Anomaly mediated supersymmetry breaking in four dimensions, naturally

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We present a simple four-dimensional model in which anomaly mediated supersymmetry breaking naturally dominates. The central ingredient is that the hidden sector is near a strongly coupled infrared fixed point for several decades of energy below the Planck scale. Strong renormalization effects then sequester the hidden sector from the visible sector. Supersymmetry is broken dynamically and requires no small input parameters. The model provides a natural and economical explanation of the hierarchy between the supersymmetry-breaking scale and the Planck scale, while allowing anomaly mediation to address the phenomenological challenges posed by weak scale supersymmetry. In particular, flavor-changing neutral currents are naturally near their experimental limits.

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I. INTRODUCTION

Anomaly mediated supersymmetry breaking (AMSB) [1,2] is a general supergravity mechanism that is tightly constrained by local supersymmetry. AMSB may play an important role in solving the major phenomenological problems of weak scale supersymmetry (SUSY): the flavor, μ, and gaugino mass problems. For example, Ref. [3] describes a complete and very plausible extended supersymmetric standard model where anomaly mediation is the main ingredient in solving these problems and leads to a realistic and distinctive spectrum. References [4] describe other proposals for weak scale AMSB.

In order for AMSB to dominate in the observable sector, SUSY breaking must originate in a special type of hidden sector. A general hidden sector model has the form

$$\mathcal{L}=\mathcal{L}_{\text{SUGRA}}+\mathcal{L}_{\text{visible}}+\mathcal{L}_{\text{hidden}}+\mathcal{L}_{\text{mixed}}, \tag{1.1}$$

where the first three terms are self-explanatory (SUGRA indicates supergravity), while $\mathcal{L}_{\text{mixed}}$ contains Planck-suppressed terms involving both visible and hidden fields that cannot naturally be forbidden by symmetries. AMSB in the visible sector arises from minimal coupling to supergravity, in particular the auxiliary scalar field in the minimal formulation. By supercovariance, this scalar couples via visible mass scales, in particular the renormalization scale associated with the scale anomaly in radiative corrections (hence "anomaly mediation"). Therefore, supersymmetry-breaking effects arising from AMSB are suppressed by loop factors. In general hidden sector models, larger visible SUSY breaking can arise directly from the hidden sector through terms in $\mathcal{L}_{\text{mixed}}$. Therefore, in order for AMSB to dominate $\mathcal{L}_{\text{mixed}}$ must be strongly suppressed.[1] That is, the hidden and visible sectors are "sequestered."

In Refs. [1,5] it was shown that sequestering can be achieved if the visible and hidden sectors are localized on different three-branes separated in extra dimensions. Recently, we demonstrated that highly warped supersymmetric anti–de Sitter (AdS) space compactifications could be stabilized with sufficient sequestering [6]. AdS/conformal field theory (CFT) duality [7] applied to such compactifications [8] then suggests that sequestering can also arise in a purely 4D context with the help of strongly coupled conformal dynamics. In Ref. [9], we showed that sequestering in fact occurs in a large class of supersymmetric CFT's. We also presented a specific model incorporating SUSY breaking of the required type. This model is technically natural, but it requires several unexplained small input parameters. In this paper, we will present a very simple and plausible model of conformal sequestering, in which all large hierarchies are dynamically generated. Using "naïve dimensional analysis" to estimate the strong interaction coefficients, we find that the model easily gives enough sequestering so that anomaly mediation dominates, and flavor-changing neutral currents are near their experimental limits.

The basic structure of our model is as follows. The central component of $\mathcal{L}_{\text{hidden}}$ is a SUSY theory that is near a strongly coupled conformal fixed point below the Planck scale. The infrared approach to the fixed point is governed by an order 1 critical exponent β_*'. Imposing certain exact hidden symmetries restricts the hidden sector factors in $\mathcal{L}_{\text{mixed}}$ to have the same form as the operators in $\mathcal{L}_{\text{hidden}}$. Because of this, the operators in $\mathcal{L}_{\text{mixed}}$ can be viewed as perturbations of hidden sector couplings with coefficients that depend on visible sector fields. All such perturbations are suppressed by $(\mu/M)^{\beta_*'}$ as the hidden sector approaches the fixed point, where μ is the renormalization scale and M is the Planck scale. This is the conformal sequestering mechanism.

