Uplifted metastable vacua and gauge mediation in SQCD

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ABSTRACT: Anomalously small gaugino masses are a common feature of various models of direct gauge mediation. This problem is closely related to the vacuum structure of the theory. In this paper we show that massive SQCD can have SUSY-breaking vacua which are qualitatively different from the ISS vacuum. These novel vacua are metastable with respect to decay to the ISS vacuum. We demonstrate the possibility of addressing the gaugino mass problem in this framework. We study the properties of these vacua and construct an example of a model of direct gauge mediation.

KEYWORDS: Supersymmetry Breaking, Supersymmetry and Duality

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1   Introduction and summary of results

Gauge mediation of supersymmetry (SUSY) breaking is a compelling scenario [1–4]. It addresses the flavor puzzle and often has the virtue of being calculable. Its distinctive footprints were summarized in the form of sum rules in [5] and further explored in [6–9]. Direct gauge mediation [10] is especially appealing since there is no need to add a separate sector of messengers; they emerge from the dynamics.

However, we are still far from having satisfactory complete models. Gauge mediation is generally afflicted by the μ/B_μ problem[1] and by the Landau pole problem.[2] A somewhat less well-known (but still very common) problem is the unexpected smallness of the visible gaugino masses relative to the masses of the scalars. Since the gauginos cannot be lighter than the electroweak scale the scalars are rendered heavy, exacerbating the need for fine-tuning.

Recently the problem of gaugino masses has been studied in [19]. It has been shown that the smallness of gaugino masses is closely related to global properties of the vacua of the theory. In particular, it has been shown that models which break SUSY in the lowest energy state of the low-energy effective theory (namely, the renormalizable theory around the SUSY-breaking vacuum) necessarily have anomalously small gaugino masses. More precisely, the contributions to the gaugino masses vanish at leading order in SUSY breaking.

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[1]This has recently been discussed in [11–13].
[2]This is ameliorated in theories with a sector of messengers [14–18].

The problem of gaugino masses is rather pervasive in models of direct gauge mediation. For example, in models based on massive SQCD [20] the absence of gaugino masses at leading order in SUSY breaking (and the resulting split-SUSY-like spectrum) has been observed in a few examples, e.g. [21–24].[3] In light of the discussion above, the reason is clear. The ISS vacuum is the ground state of the renormalizable theory at the IR. The existence of a SUSY vacuum due to non-perturbative effects very far in field space is not sufficient to generate gaugino masses. Given the arguments above we are therefore motivated to study different vacua of massive SQCD; such vacua have a completely different kind of metastability than the original example of ISS (and many deformations of it).

To find vacua of this sort we abandon the usual strategy of looking for SUSY-breaking minima. Instead, the approach we pursue is to look for states in massive SQCD with a higher vacuum energy than the ISS vacuum.[4]

In this paper we focus on studying the gaugino mass problem in the context of massive SQCD. We show that the theory contains pseudomoduli spaces (i.e. classical flat directions in field space) with a higher vacuum energy than the ISS vacuum. The problem then reduces to constructing vacua on these pseudomoduli spaces. This can be done in variety of ways and for the sake of concreteness and providing an existence proof we focus on certain deformations of massive SQCD. The result is that we find new SUSY-breaking vacua. An important property of these vacua is that they can decay to the ISS vacuum as well as to supersymmetric vacua. These decays are visible in the renormalizable approximation around our SUSY-breaking vacuum. Therefore, this vacuum has the right properties to generate sizeable gaugino masses. We use this SUSY-breaking vacuum to exhibit a generic[5] model of gauge mediation based on massive SQCD.

Of course, it should be mentioned that we do not attempt to construct a model that solves all the difficulties of gauge mediation. The goal of this paper is to demonstrate that the problem of gaugino masses can be addressed by reconsidering the vacuum structure of SUSY-breaking models and more generally the strategy of model building. We hope that these general guidelines will lead to further progress in realizing SUSY breaking and its transmission to the visible sector.

The paper is organized as follows. In section 2 we review the relation between gaugino masses and the vacuum structure. In section 3 we describe our construction and comment on some phenomenological issues. Finally, we conclude in section 4 where we also discuss open problems. Some technical details on Wess-Zumino models are relegated to an appendix.

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[3]Other classes of examples where this phenomenon occurs are, for instance, [17, 18, 25] and in a holographic realization of gauge mediation [26].

[4]Interestingly, in this context [27] provided a non-generic model with a conventional spectrum of soft scales. Indeed, the solution of [27] is not a ground state of the IR theory and it falls to the class of theories we are describing. Here we emphasize the generality of these models and the inevitability of such metastable structures. We also argue that such vacua are easy to find in massive SQCD and are easily rendered generic.

[5]By "generic" we mean that one can add any operator consistent with the symmetries whose coefficient is of the same order of magnitude as existing operators, without spoiling the properties of the metastable vacuum.

2 Gaugino masses in gauge mediation

This section is a review of [19]. For references and more details the reader is advised to consult the original paper. Our discussion here is sufficient for the purposes of this work.

One perplexing property of many models of gauge mediation is that they predict anomalously small gaugino masses. In other words, if the gaugino masses are set to be around the expected soft scale, the scalar masses are at the multi-TeV scale. Such heavy scalar masses are phenomenologically undesirable.

