SUSY breaking mediation by throat fields

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Abstract

We investigate, in the general framework of KKLT, the mediation of supersymmetry breaking by fields propagating in the strongly warped region of the compactification manifold (‘throat fields’). Such fields can couple both to the supersymmetry breaking sector at the IR end of the throat and to the visible sector at the UV end. We model the supersymmetry breaking sector by a chiral superfield which develops an $F$ term vacuum expectation value (also responsible for the uplift). It turns out that the mediation effect of vector multiplets propagating in the throat can compete with modulus-anomaly mediation. Moreover, such vector fields are naturally present as the gauge fields arising from isometries of the throat (most notably the SO(4) isometry of the Klebanov-Strassler solution). Their mediation effect is important in spite of their large 4d mass. The latter is due to the breaking of the throat isometry by the compact manifold at the UV end of the throat. The contribution from heavy chiral superfields is found to be subdominant.
1 Introduction

The KKLT scenario [1] suggests a way in which all moduli of a string compactification could be stabilized in a de Sitter vacuum with small cosmological constant. The construction is based on the supergravity approximation to type IIB superstring theory. The complex structure moduli of the internal manifold and the dilaton are stabilized by fluxes (see e.g. [2] and references therein). There is always at least one Kähler modulus $T$ which does not appear in the flux superpotential and which KKLT take to be stabilized nonperturbatively (e.g. by gaugino condensation). This results in a supersymmetric AdS vacuum. SUSY may be broken and the cosmological constant uplifted to a positive value by introducing $\overline{D}3$ branes.

The flux-stabilized compactification geometry for type IIB models was worked out in [3], where it was found that fluxes lead to a warped product structure between the four non-compact spacetime dimensions and the internal manifold. Strongly warped regions or ‘throats’ occur naturally, with a large hierarchy of scales between the bottom of the throat (‘IR end’) and the weakly warped remainder of the internal space (‘UV end’ or ‘compact manifold’). This hierarchy is crucial for the KKLT construction since the $\overline{D}3$ branes are dynamically confined to the IR end of the throat, where their contribution to the vacuum energy density is maximally redshifted. Thus it is possible to start with a pre-uplift AdS vacuum at parametrically large volume, and still (nearly) cancel its parametrically small cosmological constant.

In this paper we study the mediation of SUSY breaking between the bottom of the throat and the UV region. Taking the visible sector to be localized in the UV and the SUSY breaking sector in the IR, we focus on the mediation effects of fields propagating in the throat and coupling to both sectors. In KKLT, supersymmetry breaking is induced by the $\overline{D}3$ branes, which represents a hard breaking from the point of view of 4d supergravity. The systematic analysis of soft terms in this setup was pioneered in [4], employing non-linearly realized SUSY for the description of the uplifting sector in the throat. Here, we will instead introduce a dynamical chiral superfield $X$, confined to the bottom of the throat. The superpotential and Kähler potential of $X$ are assumed to ensure a non-zero vacuum expectation value of its auxiliary component $F_X$ (see e.g. [8,9] for related examples involving $F$ term uplifts). It will turn out that much of the SUSY breaking dynamics as seen by the visible sector is independent of the details of the $X$ Lagrangian. Our SUSY breaking $X$ sector may be seen in two ways: On the one hand, it can model the $\overline{D}3$ brane sector (including, in the limit of an extremely heavy $X$, much of what can be done using non-linearly realized SUSY). On the other hand, it can be taken seriously in its own right, expecting that, in the future, a stringy realization of such a sector will be found and used to replace the $\overline{D}3$ brane sector of KKLT. For example, it could be realized by branes at singularities at the bottom of the throat [10].

Our main result is that, under certain quite natural conditions, the sequestering assumption (see [11] and the critical discussion of [12]) is generically violated by effective

\footnote{1 See also [5]. See [6] for some other work on soft terms in KKLT, and [7] for a selection of earlier papers on soft terms in type IIB flux compactifications.}
5d gauge fields representing the SO(4) isometry of the Klebanov-Strassler throat [13, 14]. More specifically, we assume that our SUSY-breaking field $X$ is charged under the isometry of the throat. This is natural given that, for example, $\overline{D3}$ branes at the bottom of the throat break this isometry. Furthermore, we assume that the compact space at the UV end of the throat breaks the isometry softly. Again, this is natural since one can clearly imagine a compact manifold respecting (part of) the throat isometries and view the actual Calabi-Yau manifold of a realistic model as a deformation thereof. Under these conditions, it is easy to see that the massive vector fields originating in the isometry develop a $D$ term and that this $D$ term induces SUSY-breaking scalar masses in the UV sector of the model which can compete with those of the sequestered case.