Superconformal field theories naturally have a moduli space. They are exactly superconformal only at the origin of moduli space, but away from the origin superconformal invariance is spontaneously broken. The degeneracy of these vacua can be lifted by weak, even technically irrelevant per-

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[1]Note that this is true in superspace, but not after component level field redefinitions to go to the Einstein frame. See Ref. [3].

turbations to the fixed point theory. We use such effects to generate an effective potential for the hidden moduli space which stabilizes the moduli away from the origin with a SUSY-breaking vacuum energy. The small numbers needed to ensure that the SUSY-breaking scale and the moduli vacuum expectation values (VEV's) are hierarchically smaller than the Planck scale are naturally generated by nonperturbative effects.

For an earlier application of strong conformal dynamics (in the visible sector) to supersymmetric model building see Ref. [10].

II. THE MODEL

In this paper, we focus on the hidden sector and the mechanism for sequestering from the visible sector. The visible sector can be any theory for which AMSB yields an acceptable phenomenology. For now, we will restrict ourselves to global SUSY. In Sec. IV, we will consider the (SUGRA) corrections to the effective potential, which are important for modulus stabilization and canceling the cosmological constant.

Our model of the hidden sector consists of two supersymmetric QCD (SQCD) subsectors: a SU(2) gauge theory with four flavors (eight fundamentals) T^{Ja} $(J=1,...,4;a=1,2)$, denoted as SQCD_2; and a SU(3) gauge theory with two flavors P^a, $\overline{P}_a$ $(a=1,2)$, denoted by SQCD_3. Throughout the paper we will suppress all gauge indices, and we will suppress the $a=1,2$ index when the meaning is clear. We impose the following symmetries on the hidden sector: permutations of the T^J, multiplication of any of the T^J by -1, charge conjugation for SQCD_3, and a global SU(2) symmetry acting on the $a=1,2$ index.[2] The theory has a superpotential invariant under these symmetries:

$$W=\frac{\lambda}{M}\sum_J (T^JT^J)(\overline{P}P)+\frac{\lambda'}{M}\sum_{J\neq K}(T^JT^J)(T^KT^K), \tag{2.1}$$

where M is the Planck scale. We will show that this simple model sequesters itself from the visible sector and has a local minimum that dynamically breaks SUSY.

The SQCD_2 sector is at the self-dual point of Seiberg's conformal window [11] and we will assume that it starts near its ir fixed point coupling at the Planck scale. It is therefore strongly coupled. We assume the SQCD_3 sector is weakly coupled at the Planck scale. We also assume that the superpotential couplings λ and λ' are sufficiently small that they can be treated as perturbations of the SQCD_2 fixed point.

The leading dangerous terms in $\mathcal{L}_{\mathrm{mixed}}$ compatible with the hidden sector symmetries are

$$\mathcal{L}_{\mathrm{mixed}}(M)=\int d^4\theta\left[\frac{c^j{}_k}{M^2}Q_j^\dagger Q^k\sum_J T_J^\dagger T^J + \frac{(c_P)^j{}_k}{M^2}Q_j^\dagger Q^k(P^\dagger P+\overline{P}^\dagger\overline{P})\right], \tag{2.2}$$

where the Q_j are visible chiral superfields. The danger is that, for c,c_P of order unity and containing standard model flavor violation, flavor-violating visible scalar masses will be generated upon SUSY breaking in the hidden sector that will dominate over the flavor-blind AMSB contributions. We will show that the SQCD_2 conformal dynamics naturally suppresses the effects of c,c_P at low energies, allowing AMSB to dominate the visible sector. Supergravity loops can contribute to mixed couplings, but they are dominant in the ultraviolet, so their leading effects can be absorbed into the c coefficients.