The key point is that the smallness of gaugino masses has to do with global properties of the theory. We can understand this in a somewhat simplified setup. Consider the most general theory of messengers in the $5 \oplus \overline{5}$ representation of SU(5),

$$W = \lambda_j^i X \psi_i \bar{\psi}^j + m_j^i \psi_i \bar{\psi}^j \,, \tag{2.1}$$

where X is a spurion, $\langle X \rangle = x + \theta^2 F$, and $\psi, \bar{\psi}$ are chiral superfields of messengers in $5 \oplus \overline{5}$. This has been first discussed in full generality in [28]. The well known paradigm of minimal gauge mediation is the special case $m_j^i = 0$.

As has been shown in [29], in renormalizable theories the direction labeled by the bottom component of the spurion X is flat at tree-level.[6] This is related to a general result: in vacua without tachyons, the scalar in the same chiral multiplet with a massless Weyl fermion is also massless. Since the fermion in the X multiplet is the massless Goldstino we see that x must be massless.

Therefore, when discussing predictions of the model (2.1) it is natural to discuss them as a function of x, since we do not know a-priori where quantum effects stabilize it. If supersymmetry breaking is small for every x (i.e. the splittings in the messenger multiplets are small) we can calculate the gaugino mass via [28]

$$m_\lambda \sim \partial_x \log \det \left(x \lambda_j^i + m_j^i \right) \,. \tag{2.2}$$

We see that if $m_\lambda \neq 0$ then $\det \left(x \lambda_j^i + m_j^i \right)$ is x dependent. However, since it is a polynomial in x, this implies that there is at least one zero of the polynomial, x_0,

$$\det \left(x_0 \lambda_j^i + m_j^i \right) = 0 \,. \tag{2.3}$$

Since $x \lambda_j^i + m_j^i$ is the mass matrix of the fermion messengers, (2.3) means that there is a massless fermion messenger at x_0. As we have remarked, massless fermions also lead to massless bosons (or tachyons) and therefore there is also a massless (or tachyonic) spin-0 messenger. With a little more work it can be shown that following the direction of this scalar the energy can always be decreased (or the messenger is decoupled from the theory (2.1) and can be therefore disregarded).

We conclude that in order to get $m_\lambda \neq 0$ there must be states with lower energy in the system, visible already at the renormalizable level. Equivalently, in order to get non-vanishing gaugino masses one is led to look for theories where the pseudomodulus space

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[6]One may need to perform a holomorphic change of variables to make this manifest. We assume that this has been done and write (2.1) without loss of generality.

spanned by x is *not locally-stable everywhere.* There must be points on the pseudomodulus space where some messengers are unstable. (However, the true vacuum lies elsewhere and is of course stable.) One context in which such theories have arisen in the past is the inverted hierarchy mechanism [30–34].

One important assumption in our argument was that SUSY breaking is a small effect. In fact, in all the known examples it turns out that introducing large splittings inside the multiplets does not help much. In spite of the fact that now m_λ are non-zero, a sizable hierarchy between the gauginos and the scalars remains. This is reminiscent of the known (in)dependence of gaugino masses on SUSY breaking in minimal gauge mediation. The change from small to large SUSY breaking is very mild [35, 36].

In more detail, suppose the leading order contribution to the gaugino masses vanishes. Then, in many examples, it turns out that even if SUSY breaking is fine-tuned to be the maximal allowed one, it is possible to achieve $m_\lambda/m_{\mathrm{sfermion}} \sim 1/10$. However, in most of the parameter space the hierarchy is much more significant (see e.g. [37]). It would be interesting to understand this phenomenon more generally.

A central theme of this work is the comparison between gaugino masses and scalar masses. This can be quantified [28] using the concept of "effective number of messengers," which we review here for completeness. We can parameterize the gaugino mass as

$$m_{\lambda_r} = \frac{\alpha_r}{4\pi} \Lambda_G \,, \tag{2.4}$$

with the label $r = 1, 2, 3$ standing for $\mathrm{U}(1), \mathrm{SU}(2), \mathrm{SU}(3)$, respectively. The mass of the superpartner $\tilde{f}$ is given by

$$m_{\tilde{f}}^2 = 2 \sum_r C_{\tilde{f}}^r \left(\frac{\alpha_r}{4\pi}\right)^2 \Lambda_S^2 \,, \tag{2.5}$$

where $C_{\tilde{f}}^r$ is the quadratic Casimir of $\tilde{f}$ in the gauge group corresponding to r. With this we can define the effective number of messengers to be

$$\mathcal{N}_{\mathrm{eff}} = \frac{\Lambda_G^2}{\Lambda_S^2} \,. \tag{2.6}$$

In the case of minimal gauge mediation $\mathcal{N}_{\mathrm{eff}}$ coincides with the actual number of messenger superfields. In a more general setup it does not have to be an integer. It is possible to generalize the notion of effective number of messengers to theories that respect just the SM gauge symmetry, $\mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$, but we will not need it here.

