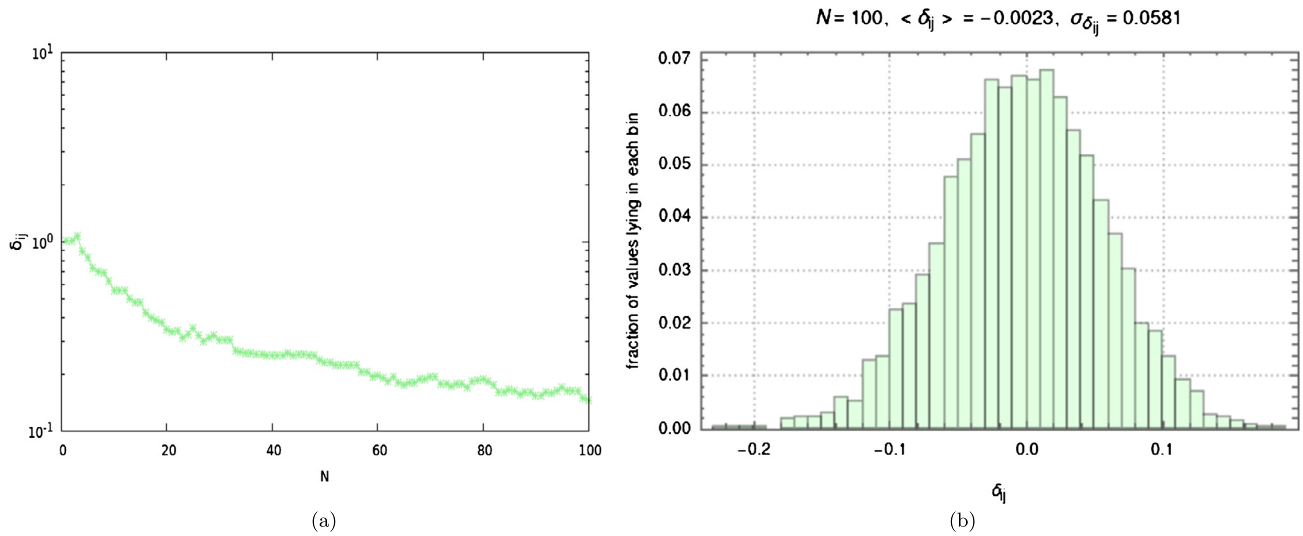


Fig. 1. A scatter plot showing the variation of maximum value of δ_{ij} of the type LL/RR with respect to number of hidden sectors, N at the high scale (left side). A histogram of the δ_{ij} is presented. As expected the mean is very close to zero and the variance is as computed in the text.