III. SEQUESTERING

We first consider the limit $\lambda=\lambda'=0$. In this limit the SQCD_3 sector completely decouples and we can omit it from the discussion. The leading terms in $\mathcal{L}_{\mathrm{mixed}}$ [see Eq. (2.2)] can be viewed as perturbations to the wave function of the hidden fields renormalized at the Planck scale:[3]

$$(\mathcal{L}_{\mathrm{hidden}}+\mathcal{L}_{\mathrm{mixed}})(M)=\int d^4\theta\, Z_0T^\dagger T + \left(\int d^2\theta\,\tau_{\mathrm{hol},0}\mathrm{tr}W^\alpha W_\alpha+\mathrm{H.c.}\right), \tag{3.1}$$

where

$$Z_0=z_0+\frac{c^j{}_k}{M^2}Q_j^\dagger Q^k. \tag{3.2}$$

We will explain the role of z_0 below. $\tau_{\mathrm{hol},0}$ is the holomorphic SU(2) gauge coupling. The theory defined by Eq. (3.1) has only one physical coupling, namely, the physical gauge coupling $\tau=1/g^2$, given by[4]

$$\tau=\mathrm{Re}(\tau_{\mathrm{hol}})-\frac{F}{8\pi^2}\ln Z+\frac{N}{8\pi^2}\ln\tau+f(\tau), \tag{3.3}$$

where N is the number of colors and F is the number of flavors; in our theory, $N=2$, $F=4$. Here $f(\tau)=\mathrm{const}+\mathcal{O}(\tau^{-1})$ parametrizes the scheme dependence. The perturbation Eq. (3.2) (holding $\tau_{\mathrm{hol},0}$ fixed) therefore gives rise to a

[2]For readers concerned by quantum gravity violation of global symmetries: the SU(2) group can be weakly gauged, or can be replaced by a suitable discrete subgroup.

[3]Note that, without imposing the hidden-flavor symmetries discussed in Sec. II, the mixed terms could be more general than this form. In this case we would encounter the difficulties discussed in Ref. [9].

[4]Note that Z_0 in Eq. (3.2) is a vector superfield. By "analytic continuation into superspace" Eq. (3.3) can be interpreted as an equality of vector superfields [13].

perturbation of the physical gauge coupling. Taking the derivative d/dt of Eq. (3.3), where $t \equiv \ln \mu/M$, we obtain

$$\beta(\tau) = \frac{b/8\pi^2 - (F/8\pi^2)\gamma(\tau)}{1-(N/8\pi^2)(1/\tau) - f'(\tau)}, \quad b = 3N - F, \quad (3.4)$$

where $\beta \equiv d\tau/dt$, $\gamma \equiv d\ln Z/dt$, and we have used $d\tau_{\rm hol}/dt = b/8\pi^2$. In the Novikov-Shifman-Vainshtein-Zakharov (NSVZ) scheme $f \equiv 0$, Eqs. (3.3) and (3.4) are the famous formulas of Refs. [12].

Because Eq. (3.2) is a perturbation to the UV gauge coupling, it is clear that it is irrelevant near the IR fixed point. This means that the effects of the perturbation $c^j{}_k Q_j^\dagger Q^k/M^2$ are suppressed in the IR. This is the underlying mechanism for sequestering in this class of models [9].