A useful and simple corollary which will be relevant for our specific example is the following. Consider a theory with a set of messengers $\psi, \bar\psi$ which are stable (i.e. massive) on the entire pseudomodulus space, and some other set of messengers $\varphi, \bar\varphi$ which are tachyonic at some points on the pseudomodulus space. Then only $\varphi, \bar\varphi$ can contribute to the gaugino masses, while generally all of the messengers contribute to the scalar masses. Thus, the addition of $\psi, \bar\psi$ does not affect Λ_G in (2.4) but increases Λ_S in (2.5). In this case one finds that the total effective number of messengers is smaller than the effective number of $\varphi, \bar\varphi$ pairs.

3 Constructing a new solution of massive SQCD

3.1 The ISS solution and beyond

Consider SQCD with an $SU(N_c)$ gauge group and N_f (anti-)fundamental electric quarks $(\bar{Q}_i)\, Q^i$ where $i = 1 \ldots N_f$. We suppress the electric color index. Our interest is in the regime $N_c < N_f < \frac{3}{2}N_c$ where the theory has a known weakly coupled dual description in the IR [38]. The authors of [20] considered this theory deformed by a mass term in the UV,

$$W = mQ^i\bar{Q}_i \ . \tag{3.1}$$

Denoting by Λ the strong coupling scale of the theory, we have to assume that $m \ll \Lambda$ for the analysis below to be valid. Near the origin of field space, the dynamics of this theory at low energies is captured by an $SU(N)$ gauge theory, $N \equiv N_f - N_c$, with N_f (anti-) fundamental magnetic quarks $(\bar{q}^i)\, q_i$ where $i = 1...N_f$. There is also an $N_f \times N_f$ meson field Φ^i_j transforming as a singlet under the gauge group. The superpotential at low energies is

$$W = hq_i\Phi^i_j\bar{q}^j - h\mu^2\Phi^i_i , \tag{3.2}$$

where $\mu^2 = -m\Lambda$. The kinetic terms for these emergent degrees of freedom are canonical.

It has long been known that this theory has no SUSY vacuum at $\langle\Phi\rangle = 0$. This can be seen due to the rank condition: looking at the F-terms of the meson components,

$$F^j_i = hq_i\bar{q}^j - h\mu^2\delta^j_i , \tag{3.3}$$

we see that the two terms cannot cancel each other due to the fact that the rank of the first matrix is at most $N_f - N_c$. The remarkable discovery of [20] is that in fact the origin is a SUSY-breaking solution. In this vacuum the magnetic quarks q obtain expectation values of the order μ. This is much smaller than Λ and therefore does not mix with the intricate physics at the strong coupling scale.

Let us use the following notation for the meson field and the magnetic quarks:

$$\Phi = \begin{pmatrix} (V)_{N \times N} & (Y)_{N \times (N_f - N)} \\ (\bar{Y})_{(N_f - N) \times N} & (Z)_{(N_f - N) \times (N_f - N)} \end{pmatrix} ,$$

$$q = \begin{pmatrix} \chi_{N \times N} & \rho_{N \times (N_f - N)} \end{pmatrix} , \qquad \bar{q} = \begin{pmatrix} \bar{\chi}_{N \times N} \\ \bar{\rho}_{(N_f - N) \times N} \end{pmatrix} . \tag{3.4}$$

The solution found by ISS is constructed to cancel as many of the F-terms (3.3) as possible. Namely, we choose the matrix $q_i\bar{q}^j$ to have the maximal possible rank, N. It turns out that by doing so the components of the matrix Z in (3.4) remain undetermined. In other words, the solution is given by

$$\Phi = \begin{pmatrix} 0_{N \times N} & 0 \\ 0 & Z_{(N_f - N) \times (N_f - N)} \end{pmatrix} , \qquad q_i\bar{q}^j = \begin{pmatrix} \mu^2\mathbb{I}_{N \times N} & 0 \\ 0 & 0_{(N_f - N) \times (N_f - N)} \end{pmatrix} . \tag{3.5}$$

As we have mentioned in section 2, classical solutions of such theories are always accompanied by at least one complex flat direction. Indeed, there are a few massless particles of the

theory encountered upon expanding around (3.5). The most important one for us is the $N_c \times N_c$ bottom block of the meson field, Z. The one-loop Coleman-Weinberg potential for all the massless modes can be calculated and one gets that the metastable state is located at

$$\Phi = 0, \qquad q = \mu \left(\mathbb{I}_{N \times N} \ 0_{N \times (N_f - N)} \right), \qquad \bar{q} = \mu \begin{pmatrix} \mathbb{I}_{N \times N} \\ 0_{(N_f - N) \times N} \end{pmatrix}. \qquad (3.6)$$

This Higgses the magnetic gauge group completely.

For us it will be important to understand some global properties of the theory (3.5). At one-loop, the theory around the pseudomoduli directions Z_i^j can be described as a set of decoupled O'Raifeartaigh-like models. In the variables of (3.4) this is given by

$$\frac{1}{h} W = -\mu^2 Z_i^i + Z_i^j \rho_j^c \bar{\rho}_c^i + \mu \rho_i^c \bar{Y}_c^i + \mu Y_i^c \bar{\rho}_c^i. \qquad (3.7)$$

$SU(N_f - N)$ flavor symmetry together with the fact that we consider only one-loop diagrams guarantees that the result is only a function of single-trace combinations of the form $\text{Tr} \left((ZZ^\dagger)^n \right)$ and therefore it is enough to consider only one component of the matrix Z_i^j together with the fields coupled to it. For example, consider Z_1^1. Then the model we get is just N copies of the theory:

$$\frac{1}{h} W = -\mu^2 Z + Z \rho \bar{\rho} + \mu \rho \bar{Y} + \mu Y \bar{\rho}, \qquad (3.8)$$

where we have stripped all the indices for simplicity. As we will see, to understand many of the interesting features of massive SQCD it is enough to consider (3.8).

