Supersymmetry breaking and composite extra dimensions

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(Received 21 May 2001; published 28 February 2002)

We study supergravity models in *four dimensions* where the hidden sector is superconformal and strongly coupled over several decades of energy below the Planck scale, before undergoing spontaneous breakdown of scale invariance and supersymmetry. We show that large anomalous dimensions can suppress Kähler contact terms between the hidden and visible sectors, leading to models in which the hidden sector is ''sequestered'' and anomaly-mediated supersymmetry breaking can naturally dominate, thus solving the supersymmetric flavor problem. We construct simple, explicit models of the hidden sector based on supersymmetric QCD in the conformal window. The present approach can be usefully interpreted as having an extra dimension responsible for sequestering replaced by the many states of a (spontaneously broken) strongly coupled superconformal hidden sector, as dictated by the anti–de Sitter conformal field theory correspondence.

DOI: 10.1103/PhysRevD.65.066004 PACS number(s): 11.25.Mj, 04.65.+e, 11.10.Kk

I. INTRODUCTION

The anti–de Sitter (AdS) conformal field theory (CFT) correspondence [1] asserts that a gravity theory on 5D AdS space is ''dual'' to a 4D CFT. The duality takes the form of an equality of generating functionals depending on some 4D field h_0 that act as boundary values for the gravity fields on AdS and source terms for operators in the CFT:

$$\int_{h|_{\rm bdy}=h_0} d[h] e^{iS_{\rm grav}[h]} = \langle e^{i\int h_0 \cdot \mathcal{O}} \rangle_{\rm CFT}. \tag{1.1}$$

This correspondence remains at present an unproven conjecture, but in string theory realizations it passes an impressive number of quantitative and qualitative consistency checks, and has proved extremely fruitful in suggesting new connections [2].

It has been argued in Refs. [3,4,5] that this duality can be extended to the equivalence between the 5D ''brane world'' scenario of Randall and Sundrum (RS) [6] and a conformal field theory perturbed by four-dimensional gravity. In this duality the ''UV brane'' in the RS model where gravity is localized is mapped to a UV cutoff for the CFT and the redshifted ''IR brane'' is mapped to spontaneous breaking of conformal invariance [4,5], which provides an IR cutoff of the CFT. As stressed in Ref. [4], both sides of this perturbed duality are macroscopically 4D theories with a discrete spectrum, and hence the equivalence reduces simply to the statement that both theories give identical predictions for all physical quantities such as the S matrix. This conjecture also passes several quantitative checks [7] and many qualitative ones [4,5].

In this paper, we will study 4D CFT's that can be viewed as being dual to the RS model with supersymmetry (SUSY) [8,9,10]. The phenomenological motivation for this class of models is quite different from the original RS model, where the redshift factor between the UV and IR branes was used to explain the hierarchy between the Planck and weak scales. In the SUSY RS model, it is SUSY that solves the hierarchy problem. The motivation for the extra dimension comes from the SUSY flavor problem. If we assume that there is no flavor symmetry at the Planck scale, then the low-energy effective theory necessarily includes contact terms of the form

$$\Delta \mathcal{L}_{\rm eff} \sim \int d^4\theta \frac{1}{M_{\rm Pl}^2} \Sigma^\dagger \Sigma Q^\dagger Q, \tag{1.2}$$

where Σ is the hidden sector field that breaks SUSY and Q is a visible sector field. Such terms cannot be forbidden by any symmetry, and give a contribution to the squark masses of order $m_{3/2} \sim \langle F_\Sigma \rangle / M_{\rm Pl}$ that has no reason to be flavor diagonal. In hidden sector models where soft SUSY breaking parameters are of order $m_{3/2}$, it is therefore difficult to understand why the squark masses are nearly flavor independent, as required by constraints on flavor-changing neutral currents. Reference [11] showed that this problem can be solved if the visible and hidden sectors are localized on different branes, separated in an extra dimension. The spatial separation suppresses contact terms between the hidden and visible sectors [11,12]. In Ref. [11], this was referred to as ''sequestering'' the hidden sector. If there are no massless fields in the bulk other than supergravity, SUSY breaking is communicated from the supergravity sector to the visible sector by anomaly mediation [11,13] (after radion stabilization [12]).[1] Assuming that the visible sector contains only the minimal supersymmetric standard model (MSSM) and assuming no

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[1]If the visible sector gauge fields propagate in the bulk, this scenario leads to gaugino mediated SUSY breaking [14] or radion mediated SUSY breaking [15]. The dual CFT description of these mechanisms will be discussed elsewhere [16].

other couplings between the visible and hidden sectors implies that slepton mass-squared terms are negative, but by relaxing these assumptions realistic and predictive models have been constructed that preserve the attractive features of the scenario [17].