We do not consider this a negative result for the following reason: Our findings suggest that the above effect of ‘vector mediation’ is the only one competing with mixed modulus-anomaly mediation. Thus, it is conceivable that, in specific models, the impact on soft scalar masses will turn out to be small or, more interestingly, will be large and calculable. Clearly, this will require a better understanding of both the SUSY breaking sector at the bottom of the throat as well as the geometry of the compact space.

While this work was being finalized, Ref. [15] appeared, which uses the formalism of non-linear realizations and focuses on the effect of anomalous U(1) gauge symmetries. Nevertheless, in its discussion of sequestering it has a significant overlap with this analysis. Our perspective differs in that we identify throat vector fields (which are also discussed in [15]) as an intrinsic feature of the KKLT setup, forced upon us by the symmetry of the supergravity solution. Furthermore, our dynamical description of the SUSY breaking sector based on the chiral superfield $X$ allows us to specify its effect on isometry vector fields in a very direct and physical way. We will comment on some additional fine points in which our analyses differ as we go along.

The present paper is organized as follows. We begin with a streamlined discussion of the basic KKLT setup and its SUSY breaking dynamics in Sect. 2. The use of the chiral compensator formalism and the realization of the SUSY breaking and uplifting sector as conventional $F$ term breaking make the discussion of energy scales and sequestering particularly transparent.

In Sect. 3 we introduce, motivated by the isometry of the Klebanov-Strassler throat, bulk vector fields and calculate the $D$ terms they acquire if the SUSY-breaking chiral superfield $X$ is charged. It turns out that their effect on the scalar masses of Standard Model (or other chiral) superfields in the UV sector is potentially large.

For completeness, we analyse in Sect. 4 the effect of throat fields whose zero modes are 4d chiral superfields. Assuming that those chiral superfields acquire large masses due to some dynamics in the UV region and that they do not mix with $T$ at the perturbative level, we find that their effect is small and sequestering is respected.

Section 5 contains our conclusions and a brief discussion of open issues. Two technical calculations related to a concrete SUSY breaking model and to the mediation effect of massive chiral superfields are given in the Appendix.
2 The minimal scenario

Following [1], we consider a type IIB compactification with all complex structure moduli and the dilaton stabilized by fluxes at a supersymmetric minimum, at which the superpotential is $W_0$. We assume that there is only a single Kähler modulus $T$ with no-scale kinetic function and Kähler potential

$$\Omega = -(T + \bar{T})$$

and

$$K = -3 \log(-\Omega) = -3 \log(T + \bar{T}). \tag{1}$$

It is stabilized, e.g., by gaugino condensation, such that the total superpotential becomes

$$W = W_0 + e^{-T}. \tag{2}$$

We suppress $O(1)$ numerical coefficients here and below and work in units where $\alpha' \sim 1$.

The scalar potential (in a frame where $\Omega$ is the coefficient of the Einstein-Hilbert term) is most easily obtained from the standard $D = 4$, $N = 1$ Lagrangian

$$\mathcal{L} = \int d^4 \theta \varphi \Omega + \left( \int d^2 \theta \varphi^3 W + \text{h.c.} \right), \tag{3}$$

where $\varphi = 1 + \theta^2 F_\varphi$ is the chiral compensator [16].

With the above kinetic function and superpotential (and using the same symbol for a chiral superfield and its lowest component), the scalar-potential-part of Eq. (3) is

$$\mathcal{L} \supset -(T + \bar{T}) |F_\varphi|^2 - (F_T F_\bar{\varphi} + \text{h.c.}) + \left[ (3(W_0 + e^{-T}) F_\varphi - e^{-T} F_T) + \text{h.c.} \right]. \tag{4}$$

The resulting equations of motion are

$$F_\bar{\varphi} : \quad -(T + \bar{T}) F_\varphi - F_T + 3(W_0 + e^{-T}) = 0, \tag{5}$$

$$F_T : \quad - F_\varphi - e^{-T} = 0, \tag{6}$$

$$\bar{T} : \quad - |F_\varphi|^2 - 3 e^{-T} F_\varphi + e^{-T} F_T = 0. \tag{7}$$

Taking $W_0$ to be parametrically small (which may be justified by the exponentially large number of flux choices), it is easy to see that the above equations are solved for

$$F_T \sim F_\varphi \sim e^{-T} \sim W_0. \tag{8}$$

Note that here and below we focus on parametrically small factors $\sim e^{-T}$ but ignore factors $\sim 1/T$ (which are strictly speaking also parametrically small, but to a much lesser degree). The vacuum energy density is negative and $\sim W_0^2$.