Setting $\rho = Y = \bar{\rho} = \bar{Y} = 0$ we see that $Z \in \mathbb{C}$ is a complex line of degenerate classical solutions with vacuum energy $V = h^2 |\mu|^4$. It can be easily seen that, as a function of Z, the particles $\rho, Y, \bar{\rho}, \bar{Y}$ are always non-tachyonic. An equivalent property is that the fermions in these chiral multiplets are strictly massive for any Z. As explained in [19] these two properties result from the complex line Z being a local ground state of the theory, or in other words, *"a locally stable pseudomodulus space."* We have reviewed in section 2 the reason for why these theories lead to anomalously small values for the gaugino masses.

Suppose the visible sector gauge group is embedded in the flavor symmetry group of (3.7). Then, $\rho, \bar{\rho}, Y, \bar{Y}$ transform in some vector-like representation while $\text{Tr} Z$ is neutral. Regardless of the value of $\langle \text{Tr} Z \rangle$ at which the effective potential stabilizes the modulus, the leading order contribution to the gaugino mass vanishes. Many deformations of the basic starting point (3.1) were constructed and the dynamics essentially boiled down to shifting the original ISS vacuum $\langle \text{Tr} Z \rangle = 0$ to some non-zero value of $\langle \text{Tr} Z \rangle$. Of course, in all these cases it was obtained that the gaugino mass vanishes at leading order [21–24].[7]

We see here very explicitly that to obtain comparable masses for the scalars and the gauginos, a different strategy has to be invoked. In the next subsection we describe a different pseudomoduli space of the theory (3.1). Along this pseudomoduli space the

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[7]In fact, for $\langle \text{Tr} Z \rangle = 0$ there is a restoration of an R-symmetry under which $R(Z) = 2$ and therefore the gaugino masses are exactly zero. For nonzero $\langle \text{Tr} Z \rangle$ there is no unbroken R-symmetry but the leading contributions to gaugino masses still vanish.

vacuum energy is higher than the vacuum energy of the ISS solution. In addition, this pseudomoduli space is not locally stable everywhere. At some regions of the pseudomoduli space there are tachyons, but the metastable state could be in a different place. This "global" property of the space of classical solutions is essential for obtaining a leading order contribution to gaugino masses, and indeed we will see that upon finding a local minimum on our lifted pseudomoduli space, gaugino masses are generated.

3.2 A different pseudomoduli space of massive SQCD

Consider again the theory defined by (3.2). The solution of ISS is constructed to cancel as many of the F-terms (3.3) as possible. Indeed, the rank of the quark matrix $q_i \bar{q}^j$ is at most N, which is exactly saturated by (3.5). Now let us consider the case where the rank of the matrix $q_i \bar{q}^j$ is

$$\mathrm{rank}(q_i \bar{q}^j) = k, \qquad 0 < k \le N \ . \tag{3.9}$$

For $k = N$ we recover ISS, but here we are interested in $k < N$. Related ideas were discussed in [24, 39, 40].

From now on it will be convenient to use the following parametrization for the meson superfield and the magnetic quarks:

$$\Phi = \begin{pmatrix} (V)_{k \times k} & (Y)_{k \times (N_f - k)} \\ (\bar{Y})_{(N_f - k) \times k} & (Z)_{(N_f - k) \times (N_f - k)} \end{pmatrix},$$

$$q = \begin{pmatrix} (\chi_1)_{k \times k} & (\rho_1)_{k \times (N_f - k)} \\ (\chi_2)_{(N-k) \times k} & (\rho_2)_{(N-k) \times (N_f - k)} \end{pmatrix}, \quad \bar{q} = \begin{pmatrix} (\bar{\chi}_1)_{k \times k} & (\bar{\chi}_2)_{k \times (N-k)} \\ (\bar{\rho}_1)_{(N_f - k) \times k} & (\bar{\rho}_2)_{(N_f - k) \times (N-k)} \end{pmatrix}. (3.10)$$

To realize (3.9) we look at the following classical solution:

$$\Phi = \begin{pmatrix} 0_{k \times k} & 0 \\ 0 & Z_{(N_f - k) \times (N_f - k)} \end{pmatrix} = 0 \,,$$

$$q = \mu \begin{pmatrix} \mathbb{I}_{k \times k} & 0_{k \times (N_f - k)} \\ 0_{(N-k) \times k} & 0_{(N-k) \times (N_f - k)} \end{pmatrix} \,, \quad \bar{q} = \mu \begin{pmatrix} \mathbb{I}_{k \times k} & 0_{k \times (N-k)} \\ 0_{(N_f - k) \times k} & 0_{(N_f - k) \times (N-k)} \end{pmatrix} \ . \tag{3.11}$$

The magnetic gauge group is now only partly Higgsed to $SU(N - k)$. All the χ quarks acquire mass either from tree-level terms or from the super-Higgs mechanism. The massless particles in the vacuum (3.11) are thus only the components of the $(N_f - k) \times (N_f - k)$ matrix Z. The energy of the degenerate classical solutions labelled by Z is higher than the one of ISS,

$$\Delta V = h^2 (N - k) |\mu|^4 \ . \tag{3.12}$$

Near the origin of Z, there are some tachyonic magnetic quarks coming from combinations of ρ_2, $\bar{\rho}_2$. Following these tachyons we reach the ISS vacuum. However, as we survey the pseudomoduli space spanned by Z we find that, schematically, for $Z > \mu$ all the magnetic quarks are massive. We shall next discuss it in more detail.

