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Fluxed minimal supersymmetric standard model

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Recent developments in string compactifications in the presence of antisymmetric field backgrounds suggest a new simple and predictive structure for soft terms in the MSSM depending only on two parameters. They give rise to a positive definite scalar potential, a solution to the μ-problem, flavor universality and absence of a SUSY-*CP* problem.

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I. INTRODUCTION

The minimal supersymmetric standard model (MSSM) [1] is one of the most prominent candidates for an extension of the SM addressing the hierarchy problem. Gauge coupling unification is an important success in this scheme [2] and electroweak symmetry breaking is naturally induced as a consequence of SUSY breaking and a large t-quark mass [3]. The weakest point of the MSSM is the origin and structure of SUSY breaking. We know that after the dust settles one can parametrize our ignorance in terms of dim≤ 3 SUSY-breaking soft terms like gaugino and scalar masses. Before SUSY breaking the globally SUSY scalar potential has the form

$$\begin{aligned} V_{\rm SUSY} = &|-\mu H_d + h_U^{ij} Q_i U_j|^2 + |-\mu H_u + h_D^{ij} Q_i D_j \\ &+ h_L^{ij} L_i E_j|^2 + \sum_i (|h_U^{ij} U_j H_u + h_D^{ij} D_j H_d|^2 \\ &+ |h_U^{ij} Q_j H_u|^2 + |h_D^{ij} Q_j H_d|^2 + |h_L^{ij} L_j H_d|^2 \\ &+ |h_L^{ij} E_j H_d|^2) + V_{D-\rm terms} \end{aligned} \tag{1.1}$$

where μ is the SUSY Higgs mass and the Yukawa couplings are complex matrices in generation space. Q and U, D are the left and right-handed squarks, respectively, whereas L, R are left- and right-handed sleptons. Upon SUSY breaking the most general form of the SUSY breaking yields terms

$$L_g = \frac{1}{2}\sum_a M_a \lambda_a \lambda_a + \text{h.c.}, \tag{1.2}$$

$$\begin{aligned} L_{m^2} = &-m_{H_d}^2 |H_d|^2 - m_{H_u}^2 |H_u|^2 - m_{Q_{ij}}^2 Q_i Q_j^* - m_{U_{ij}}^2 U_i U_j^* \\ &- m_{D_{ij}}^2 D_i D_j^* - m_{L_{ij}}^2 L_i L_j^* - m_{E_{ij}}^2 E_i E_j^*, \\ L_{A,B} = &-A_{ij}^U Q_i U_j H_u - A_{ij}^D Q_i D_j H_d - A_{ij}^L L_i E_j H_d \\ &- B H_d H_u + \text{h.c.} \end{aligned}$$

It is well known that the most general form of soft terms has a variety of problems which include

(1) Lack of flavor universality (at least for the lightest generations) may induce too large flavor violating neutral currents (FCNC).

(2) Arbitrary complex phases for the A, B, μ and gaugino masses may lead to large *CP*-violating electric dipole moments [4].

(3) For arbitrary soft terms the scalar potential is unbounded below and may lead to $SU(2) \times U(1)$ breaking at the unification/string scale and/or charge- and color-breaking minima.

(4) The μ-problem [5]. We would like to understand why a SUSY mass parameter like μ turns out to be of the same order of magnitude as the SUSY-breaking mass terms.

The first problem is often solved by hand by postulating universal scalar masses and $A_{ij} = h_{ij} A$ parameters at the large scale. This is what is done, for example, in the popular mSUGRA scenario [6] in which there are just five parameters m, M, A, B, μ. In this scenario SUSY is broken in some hidden sector of the theory at a scale of order $M_{SB} \simeq 10^{10}$ GeV and it is transmitted to the observable sector of particle theory by supergravity interactions. The soft terms are then of order $\simeq M_{SB}^2/M_{\rm Planck} \simeq 10^2$ GeV. Concerning the second problem, again one can postulate that all soft terms are real or else that there is some (so far not very well motivated) phase alignment taking place. Finally, in order to get a stable scalar potential one usually restricts oneself to certain regions of soft parameter space. Concerning the μ problem, a natural mechanism in the context of $N =$ supergravity was proposed in Ref [7].

It is clear that in order to find a satisfactory solution for these problems we need a theory of supersymmetry breaking. It would be particularly interesting to have a well motivated underlying theory in which SUSY breaking takes place naturally and in which all the above problems are addressed.

Here we are interested in gravity mediated models [6,8] which naturally appear when combining $N = 1$ supersymmetry with gravitational interactions. In this case the scalar potential of the massless chiral fields is given by the general expression [9] (we are using here Planck mass units)

$$V = e^K (g^{i\bar{j}} (D_i W)(\bar{D}_{\bar{j}} \bar{W}) - 3|W|^2) + D\text{-term} \tag{1.3}$$

where the index run over all the chiral fields and K and W are the Kahler potential and superpotential of the theory. One also has for the Kahler covariant derivative

$$D_i W = \partial_i W + W K_i. \quad (1.4)$$

The general idea is that spontaneous supersymmetry breaking takes place in some hidden sector of the theory so that the gravitino gets a mass $m_{3/2} = \exp(K/2)W$. By taking the limit $M_{\text{Planck}} \to \infty$ while maintaining $m_{3/2}$ fixed one generically obtains bosonic SUSY-breaking soft terms. The gaugino masses are determined by the first derivative of the gauge kinetic function, i.e. f_a by

$$M_a = (2Ref_a)^{-1} F^i \partial_i f_a, \quad (1.5)$$

where the F^i is the vev of the auxiliary field corresponding to the chiral multiplet ϕ_i. The form of the so obtained soft terms thus depend on the structure of K, W and f_a, as well as on which chiral (hidden sector) fields are involved in the process of SUSY breaking ($F^i \neq 0$). Thus within this scheme a theory of soft terms correspond to a choice for these functions and a minimization of the potential leading to SUSY breaking.