At this point, it is crucial to recall that a solution of the above equations does not represent a true vacuum of the model unless the curvature scalar (which is multiplied by $\Omega$) vanishes. This shortcoming will now be corrected.

Assume that, due to the fluxes, the compactification manifold has developed a strongly warped region or throat. We normalize the warp factor such that it is $O(1)$.
at the UV end (and throughout the compact space) and equals $\omega$ (with $\omega \ll 1$) deep in the IR. Thus, mass scales at the bottom of the throat are redshifted by a factor $\omega$.

Now let us add a SUSY-breaking sector in the throat, i.e. some physical degrees of freedom localized at the bottom of the throat which can, in 4d language, be described by a chiral superfield $X$. We assume that $X$ is sequestered from all the light fields on the compact space in the UV. (At the present stage of our analysis, the only relevant light field is the universal Kähler modulus $T$.) As discussed in some detail in [18], this setup can be viewed as a Goldberger-Wise stabilized Randall-Sundrum model with $T$ being a no-scale field localized at the UV brane (see [19] for general analysis of deformations of warped models). Thus, the sequestering assumption can indeed be made and the $X$- or uplifting sector leads to corrections of the type [20]

$$\Omega_{up} = \Omega + \omega^2 \Delta \Omega(X, \overline{X})$$
$$W_{up} = W + \omega^3 \Delta W(X) = W_0 + e^{-T} + \omega^3 \Delta W(X).$$

(9)

Here the warp factor dependence (which follows on dimensional grounds) has been given explicitly. Therefore the functions $\Delta W$ and $\Delta \Omega$ are naturally of the order of the string scale (i.e. they are $O(1)$ in our units and contain no small or large parameters).

Note that we could in principle have performed a Kähler-Weyl rescaling by $T^\alpha$ before imposing sequestering, which would correspond to using

$$\Omega = -(T + \overline{T})(T \overline{T})^\alpha, \quad W = T^{3\alpha}(W_0 + e^{-T})$$

(10)

in Eq. (9). The value of $\alpha$ can be fixed by requiring the uplift energy to scale as $1/(T + \overline{T})^2$ [21]. It is easily checked that $\alpha = 0$ is the correct value, which shows that, in Eq. (9), we should indeed use the unrescaled form of $\Omega$ and $W$ as given in Eqs. (1) and (2).

Neglecting for the moment the influence of $F_\phi$ on the $X$ sector (this will be easy to justify a posteriori), the equation of motion for $F_X$ reads

$$\omega^2 \Delta \Omega_{XX} F_X + \omega^3 \Delta W_X = 0,$$

(11)

where the indices of $\Delta W$ and $\Delta \Omega$ denote partial derivatives.

Obviously, if $\Delta W$ and $\Delta \Omega$ are such that, in the absence of warping, $F_X$ would break SUSY at the string scale, then

$$F_X \sim \omega$$

(12)

in the present context. The vacuum energy density induced by the $X$ sector is $\sim \omega^4$, with comparable contributions coming from $\Delta W$ and $\Delta \Omega$.

It now becomes clear that to uplift the previously found negative vacuum energy density $\sim W_0^2$ to a realistic positive value (i.e. to zero, for all practical purposes), we need $W_0 \sim \omega^2$. Thus, there is in fact only one small parameter in the model, which we

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2Alternatively, a dynamical SUSY-breaking can be introduced via $D$ terms [17] or non-sequestered $F$ terms [9].
choose to be $\omega$. It is also immediately clear that, in this situation, the influence of the $X$ sector on the previously found solution for $F_\varphi$ (and hence on $T$ and $F_T$) is of higher order in $\omega$. Thus, Eqs. (5) - (7) continue to be the right equations to solve. The $X$ sector simply adds the necessary positive vacuum energy to promote the solutions of these equations to a physical vacuum with

$$ F_T \sim F_\varphi \sim W_0 \sim \omega^2 \quad \text{and} \quad F_X \sim \omega. \quad (13) $$

This lies at the basis of ‘mixed modulus-anomaly mediation’ [4]. Recall that, due to sequestering, $F_X$ has no direct effect on soft terms in the visible sector.

## 3 Vector superfields

In this section we study the effect of a massive vector superfield on sequestering. Such a field can emerge from an isometry of the throat, which becomes a gauge symmetry in the corresponding 4d field theory. If the isometry is not a symmetry of the entire internal manifold (in particular, of the UV end), this gauge symmetry is spontaneously broken, and the gauge field acquires a UV-scale mass.