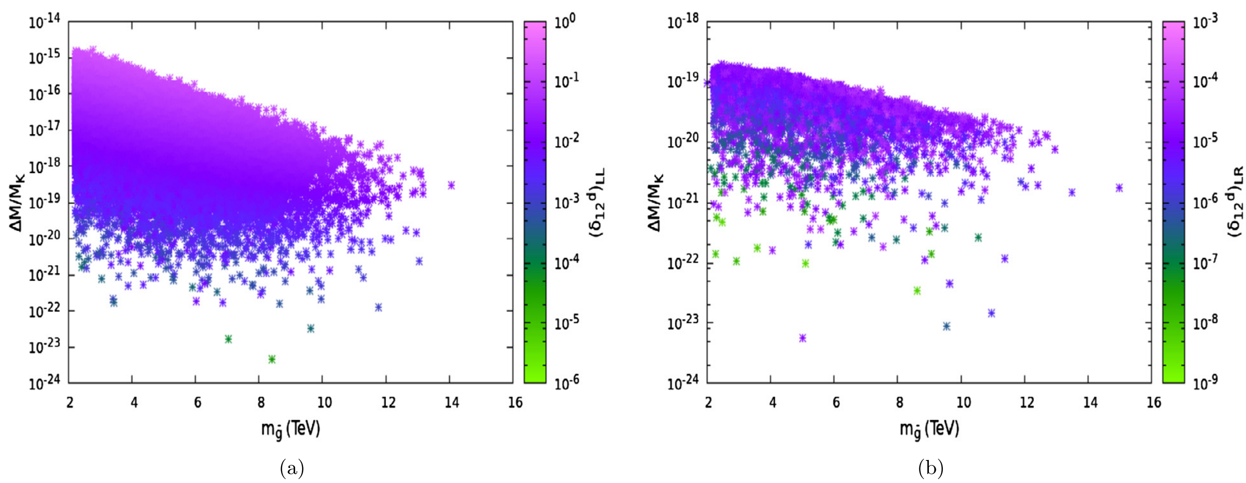


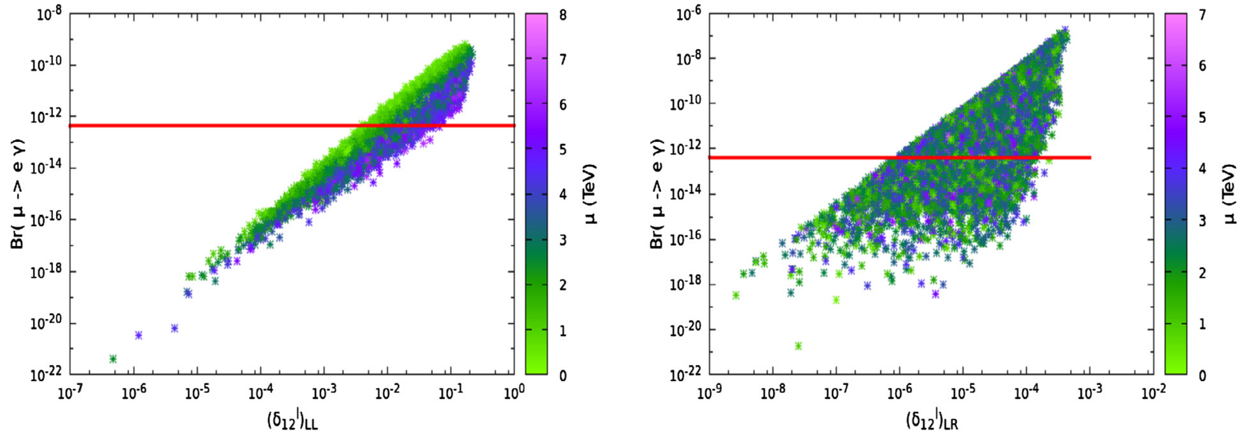
Fig. 2. Regions of the parameter space which satisfy the bounds from LHC, Higgs mass and other phenomenological bounds. We have chosen N to be 100 and $\tilde{m}$ to be 20 GeV. The distribution of δ values is as per the Eqs. (27) and (32). The left side plot is for LL type mass insertion whereas the right hand side is for LR type mass insertion. All points satisfy the experimental constraint from ΔM_K.

discretely and randomly from -100 to 100. We explored taking $\tilde{m}$ to be 20 GeV. The maximum value of $m_{3/2}$ is roughly about 6 TeV. The parameters $d_\alpha, s_\alpha, a_\alpha, g_\alpha$ are varied between {-1/4,1/4}. A larger value for the d_0 parameters would lead to significant number of the points ruled out due to tachyonic masses at the weak scale. We believe that larger values $d_0 \sim \mathcal{O}(1)$ would not significantly alter the results presented here. Finally we set $\tan\beta = 10$. We show that these values of N are enough to demonstrate the $1/\sqrt{N}$ suppression on the flavor violating off-diagonal entries. We use Suseflav [62] for computation of the spectrum and computing the flavor observables.

The variation of the off-diagonal flavor entries in the sfermion mass matrices is presented in Figs. 1(a) and 1(b). In Fig. 1(a) we present the scatter plot of a typical δ_{ij} as defined in eq. (27). From the plot it is clear that the δ_{ij} does fall off as $1/\sqrt{N}$. The second figure show the same data in terms of a histogram, where as we can see the mean value is close to zero and the variance is as expected from the formulae in eqs. (26, 27).