We now make this quantitative. Exactly at the fixed point, $\tau = \tau_* = \text{const}$, so $\gamma(\tau) = \gamma(\tau_*) \equiv \gamma_* = \text{const}$. Therefore [taking $Z_*(t=0) = 1$]

$$Z_*(t) = e^{\gamma_* t}. \quad (3.5)$$

The theory is at a fixed point despite the running of Z because the running of $\tau_{\rm hol}$ compensates so that $\beta(\tau_*) = 0$. From Eq. (3.4) we see that this requires [11]

$$\gamma_* = \frac{b}{F}. \quad (3.6)$$

We now consider the perturbations about the fixed point. We expand the renormalization group (RG) functions in $\Delta\tau \equiv \tau - \tau_*$ to first order to define critical exponents

$$\beta(\tau) \simeq \beta'_* \cdot \Delta\tau, \quad (3.7)$$

$$\gamma(\tau) \simeq \gamma_* + \gamma'_* \cdot \Delta\tau. \quad (3.8)$$

We factor out the fixed point running by defining

$$\Delta \ln Z \equiv \ln Z - \gamma_* t. \quad (3.9)$$

Then we have

$$\frac{d(\Delta \ln Z)}{dt} = \gamma'_* \cdot \Delta\tau, \quad \frac{d(\Delta\tau)}{dt} = \beta'_* \cdot \Delta\tau. \quad (3.10)$$

Because of the relation Eq. (3.3), these equations are not independent. Using Eq. (3.3) we can write a RG equation for $\Delta \ln Z$ alone:

$$\frac{d(\Delta \ln Z)}{dt} = \beta'_* \left[\Delta \ln Z - \frac{8\pi^2}{F} \Delta\tau_{\rm hol,0} \right]. \quad (3.11)$$

Here $\Delta\tau_{\rm hol} \equiv \tau_{\rm hol} - \tau_{\rm hol,*}$, and $\tau_{\rm hol,*}$ satisfies Eq. (3.3) for $\tau = \tau_*$, $Z = Z_*$. The deviation from the fixed point in the UV is parametrized by $Z_0 \neq 1$ [see Eqs. (3.2) and (3.5)] and $\Delta\tau_{\rm hol,0} \neq 0$. From Eq. (3.3) we can see that these are not independent perturbations, so we can choose $\Delta\tau_{\rm hol,0} = 0$ and parametrize the perturbation by Z_0 alone. The solution to Eq. (3.11) is then simply

$$\Delta \ln Z = e^{\beta'_* t} (\Delta \ln Z)_0. \quad (3.12)$$

β'_* is a strong interaction critical exponent of order 1. Since the dangerous terms in $\mathcal{L}_{\rm mixed}$ are contained in $\Delta \ln Z_0$, this clearly shows the sequestering.

Now we include the effects of λ, λ', and τ_3. From now on we specialize to the case $F = 2N$ for the $SQCD_2$ sector, so that $\gamma_* = 1/2$. We must now include the additional mixed terms

$$Z_{P,0} = 1 + \frac{(c_P)^j{}_k}{M^2} Q_j^\dagger Q^k. \quad (3.13)$$

Because the SU(3) sector is not a CFT, we expect at most an order 1 renormalization of Z_P. Since we are only interested in the order of magnitude of Z_P, we will simply use the approximation $Z_P \simeq Z_{P,0}$. The mixed terms in Eq. (3.13) do not directly give rise to large visible soft masses because in our model the dominant source of SUSY breaking is in the $SQCD_2$ sector.

However, we must determine the leading effects of the perturbation Eq. (3.13) on the $SQCD_2$ sector. These can be studied in the RG equation for $\Delta \ln Z$:

$$\frac{d(\Delta \ln Z)}{dt} = \gamma'_* \cdot \Delta\tau + \Delta\gamma(\tau, \tau_3, \lambda_{\rm phys}, \lambda'_{\rm phys}), \quad (3.14)$$

where

$$\lambda_{\rm phys} = \frac{\lambda\mu}{MZZ_P} = \frac{\lambda e^{t/2} e^{-\Delta \ln Z}}{Z_P},$$

$$\lambda'_{\rm phys} = \frac{\lambda'\mu}{MZ^2} = \lambda' e^{-2\Delta \ln Z}. \quad (3.15)$$