A well motivated underlying theory would be string theory. From the very early days of string theory phenomenology attempts were made to understand the possible origin of SUSY-breaking soft terms [10]. It was also soon realized that the Kahler moduli T_i, dilaton S and complex-structure M_i fields which appear in string compactifications are natural candidates to constitute the hidden sector of the theory. For certain classes of heterotic compactifications [i.e. Abelian orbifolds and certain large volume limits of Calabi-Yau (CY) compactifications] it was possible to compute in perturbation theory the form of the functions K, W and f_a [11]. Two natural sources of SUSY breaking were put forward: gaugino condensation and a nonvanishing flux H_{ijk} in compact dimensions for the two index antisymmetric field B_{ij} appearing in the heterotic [10]. However the latter source did not look very promising since those fluxes are quantized and give rise to too large SUSY breaking of order the Planck scale. Gaugino condensation [12] may lead to hierarchically small SUSY breaking, however specific models had two generic problems: first, there were to many moduli/dilaton fields to be determined by the gaugino condensation potential; second, the vacuum energy at the minima of the scalar potentials was in general large and negative, leading to AdS space.

Another slightly more model independent proposal was made in order to be able to compute soft terms [13–16]. The idea was to assume that the source of SUSY breaking resides in the auxiliary fields of the dilaton/moduli fields, e.g. , F_S, F_{T_i}. Even without knowing of what could be the source of these nonvanishing auxiliary fields, knowledge of the Kahler potential and gauge kinetic function in some simple heterotic compactifications allowed (under the assumption of a vanishing cosmological constant) for the computation of soft terms as a function of the auxiliary fields. Two limits were particularly simple: in the case in which the auxiliary field of the overall volume field T dominates ($F_T \neq 0$) one gets a "no-scale structure" [17] leading to a leading order vanishing cosmological constant. This looks like a quite interesting starting point, however in that limit no soft terms whatsoever were generated (to leading order) [15]. In the dilaton domination case ($F_S \neq 0$) one obtains a set of appealing flavor-independent SUSY-breaking soft terms [14,15]. However no microscopic source for such $F_S \neq 0$ was found.

In the last nine years a number of developments have taken place which suggests to revisit these problems. First, it has been realized that type II and type I strings offer quite promising possibilities for the construction of string vacua close to the structure of the SM. A crucial ingredient in this new model building are Dp branes, nonperturbative configurations in string theory corresponding to $(p+1)$ dimensional subspaces inside the full 10-dimensional theory. The crucial property of Dp branes is that open strings are forced to have their boundaries on them. String excitations of open strings on the Dp branes give rise to massless gauge fields as well as fermions and scalars. Those fields are then to be identified with the fields of the SM. In fact a number of D-brane string configurations have been constructed using e.g. D branes at singularities [18] and/or intersecting D branes [19] with a massless spectrum remarkably close to the SM.[1]

A second new ingredient whose importance has only recently been realized is the role played by antisymmetric field fluxes in generic string compactifications [20–22] . The case of 3-form fluxes in type IIB CY (orientifold) compactifications has been studied with particular intensity in the last couple of years. It was realized in [21] that such kind of fluxes in type IIB orientifold theories give rise to a scalar potential which fixes both the dilaton and the complex-structure moduli M_i. Furthermore the hope exists that, when including nonperturbative effects depending on the volume moduli T_i all the moduli in these compactifications could be determined [22]. This would be an important result since the proliferation of undetermined scalar moduli vevs has been for many years one of the outstanding problems of string theory.

In a different development it has been recently shown [23–27] that fluxes of this type give also rise to SUSY-breaking soft terms on the world volume of $D3$ branes and $D7$ branes. In particular it was noted [24,25] that fluxes induce nonvanishing expectation values for the auxiliary fields of the dilaton S and/or moduli T_i, in this way making

[1]One can argue that the semirealistic perturbative heterotic models studied up to now are S-dual to orientifold type IIB compactifications with $D9$ branes. It is thus not surprising that considering more general configurations with different Dp branes lead to new model-building possibilities not previously envisaged in perturbative heterotic compactifications.

contact with the approach followed in Refs. [13–16] and providing a microscopic explanation for the vevs of the auxiliary fields. In particular in Ref. [26] certain classes of soft terms for matter fields in the world volume of $D7$ branes and with potential phenomenological interest have appeared. They are particularly interesting since, unlike other previous attempts to compute soft terms from string theory, they correspond to type IIB orientifold compactifications which solve the classical equations of motion. In the present paper we try to obtain general patterns of MSSM SUSY-breaking soft terms based on those recent results.

II. A BOTTOM-UP MOTIVATION

The results suggested by flux-induced SUSY breaking may be also motivated from a bottom-up approach. The first (FCNC) problem of the MSSM suggests to start with flavor-independent mass and trilinear terms for squarks and leptons. Let us consider now the second of the MSSM problems listed above which concerns complex phases in soft terms. In a universal setting complex phases may appear from μ, B, M and A parameters. Physical phases actually depend on the two linear combinations [1]

$$\phi_1 = \phi_\mu + \phi_A - \phi_B, \qquad \phi_2 = \phi_\mu + \phi_M - \phi_B, \tag{2.1}$$

where $\mu = |\mu|e^{i\phi_\mu}$, $M = |M|e^{i\phi_M}$, $A = |A|e^{i\phi_A}$ and $B = |B|e^{i\phi_B}$. For $\phi_{1,2}$ to vanish one needs to have

$$\phi_A = \phi_M; \qquad \phi_B = \phi_\mu + \phi_M. \tag{2.2}$$

The simplest way to achieve this is having soft terms related by

$$A = aM; \qquad B = bM\mu, \tag{2.3}$$

with a, b constant real parameters.