The AdS-CFT correspondence extended to the SUSY RS setup asserts that anomaly-mediated models can be realized in a purely 4D theory. The necessary ingredients in the 4D theory are determined simply by following the correspondence. The bulk supergravity modes in the 5D description are mapped to a strongly coupled superconformal field theory (SCFT) in the 4D theory. Visible sector fields localized on the UV brane in the 5D description are mapped to elementary fields in the 4D theory that are coupled to the SCFT only through Planck-suppressed operators. Hidden sector fields localized on the IR brane in the 5D description are mapped to composites of the SCFT that arise from the spontaneous breaking of conformal invariance. Stabilization of the extra-dimensional radius is mapped to the stabilization of the modulus in the CFT responsible for spontaneous breakdown of scale invariance. Finally, the condition on the 5D theory that there are no light bulk modes other than supergravity responsible for transmitting supersymmetry breaking is mapped to the condition that only irrelevant SCFT operators couple to the visible fields.

We will be interested in strongly coupled SCFT's with no expansion parameters, such as large N or large 't Hooft parameter. AdS duals of such theories are not known, presumably because there is no parametric separation between the string length, the AdS radius, and the Planck length. In order to achieve sequestering in such a theory by decoupling the effects of massive bulk states, we require an extra dimension that extends over several AdS radii. At the level of effective field theory, the existence of such strongly warped SUSY models was demonstrated in Refs. [8]. In Ref. [9] it was shown that radius stabilization and anomaly-mediated SUSY breaking can be realized in this scenario.

The purpose of this paper is to show that sequestering and anomaly mediation are indeed realized in a large class of 4D SUSY theories. The theories can be explicitly constructed and understood from a purely 4D perspective, and demonstrate that sequestering can be realized without positing extra dimensions or branes. However, we find the dual 5D description, where sequestering has a simple geometrical origin, very illuminating. We therefore refer to this class of models as "composite extra dimensions."

Reference [18] considered SUSY models where the MSSM has superpotential couplings to a strong SCFT and studied implications for flavor and SUSY breaking. References [19] constructed non-SUSY gauge theories whose low-energy dynamics mimics that of a theory with an extra dimension. We will briefly discuss the relation of these papers to our work in the conclusions.

II. CFT SUPPRESSION OF HIDDEN-VISIBLE CONTACT TERMS

The prototype of the kind of SCFT to which our results apply is the strongly coupled fixed point of SU(N) SUSY QCD with F flavors found by Seiberg [20]. For $\frac{3}{2}N<F<3N$, this theory is asymptotically free in the UV but has a nontrivial conformal fixed point in the IR. For $F\simeq 3N$ the IR fixed point is weakly coupled [21], and for $F\simeq\frac{3}{2}N$ the IR fixed point has a weakly coupled dual description [20], but in the middle of the range the IR fixed point has no known weakly coupled Lagrangian description. In what follows, we will write our results for the special case $N=2$, $F=4$ for simplicity. In that case there are eight SU(2) fundamentals T^J, $J=1,\ldots,8$. We define $\Lambda_{\rm CFT}$ to be the scale below which the theory is in the IR conformal regime, and above which the theory rapidly runs to its asymptotically free regime.

The crucial question is the size of scalar masses induced by flavor-violating couplings of the form

$$\Delta\mathcal{L}_{\rm eff}=\int d^4\theta\,\frac{c^j{}_k}{M_{\rm Pl}^2}T_J^\dagger T^J Q_j^\dagger Q^k, \tag{2.1}$$

where T is a hidden sector SUSY QCD field. Note that we have assumed that the coupling is diagonal in hidden flavor, which can be made natural by imposing (discrete and/or gauged) flavor symmetries on the hidden sector. In order for anomaly mediation to dominate, we require that this term contribute visible scalar masses $\Delta m_{\tilde{Q}}^2 \lesssim 10^{-7}V_{\rm hid}/M_{\rm Pl}^2$ (see Sec. III), where $V_{\rm hid}$ is the SUSY breaking vacuum energy. We will discuss the mechanism of this SUSY breaking in the hidden sector in the next section. The suppression factor of 10^{-7} must arise from the nontrivial CFT scaling of the operator Eq. (2.1).