Similarly to the analysis around (3.7) and (3.8), the theory on this pseudomoduli space effectively factorizes into a sum of generalized O'Raifeartaigh models. As before,

SU$(N_f - k)$ flavor symmetry constrains the dependence on Z so it is enough to look at a single entry of the matrix, e.g. Z_1^1. Doing so we get the following theory (where we omit the indices from Z, as in (3.8)):

$$\frac{1}{h}W = -\mu^2 Z + \sum_{c=1}^{N-k} Z\rho_2^c(\bar{\rho}_2)_c + \sum_{c=1}^{k} \left(Z\rho_1^c(\bar{\rho}_1)_c + \mu\rho_1^c\bar{Y}_c + \mu Y^c(\bar{\rho}_1)_c \right) . \tag{3.13}$$

Thus, we have k copies of models similar to those in (3.8). However, now there are also $N-k$ copies of sectors with the magnetic quarks $\rho_2, \bar{\rho}_2$. These sectors are qualitatively different for reasons we will now explain. Contrary to the theory in (3.7), (3.8), as mentioned above, here for $Z < \mu$ some combinations of the $\rho_2, \bar{\rho}_2$ particles are tachyonic.

Suppose some dynamics stabilizes the pseudomoduli Z farther away from μ. Then this does give rise to a well defined metastable state and because of the global properties of the pseudomoduli space, we do have a leading order contribution to the gaugino masses. The way to see this is to note that the $\rho_2, \bar{\rho}_2$ fields look like the messengers of minimal gauge mediation (see [36] for review) and would give the well known $\mu^2/\langle Z \rangle$ contribution to m_λ, rendering the scalar and gaugino masses comparable parameterically.

In general, theories of the form (3.13) are of the type classified and discussed in [28], so this branch of massive SQCD also gives a dynamical realization for some particular example of (Extra)-Ordinary Gauge Mediation (EOGM).

The pure massive SQCD theory (3.1) contains at the tree level of the low energy description the pseudomoduli space (3.11) as well as the accompanying structure (3.13), however, the dynamics of the theory does not naturally give rise to a solution at some permissible value of Z (where there are no tachyonic modes). Rather, the effective potential as a function of Z pushes the theory towards the region with tachyons and eventually to the ISS state.

Thus, as a model building quest, the situation is completely analogous to the question of breaking the R-symmetry in ISS. The theory has to be deformed or modified in one out of many possible ways. However, achieving this goal here will also give rise to a conventional spectrum of superpartners, so the typical split-SUSY spectrum encountered so far can be avoided.

The rest of the note is dedicated to a description of one particular choice of a deformation that allows to stabilize at permissible values of Z.

3.3 Deforming the model

As we have mentioned, the theory (3.1) on the pseudomoduli space (3.11) with $k < N$ has no metastable SUSY-breaking states. This conclusion remains true even if (3.1) is deformed by some small non-renormalizable operators in the UV (which may flow to renormalizable operators in the IR with small coefficients).

Let us explain why this is the case. For $Z < \mu$ there are tree-level tachyons, hence there are no metastable states in this regime. On the other hand, the result of appendix A implies that for $Z \gg \mu$ there cannot be SUSY-breaking vacua as well. In fact, the log approximation discussed in appendix A becomes predominant very quickly past μ and

therefore these two arguments, combined, explain the absence of metastable states in this theory. This has been verified in [24] for a particular deformation, but we see that this is a general result.

These arguments suggest that we should separate the mass scales of the theory. Indeed, there is no reason to have all the mass parameters of (3.1) degenerate. Suppose there are two mass scales, $m_1 > m_2$, such that k of the N_f electric quarks have mass m_1 and the rest $N_f - k$ quarks have mass m_2:

$$W = m_1 \sum_{i=1}^{k} Q^i \bar{Q}_i + m_2 \sum_{i=k+1}^{N_f} Q^i \bar{Q}_i \,. \tag{3.14}$$

The global non-Abelian flavor symmetry is $\mathrm{SU}(k) \times \mathrm{SU}(N_f - k)$. As before, we consider $k < N$ for the analysis below to be valid.