It is remarkable that there is a very simple modification of the SUSY scalar potential Eq. (1.1) which solves the first three problems listed above. This amounts to making the replacements[2]

$$W_{H_u} \longrightarrow W_{H_u} - a_u M^* H_u^*, \qquad W_{H_d} \longrightarrow W_{H_d} - a_d M^* H_d^*, \tag{2.4}$$

where W_i indicates derivative with respect to the ith scalar and a_u, a_d are real parameters. The superpotential here includes the bilinear $-\mu H_u H_d$ as usual. Note that by making these replacements one obtains a positive definite scalar potential with soft terms

$$\begin{aligned} A_u &= -a_u M; \; A_d = A_L = -a_d M, \\ m_{H_u}^2 &= a_u^2 |M|^2; \; m_{H_d}^2 = a_d^2 |M|^2; \; m_{\tilde{f}}^2 = 0, \\ B &= (a_u + a_d) M\mu, \end{aligned} \tag{2.5}$$

where $m_{\tilde{f}}^2$ are the masses of scalar partners of quarks and leptons. In addition all phases in soft terms may be rotated away. Particularly simple boundary conditions are obtained in the case with $a_u = a_d = 1$. In this situation all soft terms are determined by a couple of parameters M and μ:

$$m_{H_u}^2 = m_{H_d}^2 = |M|^2; m_{\tilde{f}}^2 = 0, \quad A = -M, \quad B = 2M\mu. \tag{2.6}$$

Note that in principle one could make a similar substitution for the rest of the chiral fields of the MSSM

$$V_{\rm SUSY} = \sum_i |\partial_i W|^2 \longrightarrow V_{SB} = \sum_i |\partial_i W - a_i M^* \phi_i^*|^2 \tag{2.7}$$

with $i = H_u, H_d, \tilde{Q}, \tilde{U}_R, \tilde{D}_R, \tilde{L}, \tilde{E}_R$. Such a procedure would give rise to universal mass terms for all squarks/sleptons and Higgs fields, as well as trilinears. Indeed we will see below that flux-induced SUSY breaking suggests to make such universal replacement with all $a_i = 1$.

III. FLUXES, D-BRANES AND SUSY-BREAKING SOFT TERMS

The kind of string context that we are going to work on here is that of type IIB orientifold compactifications. This is one of the simplest contexts in which in the last few years a number of chiral models with a particle content close to the SM have been constructed [18,19]. In these theories one compactifies type IIB strings on a CY manifold (or a toroidal orbifold) and further modes out the theory by an order-2 twist which includes the Ω operation which corresponds to world-sheet parity. Consistency of the compactification (RR tadpole cancellation conditions) requires the presence of some particular set of Dp branes with $p = 3, 5, 7, 9$ in the setting. Depending on the particular form of the orientifold operation one type or other of Dp branes is required. If we want to preserve one unbroken $D = 4$ SUSY either $D3, D7$ or alternatively $D9, D5$ sets may be added. In the case in which the orientifold operation is just Ω only $D9$ branes are required and the result is just a standard compactification of type I string theory. Since the latter is known to be S dual to the Heterotic string, the effective actions are quite similar and the phenomenology also is, so one does not expect to obtain results very different from those previously found in heterotic compactifications. We will rather focus in the case of orientifolds in which $D7$ branes (and possibly additional $D3$ branes) are present. The world volume of $D7$ branes is 8-dimensional and it is supposed to include Minkowski space

[2] As we will see below, from the $N = 1$ supergravity point of view this replacement will correspond to going from the SUSY auxiliary fields to the SUGRA ones.

and a 4-cycle inside the compact CY manifold. If the position of the $D7$ branes in the transverse dimensions is sitting on a smooth point of the CY, there appears an $N = 4$ Yang-Mills theory in the effective 4-dimensional Lagrangian. If $D7$ branes sit on top of some (e.g., orbifold) singularity the symmetry is reduced and one may get chiral $N = 1$ theories (see e.g. Ref. [18] for a description of this type of models) of phenomenological interest.[3]

As we mentioned, in the last few years it has been realized the importance of fluxes of antisymmetric fields on the structure of type IIB orientifold compactifications [20–22]. Ten-dimensional type IIB theory has a couple of antisymmetric tensors B_{NM} and A_{NM} coming, respectively, from the so called NS and RR sectors of the theory. They can have (quantized) fluxes H_{ijk}, F_{ijk} along the compact complex dimensions $i, j, k = 1, 2, 3$

$$\frac{1}{2\pi\alpha'}\int_{C_3} F_3 \in 2\pi ZZ; \qquad \frac{1}{2\pi\alpha'}\int_{C_3} H_3 \in 2\pi ZZ, \tag{3.1}$$

where C_3 is any 3-cycle inside the CY. In SUSY compactifications it is actually the complex flux combination

$$G_3 = F_3 - iSH_3 \tag{3.2}$$

which naturally appears. Here S is the complex axidilaton field. As long as the supergravity equations of motion are obeyed this is a degree of freedom which is generically there and should be considered.