This term can be viewed as a correction to the kinetic term for the T fields in the UV Lagrangian:

$$\mathcal{L}_{\rm UV}=\int d^4\theta Z_0 T_J^\dagger T^J+\cdots,\quad Z_0=1+\frac{c^j{}_k}{M_{\rm Pl}^2}Q_j^\dagger Q^k. \tag{2.2}$$

This contributes to a perturbation of the physical gauge coupling g^2, given by [22]

$$\frac{1}{g^2}=\frac{1}{g_{\rm hol}^2}-\frac{N}{8\pi^2}\ln g^2-\frac{F}{8\pi^2}\ln Z+\text{const}+\mathcal{O}(g^2), \tag{2.3}$$

where $1/g_{\rm hol}^2$ is the holomorphic gauge coupling that appears as the coefficient of the gauge kinetic term in the Lagrangian. Because Eq. (2.2) is a perturbation to the UV gauge coupling, it is necessarily irrelevant near the fixed point. This is simply because the theory near a fixed point must be insensitive to UV couplings in order to be IR attractive, as Seiberg argued is the case in SUSY QCD.

To make this quantitative, let us consider how the operator Eq. (2.2) runs down to the IR in two stages. First, the running down to the scale $\Lambda_{\rm CFT}$ where the theory becomes strong is the standard logarithmic running in the far UV, together with an order unity strong-interaction renormalization near $\Lambda_{\rm CFT}$, $\Lambda_{\rm CFT}$ is defined such that $g(\Lambda_{\rm CFT})$ is a *fixed number* close enough to the fixed point coupling g_* that below this scale we can expand about the fixed point:

$$\beta = \beta'_* \cdot (g^2 - g_*^2) + \cdots, \quad \gamma = \gamma_* + \gamma'_* \cdot (g^2 - g_*^2) + \cdots, \tag{2.4}$$

where $\beta \equiv dg^2/d\ln\mu$, $\gamma \equiv d\ln Z/d\ln\mu$. The anomalous dimension at the fixed point is determined by the (nonanomalous) $\mathrm{U}(1)_R$ symmetry to be

$$\gamma_* = \tfrac{1}{2}. \tag{2.5}$$

Integrating these renormalization group equations from $\Lambda_{\rm CFT}$ down to μ then gives

$$Z(\mu) = Z(\Lambda_{\rm CFT}) \left(\frac{\mu}{\Lambda_{\rm CFT}} \right)^{\gamma_*} \left\{ 1 + \frac{\gamma'_*}{\beta'_*} [g^2(\Lambda_{\rm CFT}) - g_*^2] \left[\left(\frac{\mu}{\Lambda_{\rm CFT}} \right)^{\beta'_*} - 1 \right] + \cdots \right\}. \tag{2.6}$$

We can rewrite this using Eq. (2.3) evaluated at $\Lambda_{\rm CFT}$ and the fact that $g_{\rm hol}$ has exact one-loop running,

$$Z(\mu) = \mathrm{const} \times \left(\frac{\mu}{|\Lambda_{\rm hol}|} \right)^{\gamma_*} \left\{ 1 + \frac{\gamma'_*}{\beta'_*} [g^2(\Lambda_{\rm CFT}) - g_*^2] \left[\left(\frac{\mu}{\Lambda_{\rm CFT}} \right)^{\beta'_*} - 1 \right] + \cdots \right\}, \tag{2.7}$$

where

$$\Lambda_{\rm hol} \equiv \mu e^{-4\pi^2/g_{\rm hol}^2(\mu)} \tag{2.8}$$

is the holomorphic one-loop strong-interaction scale.

Equation (2.7) is useful because it shows that the leading dependence on Z_0 [contained in $Z(\Lambda_{\rm CFT})$] has disappeared in the IR below $\Lambda_{\rm CFT}$. The subleading dependence on Z_0 is implicit in the dependence on $\Lambda_{\rm CFT}$, which is suppressed by the power $\beta'_* > 0$. [Note that $\Lambda_{\rm CFT}$ depends on Z_0 through its definition, $g^2(\Lambda_{\rm CFT}) = \mathrm{const}$, and also Eq. (2.3).] The fixed point behavior is cut off at the scale $\mu = u_{\rm CFT}$ where the conformal symmetry is spontaneously broken, so we obtain the required suppression for $(v_{\rm CFT}/\Lambda_{\rm CFT})^{\beta'_*} \lesssim 10^{-7}$.