The prime example of a throat is the warped deformed conifold of Klebanov and Strassler [13]. Its full 10d geometry is a deformation of $\text{AdS}_5 \times T^{1,1}$, where $T^{1,1}$ is a 5d compact Einstein space with isometry $G = \text{SU}(2) \times \text{SU}(2) \times \text{U}(1)$. The massless 5d spectrum of type IIB supergravity compactified on $T^{1,1}$ contains seven vector multiplets from the seven generators of $G$ [22]. In the warped deformed conifold the $\text{U}(1)$ is broken, and the isometry group is $\text{SU}(2) \times \text{SU}(2) = \text{SO}(4)$ [14]. Hence the 5d solution should contain six massless vector multiplets. This solution can be interpreted as the bulk of a Randall-Sundrum-like model, with the Calabi-Yau playing the role of the UV brane. Since Calabi-Yaus admit no isometries, the zero modes of the above vector multiplets acquire a large mass which, from the point of view of the Randall-Sundrum model, can be ascribed to a UV-brane-localized mass operator.

Independently of the above UV-scale breaking of the SO(4) gauge symmetry, we assume that the SUSY breaking sector at the bottom of the throat by itself also breaks this symmetry. In particular, a $\overline{D3}$ brane at the bottom of the throat already breaks part of the isometry. Clearly, as far as the mass of the 4d vector states is concerned, this IR-scale breaking can not compete with the UV-scale breaking discussed earlier. We model this situation in the following by assuming that our SUSY-breaking sector contains fields charged under the gauge symmetry, but ignoring any possible symmetry breaking $A$ term vevs from this sector.

To derive the main qualitative SUSY breaking effects of the above setup, we introduce a single 4d vector superfield $V$ (although the actual symmetry is non-abelian and several

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3In detail, the small hierarchy between $F_T$ and $F_\varphi$ based on the $\beta$-function coefficient of the condensing gauge group and ignored in this paper also plays an important rôle.

4 The model developed in the last two paragraphs differs from the 5d massive vector fields discussed in [15], yet the following technical analysis will lead to very similar results.
such fields are expected). \( V \) has the following component expansion:

\[
V = C + \theta \sigma^\mu \overline{\theta} A_\mu + \frac{1}{2} (F_V \theta \theta + \text{h.c.}) + \frac{1}{2} \theta \theta \theta \theta D + \text{fermions.}
\]

(14)

Here \( F_V \) is complex, while \( A_\mu, C, D \) are real. The UV-brane symmetry breaking (or non-linear realization of the gauge symmetry) is modelled by simply giving this vector superfield a string-scale mass term. A massive vector superfield can give rise to soft terms in two ways: it may develop \( F \) or \( D \) terms in the vacuum.\(^5\)

The dominant effect is easy to guess: Focus on the term \( CD \) (coming from the superfield mass term \( \sim V^2 \)) and the term \( \omega^2 C |F_X|^2 \) (coming from the gauge coupling \( \omega^2 X V X \) motivated above). Varying these terms with respect to \( C \) one immediately finds

\[
D \sim \omega^2 |F_X|^2 \sim \omega^4,
\]

(15)

which induces scalar masses \( \sim \omega^4 \) for standard model fields \( Q \) in the visible sector if a gauge coupling \( \bar{Q}VQ \) exists. The presence of such a gauge coupling represents, of course, a crucial assumption to which we will return at the end of this section. We note that the \( D \) term contribution to the vacuum energy density is negligible compared with \( |F_X|^2 \), which is responsible for the uplift.

To derive the above in more detail, we start with the Lagrangian

\[
\mathcal{L} = \int d^4 \theta \overline{\varphi} \left[ \Omega(T, \overline{T}) + \omega^2 \Delta \Omega(X, e^V \overline{X}) + V^2 \right]
\]

\[
+ \left( \int d^2 \theta \left[ \varphi^3 \{ W(T) + \omega^3 \Delta W(X) \} + \frac{1}{4} W^\alpha W_\alpha \right] + \text{h.c.} \right),
\]

(16)

(17)

where \( W_\alpha \) is the field strength chiral superfield corresponding to \( V \). The most relevant terms in the Lagrangian are the mass term for \( V \), the gauge-kinetic term, as well as the terms

\[
\Omega = -(T + \overline{T}) + \omega^2 X(1 + V)\overline{X} + \ldots \quad W = W_0 + e^{-T} + \ldots
\]

(18)

(as before we suppress any coefficients that are generically of order one). In components, the mass term contributes

\[
\varphi \overline{\varphi} V^2 \big|_{\theta^4} = CD + F_\phi F_{\overline{\phi}} + A_\mu A^\mu + C \left( F_\phi F_{\overline{\phi}} + \text{h.c.} \right) + C^2 F_\phi F_{\overline{\phi}},
\]