It was realized in [21] that such type of fluxes give rise to a scalar potential which fixes the vev of the dilaton and all complex-structure (shape) moduli. Specifically, G_3 backgrounds of a certain class (i.e. imaginary self-dual fluxes, ISD[4]) solve the equations of motion with a vanishing c.c. to leading order. The origin of this dynamics is the generation of a flux-induced superpotential of the form [29,30]

$$W = \kappa_{10}^{-2}\int_{M_6} G_{(3)} \wedge \Omega, \tag{3.3}$$

where $\kappa_{10}^2 = \frac{1}{2}(2\pi)^7\alpha'^4$ is the $D = 10$ gravitational constant and Ω the Calabi-Yau holomorphic 3-form (see e.g. Ref. [21] for details). This superpotential depends on the dilaton complex field S and the complex-structure moduli (through Ω) but not on the kahler moduli (T_i fields). It can be shown that upon minimization of the resulting $D = 4$ scalar potential the dilaton S and complex-structure moduli fields are fixed with a vanishing c.c. (to leading order in both the string coupling constant g_s and the inverse string tension α').

Furthermore, in [22] it was pointed out that, when combined with nonperturbative effects on the gauge couplings (like e.g., gaugino condensation) fluxes may potentially lead to a determination also of all the T-like volume moduli. As we said, if true this would be important progress, since fixing the dilaton moduli and complex-structure fields in string theory has always been one of the most outstanding problems.

As we mentioned it has also been recently shown that fluxes of this type give rise to SUSY-breaking soft terms on the world volume of $D3$ branes and $D7$ branes. In particular in Ref. [24] it was realized that certain particular choices of G_3 backgrounds give rise to soft terms corresponding to the dilaton dominance or modulus-dominance limits discussed in the heterotic literature [16]. This is important since it provides for a microscopic explanation of nonvanishing auxiliary fields for dilaton and moduli. The obtained soft terms are proportional to the flux densities $G_{(3)}$ which have a dependence for large radius $G_3 = f\frac{\alpha'}{R^3}$, with f an R-independent constant measuring the amount of quantized flux. Thus one typically obtains SUSY-breaking terms of order [24]

$$m_{\rm soft} = \frac{g_s^{1/2}}{\sqrt{2}} G_{(3)} = \frac{f g_s^{1/2}}{\sqrt{2}}\frac{\alpha'}{R^3} = \frac{fM_s^2}{M_p} \tag{3.4}$$

with M_s the string scale and $M_p = M_s^4 R^3$ the Planck scale. Thus a way to get soft terms of order the electroweak scale is having the string scale at an intermediate scale $M_s \simeq 10^{10}$ GeV. However it would be consistent to have a high string scale with $M_s \simeq M_p$ if the factor f in Eq. (3.4) is sufficiently small, i.e., if the local flux in the brane position is for some reason diluted. That is, for example, the case in the presence of a large warping suppression [21,22]. In what follows we will not deal with these issues but assume that the resulting soft terms are of order the electroweak scale, as phenomenologically required.

It turns out that the kind of fluxes which solve the type IIB equations of motion (i.e., ISD ones) do not lead to any soft terms to leading order for the fields on the world volume of $D3$ branes. Thus if we try to embed the MSSM on $D3$ branes we find no soft terms at all. From the effective field theory point of view this happens because the ISD fluxes considered correspond to "modulus-dominance" SUSY breaking which has a no-scale structure leading to no soft terms to leading order [17].

The prospects change completely if one considers the embedding of the SM inside $D7$ branes. As we said, this is a natural thing to do in the context of type IIB F-theory compactifications As recently emphasized in Ref. [26] ISD fluxes do give rise to interesting SUSY-breaking soft terms for the fields on $D7$ branes. We are not giving any details here but we can summarize the results and later we will motivate it from the effective field theory point of view.

[3]More general type IIB compactifications of this type are more efficiently described in terms of F-theory compactifications on CY 4-folds [28].

[4]Imaginary self-dual fluxes verify $G_{(3)} = -i *_6 G_{(3)}$, where $*_6$ means Hodge dual in the compact six dimensions.

A stack of $D7$ branes gives rise at low energies to charged chiral multiplets ϕ_i upon KK compactification. A large class of those admit a geometric interpretation in the sense that vevs for them parametrize the position of $D7$ brane in transverse space inside the CY manifold. What we are going to discuss now refers to that particular class of $D7$-brane charged fields.[5]

Two types of ISD G_{mnp} flux densities (which we take for simplicity to be constant over the CY) are particularly relevant. The first of them corresponds to $(0, 3)$ forms (e.g., a tensorial structure $G_{\bar{1}\bar{2}\bar{3}}$ in tori) and gives rise to SUSY-breaking soft terms. The second corresponds to $(2, 1)$ forms (e.g., a tensorial structure $G_{12\bar{3}}$ in the toroidal case) and does not break SUSY but may give rise to a μ term if the symmetries of the CY compactification allow for it (see Refs. [24,26]). In Ref. [26] (section 5.1) the soft terms induced by these types of fluxes where computed in some simple $D7$-brane settings. It was found that all the bosonic soft terms arise from positive definite contributions to the scalar potential given by

$$V_{\text{flux}} = \sum_i |-M^* \phi_i^* + \partial_i W|^2, \tag{3.5}$$

where ∂_i is the derivative with respect to ϕ_i and W is the superpotential involving the field ϕ_i. Here M is the gaugino mass which is given in terms of the fluxes by

$$M = c(G_{(0,3)})^* \tag{3.6}$$

where $c = (g_s^{1/2}/3\sqrt{2})$, with g_s the string coupling constant. In addition, if the chiral field ϕ_i is vectorlike a possible SUSY mass term appears given by the flux