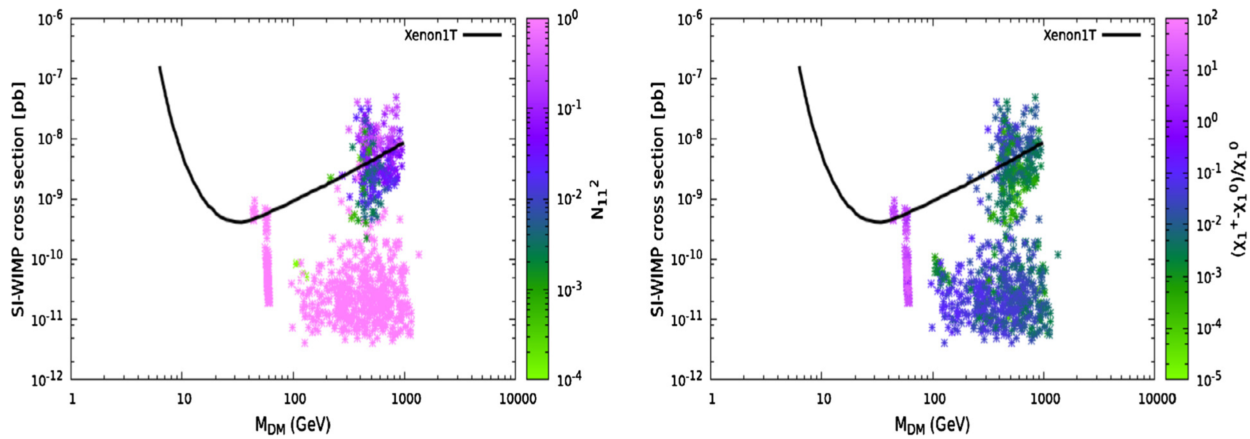
The high scale distributions are then evolved to the weak scale where the full soft supersymmetric spectrum is computed. Radiative electroweak symmetry breaking conditions are imposed. Experimental constraints from LHC and the Higgs mass are also taken in consideration. As is standard practice we consider one δ_{ij} at a time. In the present letter, we consider the two of the strongest constraints, *i.e.* the mass difference between the neutral K-mesons, ΔM_K and the leptonic rare decay $\mu \to e + \gamma$. A more detailed analysis with rest of the flavor processes will be presented elsewhere [63].

At the weak scale, the diagonal entries would be enhanced due to renormalisation group equation running, while the inter-generational entries of the squark matrices would only receive corrections suppressed by the product of Yukawa couplings and CKM angles [64]. Due to this the δ_{ij} would be further suppressed roughly by an additional factor which is proportional to the gluino mass corrections and roughly independent of the number of hidden sector fields. In Figs. 2(a), 2(b) we present the regions of the parameter space allowed by ΔM_k constraint as a function of the gluino mass. It should be noted here that we have taken the weak scale values of the mass insertions of eq. (27), where all the parameters appearing on the RHS are computed at the weak scale. The left figure is for the LL mass insertion where as the right figure is for the LR mass insertion. As can be seen from the figure, all the points lie below the experimentally measured value of ΔM_K [65]. The spectrum at the weak scale for the first two generations



(a) (b)

Fig. 3. Regions of the parameter space constrained by the leptonic rare decay $\mu \to e + \gamma$ for $(\delta_{12})_{LL}$ (on the left) and for $(\delta_{12})_{LR}$ (on the right). The horizontal (red) line is the present experimental limit from MEG experiment. As can be seen, a large part of the parameter space survives the experimental limit for N=100.



(a) (b)

Fig. 4. Regions of the parameter space which satisfy the relic density and the direct detection results from Xenon 1T. On the left we have shown the spin independent cross-section with respect to the lightest neutralino mass and the Bino component of the lightest neutralino. On the right, we show the same, with the mass difference between chargino and Bino.

is about 5-6 TeV and the gluino mass is shown in the figure after taking into consideration the limits from LHC. For this spectrum and a diluted $\delta \lesssim 10^{-1}$ the constraint from ΔM_K is satisfied.

The leptonic rare process $\mu \to e + \gamma$ is however more strongly constraining for the same set of parameters, *i.e.*, $N = 100$ and $\tilde{m} = 20$ GeV. In Figs. 3(a) and 3(b) we present results of the scanning as a function of the δ parameter and μ. The left figure is for a LL type mass insertion whereas the right figure is for RR type mass insertion. As one can see from the figures, a significantly large region of the parameter space is still compatible with the latest result from the MEG experiment [66], but the constraints from LR are significantly stronger, as expected. A larger N value $\gtrsim 10^5$ would lead to complete dilution of the δ.