While $\Delta\gamma$ is a small perturbation in Eq. (3.14), it becomes comparable to the first term on the right-hand side in the IR. We must show that this does not spoil sequestering. Since we are expanding around the fixed point we can set $\tau = \tau_*$ in $\Delta\gamma$. We will use Eq. (3.14) only in the regime where the $SQCD_3$ sector is unbroken and weakly coupled. In this regime, we can neglect the running due to τ_3. The leading terms are therefore

$$\frac{d(\Delta \ln Z)}{dt} = \gamma'_* \cdot \Delta\tau + \frac{|\lambda_{\rm phys}|^2}{\rho^4} + \frac{|\lambda'_{\rm phys}|^2}{\rho^4}. \quad (3.16)$$

Because of the $SQCD_2$ strong interaction uncertainties, we cannot compute the coefficients of the last two terms precisely, but we have estimated their order of magnitude using "naive dimensional analysis" (NDA) [14,15]. Here, and later in the paper, we will give our NDA estimates in terms of

$$\rho \sim 4\pi. \quad (3.17)$$

Separate order 1 uncertainties should then be ascribed to different terms, but these will not be written explicitly.

Once again, we would like to use Eq. (3.3) to eliminate $\Delta\tau$ on the right-hand side of Eq. (3.16) in favor of $\Delta \ln Z$. In the presence of the additional couplings λ, λ', and τ_3 Eq. (3.3) remains true, but the scheme dependent function f is in general a function of all the couplings. However, we can always choose a scheme where f is a function of τ alone. In such a scheme we have

$$\frac{d(\Delta \ln Z)}{dt} = \beta'_* \Delta \ln Z + \frac{|\lambda_{\mathrm{phys}}|^2}{\rho^4} + \frac{|\lambda'_{\mathrm{phys}}|^2}{\rho^4}. \tag{3.18}$$

The last two terms on the right-hand side are subdominant perturbations compared to the first term unless $\Delta \ln Z$ is small. Therefore we can approximate the last two terms using Eq. (3.15) in the limit $\Delta \ln Z \to 0$. Also Z_P runs only perturbatively, so we can approximate $Z_P \simeq Z_{P,0}$. We then obtain the approximate solution

$$\Delta \ln Z \simeq e^{\beta'_* t}(\Delta \ln Z)_0 + \frac{|\lambda|^2}{\rho^4 Z_{P,0}^2}\frac{e^{\beta'_* t} - e^t}{\beta'_* - 1} + \frac{|\lambda'|^2}{\rho^4}\frac{e^{\beta'_* t} - 1}{\beta'_*}. \tag{3.19}$$

The first two terms contain mixed terms, but are sequestered, while the third term is not sequestered, but contains no mixed terms. Therefore, all mixed terms are suppressed in this model provided that there is a sufficiently large range of scales for which the SQCD_2 sector is near the fixed point. In fact, the above perturbations due to λ, λ' have effects subdominant to others we will later identify and to $\Delta \ln Z$.

It is convenient to summarize the RG near the fixed point by writing the effective Lagrangian

$$\begin{aligned}\mathcal{L} \simeq & \int d^4\theta [\mu^{1/2}(1 + \Delta \ln Z)\tilde{T}^\dagger \tilde{T} + Z_{P,0}(P^\dagger P + \bar{P}^\dagger \bar{P})] \\ & + \int d^2\theta \left[\frac{\lambda}{M^{1/2}} \sum_J (\tilde{T}^J \tilde{T}^J)(\bar{P}P) + \lambda' \sum_{J \neq K} (\tilde{T}^J \tilde{T}^J) \right. \\ & \left. \times (\tilde{T}^K \tilde{T}^K) \right] + \mathrm{H.c.} + \text{gauge kinetic terms,}\end{aligned} \tag{3.20}$$

where we have defined the rescaled fields [9]

$$\tilde{T} \equiv \frac{T}{M^{1/4}}. \tag{3.21}$$

This rescaling removes the leading M dependence of the Lagrangian, and makes the canonical dimension of the $\tilde{T}$ fields the same as their fixed point scaling dimension in chiral operators.

IV. SUPERSYMMETRY BREAKING

We now determine the vacuum in this theory. We will show that there is a locally stable vacuum with broken SUSY at $T \neq 0$.