The dual description is:

$$W = h q_i \Phi^i_j \bar{q}^j - h \mu_1^2 \sum_{i=1}^{k} \Phi^i_i - h \mu_2^2 \sum_{i=k+1}^{N_f} \Phi^i_i \,, \tag{3.15}$$

with $\mu_1^2 = -m_1 \Lambda$, $\mu_2^2 = -m_2 \Lambda$. Of course, this has an $\mathrm{SU}(N)$ gauge symmetry whose indices we suppress. One may expand around the pseudomoduli space of classical solutions

$$q_i \bar{q}^j = \begin{pmatrix} \mu_1^2 \mathbb{I}_{k \times k} & 0 \\ 0 & 0_{(N_f - k) \times (N_f - k)} \end{pmatrix} \,, \qquad \Phi = \begin{pmatrix} 0_{k \times k} & 0 \\ 0 & Z_{(N_f - k) \times (N_f - k)} \end{pmatrix} = 0 \,. \tag{3.16}$$

The gauge symmetry is Higgsed to $\mathrm{SU}(N-k)$. The surviving non-Abelian global symmetry is $\mathrm{SU}(k)_d \times \mathrm{SU}(N_f - k)$, where $\mathrm{SU}(k)_d$ is a diagonal mixture of the original $\mathrm{SU}(k)$ and some broken gauge generators.

Studying the fluctuations around (3.16) we find a structure very similar to (3.13), but now the mass scales are not all the same:

$$\frac{1}{h} W = -\mu_2^2 Z + \sum_{c=1}^{N-k} Z \rho_2^c (\bar{\rho}_2)_c + \sum_{c=1}^{k} \left(Z \rho_1^c (\bar{\rho}_1)_c + \mu_1 \rho_1^c \bar{Y}_c + \mu_1 Y^c (\bar{\rho}_1)_c \right) \,. \tag{3.17}$$

The pseudomodulus space is again characterized by Z, with vanishing expectation values for all the other fields.

The theorem of appendix A is still true, regardless of how the theory is deformed: there cannot be a metastable vacuum for Z bigger than all the mass scales of the theory, namely, $Z > \mu_1$. However, now there are tachyons only when $Z < \mu_2$. (Some combinations of the fields $\rho_2, \bar{\rho}_2$ become tachyonic in this region.) We therefore have the nontrivial region

$$\mu_2 < Z < \mu_1 \,, \tag{3.18}$$

where there are no tachyons and there is no argument for the absence of metastable solutions. Indeed, we will see that it is possible to deform the theory by small non-renormalizable operators in the UV and obtain SUSY-breaking solutions in this interval.

As we have mentioned, the problem at this stage is very closely related to the question of how to break the R-symmetry in the original ISS solution. Our choice for the solution of the problem (which is by no means unique) is reminiscent of the work of [24] in the context of R-symmetry breaking in the ISS vacuum.

We deform the UV theory (3.14) by gauge invariant operators quartic in the electric quarks. For concreteness we consider the operator

$$\delta W = \frac{1}{M_*} \left(\sum_{i=k+1}^{N_f} Q^i \bar{Q}_i \right)^2 , \tag{3.19}$$

where M_* is some high energy scale. In the IR this becomes

$$\delta W = \frac{\epsilon h \mu_2}{2} (\mathrm{Tr}\, Z)^2 , \qquad \epsilon \sim \frac{\Lambda^2}{M_* \mu_2} . \tag{3.20}$$

With this deformation we are led to consider the theory

$$\frac{1}{h} W = -\mu_2^2 Z + \sum_{c=1}^{N-k} Z \rho_2^c (\bar{\rho}_2)_c + \sum_{c=1}^{k} \left(Z \rho_1^c (\bar{\rho}_1)_c + \mu_1 \rho_1^c \bar{Y}_c + \mu_1 Y^c (\bar{\rho}_1)_c \right) + \frac{\epsilon \mu_2}{2} Z^2 . \tag{3.21}$$

We will see that ϵ is a parameterically small number and therefore the scale $\epsilon \mu_2$ is much smaller than all the other low energy scales in the problem.[8] Such deformations may not leave any R-symmetry in the problem, but as we have explained, the difficulty of getting gaugino masses is not related to R-symmetry.

3.4 The dynamics of the deformed model

The basic dynamics of (3.21) can be understood from general considerations which we outline here. These considerations are enough to establish the existence of metastable states in the theory (3.21) and in turn in the complete theory.

Without the deformation $\delta W = \frac{\epsilon h \mu_2}{2} Z^2$, the Z field is flat at tree level. Since ϵ is assumed to be small, the only relevant effect of this deformation at tree level is to introduce a tadpole for Z:

$$V_{\text{tree}} = -\epsilon h^2 \mu_2^3 Z + c.c. . \tag{3.22}$$

The basic idea is to balance this tree-level tadpole, which is parameterically suppressed by ϵ, with the one-loop effects, which are ϵ independent. There are two different types of one-loop contributions. One is from the fields $\rho_1, \bar{\rho}_1$ and $Y, \bar{Y}$ which are coupled to each other. This is denoted by $V_{\text{one-loop}}^{(1)}(Z)$. The other is from the fields $\rho_2, \bar{\rho}_2$, which we denote $V_{\text{one-loop}}^{(2)}(Z)$. The total scalar potential is thus

$$V(Z) = V_{\text{tree}}(Z) + V_{\text{one-loop}}^{(1)}(Z) + V_{\text{one-loop}}^{(2)}(Z) . \tag{3.23}$$

The function $V_{\text{one-loop}}^{(2)}(Z)$ does not exist for $Z < \mu_2$ due to the fact that some components of $\rho_2, \bar{\rho}_2$ become tachyonic and cannot be integrated out. We therefore consider the potential

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[8]Replacing $(\mathrm{Tr}\, Z)^2$ in (3.20) by Z^2 in (3.21) is not quite precise, but it is done for simplicity. One can analyze the complete system and arrive at similar conclusions.