(19)

and the gauge kinetic term gives

\[
W^\alpha W_\alpha \big|_{\theta^2} = -\frac{1}{2} F_{\mu\nu} F^{\mu\nu} + D^2.
\]

(20)

From the coupling of the gauge field to the SUSY breaking field \( X \) we get

\[
\overline{\varphi} \varphi X V \overline{X} \big|_{\theta^4} = F_\varphi F_{\overline{\varphi}}XC \overline{X} + F_{\overline{\varphi}} F_X C \overline{X} + \text{h.c.} + \frac{1}{2} F_{\overline{\varphi}}F_V X \overline{X} + \text{h.c.}
\]

\[
+ F_X F_{\overline{\varphi}} + \frac{1}{2} F_X F_V X \overline{X} + \text{h.c.} + \frac{1}{2} XD \overline{X}.
\]

(21)

\(^5\) Note that for a massless vector superfield \( V \), \( F \) terms are unphysical because the \( \theta^2 \)-components of \( V \) can be gauged away using Wess-Zumino gauge. This is no longer the case when \( V \) is massive.
For $X = 0$ in the vacuum (we can define $X$ such that any nonvanishing vev is already absorbed in the gauge symmetry breaking mass term for $V$), the equations of motion are

$$C : \quad D + (F_{\bar{\phi}} F_{\phi} + \text{h.c.}) + 2 C F_{\nu} F_{\bar{\nu}} + \omega^2 F_X F_{\bar{X}} = 0,$$

(22)

$$F_{\bar{\nu}} : \quad F_V + 2 C F_{\nu} = 0,$$

(23)

$$D : \quad C + D = 0,$$

(24)

giving

$$C \sim \omega^4, \quad D \sim \omega^4, \quad F_V \sim \omega^6.$$  

(25)

It is obvious that $F_V$ is irrelevant for SUSY breaking mediation. $D$, however, will contribute significantly, because a possible coupling $\sim Q e^V \bar{Q}$ to the visible sector will clearly induce soft scalar masses $m^2 \sim D \sim \omega^4$. This is just the same order of magnitude as we get from mixed modulus-anomaly mediation, so ‘vector mediation’ will compete with these effects. Of course, this can also be easily seen by focusing on the couplings $\sim \omega^2 X e^V \bar{X}$ and $\sim Q e^V \bar{Q}$ and integrating out the heavy vector. The induced operator $\omega^2 X \bar{X} Q \bar{Q}$ provides soft masses $m^2 \sim \omega^2 |F_X|^2 \sim \omega^4$.

Note that the sign of the operator $\omega^2 X \bar{X} Q \bar{Q}$ depends on the relative sign of the $X$ and $Q$ charges. In the case that it is negative, the contribution to the $Q$ mass term is tachyonic. However, this effect need not lead to dangerous instabilities since other contributions, e.g. from mixed modulus-anomaly mediation, are of the same order of magnitude and can thus cancel the vector-mediated contributions. On the other hand, it can be avoided altogether by a suitable choice of charges — the usual constraints on gauge charges from anomaly cancellation do not apply here, because $V$ is massive and the gauge symmetry is broken. In summary, we can always arrange for $Q$ to be non-tachyonic, so that in particular the $D$ term in the vacuum is nonzero.

Finally, we want to argue in favour of the assumed coupling of the visible sector fields $Q$ to $V$. Imagine, for example, that a stack of D branes wrapped appropriately on the Calabi-Yau space in the UV gives rise to a higher-dimensional gauge theory. Let the scalars of this SUSY gauge theory be the superpartners of standard model fermions (‘matter from gauge’). Since the throat isometry may, among other effects, be broken in the UV by the position of this brane stack, it is natural to expect that the matter superfields are charged under $V$. This provides a strong motivation for the coupling $Q e^V \bar{Q}$ used above.