In the absence of the superpotential couplings Eq. (2.1), the SQCD_2 theory has 13 independent moduli, which can be parametrized by the SU(2) gauge invariant "meson" operators of the form $T^{Ja}T^{Kb}$ subject to classical constraints. Away from the origin of moduli space the superpotential couplings proportional to λ' reduce the moduli space to a single flat direction, which we assume is in the direction[5]

$$TT \propto \begin{pmatrix} X^{3/4}\epsilon & 0 \\ 0 & 0 \end{pmatrix}, \quad \epsilon = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \tag{4.1}$$

where we use the basis

$$T = \begin{pmatrix} T^{11} \\ T^{12} \\ \vdots \\ T^{42} \end{pmatrix}. \tag{4.2}$$

The field X parametrizes the flat direction. A VEV for X breaks the conformal symmetry, so X is the Nambu-Goldstone mode for spontaneous breaking of scale symmetry. We have defined X so that it has dimensions of mass.

The first threshold in this theory is given by the VEV $\langle X \rangle$, where the conformal symmetry is spontaneously broken. NDA tells us that the physical threshold is at a scale $\sim (\rho \langle \tilde{T} \rangle)^{4/3}$, and that the canonically normalized modulus field is $X \sim \rho^{1/3}\tilde{T}^{4/3}$. The effective Lagrangian below the scale of conformal symmetry breaking is written in terms of the modulus X and the SQCD_3 fields:

$$\begin{aligned}\mathcal{L}_{\mathrm{eff}}(\mu \lesssim \rho|X|) & \\ \sim & \int d^4\theta \{[1 + \Delta \ln Z(\mu \sim \rho|X|)]X^\dagger X \\ & + Z_{P,0}(P^\dagger P + \bar{P}^\dagger \bar{P})\} + \int d^2\theta \frac{\lambda}{\rho^{1/2}M^{1/2}} \\ & \times X^{3/2}\bar{P}P + \mathrm{H.c.} + \mathrm{SU}(3) \text{ gauge kinetic terms.}\end{aligned} \tag{4.3}$$

The superpotential in Eq. (4.3) gives rise to a mass for the P fields

$$m_P \sim \frac{\lambda}{\rho^{1/2}M^{1/2}} \langle X \rangle^{3/2}. \tag{4.4}$$

We consider the case $m_P > \Lambda_3$, where Λ_3 is the scale where the SQCD_3 gauge theory with two flavors becomes strong, and will check the self-consistency of this choice later. In this case, we can integrate out the P fields perturbatively at the scale m_P, and the effective theory is

$$\begin{aligned}\mathcal{L}(\mu \lesssim m_P) \sim & \int d^4\theta \left[1 + \Delta \ln Z(\mu \sim \rho|X|) \right. \\ & \left. + \frac{\lambda^2}{\rho^3 Z_{P,0}^2 M}|X| \ln\left(\frac{\rho|X|}{m_P}\right) \right] X^\dagger X \\ & + \mathrm{SU}(3) \text{ gauge kinetic terms.}\end{aligned} \tag{4.5}$$

[5] For more detail on the moduli space of this theory, see Ref. [9].

The $\ln|X|$ term gives the leading effect of P loops between the scales $\rho\langle X\rangle$ and m_P.[6]

The pure SU(3) gauge theory becomes strong at a scale

$$\lambda_{3,\mathrm{eff}} \sim m_P^{2/9}\Lambda_3^{7/9}. \tag{4.6}$$

Gaugino condensation gives rise to an effective superpotential

$$W_{\mathrm{dyn}} \sim \frac{\Lambda_{3,\mathrm{eff}}^3}{\rho^2} \sim \Lambda_{\mathrm{int}}^2 X, \tag{4.7}$$

where

$$\Lambda_{\mathrm{int}}^2 \sim \frac{\lambda^{2/3}\Lambda_3^{7/3}}{\rho^{7/3}M^{1/3}}. \tag{4.8}$$

Equation (4.7) is the superpotential of a Polonyi model, which breaks SUSY provided that the Kähler terms stabilize the field X. The vacuum energy is then of order Λ_{int}^4, and therefore $m_{3/2} \sim \Lambda_{\mathrm{int}}^2/M$.