$V(Z)$ only in the regime $Z > \mu_2$. In this regime, $V_{\text{one-loop}}^{(2)}(Z)$ can be crudely approximated by the leading log. (This is rigorously justified if $Z \gg \mu_2$.)

Similarly, in the regime (3.18), $V_{\text{one-loop}}^{(1)}(Z)$ is roughly approximated by a quadratic function. In fact, the quadratic approximation is formally valid in the limit $Z \ll \mu_1$, but our purpose here is to obtain a gross understanding of the physics so this is good enough. We therefore get that the potential we should study is

$$\frac{1}{h^2}V(Z) = -\epsilon\mu_2^3 Z - \epsilon\mu_2^3 Z^\dagger + (N-k)\frac{\alpha_h}{4\pi}\mu_2^4 \log(|Z|^2) + k\frac{\alpha_h}{4\pi}\frac{\mu_2^4}{\mu_1^2}|Z|^2 \,, \tag{3.24}$$

where $\alpha_h = \frac{h^2}{4\pi}$. We do not include here various order one numerical coefficients from the one-loop potential in order not to clutter the expressions, however, all the signs are correct.

The potential (3.24) has a minimum if[9]

$$\frac{N-k}{k} \lesssim \left(\frac{Z}{\mu_1}\right)^2 \,, \tag{3.25}$$

and if

$$\sqrt{k(N-k)}\frac{\alpha_h}{4\pi}\frac{\mu_2}{\mu_1} \lesssim \epsilon \lesssim k\frac{\alpha_h}{4\pi}\frac{\mu_2}{\mu_1} \,. \tag{3.26}$$

Since $Z/\mu_1 < 1$ for our approximation to make qualitative sense, from (3.25) we see that we need $k > N-k$, which means that we should expect to have more of the $\rho_1 - Y$-type sectors than the ρ_2-type fields. From (3.26) we see that, self-consistently with our assumption, ϵ is parameterically smaller than any other tree-level quantity in the Lagrangian and, therefore, the backreaction from the quadratic deformation is indeed negligible. The value of ϵ is almost determined in terms of h and μ_2/μ_1; if ϵ is not in the range (3.26), then either there is no minimum or it is outside of the range of validity of our approximation. (An exact treatment shows that there is typically no minimum in this case.)

Of course, the model (3.21) can also be solved exactly at one-loop and it has been verified that all its qualitative features agree remarkably well with the analysis above. The theory (3.21) also contains solutions with lower vacuum energies, exactly like the parent SQCD theory. In one of them the $\rho_2, \bar{\rho}_2$ fields obtain nonzero VEVs. This corresponds to the ISS solution. The other is where $Z = \frac{\mu_2}{\epsilon}$, where we have a new SUSY vacuum. In both cases the separation in field space is at least of order μ_1 and the energy difference is μ_2, therefore the parametric separation of scales guarantees longevity.

3.5 Direct mediation of SUSY breaking

Imagine that we embed the SU(5) GUT group in the flavor symmetry group $SU(N_f - k)$. Then, some components of ρ_1, ρ_2, Y and $\bar{\rho}_1, \bar{\rho}_2, \bar{Y}$ are in the 5 and $\bar{5}$ of SU(5), respectively. The matrix Z decomposes to a matrix in the adjoint, 5, $\bar{5}$ and the singlet representation of SU(5). Let us study the branch where $\text{Tr}\, Z \neq 0$ and the GUT group remains unbroken.

We can then calculate the scalar and gaugino masses of the model (3.17). The various components of the matrix Z do not have a supersymmetric spectrum but their contributions

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[9]For simplicity, we take all the parameters to be real and positive.

to the soft terms in the MSSM are negligible. The reason is that the typical mass of the scalars in Z, hence the typical mass splitting in the supermultiplet, is $\sqrt{\frac{\alpha_h}{4\pi}}\frac{\mu_2^2}{\mu_1}$. This has a negligible effect on the visible soft masses.

Therefore, for our purposes we may forget about the matrix Z. The messengers are only the magnetic quarks $\rho_{1,2}$, their conjugates $\bar{\rho}_{1,2}$ and the meson components Y, $\bar{Y}$. These types of models have been analyzed in EOGM and formulae for the gaugino and scalar masses in the regime $Z \gg \mu_2$ were presented (see equations (2.5), (2.6) in [28]). We can parameterize the ratio of gaugino to scalar masses with the effective number of messengers $\mathcal{N}_{\text{eff}}$, defined in (2.6). The result is:

$$(\mathcal{N}_{\text{eff}})^{-1} = \frac{|z|^2}{(N-k)^2}\left(\frac{N-k}{|z|^2} + \frac{2k}{|z|^2+4} + \frac{2k\log\left(\frac{|z|^2+2+\sqrt{|z|^4+4|z|^2}}{|z|^2+2-\sqrt{|z|^4+4|z|^2}}\right)}{(|z|^2+4)\sqrt{|z|^4+4|z|^2}}\right), \qquad (3.27)$$

where we denote $z = Z/\mu_1$. Note that the $\rho_1 - Y$-type messengers do not generate gaugino masses at leading order (which is in accord with the discussion in section 2) but do contribute to the scalar masses. This is the reason for the inequality $\mathcal{N}_{\text{eff}} \leq N - k$. We therefore end up with theories that have no parametric suppression of gaugino masses, but the effective number of messengers is bounded by $N - k$.