# 4 Chiral superfields

Consider now a chiral superfield $Y$ with a string-scale mass term, such as might be produced by flux stabilization. We will show that, even if $Y$ has direct couplings to both the visible and the SUSY breaking sector, sequestering is not violated to leading order, as the contributions of $Y$ to SUSY breaking mediation are subdominant. We will estimate the $F$ term of $Y$ in the vacuum, since it may give SUSY breaking soft masses to visible sector fields via terms like $\bar{Y} Y \bar{Q} Q$.
Suppressing $O(1)$ coefficients, the dominant terms in the kinetic function and superpotential are

\begin{align}
\Omega &= -(T + \overline{T}) + Y\overline{Y} + \omega^2 (X\overline{Y} + \text{h.c.}) + \ldots \\
W &= W_0 + Y^2 + (1 + Y)e^{-T} + \omega^3 XY + \ldots .
\end{align}

(26)

Since we imagine that $Y$ contains fields propagating in the throat, we have allowed for the strongest possible couplings to the $X$ sector. Furthermore, since $Y$ does not represent a modulus of the fluxed Calabi-Yau, we have allowed for an unsuppressed mass term $\sim Y^2$ but excluded any leading-order linear term in $Y$ or a mixing of $Y$ and $T$. However, once non-perturbative effects (e.g. gaugino condensation) are incorporated, the clear separation between $Y$ and $T$ may be blurred, which motivates us to include the term $\sim Ye^{-T}$, as an example for such effects.\[\text{Note that we could have replaced } Y^2 \text{ by } (T + \overline{T})Y \text{ without affecting the results of the following analysis.}\]

Since $F_\phi$ would by itself not generate a non-zero $F_Y$, we will neglect its influence for the moment. Afterwards we will show that the backreaction of $Y$ on $F_\phi$ is indeed negligible, hence this ansatz is fully self-consistent.

We obtain the following $Y$- and $F_Y$-dependent terms in the bosonic Lagrangian:

\begin{align}
\mathcal{L} &\supset F_Y F_{\overline{Y}} + \omega^2 (F_X F_{\overline{Y}} + \text{h.c.}) + [2Y F_Y + F_Y e^{-T} - Ye^{-T} F_T + \text{h.c.}] \\
&\quad + \omega^3 (F_X Y + XF_Y + \text{h.c.}).
\end{align}

(27)

Recall that $e^{-T} \sim \omega^2$ by assumption. This leads to the equation of motion for $Y$

\[2F_Y - \omega^2 F_T + \omega^3 F_X = 0 ,\]

(28)

hence

\[F_Y \sim \omega^2 F_T \sim \omega^3 F_X \sim \omega^4 .\]

(29)

The equation of motion for $F_{\overline{Y}}$ reduces to

\[F_Y + \omega^2 F_X + 2\overline{Y} + \omega^2 + \omega^3 \overline{X} = 0 ,\]

(30)

thus

\[Y \sim \omega^2 .\]

(31)

To ensure that this estimate is correct, we now need to prove that there are no contributions to $F_\phi$ and $F_T$ of order $\omega^2$. This is fairly obvious, however, since what we are adding to the Lagrangian by including $Y$ is, in the vacuum, suppressed by sufficiently high powers of $\omega$. For example, we can check that Eq. [6], the equation of motion for $F_T$, now becomes

\[F_\phi + (1 + Y)e^{-T} = 0 ,\]

(32)

inducing a negligible correction to $F_\phi \sim \omega^2$. (This remains correct if $Y\overline{Y}$ is replaced by $(T + \overline{T})Y\overline{Y}$. Similarly, it is easy to check that the vacuum values of $T$ are not affected at leading order in $\omega$.

\[\text{Note that we could have replaced } Y^2 \text{ by } (T + \overline{T})Y \text{ without affecting the results of the following analysis.}\]

\[\text{This is a slight generalization of the otherwise similar analysis of [15]}\]
In summary, we have seen that throat fields which are described by heavy chiral superfields in the 4d effective theory cannot contribute sizeably to SUSY breaking mediation because their $F$ terms are always subdominant compared to $F_{\varphi}$ and $F_T$.

We will give an alternative (and somewhat more general) derivation of the warp factor dependence in appendix B, using standard supergravity relations.

5 Conclusions

In this paper we have studied SUSY breaking mediation in the KKLT setup, in particular the contributions from vector and chiral fields propagating in the throat. Vector fields are naturally present in warped throat constructions. For instance, the currently best-understood example of a warped throat, the Klebanov-Strassler solution, has an SO(4) global symmetry acting on the compactification manifold and the flux background, hence there will be six vector fields characterizing the corresponding gauge symmetry in the compactified theory. These will become massive because at the UV end of the throat the symmetry will in general be broken. We find that, despite their string-scale mass term, they can make a contribution to SUSY breaking mediation which is equally important as that from modulus-anomaly mediation. In detail, from earlier analyses [4] one can see that the soft masses induced by modulus-anomaly mediation are of the order $m^2 \sim \omega^4 M_P^2$, where $\omega \ll 1$ is the relative redshift between the UV and IR ends of the throat. This is the same order of magnitude as we find for additional vector fields. Therefore in any realistic scenario it will be essential to take ‘vector mediation’ into account as one of the leading effects. By contrast, the contribution from heavy chiral superfields is suppressed by another factor $\omega^4$, so these are truly irrelevant for communicating SUSY breaking.