$$\mu = -c(G_{(2,1)})^*. \tag{3.7}$$

As explained in Refs. [24,26] (see also [25]) these results may be understood also from the effective $N = 1$ supergravity point of view. Indeed it may be seen that a nonvanishing value for $G_{(0,3)}$ induces a nonvanishing expectation value for F_T, the auxiliary field of the overall modulus field T. A constant superpotential proportional to $G_{(0,3)}$ is also induced. Now on $D7$ branes the gauge kinetic function is simply given by $f_a = T$ and hence from Eq. (1.5) gauginos get a mass proportional to F_T [and hence to $G_{(0,3)}$]. The SUSY-breaking terms above may be understood as arising from a ''modulus domination'' scheme. This may sound surprising for readers familiar with heterotic compactifications, since in that case it is well known that modulus domination leads to no soft terms at all to leading order. Indeed that would be the case also in type IIB orientifolds if e.g., one considers the charged fields arising from $D3$ branes. Those do not get any soft terms either from a $G_{(0,3)}$ background. It is the fact that our charged fields are living on $D7$ branes (which is the natural situation in F theory) that makes the difference. The fact that modulus dominance may lead to nontrivial soft terms for branes different than $D9$ or $D3$ already appeared in Ref. [31] in which the approach of Ref. [15] was applied to Dp-brane type I systems (see [26] for a description of its connection with flux-induced soft terms).

From the $N = 1$ supergravity point of view it is easy to understand the appearance of this positive definite SUSY-breaking scalar potential. It is well known that if only the auxiliary field of the overall modulus T is breaking SUSY, the negative piece of the scalar potential Eq. (1.3), $-e^K 3|W|^2$ is canceled by a positive contribution coming from the T-field auxiliary field $|D_T W|^2$ leading to a vanishing vacuum energy, this is the no-scale structure. On the other hand one observes that it remains a positive definite piece which is uncanceled and is given by

$$V = e^K (g^{i\bar{j}} (D_i W)(\bar{D}_{\bar{j}} \bar{W})) \tag{3.8}$$

where the sum now only runs over the matter fields (the contribution of S vanishes identically since the fluxes considered do not contribute to F_S). Now substituting $D_i W = \partial_i W + W K_i$, normalizing canonically the fields and recalling that both the constant superpotential and gaugino masses are proportional to $G_{(0,3)}$, one obtains the result (3.5). Note that this has the same form as the SUSY scalar potential for matter fields with the only difference that one replaces the usual derivative ∂_i by the covariant Kahler derivative $D_i W$. This is precisely the kind of substitution required from a bottom-up argumentation in the previous section.

Note that sometimes some $D7$ charged fields may not appear in this scalar potential. As we said, this is the case of $D7$ chiral fields not corresponding to geometric $D7$-brane moduli. A particular example appears in toroidal or orbifold orientifolds in which some massless $D7$-scalars parametrize possible continuous Wilson lines. Those fields have kinetic terms which depend only on T and then the standard cancellation for soft masses characteristic of no-scale models takes place. Thus in this toroidal case, scalars ϕ_1, ϕ_2 corresponding to Wilson lines in the first and second complex planes remain massless whereas ϕ_3, which parametrizes the $D7$ position is transverse space get masses in the form described above (see [26]). However, in generic CY compactifications at most discrete (not continuous) Wilson lines may be added, so the presence of these type of $D7$-brane moduli is ungeneric.

IV. THE FLUXED MSSM

Although there has been important recent progress in obtaining realistic models from type II orientifolds [18,19], most of the examples considered assume vanishing antisymmetric fluxes. Some preliminary steps on realistic models with fluxes have however been given

[5] In many F-theory compactifications all charged $D7$ zero modes have this geometric character. In simpler less generic compactifications (e.g. toroidal or orbifold orientifolds) some $D7$ zero modes rather parametrize e.g. values of continuous Wilson lines. We will refer later to those.

[24,26,32,33] . Still, although the $D7$ brane flux configurations considered up to now are very simplified, it could well be that similar structures could appear in more realistic type IIB orientifolds or F-theory compactifications. In particular, the fact that the scalar potential including soft terms is positive definite and involves the scalars parametrizing the $D7$-brane positions seems to be a general property of SUSY breaking induced by ISD antisymmetric backgrounds.

It seems then worth considering in which way such a structure could appear in a theory including the spectrum and interactions of the MSSM. A first simple option is to assume that all the fields of the MSSM correspond to geometric $D7$-brane moduli in some F-theory compactification. We then assume that ISD fluxes of type $G_{(0,3)}$ are present leading to modulus dominated SUSY breaking. In addition, if a $G_{(2,1)}$ background is present, the Higgs multiplets may generically get a μ term, as explained above. Then the full SUSY-breaking scalar potential will have the form

$$
\begin{aligned}
V_{\rm FMSSM} = \; & \Big| -\mu H_d - M^* H_u{}^* + \sum_{ij} h_U^{ij} Q_i U_j \Big|^2 + \Big| -\mu H_u \\
& - M^* H_d^* + \sum_{ij} h_D^{ij} Q_i D_j + \sum_{ij} h_L^{ij} L_i E_j \Big|^2 \\
& + \sum_i (| - M^* Q_L^{i*} + h_U^{ij} U_j H_u + h_D^{ij} D_j H_d|^2 \\
& + | - M^* U^{i*} + h_U^{ij} Q_j H_u|^2 + | - M^* D^{i*} \\
& + h_D^{ij} Q_j H_d|^2 + | - M^* E^{i*} + h_L^{ij} L_j H_d|^2 \\
& + | - M^* L^{i*} + h_L^{ij} E_j H_d|^2)) + V_{D\text{-terms}}. \qquad (4.1)
\end{aligned}
$$

This potential gives rise to the following set of bosonic SUSY-breaking soft terms for the MSSM:

$$
\begin{aligned}
& m_{H_u}^2 = m_{H_d}^2 = m_{\tilde{q}}^2 = m_{\tilde{l}}^2 = |M|^2, \\
& A_U = A_D = A_L = -3M, \qquad B = 2M\mu.
\end{aligned} \qquad (4.2)
$$

Note that all soft terms are universal and given by only two free parameters M and μ which are determined by the ISD fluxes $G_{(0,3)}$ and $G_{(2,1)}$ respectively. As a general remark note that, since both the SUSY-breaking parameter M and the μ term arise from fluxes, it is natural for them to be of the same order of magnitude. Thus flux SUSY breaking solves naturally the μ problem.