Finally we have also looked for regions with neutralino dark matter which could lead to correct relic density while satisfying the constraints from direct detection and flavour. As can be seen from Fig. 4(a), there are two branches which satisfy relic density as well as the direct detection result. The first branch has dark matter masses $\lesssim$ 100 GeV and the lightest neutralino is a pure bino. In the second branch the neutralino has a region in which it is a pure bino and another region where there is a significant admixture from wino and higgsino. The regions where the neutralino are pure bino have significant co-annihilations with the chargino as can be seen from the Fig. 4(b). These regions arise due to the non-universality in the gaugino masses at the high scale due to the s_α parameters. On the other hand, regions with bino-higgsino mixing arise due to cancellations in $M_{1/2}$ in contributions from various fluxes of different spurion fields. As the charges/fluxes m_α take both signs, for significantly large N there is an enhanced probability of cancellations between the charges leading to small $M_{1/2}$ at the high scale. We numerically found that this probability is significantly high for $N \gtrsim 30$. Due to the universal nature of the gravitino mass, such cancellations do not occur in the soft scalar mass terms. A low value for μ is very probable in these regions leading to significant bino-higgsino mixing. Together they lead to regions with physically viable regions of neutralino dark matter. More details of these regions will be presented elsewhere [63].

5. Conclusions and outlook

We presented a novel solution to the supersymmetric flavor problem in the presence of large number of hidden sector (spurion) fields. Such a scenario naturally arises in the string landscape. The result does not depend on the explicit details of the string construction, but crucially on the form of the soft terms in the supergravity potential in the presence of a large number of hidden sector fields, eqs. (20). They naturally lead to a suppression of the flavour violating entries as $1/\sqrt{N}$. At the weak scale, there is further suppression due to the renormalisation group running, especially for the hadronic mass insertions. We have shown that numerically $N = 100$ is sufficient to remove the constraints from ΔM_K, whereas a much larger N would be required to eliminate completely the constraints on the leptonic sector from $\mu \to e + \gamma$. Conversely, a discovery of such leptonic processes in forthcoming experiments could be a smoking gun of such a scenario. The four fluxes contribution to the soft terms presented here provides an interesting framework to further study the implications for low energy phenomenology.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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References