The effective potential for X including the Kähler terms of Eq. (4.5) is

$$V_{\mathrm{eff}} = \Lambda_{\mathrm{int}}^4 \Bigg/ \frac{\partial^2 K_{\mathrm{eff}}}{\partial X\, \partial X^\dagger} + \Delta V_{\mathrm{SUGRA}} \sim \Lambda_{\mathrm{int}}^4 \left[1 + (\ln z_0)\left(\frac{\rho|X|}{M}\right)^{\beta'_*} - \frac{\lambda^2}{\rho^3}\frac{|X|}{M}\ln\left(\frac{\rho|X|}{m_P}\right)\right] - \Lambda_{\mathrm{int}}^4[1 + \mathrm{Re}(X)/M]. \tag{4.9}$$

Here we have written out the leading terms in the solution for $\Delta \ln Z$ [see Eq. (3.19)] and used $(\Delta \ln Z)_0 = \ln z_0$ [see Eq. (3.2)]. We have dropped terms comparable to the $\ln|X|$ term that are not log enhanced. The coefficient of the $\ln z_0$ term depends on strong interactions, but we can choose the sign of $\ln z_0$ so that the sign of this term is positive. The last line contains the leading SUGRA corrections once we add a constant superpotential so as to cancel the Λ_{int}^4 contribution to the cosmological constant. We will demand that the supergravity corrections to the potential dominate over the $|X|\ln|X|$ term. This gives the restriction

$$\frac{\lambda^2}{\rho^3}\ln\left(\frac{\rho\langle X\rangle}{m_P}\right) \lesssim 1. \tag{4.10}$$

We then find a stable minimum at

$$\text{sequestering} \equiv \left(\frac{\rho\langle X\rangle}{M}\right)^{\beta'_*} \sim \left[\frac{1}{\rho \ln z_0}\right]^{\beta'_*/(\beta'_*-1)}, \tag{4.11}$$

where we have solved for the sequestering factor for the mixed terms in Z_0. [By Eq. (4.10), the dangerous mixed terms arising from $Z_{P,0}$ are even more suppressed.] The term $\ln z_0$ parametrizes the deviation of SQCD_2 from the fixed point at the Planck scale, and must be small enough that we can trust the fixed point expansions Eq. (3.7). NDA yields tells us that this requires $\ln z_0 \lesssim 1$. The anomalous dimension β'_* is order 1 (and positive), and therefore the sequestering factor is an order 1 power of a loop suppression factor (up to a logarithmic correction).

In fact, there is an adjustable parameter that controls the amount of sequestering in our model. It is completely natural for the SQCD_2 sector to enter the strong coupling conformal regime at a sub-Planckian scale, $\tilde{M} < M$, although we have taken the two scales to be equal. In this more general case, we must substitute $1/\rho \to \tilde{M}/(\rho M)$ on the right-hand side of Eq. (4.11). We can therefore obtain any desired amount of sequestering by taking $\tilde{M} \ll M$. Our analysis assumed that $\Lambda_3 < m_P$, so that the P's were integrated out of the theory before the SQCD_3 subsector became strongly coupled. This naturally occurs for sufficiently small Λ_3, which also sets the SUSY-breaking scale according to Eq. (4.8). At the qualitative level, these observations show that the model naturally breaks SUSY far below the Planck scale and generates a large amount of sequestering. In the next section we will see that quantitatively we must saturate the inequalities Eq. (4.10), $\Lambda_3 < m_P$, and $\ln z_0 \lesssim 1$ in order to get maximal sequestering for the real world. It is also optimal to take $\tilde{M} \sim M$ as we have throughout the paper.