We thus see that under the assumption that (1) our SM fields are embedded as geometric $D7$-brane fields in a general type IIB orientifold (or more generally, F-theory) compactification and (2) that ISD fluxes are present we obtain a rather simple structure of soft terms addressing the four MSSM problems listed at the beginning of the paper.[6]

Most of the features of the above simplest choice of soft terms may be also obtained from a simple $N = 1$ supergravity toy model. Indeed, consider the following string motivated type of gauge kinetic function f_a and Kahler potential

$$
f_a = T, \qquad (4.3)
$$

$$
\begin{aligned}
K = & -\log(S + S^* - |H_u + H_d^*|^2 - \sum_i |\phi_i|^2) \\
& - 3\log(T + T^*), \\
W(S) = & \; aM_p^2 S + bM_p^2,
\end{aligned}
$$

where ϕ_i represents the squark and slepton fields. The superpotential $W(S)$ is modeling the general flux-induced superpotential Eq. (3.3) and a, b are complex constants related to the flux densities $H_{(3)}, F_{(3)}$ integrated over the CY space.[7] We will however treat a, b as free constant parameters. Readers familiar with string derived $N = 1$ models will note a number of differences from the standard perturbative heterotic lore. First, the gauge kinetic functions are not given by the axidilaton field S but rather by the complexified volume field T. A second difference is that the MSSM fields appear in the Kahler potential combined with the complex dilaton S field rather than the overall modulus T. However this is precisely the form of the gauge kinetic functions and Kahler potential which appear when considering simple toroidal/orbifold compactifications of type IIB orientifolds with geometric (no-Wilson-line) $D7$-brane matter fields [31,34]. In this toy example it is easy to see that upon minimization of the scalar potential one has $F_S = 0$, which fixes the value of the complex dilaton field at $S = (b^*/a^*)$. At the minimum a constant superpotential is obtained at $W_0 = M_p^2[(a/a^*)b^* + b]$. Because of the no-scale structure of the T-field SUSY is broken by a nonvanishing F_T, with a vanishing vacuum energy, leaving the T-vev undetermined. Using standard supergravity formulas (see e.g. Ref. [16]) it is an easy exercise to show that precisely the simple choice of soft terms (4.2) are obtained. This particular form of Kahler potential leads to a contribution to the μ term generated a la Giudice-Masiero $\mu_0 = M$ [7]. More generally the flux analysis shows that the complete μ term and gaugino masses M depend on different fluxes and hence are in general independent parameters. Thus one can reproduce this more general case by considering an explicit μ term in the original superpotential.

Irrespective of its string theory motivation, the above simple choice of Kahler potential, gauge kinetic term and S-dependent superpotential constitute an interesting $N = 1$

[6]In the addendum we address a possible generalization which includes the case in which the MSSM particles live at the intersections of different stacks of $D7$ branes.

[7]Note that this flux inspired superpotential has the form of a simple Polony superpotential for the dilaton complex field. However the "Polony field" here is S which does not have a canonical metric as in the old supergravity models.

supergravity model containing the MSSM spectrum. No fine tuning is required to get a (tree-level) vanishing cosmological constant, still it gives rise to universal SUSY-breaking soft terms.

Let us end with a number of comments. The above simple universal result is obtained under the (reasonable) assumption that all the MSSM fields are $D7$-geometric moduli. We have mentioned, however that in certain cases some $D7$ zero modes do not have a geometric meaning but rather correspond to, e.g., the possible existence of continuous W.L. backgrounds on the $D7$-brane world volume. Those may remain massless even in the presence of fluxes. Thus one may perhaps consider other ways of embedding the MSSM inside type IIB orientifolds in which some of the MSSM chiral fields do not get masses. That may lead to nonuniversal scenarios in which some of the MSSM fields get soft masses and others do not. These mixed scenarios seem however less natural that the universal one described above, since, as we mentioned, the presence of continuous Wilson lines is ungeneric in CY compactifications. Also, from the phenomenological point of view many of those nonuniversal possibilities are problematic due to FCNC constraints. One possibility which would be quite simple and universal would be one in which only the Higgs multiplets correspond to $D7$ geometric moduli. In this case the only source of scalar SUSY-breaking soft terms would be those arising from the first two terms in Eq. (4.2) which lead to soft terms of the form $m^2_{H_u} = m^2_{H_d} = |M|^2$; $m^2_{\tilde{f}} = 0$; $A_U = A_D = A_L = -M$, $B = 2M\mu$. There is finally a third possibility also leading to universal soft terms, which is to assume that all chiral fermions correspond to $D7$ geometric moduli, but not the Higgs fields. In that case the obtained soft terms have $m^2_{\tilde{f}} = |M|^2$, $A_U = A_D = A_L = -2M$, and $m^2_{H_u} = m^2_{H_d} = \mu^2 = B = 0$. Note that soft terms like these two may be obtained from the toy supergravity model above by having the squark and slepton fields combined with the T field (S field) in the no-scale fashion in Eq. (4.3) while maintaining the Higgs (squark/slepton) fields combined with S respectively. Let us however emphasize again that these other nonuniversal possibilities look less generic in the context of F theory.