[1] F. Gabbiani, E. Gabrielli, A. Masiero, L. Silvestrini, A complete analysis of FCNC and CP constraints in general SUSY extensions of the standard model, Nucl. Phys. B 477 (1996) 321–352.
[2] S. Dimopoulos, D.W. Sutter, The supersymmetric flavor problem, Nucl. Phys. B 452 (1995) 496–512.
[3] M. Misiak, S. Pokorski, J. Rosiek, Supersymmetry and FCNC effects, Adv. Ser. Dir. High Energy Phys. 15 (1998) 795–828. 795 (1997).
[4] A. Masiero, S.K. Vempati, O. Vives, Flavour physics and grand unification, in: Particle Physics Beyond the Standard Model. Proceedings, Summer School on Theoretical Physics, 84th Session, Les Houches, France, August 1–26, 2005, 2005, pp. 1–78.
[5] M. Dine, A.E. Nelson, Dynamical supersymmetry breaking at low-energies, Phys. Rev. D 48 (1993) 1277–1287.
[6] M. Dine, A.E. Nelson, Y. Shirman, Low-energy dynamical supersymmetry breaking simplified, Phys. Rev. D 51 (1995) 1362–1370.
[7] M. Dine, A.E. Nelson, Y. Nir, Y. Shirman, New tools for low-energy dynamical supersymmetry breaking, Phys. Rev. D 53 (1996) 2658–2669.
[8] G.F. Giudice, R. Rattazzi, Theories with gauge mediated supersymmetry breaking, Phys. Rep. 322 (1999) 419–499.
[9] E. Bagnaschi, et al., Supersymmetric models in light of improved Higgs mass calculations, Eur. Phys. J. C 79 (2) (2019) 149.
[10] P. Draper, P. Meade, M. Reece, D. Shih, Implications of a 125 GeV Higgs for the MSSM and low-scale SUSY breaking, Phys. Rev. D 85 (2012) 095007.
[11] A. Brignole, L.E. Ibanez, C. Munoz, Towards a theory of soft terms for the supersymmetric Standard Model, Nucl. Phys. B 422 (1994) 125–171, Erratum: Nucl. Phys. B 436 (1995) 747.
[12] A. Brignole, L.E. Ibanez, C. Munoz, Soft supersymmetry breaking terms from supergravity and superstring models, Adv. Ser. Dir. High Energy Phys. 18 (1998) 125–148.
[13] V.S. Kaplunovsky, J. Louis, Model independent analysis of soft terms in effective supergravity and in string theory, Phys. Lett. B 306 (1993) 269–275.
[14] A.G. Cohen, D.B. Kaplan, A.E. Nelson, The more minimal supersymmetric standard model, Phys. Lett. B 388 (1996) 588–598.
[15] S. Antusch, S.F. King, M. Malinsky, G.G. Ross, Solving the SUSY flavour and CP problems with non-Abelian family symmetry and supergravity, Phys. Lett. B 670 (2009) 383–389.
[16] A. Pomarol, D. Tommasini, Horizontal symmetries for the supersymmetric flavor problem, Nucl. Phys. B 466 (1996) 3–24.
[17] M. Dine, R.G. Leigh, A. Kagan, Flavor symmetries and the problem of squark degeneracy, Phys. Rev. D 48 (1993) 4269–4274.
[18] E. Dudas, S. Pokorski, C.A. Savoy, Soft scalar masses in supergravity with horizontal U(1)-x gauge symmetry, Phys. Lett. B 369 (1996) 255–261.
[19] C. Cheung, Y. Nomura, J. Thaler, Goldstini, J. High Energy Phys. 03 (2010) 073.
[20] C. Cheung, J. Mardon, Y. Nomura, J. Thaler, A definitive signal of multiple supersymmetry breaking, J. High Energy Phys. 07 (2010) 035.
[21] K. Benakli, C. Moura, Brane-worlds pseudo-goldstinos, Nucl. Phys. B 791 (2008) 125–163.
[22] F. Denef, M.R. Douglas, Distributions of flux vacua, J. High Energy Phys. 05 (2004) 072.
[23] H. Baer, V. Barger, D. Sengupta, Mirage Mediation from the Landscape, 2019.
[24] H. Baer, V. Barger, S. Salam, H. Serce, K. Sinha, LHC SUSY and WIMP dark matter searches confront the string theory landscape, J. High Energy Phys. 04 (2019) 043.
[25] H. Baer, V. Barger, H. Serce, K. Sinha, Higgs and superparticle mass predictions from the landscape, J. High Energy Phys. 03 (2018) 002.
[26] H. Baer, V. Barger, M. Savoy, H. Serce, X. Tata, Superparticle phenomenology from the natural mini-landscape, J. High Energy Phys. 06 (2017) 101.
[27] H. Baer, V. Barger, D. Sengupta, A landscape solution to the SUSY flavor and CP problems, Phys. Rev. Res. 1 (3) (2019) 033179, Phys. Rev. Res. 1 (2019) 033179.
[28] R. Bousso, J. Polchinski, Quantization of four form fluxes and dynamical neutralization of the cosmological constant, J. High Energy Phys. 06 (2000) 006.
[29] A. Aurilia, H. Nicolai, P.K. Townsend, Hidden constants: the theta parameter of QCD and the cosmological constant of N=8 supergravity, Nucl. Phys. B 176 (1980) 509–522.
[30] E. Witten, Fermion quantum numbers Kaluza-Klein theory, in: 1983. Conf. Proc. C8306011, 1983, p. 227.
[31] M. Henneaux, C. Teitelboim, The cosmological constant as a canonical variable, Phys. Lett. B 143 (1984) 415–420.
[32] J.D. Brown, C. Teitelboim, Dynamical neutralization of the cosmological constant, Phys. Lett. B 195 (1987) 177–182.
[33] J.D. Brown, C. Teitelboim, Neutralization of the cosmological constant by membrane creation, Nucl. Phys. B 297 (1988) 787–836.
[34] M.J. Duff, The cosmological constant is possibly zero, but the proof is probably wrong, Phys. Lett. B 226 (1989) 36, in: Conf. Proc. C8903131, 1989, p. 403.
[35] J.L. Feng, J. March-Russell, S. Sethi, F. Wilczek, Saltatory relaxation of the cosmological constant, Nucl. Phys. B 602 (2001) 307–328.
[36] G.R. Dvali, A. Vilenkin, Field theory models for variable cosmological constant, Phys. Rev. D 64 (2001) 063509.
[37] G. Dvali, A. Vilenkin, Cosmic attractors and gauge hierarchy, Phys. Rev. D 70 (2004) 063501.
[38] G. Dvali, Large hierarchies from attractor vacua, Phys. Rev. D 74 (2006) 025018.
[39] A. Herraez, L.E. Ibanez, An axion-induced SM/MSSM Higgs landscape and the weak gravity conjecture, J. High Energy Phys. 02 (2017) 109.
[40] G.F. Giudice, A. Kehagias, A. Riotto, The selfish Higgs, J. High Energy Phys. 10 (2019) 199.
[41] H.M. Lee, Relaxation of Higgs Mass and Cosmological Constant with Four-Form Fluxes and Reheating, 2019.
[42] G. Dvali, Three-Form Gauging of Axion Symmetries and Gravity, 2005.
[43] G. Dvali, A vacuum accumulation solution to the strong CP problem, Phys. Rev. D 74 (2006) 025019.
[44] N. Kaloper, L. Sorbo, A natural framework for chaotic inflation, Phys. Rev. Lett. 102 (2009) 121301.
[45] N. Kaloper, A. Lawrence, L. Sorbo, An ignoble approach to large field inflation, J. Cosmol. Astropart. Phys. 1103 (2011) 023.
[46] N. Kaloper, A. Lawrence, Natural chaotic inflation and ultraviolet sensitivity, Phys. Rev. D 90 (2) (2014) 023506.
[47] E. Dudas, Three-form multiplet and inflation, J. High Energy Phys. 12 (2014) 014.
[48] F. Farakos, A. Kehagias, D. Racco, A. Riotto, Scanning of the supersymmetry breaking scale and the gravitino mass in supergravity, J. High Energy Phys. 06 (2016) 120.
[49] F. Farakos, S. Lanza, L. Martucci, D. Sorokin, Three-forms in supergravity and flux compactifications, Eur. Phys. J. C 77 (9) (2017) 602.
[50] A. Herraez, L.E. Ibanez, F. Marchesano, G. Zoccarato, The type IIA flux potential, 4-forms and Freed-Witten anomalies, J. High Energy Phys. 09 (2018) 018.
[51] S. Lanza, F. Marchesano, L. Martucci, D. Sorokin, How many fluxes fit in an EFT?, J. High Energy Phys. 10 (2019) 110.
[52] S. Lanza, Exploring the Landscape of effective field theories, PhD thesis, 2019.
[53] S.J. Gates Jr., Super P form gauge superfields, Nucl. Phys. B 184 (1981) 381–390.
[54] C.P. Burgess, J.P. Derendinger, F. Quevedo, M. Quiros, Gaugino condensates and chiral linear duality: an effective Lagrangian analysis, Phys. Lett. B 348 (1995) 428–442.
[55] C.P. Burgess, J.P. Derendinger, F. Quevedo, M. Quiros, On gaugino condensation with field dependent gauge couplings, Ann. Phys. 250 (1996) 193–233.
[56] P. Binetruy, M.K. Gaillard, T.R. Taylor, Dynamical supersymmetric breaking and the linear multiplet, Nucl. Phys. B 455 (1995) 97–108.
[57] P. Binetruy, F. Pillon, G. Girardi, R. Grimm, The three form multiplet in supergravity, Nucl. Phys. B 477 (1996) 175–202.
[58] P. Binetruy, G. Girardi, R. Grimm, Supergravity couplings: a geometric formulation, Phys. Rep. 343 (2001) 255–462.
[59] K. Groh, J. Louis, J. Sommerfeld, Duality and couplings of 3-form-multiplets in N=1 supersymmetry, J. High Energy Phys. 05 (2013) 001.