It is crucial that the perturbation Eq. (2.2) is a singlet under the hidden sector flavor symmetries. For a nonsinglet operator of the form

$$\Delta\mathcal{L} = \int d^4\theta \frac{c^{jJ}{}_{kK}}{M_{\rm Pl}^2} T_J^\dagger T^K Q_j^\dagger Q^k, \tag{2.9}$$

with $c^{jJ}{}_{kJ} = 0$, the contribution to the visible sector scalar masses is

$$(\Delta m_{\tilde{Q}}^2)^j{}_k = -\frac{c^{jJ}{}_{kK}}{M_{\rm Pl}^2} \left\langle \int d^4\theta T_J^\dagger T^K \right\rangle. \tag{2.10}$$

This is a matrix element of a conserved current supermultiplet with vanishing anomalous dimension, so there is no CFT suppression of operators of the form Eq. (2.9). A model-building requirement is therefore to insist on enough symmetry in the hidden dynamics to prohibit such operators.

In the SUSY limit the theory above has a moduli space of vacua, and away from the origin of moduli space the conformal symmetry is spontaneously broken. The light fields below the scale $v_{\rm CFT}$ where the conformal symmetry is broken are the moduli, which can be thought of as composites of the CFT. We now consider the effective field theory for these moduli. The moduli space can be parametrized by the gauge-invariant holomorphic operators [23]. In the SU(2) gauge theory we are considering, the moduli space is parametrized completely by the "meson" invariants

$$M^{JK} = T^J T^K = -M^{KJ}. \tag{2.11}$$

From Eq. (2.7) we see that the dependence on $\Lambda_{\rm hol}$ can be eliminated by working in terms of the renormalized fields

$$T'^J = \frac{T^J}{(\Lambda_{\rm hol})^{1/4}}. \tag{2.12}$$

Since $\Lambda_{\rm hol}$ parametrizes the only explicit scale in the hidden dynamics, and since the new fields eliminate dependence on this scale near the IR fixed point, the leading low-energy interactions must be given by

$$\mathcal{L}_{\rm eff} = \int d^4\theta f(M', M'^\dagger) + \mathcal{O}(\partial^4) + \mathcal{O}(\Lambda_{\rm CFT}^{-\beta'_*}), \tag{2.13}$$

where $M' = T'T'$ and f is a homogeneous function with its degree determined by dimensional analysis:

$$f(\alpha M', \alpha M'^\dagger) = \alpha^{4/3} f(M', M'^\dagger). \tag{2.14}$$

The new fields are very convenient later because they have the same canonical dimension as their scaling dimension in superpotential terms, allowing us to simultaneously nonlinearly realize the asymptotic (canonical) scale invariance in the UV and the nontrivial asymptotic scale invariance in the IR.

There is another way to derive the absence of Z_0 dependence in the leading terms of the low-energy theory, Eq. (2.13), that may be illuminating. We can regard Z_0 as a background gauge connection for an anomalous $\mathrm{U}(1)_A$ symmetry [24]. Because of the anomaly, $\Lambda_{\rm hol}$ is charged under this symmetry with a charge such that the renormalized fields are uncharged under $\mathrm{U}(1)_A$. The leading terms in the low-energy effective theory are therefore independent of the $\mathrm{U}(1)_A$ gauge connection. By contrast, perturbations of the form Eq. (2.9) can be regarded as background gauge connections for nonanomalous hidden flavor symmetries. Therefore $\Lambda_{\rm hol}$ is uncharged under these gauge connections and cannot cancel their effects in the low-energy theory, Eq. (2.13).

III. A REALISTIC MODEL

We now show how to construct a realistic 4D model in which SUSY breaking is communicated by anomaly mediation, with the suppression of contact terms explained by the mechanism described above. Our aim is to construct a model that illustrates the issues in constructing a realistic model, and separates these issues as clearly as possible. The model we discuss contains several explicit small superpotential couplings whose origin is not explained. We believe that completely natural models without fundamental small parameters are possible, but we leave their construction for future work.

The hidden sector will be taken to be SU(2) gauge theory with eight fundamentals T^J $(J=1,...,8)$, as discussed in Sec. II. The classical moduli space of this theory can be parametrized by the ''meson'' operators M^{JK} [see Eq. (2.11)]. M is an antisymmetric matrix with rank 2, which can be conveniently parametrized by

$$M=\begin{pmatrix} X\epsilon & -Y^T \\ Y & \mathcal{O}(Y^2/X)\end{pmatrix},\quad \epsilon=\begin{pmatrix} 0 & 1\\ -1 & 0\end{pmatrix}, \tag{3.1}$$

where X and Y are unconstrained. Note that this has 13 complex degrees of freedom, as required given the 16 quark fields and the three D-flatness conditions. We will expand about the vacuum

$$\langle M\rangle=\begin{pmatrix} \langle X\rangle\epsilon & 0\\ 0 & 0\end{pmatrix}. \tag{3.2}$$

It will be convenient to further parametrize

$$Y=\begin{pmatrix} \Sigma\epsilon+\Pi \\ \Pi'\end{pmatrix},\quad \mathrm{tr}(\epsilon\Pi)=0. \tag{3.3}$$

Upon adding superpotential terms we will see that SUSY is broken by $\langle F_\Sigma\rangle\neq 0$.