A further comment concerns gaugino masses and gauge coupling unification. If the relevant $D7$ branes containing the MSSM fields have all the same geometry (i.e., wrap the same 4-cycle in the CY compact space) and are located at the same point in transverse space, gauge coupling unification at the string scale is expected. In that case there will also be in general a universal gaugino mass parameter. This has been our simplifying assumption above, although generalizations without this property could be envisaged.

V. FINAL COMMENTS

We have argued that the presence of antisymmetric fluxes in type IIB orientifolds with $D7$ branes (or, in general, F-theory compactifications in complex 4-folds) give rise to an interesting class of soft SUSY-breaking soft terms, Eq. (4.2). An important point to emphasize is that the relevant class of fluxes studied (imaginary self-dual 3-form fluxes) solves the classical equations of motion for compactified type IIB string theory [21]. Thus the class of models discussed gives rise, to leading order, to consistent $N = 1$ low-energy theories with softly broken soft terms. To our knowledge, this is the first class of classical string compactifications in which that is the case.

From the low-energy $N = 1$ supergravity effective Lagrangian point of view the presence of fluxes give rise to modulus-dominance SUSY breaking. However, unlike a type of SUSY breaking studied in perturbative heterotic compactifications (or its type I dual with $D9$ branes), in the present case interesting classes of soft terms do appear on the massless fields coming from $D7$ branes after compactification.

Assuming that the MSSM fields may be embedded as geometric $D7$-brane moduli, we have argued that the relevant MSSM soft terms depend only on two free parameters, the gaugino mass M and a μ term. Those are in turn given by certain classes of flux densities. Thus one expects both parameters to be of the same order of magnitude, given the fact that they have a common origin, fluxes. The SUSY-breaking scalar potential is positive definite and it is simply obtained from the SUSY scalar potential by making the replacement $\partial_i W \to D_i W$, with i running over all the chiral multiplets of the MSSM. The set of soft terms (4.2) so obtained is universal and solves the SUSY-CP problem due to the specific relationships obtained between the A, B, M and μ parameters.

A natural question to ask is whether one could find a scheme in which M and μ were related. In that case we would be left with a single parameter describing all soft terms. In principle the flux densities $G_{(0,3)}$ and $G_{(2,1)}$ are independent parameters. In fact, in generic CY compactifications there are a number of different 3-cycles through which fluxes can exist. The integral of the corresponding RR and NS 3-forms over 3-cycles in the CY are quantized [Eq. (3.1)] although the flux densities themselves are not. Thus parametrically these densities $G_{(0,3)}$, $G_{(2,1)}$ go like $\simeq N/(\mathrm{Vol}_C)$, with N integers and Vol_C the volume of the corresponding 3-cycle. In specific compactifications the volumes of the different cycles could be related (e.g. equal) and one would expect specific relationships (e.g. $\mu = 2M$) depending on the different integers. To find this kind of relationships would require however an specific example of compactification yielding the MSSM.

A number of other interesting issues should be addressed. Of course, it would be important to have orientifolds or F-theory examples with a massless spectrum as close as possible to that of the MSSM. Also it would be important to realize specific examples (perhaps with some large dimension transverse to $D7$ branes and/or warping or other) in which the size of soft terms is naturally of order the electroweak scale. We have also ignored in our analysis the dynamics which eventually determines the volume moduli T (see Ref. [35] for a recent discussion of this issue). Another obvious point is to study the low-energy spectrum of SUSY masses as well as the generation of radiative electroweak symmetry breaking in this class of theories. The latter will be presented elsewhere [36].

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APPENDIX: INTERSECTING $D7$ BRANES

Some of the most promising semirealistic type IIA string orientifold models are based on intersecting $D6$ branes [19]. Upon T duality this class of models may be equivalently described as type IIB orientifold models with intersecting $D7$ branes with magnetic fluxes in their world volume (branes of other dimensions may also be present, in particular, models). Thus it would be interesting to see what kind of soft terms are induced by ISD fluxes in this type of intersecting models. Most of those models are toroidal (or orbifold) compactifications. So we will consider here a case with a toroidal compactification on $T^2 \times T^2 \times T^2$ in which we have three classes of $D7$ branes, $D7_i$, $i = 1, 2, 3$ which are transverse to the ith 2-torus, respectively. Thus in addition to the fields coming from the world volume of a stack of $D7$ branes $D7_i$ considered in the main text, we will now have in general new chiral fields ϕ_{ij} corresponding to open strings living at the intersections of a pair of distinct stacks of branes $D7_i - D7_j$. Let us study here what kind of soft terms appear for these extra fields living at the intersections in the presence of ISD fluxes (corresponding in field theory language to modulus domination). Using the effective Lagrangian approach for the overall modulus-dominance in Refs. [31,34] as well as [25,26] one can figure out what to expect. One finds that the bosonic SUSY-breaking soft terms for all chiral fields may be obtained as arising from a slight generalization of the results in the main text, namely

$$
\begin{aligned}
V_{\rm SUSY} &= \sum_i |\partial_i W|^2 \longrightarrow V_{SB} \\
&= \sum_i (1-\xi_i)|\partial_i W - M^*\phi_i^*|^2 + \sum_i \xi_i |\partial_i W|^2,
\end{aligned}
\tag{A1}
$$