It would be desirable to have a concrete model in which the fields propagating in the throat are identified in terms of the supergravity solution and SUSY breaking mediation effects can be explicitly calculated. A step in this direction would be to understand properly the 5d SUSY description of the Klebanov-Strassler throat. Such an intermediate-scale 5d picture, applicable before going to 4d, could capture the geometric sequestering properties of the model, and yet be much simpler to analyze than the full 10d system (see e.g. [18,24] for 5d approaches to the Klebanov-Strassler throat). Recently developed methods in 5d off-shell supergravity [23] might prove useful for this analysis. The appropriate framework should be a 5d gauged supergravity, obtained as some flux-induced deformation of the $\text{AdS}_5 \times T^{1,1}$ theory anticipated in [22].

Once the 5d theory has been understood, the next challenge will be to find descriptions for the coupling to the SUSY breaking sector and the visible sector. We anticipate that both the UV manifold and the $\overline{D3}$ brane will lead to a nonlinear realization of the SO(4). How precisely a D-brane system (containing the visible sector fields) and an anti-brane in the throat (comprising the hidden sector) can couple to such a nonlinearly realized symmetry and what the 5d description of these coupling would look like appears to be highly non-trivial at this stage. It thus seems to be a long way to finally calculating soft terms.
However, some interesting questions may perhaps be raised and answered without knowing the details of the mechanism underlying a specific model. It would be interesting to see at a more quantitative level what vector mediation has to say about the tachyonic slepton problem of anomaly mediation. An important generalization of our analysis, so far conducted at the tree-level, would be the incorporation of loop effects [25].

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Appendix A: A simple model for a sequestered hidden sector in the throat

We take

\[ \Omega = -(T + \overline{T}) + \omega^2 (X \overline{X} - (X \overline{X})^2), \quad W = W_0 + e^{-T} + \omega^3 X. \]  

in the \( D = 4, N = 1 \) Lagrangian Eq. (3), so that its potential part becomes

\[
\mathcal{L} \supset - F_{\varphi} \overline{F_{\varphi}} (T + \overline{T}) - (F_T \overline{F_T} + \text{h.c.}) + \omega^2 F_{\varphi} \overline{F_{\varphi}} (X \overline{X} - (X \overline{X})^2) + \omega^2 F_X \overline{F_X} \\
+ \omega^2 (F_T F_X \overline{X} (1 - 2X \overline{X}) + \text{h.c.}) - 4\omega^2 F_X \overline{F_X} X \overline{X} \\
+ \left[ (3F_{\varphi}(W_0 + e^{-T} + \omega^3 X) - F_T e^{-T} + \omega^3 F_X) + \text{h.c.} \right].
\]  

The equations of motion read (note that those for \( F_T \) and \( \overline{T} \) are unchanged from Eqs. (6) and (7))

\[
F_T : \quad F_{\varphi} (T + \overline{T}) + F_T - \omega^2 F_{\varphi} (X \overline{X} - (X \overline{X})^2) - \omega^2 F_X \overline{X} (1 - 2X \overline{X}) \\
- 3W_0 - 3e^{-T} - 3\omega^3 \overline{X} = 0,
\]  

\[
F_{\varphi} : \quad F_{\varphi} + e^{-T} = 0,
\]  

\[
F_X : \quad \omega^2 F_X + \omega^2 F_{\varphi} X (1 - 2X \overline{X}) - 4\omega^2 F_X X \overline{X} + \omega^3 = 0,
\]  

\[
T : \quad F_{\varphi} \overline{F_{\varphi}} + 3F_{\varphi} e^{-T} - F_T e^{-T} = 0,
\]  

\[
X : \quad \omega^2 F_{\varphi} \overline{F_{\varphi}} (1 - 2X \overline{X}) + \omega^2 F_{\varphi} F_X (1 - 4X \overline{X}) - 2\omega^2 F_{\varphi} F_X X \overline{X} \\
- 4\omega^2 F_X F_X \overline{X} + 3\omega^3 F_{\varphi} = 0.
\]  

As before, from Eq. (36) and the condition that the pre-uplift superpotential in the vacuum should be \( \sim \omega^2 \), one can immediately see that \( F_{\varphi} \sim \omega^2 \). From Eq. (38) it follows that \( F_T \sim \omega^2 \), and from Eq. (37) we can deduce that \( F_X \sim \omega \).
Appendix B: Alternative estimation of chiral superfield contribution