The low-energy effective theory below the scale determined by the VEV $\langle M\rangle$ was described in Sec. II. The contribution to the effective Kähler potential is a homogeneous function of the meson fields of degree $\frac{4}{3}$, which must also be SU(8) invariant. We will always work in terms of the ''renormalized'' primed fields of the previous section, dropping the primes. The vacuum expectation value (VEV) breaks the SU(8) global symmetry to SU(2)×SU(6), so expanding in powers of Y gives

$$K_{\rm eff}=a_0(X^\dagger X)^{2/3}\left[1+a_1\frac{\mathrm{tr}(Y^\dagger Y)}{X^\dagger X}+\mathcal{O}(|Y|^4/|X|^2)\right], \tag{3.4}$$

where $a_{0,1}$ are unknown strong interaction parameters. It is convenient to work in terms of redefined fields

$$\hat{X}=a_0^{1/2}X^{3/2},\quad \hat{Y}=a_0^{3/4}a_1^{1/2}\frac{Y}{\hat{X}^{1/2}}, \tag{3.5}$$

which have a canonical Kähler potential:

$$K_{\rm eff}=\hat{X}^\dagger\hat{X}+\hat{Y}^\dagger\hat{Y}+\mathcal{O}(|\hat{Y}|^4/|\hat{X}|^2). \tag{3.6}$$

As described so far, the model has unbroken SUSY and a moduli space of vacua. We now add superpotential terms that stabilize the moduli and break SUSY. We will take our model, and the superpotential in particular, to respect an SU(2) subgroup of the global SU(8) symmetry. [This flavor SU(2) symmetry can be weakly gauged but we will not consider this here.] For convenience we give names to the four SU(2) doublets as follows:

$$\begin{aligned} P^{1,2}&=T^{1,2},\quad \overline{P}^{1,2}=T^{3,4},\\ N^{1,2}&=T^{5,6},\quad \overline{N}^{1,2}=T^{7,8},\end{aligned} \tag{3.7}$$

so that

$$X=P^jP_j,\quad \Sigma=P^j\overline{P}_j, \tag{3.8}$$

where j, $k=1$, 2 are global SU(2) indices, and we have defined $P_j\equiv\epsilon_{jk}P^k$, etc. In addition to the global SU(2) we impose the following discrete symmetries:

$$\text{(i)}\quad P\leftrightarrow\overline{P},\quad N\leftrightarrow\overline{N}, \tag{3.9}$$

$$\text{(ii)}\quad P\leftrightarrow N,\quad \overline{P}\leftrightarrow\overline{N}, \tag{3.10}$$

$$\text{(iii)}\quad P\mapsto iP,\quad \overline{P}\mapsto -i\overline{P},\quad N\mapsto iN,\quad \overline{N}\mapsto -i\overline{N}, \tag{3.11}$$

$$\text{(iv)}\quad P\mapsto iP,\quad \overline{P}\mapsto -i\overline{P},\quad N\mapsto -iN,\quad \overline{N}\mapsto i\overline{N}. \tag{3.12}$$

These symmetries ensure that the only allowed term of the form $T^\dagger T$ is (accidentally) SU(8) invariant. Therefore the only Kähler term of the form $T^\dagger TQ^\dagger Q/M_{\rm Pl}^2$ is a singlet of the CFT flavor symmetries, and is suppressed by the renormalization group arguments of Sec. II.