where the T dependence of the metric of the field ϕ_i is $(T+T^*)^{-\xi_i}$.[8] Thus the case considered in the previous sections (geometric $D7$-brane moduli, no T dependence in the kinetic term of the field) corresponds to $\xi_i = 0$ and the case of Wilson-line $D7$-fields (no S dependence in the metric) would correspond to $\xi_i = 1$, giving rise to no soft terms. From section 7 in Ref. [31] (see also [25]) one can see that the matter fields in $D7_i - D7_j$ intersections have metric proportional to $((S+S^*)(T+T^*))^{-1/2}$ and then one has $\xi_i = 1/2$ for the fields. One can easily check that Eq. (A1) with $\xi_i = 1/2$ indeed reproduces the bosonic soft terms in [31]. It is easy to understand qualitatively these results. For fields with metric proportional to $(T+T^*)^{-1}$ (corresponding to $\xi = 1$) the usual no-scale cancellation gives zero soft terms for the fields. For fields with metric proportional to $(S+S^*)^{-1}$ there is no cancellation at all and soft terms appear as in previous sections. On the other hand the metric for the fields at intersections is in some sense in between and there is a partial no-scale cancellation of soft terms, giving rise to the $\xi_i = 1/2$ factor. Note that in more general intersecting $D7$-brane models in which there are magnetic fluxes in their world volume the T dependence of the metric of the fields at intersections will in general depend on the (magnetic) fluxes so that the ξ_i will be model dependent and have to be computed in each model.

Let us now apply these ideas to the MSSM. In order to have universality the ξ_i's corresponding to different generations of the same quark or lepton should be equal. But in fact in the class of models under discussion that is in general the case, different generations of fields have the same metric. Thus universality is a natural situation even in these more general configurations. So apart from the flux-induced parameters M and μ, there will be a set of 7 model dependent parameters ξ_i, $i = Q, U, D, L, E, H_u, H_d$. In terms of all these and using Eq. (A1) one can write down the following set of soft terms arising from ISD fluxes for the MSSM:

$$
\begin{aligned}
m_i^2 &= (1-\xi_i)|M|^2, \quad i = Q, U, D, L, E, H_u, H_d, \\
A_U &= -M(3 - \xi_{H_u} - \xi_Q - \xi_U), \\
A_D &= -M(3 - \xi_{H_d} - \xi_Q - \xi_D), \\
A_L &= -M(3 - \xi_{H_d} - \xi_L - \xi_E), \\
B &= M\mu(2 - \xi_{H_u} - \xi_{H_d}).
\end{aligned}
\tag{A2}
$$

It must be emphasized that in a given model the ξ_i are computable quantities determined by the T dependence of the metric of the corresponding field and hence only M, μ remain as free parameters. Note that for all $\xi = 0$ we recover the situation described in the previous sections, Eq. (4.2). As in the simpler case discussed in the previous

[8] This may be considered as the type IIB analogue of the modular weights of Refs. [13,15]).

sections, these soft terms may be derived from a simple $N = 1$ supergravity model with the same gauge kinetic function and S-dependent superpotential as in Eq. (4.3) but with a generalized Kahler potential in which the metric of matter fields include the mentioned $(T + T^*)^{-\xi_i}$ dependence (see [31,34,37]).

It would be interesting to compute the ISD flux-induced soft terms in specific semirealistic compactifications. An example of this is the type II intersecting D-brane configuration yielding an MSSM-like spectrum proposed in [38]. This local D-brane configuration may be embedded into a full $N = 1$ SUSY type IIB orientifold [39] $Z_2 \times Z_2$ compactification [33] with additional ISD fluxes. In the latter constructions the MSSM fields appear at the intersections of 3 sets of branes $D7_i$, $i = a, b, c$, very much as described above. A stack of 8 $D7$ branes $D7_a$ give rise to the Pati-Salam group $SU(4)$ [which may easily be broken down to $SU(3) \times U(1)_{B-L}$ in the presence of Wilson-line backgrounds]. Two parallel $D7$ branes $D7_b(D7_c)$ give rise to the gauge group $SU(2)_L(SU(2)_R)$. The $D7_a$ stack has in addition magnetic flux in its world volume, giving rise to the replication of generations. It is beyond the scope of the present paper to give a detailed description of the soft terms induced in a model like this (see Refs. [40,41] for recent analysis). It may be however illustrative to figure out simple main features of the expected structure of soft terms. In this particular model the stacks $D7_b$ and $D7_c$ have no magnetic flux in their world volume. Thus the fields at their intersection (one set of MSSM Higgs fields) will simply have $\xi_{H_u} = \xi_{H_d} = 1/2$. The quarks and leptons reside at intersections $D7_a - D7_b$ and $D7_a - D7_c$. Given the symmetries of the brane configuration in this model (and the built-in Pati-Salam symmetry) all quarks and leptons are universal, $\xi_Q = \xi_U = \xi_D = \xi_L = \xi_E = \xi$. So all in all the general form of soft terms will be

$$\begin{aligned} &m_{H_u}^2 = m_{H_d}^2 = \frac{|M|^2}{2}, \\ &m_Q^2 = m_U^2 = m_D^2 = m_L^2 = m_E^2 = (1 - \xi)|M|^2, \\ &A_U = A_D = A_L = -M(5/2 - 2\xi), \qquad B = M\mu. \end{aligned} \tag{A3}$$

For large T values the magnetic fluxes are diluted and one expects to recover the case without fluxes with $\xi \simeq 1/2$. In that limit one would have the universal result

$$\begin{aligned} &m_i^2 = \frac{|M|^2}{2};\ i = Q, U, D, L, E, H_u, H_d, \\ &A_U = A_D = A_L = -M(3/2), \qquad B = M\mu. \end{aligned} \tag{A4}$$

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