First let us collect some supergravity relations for reference: The kinetic function $\Omega$ and the Kähler potential $K$ are related by

$$K = -3 \log(-\Omega/3),$$  \hspace{1cm} (40)

where we have reinstated a factor $1/3$ suppressed in the main text. Given the Lagrangian of Eq. (3), we are able to deduce the $F$ term for any chiral superfield from its equation of motion,

$$F^I = \frac{3}{\Omega} K^{IJ} (D_J W).$$  \hspace{1cm} (41)

Notice our non-standard normalization motivated by the chiral compensator formalism. The above calculation requires the chiral compensator $F$ term

$$F_{\varphi} = -\frac{1}{\Omega} (\Omega I F^I + 3\bar{W}).$$  \hspace{1cm} (42)

Here $W_I \equiv \partial_I W$, $\Omega_I \equiv \partial_I \Omega$, $K^{IJ}$ is the inverse of the Kähler metric $K_{IJ} \equiv \partial_I \partial_J K$, and the Kähler covariant derivative is $D_I \equiv \partial_I + K_I$. As opposed to the main text, we now distinguish between upper and lower chiral multiplet indices.

As in Sect. 4, consider a model with the Kähler modulus $T$, a chiral superfield $X$ localized at the bottom of the throat (the SUSY breaking sector), and a heavy chiral superfield $Y$ which couples to both $X$ and the visible sector. We take

$$\Omega = \Omega_0(T,Y) + \omega^2 \Delta \Omega(X,Y)$$  \hspace{1cm} (43)

and

$$W = W_0 + W_{\text{np}}(T,Y) + \bar{W}(Y) + \omega^3 \Delta W(X,Y).$$  \hspace{1cm} (44)

We have absorbed the superpotential contributions from all fields which have been integrated out in $W_0$. The non-perturbative part $W_{\text{np}}$ may in principle involve $Y$ as well as $T$. In order to have a vanishing or very small cosmological constant we should have $W_0 \sim W_{\text{np}} \sim \omega^2$ including partial derivatives, while $\Omega$ is generically of order one. As we will see later, the mass term $\bar{W}(Y) \sim Y^2$ is $O(\omega^4)$ in the vacuum.

We also have $\Omega_T \sim 1$, $\Omega_X \sim \omega^2$, $W_X \sim \omega^3$, and $W_T \sim \omega^2$. Furthermore, $K_{TT} \sim 1$, $K_{YT} \sim 1$, and $K_{XX} \sim \omega^2$. Hence the only term in the inverse Kähler metric that is actually warp-enhanced (rather than suppressed or $O(1)$) is $K^{XX} \sim \omega^{-2}$. This is true even if there is mixing between $X$ and the other fields in the Kähler potential, as can easily be seen by explicitly inverting $K_{IJ}$.

From Eq. (41b) it follows that

$$F^X \sim \omega + \bar{W}_Y K^{YY}, \quad F^Y \sim \omega^2 + \bar{W}_Y K^{YY}, \quad F^\tau \sim \omega^2 + \bar{W}_Y K^{YY}. \hspace{1cm} (45)$$

Here we kept the leading terms in the $\omega$ as well as the yet unspecified $\bar{W}_Y$ contributions.
We now consider the equation of motion for $Y$, replacing $F$ terms by the expressions given in Eqs. (42) and (45). The resulting equation for $\tilde{W}_Y$ has the structure

$$\omega^2 + \tilde{W}_Y + \tilde{W}_Y^2 = 0. \quad (46)$$

Recall that, by assumption, the perturbative superpotential $\tilde{W}(Y)$ contains no linear term in $Y$ and stabilizes $Y$ at zero as long as non-perturbative effects and SUSY breaking are ignored. This implies that, in the limit $\omega \to 0$ (where these effects are switched off), the vacuum value of $\tilde{W}_Y$ vanishes. From Eq. (46) it now follows that, in this limit,

$$\tilde{W}_Y \sim Y \sim \omega^2. \quad (47)$$

Then, we conclude from Eqs. (42) and (45) that

$$F_\phi \sim \omega^2, \quad F_X \sim \omega, \quad F_T \sim \omega^2. \quad (48)$$

At first glance one might also expect $F_Y \sim \omega^2$, but the leading terms in fact cancel. Indeed, returning to the equation of motion for $Y$ (this time using the above estimates of Eqs. (47) and (48) but leaving $F_Y$ unspecified), one finds

$$F_Y \sim \omega^4. \quad (49)$$

Thus, the final result for the $F$ terms agrees with Eq. (29).

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