These symmetries allow us to add the following terms to the superpotential:

$$W=W_{\rm stab}+W_{\rm mass}+W_{\rm Polonyi}, \tag{3.13}$$

where

$$\begin{aligned} W_{\rm stab}=&\tfrac{1}{2}\lambda_1[(P^jP_j)^2+(\overline{P}^j\overline{P}_j)^2+(N^jN_j)^2+(\overline{N}^j\overline{N}_j)^2]\\ &+\tfrac{1}{4}\lambda_2[(P^jP_j)^4+(\overline{P}^j\overline{P}_j)^4+(N^jN_j)^4+(\overline{N}^j\overline{N}_j)^4],\end{aligned} \tag{3.14}$$

$$W_{\rm mass}=c_1\sum_{a=1}^{3}[(P^j\sigma^a{}_j{}^k\overline{P}_k)^2+(N^j\sigma^a{}_j{}^k\overline{N}_k)^2]$$
$$+c_2[(P^jN^k)(P_jN_k)+(\overline{P}^j\overline{N}^k)(\overline{P}_j\overline{N}_k)$$
$$+(\overline{P}^jN^k)(\overline{P}_jN_k)+(P^j\overline{N}^k)(P_j\overline{N}_k)], \tag{3.15}$$

$$W_{\rm Polonyi}=\kappa[P^j\overline{P}_j+N^j\overline{N}_j], \tag{3.16}$$

and where $\sigma^{1,2,3}$ are the Pauli matrices for flavor SU(2). Gauge-singlet operators are enclosed by parentheses. At the scale $\Lambda_{\rm CFT}$ the SU(2) gauge interactions become strong and the theory flows to a nontrivial conformal fixed point. At this point the scaling of the operators is controlled by the nontrivial fixed point. We will assume that

$$\Lambda_{\rm CFT}\sim\Lambda_{\rm UV}\sim 4\pi M_{\rm Pl}, \tag{3.17}$$

where $\Lambda_{\rm UV}$ is the scale where 4D quantum gravity becomes strong. We will work in units where $M_{\rm Pl}=1$.

Below the scale $\langle\hat{X}\rangle$ the conformal symmetry is spontaneously broken and the effective degrees of freedom of the CFT are the moduli. Writing the superpotential in terms of the moduli fields defined in Eqs. (3.1), (3.3), and (3.5) we have

$$W_{\rm stab}=\tfrac{1}{2}\lambda_1[\hat{X}^3+\mathcal{O}(\hat{\Sigma}^4/\hat{X})+\mathcal{O}(\hat{\Pi}^4/\hat{X})]$$
$$+\tfrac{1}{4}\lambda_2[\hat{X}^6+\mathcal{O}(\hat{\Sigma}^8/\hat{X}^2)+\mathcal{O}(\hat{\Pi}^6/\hat{X}^2)], \tag{3.18}$$

$$W_{\rm Polonyi}=\kappa\hat{X}^{1/2}\hat{\Sigma}, \tag{3.19}$$

while $W_{\rm mass}$ is a sum of mass terms for every component of Π and Π'. In the above, we have absorbed the unknown strong interaction coefficients of the Kähler potential into redefinitions of the superpotential couplings [see Eqs. (3.4) and (3.5)].

The CFT running of the superpotential perturbations is automatically taken into account by our working in terms of the primed fields defined in Sec. II. Our discussion assumes that the superpotential terms can be treated as linear perturbations to the CFT in the energy range from $\Lambda_{\rm UV}$ to $\langle\hat{X}\rangle$. The couplings λ_1, c_1, and c_2 are marginal and have dimensionless coefficients small compared to 1. There are nonlinear corrections suppressed by higher powers of these couplings, but these are negligible logarithmic corrections similar to those found in weak-coupling perturbation theory. The coupling λ_2 is irrelevant, and can therefore also be treated as a perturbation.

We now determine the VEVs. The stabilization term Eq. (3.18) gives rise to a local SUSY-preserving minimum with

$$\langle\hat{X}\rangle^3=-\frac{\lambda_1}{\lambda_2}. \tag{3.20}$$

(We will consider the effects of SUSY breaking on $\hat{X}$ below.) This stabilizes the modulus and determines the scale of spontaneous breaking of the conformal symmetry. The mass of the X field is of order

$$m_X\sim\lambda_1\langle\hat{X}\rangle. \tag{3.21}$$

Since $\lambda_2\lesssim 1$ and we want $\langle\hat{X}\rangle\ll 1$ (in Planck units) we must have $\lambda_1\ll 1$, and hence $m_X\ll\langle\hat{X}\rangle$. We will take $\lambda_2\sim 1$ in what follows.[2]