V. NUMERICAL ESTIMATES

We now turn to the numerical estimates in this model. Using Eq. (4.8) and Eq. (4.11) (with $\ln z_0 \sim 1$) to eliminate the dependence on Λ_3 and $\langle X\rangle$, the constraint $\Lambda_3 < m_P$ can be written

$$(\text{sequestering})^{7/2} \gtrsim \frac{\rho^{7/2}}{\lambda^3}\left(\frac{\Lambda_{\mathrm{int}}}{M}\right)^2. \tag{5.1}$$

We see that we obtain maximal sequestering by saturating the bound Eq. (4.10). We will approximate the logarithm in Eq. (4.10) as order 1. Note that the resulting $\lambda \sim \rho^{3/2}$ is smaller than the strong coupling value, $\lambda_{\mathrm{strong}} \sim \rho^2$. Substituting into Eq. (5.1) then gives a bound on the sequestering factor:

$$\text{sequestering} \gtrsim \frac{1}{\rho^{2/7}}\left(\frac{\Lambda_{\mathrm{int}}}{M}\right)^{4/7} \sim 6\times 10^{-5}. \tag{5.2}$$

We have taken $M = 2.4\times 10^{18}$ GeV and $\Lambda_{\mathrm{int}} \simeq 3\times 10^{11}$ GeV. By Eq. (4.11), this maximal level of sequestering is obtained for $\beta'_* \simeq 1.2$. The minimum is at $\langle X\rangle \sim 10^{14}$ GeV, and the mass of X is of order 5×10^6 GeV.

The amount of sequestering is sufficient for AMSB to dominate in the visible sector, and is within an order of magnitude of the sequestering factor 3×10^{-6} [9,16] required to adequately suppress *CP*-conserving flavor violation in anomaly mediated SUSY breaking if the coefficients c of Eq. (2.2) are of order 1. Given the considerable uncertainties in the strong interaction coefficients, our maximal sequestering could easily be at or below this flavor-violation bound. Of course it is also possible that the c's of Eq. (2.2) are of order 1/10.

[6]The precise coefficient of this log term is calculable but is unimportant because of the order 1 uncertainties in the other coefficients such as m_P.

We now consider briefly the cosmology of this model. In general, models of the hidden sector suffer from the Polonyi problem [17]. Briefly stated, the problem is that models with moduli generally have a cosmological epoch where coherent oscillations of the moduli dominate the energy density of the universe, and the interactions of the moduli with the visible sector are too weak to reheat the universe to a sufficiently high temperature to allow nucleosynthesis. In the present model, this problem is less severe than in standard hidden sector models because the mass of the modulus is large compared to the weak scale and the self-interactions of the moduli are much stronger than gravitational strength. We will leave a full analysis of this issue for future work.

Another cosmological issue is the fact that the minimum we have found is a false vacuum. There is a supersymmetric vacuum at the origin $T=0$, but because $\langle X\rangle \gg \Lambda_{\rm int}$, the tunneling rate is suppressed by a large exponent and is cosmologically safe [18].

VI. DISCUSSION AND CONCLUSIONS

It is remarkable that the simple four-dimensional model of the hidden sector presented here dynamically breaks supersymmetry and sequesters itself from the visible sector, naturally allowing anomaly mediation to dominate visible sector supersymmetry breaking. We believe that similar mechanisms of sequestering and dynamical supersymmetry breaking can occur in a large class of models, although it is difficult to check this outside of supersymmetric QCD because of the limited number of superconformal theories that are known explicitly.

According to our estimates, *CP*-conserving flavor-changing neutral current processes are near their experimental limits. Given the large uncertainties, it is possible that there is more sequestering than given in our estimates, so that *CP*-violating flavor violation is also sufficiently suppressed. Alternatively, suppressing *CP*-violating flavor violation may require additional structure. In any case, we expect some flavor-changing neutral current processes to be close to their experimental limits.

We hope that this work will help open new directions for constructing complete, compelling, and realistic hidden sector models of supersymmetry breaking.

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