[60] S. Ferrara, C. Kounnas, F. Zwirner, Mass formulae and natural hierarchy in string effective supergravities, Nucl. Phys. B 429 (1994) 589–625, Erratum: Nucl. Phys. B 433 (1995) 255.
[61] E. Dudas, S.K. Vempati, Large D-terms, hierarchical soft spectra and moduli stabilisation, Nucl. Phys. B 727 (2005) 139–162.
[62] D. Chowdhury, R. Garani, S.K. Vempati, SUSEFLAV: program for supersymmetric mass spectra with seesaw mechanism and rare lepton flavor violating decays, Comput. Phys. Commun. 184 (2013) 899–918.
[63] E. Dudas, P. Lamba, S.K. Vempati, in press.
[64] M. Ciuchini, A. Masiero, P. Paradisi, L. Silvestrini, S.K. Vempati, O. Vives, Soft SUSY breaking grand unification: leptons versus quarks on the flavor playground, Nucl. Phys. B 783 (2007) 112–142.
[65] M. Tanabashi, et al., Review of particle physics, Phys. Rev. D 98 (3) (2018) 030001.
[66] A.M. Baldini, et al., Search for the lepton flavour violating decay $\mu^+ \to e^+\gamma$ with the full dataset of the MEG experiment, Eur. Phys. J. C 76 (8) (2016) 434.