Below the scale m_X we integrate out X and consider the effective Lagrangian of the remaining light degrees of freedom. We chose the couplings in $W_{\rm mass}$ so that all of the modes parametrized by Π and Π' get masses $m_\Pi\lesssim m_X$. The only remaining degree of freedom below the scale m_Π is then Σ. The effective superpotential is then

$$W_{\rm eff}=\kappa\langle\hat{X}\rangle^{1/2}\hat{\Sigma}+\mathcal{O}(\kappa^2\hat{\Sigma}^2/\langle\hat{X}\rangle^5), \tag{3.22}$$

where the terms higher order in $\hat{\Sigma}$ come from the $\hat{X}$ dependence in Eq. (3.19). The effective Kähler potential is

$$K_{\rm eff}=\hat{\Sigma}^\dagger\hat{\Sigma}+c\,\frac{(\hat{\Sigma}^\dagger\hat{\Sigma})^2}{|\langle\hat{X}\rangle|^2}+\cdots, \tag{3.23}$$

where c is an order one unknown strong interaction coefficient. If $c<0$, this theory has a (local) SUSY breaking minimum with

$$\langle F_{\hat{\Sigma}}\rangle\sim\kappa\langle\hat{X}\rangle^{1/2}[1+\mathcal{O}(\kappa^2/\hat{X}^9)],\quad \langle\hat{\Sigma}\rangle=\mathcal{O}(\kappa/\hat{X}^{7/2}), \tag{3.24}$$

and Σ gets a mass

$$m_\Sigma\sim\frac{\langle F_{\hat{\Sigma}}\rangle}{\langle\hat{X}\rangle}. \tag{3.25}$$

Here we will simply make the dynamical assumption that $c<0$. The condition that the higher order terms make only a small fractional correction to the SUSY breaking order parameter $F_{\hat{\Sigma}}$ is

$$\frac{\langle F_{\hat{\Sigma}}\rangle^2M_{\rm Pl}^6}{\langle\hat{X}\rangle^{10}}\lesssim 1. \tag{3.26}$$

This discussion assumes that $\hat{X}$ is sufficiently heavy that we can integrate it out for purposes of SUSY breaking. We also assumed that SUSY breaking does not significantly shift the $\langle\hat{X}\rangle$ away from its SUSY value. It is easily checked that both of these constraints are equivalent to Eq. (3.26).

We are now ready to check the numbers. The most stringent constraints on flavor-changing neutral currents arise from $K^0-\overline{K}^0$ mixing [27]:

$$\frac{m^2_{\tilde{d}\tilde{s}}}{m^2_{\tilde{s}}}\lesssim(6\times 10^{-3})\left(\frac{m_{\tilde{s}}}{1\ \text{TeV}}\right), \tag{3.27}$$

[2]This is in fact conservative, since naive dimensional analysis [26] allows a larger coefficient.

$$\mathrm{Im}\left(\frac{m^2_{\tilde{d}\tilde{s}}}{m^2_{\tilde{s}}}\right) \lesssim (4\times 10^{-4})\left(\frac{m_{\tilde{s}}}{1\ \mathrm{TeV}}\right).$$

We assume squark masses of order 1 TeV in the following. Anomaly mediation gives a flavor-independent mass to squarks of order

$$m_{\tilde{q}} \sim 2\times 10^{-2} F_\phi, \tag{3.28}$$

where $F_\phi \sim m_{3/2} \sim \langle F_{\hat{\Sigma}}\rangle / M_{\mathrm{Pl}}$. This fixes

$$\langle F_{\hat{\Sigma}}\rangle \sim 8\times 10^{22}\ \ \mathrm{GeV}^2. \tag{3.29}$$

The required suppression of FCNC's is obtained provided

$$\left(\frac{\langle \hat{X}\rangle}{\Lambda_{\mathrm{UV}}}\right)^{\beta'_*} \lesssim 2\times 10^{-7} \tag{3.30}$$

using the stronger *CP*-violating constraint. The constraint Eq. (3.26) then gives

$$\langle \hat{X}\rangle \gtrsim 4\times 10^{15}\ \ \mathrm{GeV}. \tag{3.31}$$

The constraint Eq. (3.30) is satisfied provided

$$\beta'_* \gtrsim 1.7. \tag{3.32}$$

As discussed above, β'_* is a nonperturbatively determined exponent which we cannot calculate. Naive dimensional analysis [26] tells us that $\beta'_* \sim 1$. Extrapolations using perturbation theory valid for the Banks-Zaks fixed points [25], $1-F/(3N) \ll 1$, suggest that $\beta'_* \simeq 1$ at the self-dual point, $F=2N$. In the absence of more rigorous information, we believe that values such as this are very reasonable. In fact we are able to construct models that allow smaller values of β'_* than Eq. (3.32) by using stabilizing superpotentials with smaller powers of X. In the present model, such powers are forbidden by discrete symmetries, but we can add more fields and couplings that spontaneously break these symmetries and generate lower powers of X below the breaking scale. The analysis of such models is slightly more involved than our present model and will not be detailed here.

Another dynamical assumption required in this model is that the uncalculable strong Kähler corrections have the right sign ($c<0$) to give a local SUSY breaking vacuum at $\langle\hat{\Sigma}\rangle=0$. This dynamical assumption can be avoided by replacing W_{Polonyi} by an O'Raifeartaigh sector with additional singlet fields. Basically, the additional fields in the O'Raifeartaigh sector give larger calculable Kähler corrections than the uncalculable Kähler corrections if these singlets are sufficiently light. If some of the additional singlet fields are elementary, one must ensure that they do not get substantial F terms, since (standard model flavor violating) contact interactions between these fields and the visible fields are not suppressed. We have constructed explicit models of this type.

So far we have considered the hidden dynamics in flat spacetime, showing how to stabilize the moduli and break SUSY. Because the energy scales and VEVs in the hidden sector are much smaller than the Planck scale, the main effect of coupling the hidden sector to supergravity is that the SUSY breaking contribution to the cosmological constant can be canceled in the usual way. Supergravity has a very small effect on the hidden dynamics and vacuum stabilization. The main effect of coupling supergravity to the visible sector is that SUSY breaking is communicated to the visible sector by the mechanism of anomaly mediation.

It is straightforward to adapt proposals in the literature [17] for solving the tachyonic slepton problem and μ problem of anomaly mediation to the present framework.

Note that the stabilizing superpotential has another supersymmetric solution, $X=0$. At this point on the moduli space the theory remains superconformal and supersymmetric, and therefore has lower energy than the local minimum, Eq. (3.20). In other words our supersymmetry-breaking vacuum is only metastable. However we have checked that tunneling to the true supersymmetry-preserving vacuum is highly suppressed over cosmological time scales, just as in the SUSY breaking scenario of Ref. [28].

IV. CONCLUSIONS

The main result of this paper is that it is possible to construct 4D SUSY field theories that realize the sequestering of the hidden sector. The original sequestering mechanism of Ref. [11] had its origin in the spatial separation of the visible and hidden sectors in an extra dimension. In the four-dimensional models considered here the role of the extra dimension is played by the many states of a superconformal field theory, as dictated by the AdS-CFT correspondence. Using these ideas we have constructed an explicit realistic 4D model in which anomaly mediation dominates in the visible sector.

The higher-dimensional realization of sequestering is geometric and highly intuitive. However, the local higher-dimensional ($\mathcal{N}=2$) SUSY is a significant technical complication that makes the construction of explicit models difficult. In the 4D models considered here the extra supersymmetry is implicit in the enhanced superconformal symmetry of the fixed point, and we only need to keep track of $\mathcal{N}=1$ SUSY for model building.

There are many interesting further directions to pursue. In future work [16], we intend to extend the ideas of this paper to study 4D realizations of gaugino mediation [14] and radion mediation [15], where the hidden sector has a superconformal regime (dual to having the hidden sector on the IR brane in a SUSY RS setup). We also wish to consider the important question of constructing fully natural models with a dynamical origin for scale hierarchies.

We end with some comments on related work that has appeared recently. Reference [18] studied strong SUSY CFT's applied to the flavor problem and SUSY breaking. Although the relation to the AdS-CFT correspondence was not discussed in this paper, for purposes of SUSY breaking the models of Ref. [18] can be viewed as the CFT dual of gaugino mediation [14], with the hidden sector localized on the UV brane. There are some important differences between this work and that of the present paper, beyond the obvious difference in the SUSY-breaking mediation mechanism. In Ref. [18] the conformal symmetry is broken by relevant operators, and suppressing all soft terms requires flavor symmetries in the standard model to be completely broken. While the scenario of Ref. [18] implements a specific proposal for understanding the structure of Yukawa couplings as well as giving a solution to the SUSY flavor problem, our present work is aimed only at the SUSY flavor problem. On the other hand, realistic model building appears to be simpler in the present approach where the hidden sector originates from the CFT (dual to having the SUSY breaking on the IR brane).

Reference [19] gave a very simple and explicit construction of gauge theories whose low-energy dynamics mimics that of a flat extra dimension without gravity. In this approach the many states of the extra dimension arise from having many four-dimensional gauge sectors, while in our approach they arise from the excited states of a simple CFT. In the framework of Refs. [19] sequestering is difficult to realize because it is not clear how to maintain locality in the extra-dimensional interpretation in the presence of gravity.

ACKNOWLEDGMENTS

M.A.L. was supported by NSF Grant No. PHY-98-02551.

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