3.6 Comments on phenomenology

These theories generically suffer from the Landau pole problem. For example, if $k = N-1 = 7$, one finds 15 messengers and all the representations coming from the matrix Z. The latter can be lifted by introducing another deformation $\delta W = \epsilon_{ad}h\mu_2\text{Tr}\left(Z^2\right)$, which is linearly independent of $(\text{Tr}\,Z)^2$. This can be arranged so that all the components of the matrix Z but the trace get in the IR a large enough mass. In this way, all the contributions to the beta functions from the matrix Z can be pushed to higher energies. Consequently, for appropriately chosen ϵ_{ad}, the Landau pole can be pushed to lie well above the strong coupling scale Λ.

For $k < N - 1$, the Landau pole problem is more severe. This is simply because the overall number of messengers is scaled with the number of unbroken color generators, which is $N - k$. One could perhaps try to address the Landau pole problem as in [41].

Let us discuss some scales in the problem. Hereafter we assume that $h \sim 1$. First, the gaugino mass scale is given by

$$m_{\lambda_r} \sim \frac{\alpha_r}{4\pi}\frac{\mu_2^2}{\mu_1} \sim 100 \text{ GeV} \quad \Rightarrow \quad \frac{\mu_2^2}{\mu_1} \sim 10^5 \text{ GeV} . \qquad (3.28)$$

In order to get ϵ to be of the right order of magnitude, and if we choose $M_* \sim M_{\text{Pl}}$, we get:

$$\epsilon\mu_2 \sim \frac{h^2}{16\pi^2}\frac{\mu_2^2}{\mu_1} \sim \frac{\Lambda^2}{M_{\text{Pl}}} \Rightarrow \Lambda \sim 10^{11} \text{ GeV} . \qquad (3.29)$$

Now let us consider the constraints from calculability. A typical value of the incalculable contributions to all the masses from corrections to the Kähler potential is $\delta m^2 \sim \frac{\mu_2^4}{\Lambda^2}$.

Comparing with the one-loop contributions, which are of the order $\frac{h^2}{16\pi^2}\frac{\mu_2^4}{\mu_1^2}$, we see that we should demand $h\Lambda \gg 4\pi\mu_1$, and consequently $\mu_1 \lesssim 10^9$ GeV.

Indeed, we can take $\mu_1 \sim 10^9$ GeV and, therefore, $\mu_2 \sim 10^7$ GeV. In this case, by choosing an appropriate ϵ_{ad}, the Landau pole can be pushed above Λ. Relaxing the condition $M_* \sim M_{\rm Pl}$ one can do better by considering more general constructions. As a numerical example, for $k = 7$ and $N = 8$ we find metastable states (by choosing an appropriate value for ϵ, as in (3.26)). Calculating the value of the pseudomodulus $\mathrm{Tr}\, Z$ at the vacuum and using the formula (3.27), we can typically obtain $\mathcal{N}_{\rm eff} \sim 1/4$. Note that this means that the actual hierarchy in masses is just a factor $\sim 1/2$ on top of the usual prediction of minimal gauge mediation with one messenger. This is better than what can be obtained from studying metastable states on locally stable pseudomoduli spaces.

In spite of the fact that we analyzed a particular set of quartic operators in the UV, it is possible to see that the dynamics of the model is only slightly affected by other small perturbations consistent with the symmetries. In this sense, the mechanism above is generic.

Of course, as many of the model building attempts based on on ISS, this has a few unappealing features. One is the difficulty in attaining perturbative coupling unification. Second, like in many other examples, we need to deform the model and tie the deformation parameters with other low energy scales (which seems unnatural and unappealing). Last but not least, in the simplest form of such models, the mass scale for the electric quarks which one needs to introduce in the UV spoils the main motivation for dynamical SUSY breaking (which is to explain the smallness of the electroweak scale).

Here we did not attempt to solve these problems but rather to demonstrate a new principle in model building, which has shown to be fruitful in surmounting one pervasive problem of gauge mediation: the anomalously small gaugino masses. This problem is interesting and perhaps special as it is connected to the structure of the vacua in the theory.

4 Conclusions and outlook

In this paper we have discussed pseudomoduli spaces of metastable SUSY-breaking solutions of massive SQCD. The main purpose has been to establish an existence proof for such vacua and to demonstrate their crucial phenomenological difference from the more conventional approach. These vacua are metastable already in the renormalizable approximation and lead to sizeable gaugino masses.

This is to be contrasted with the ISS solution (and deformations thereof) which are generally afflicted by anomalously small contributions to the gaugino masses. We have repeated the argument of [19] and have explained why this is the case. We identified different possible states of massive SQCD which are metastable even within the low energy effective theory (however, longevity can be ensured) and solve the parametric suppression of the gaugino masses. The construction we presented provides a realization of a particular model of the EOGM type. It would be nice to understand whether more general models of messengers can be embedded in dynamical SUSY breaking.

Our discussion on how to find these solutions is general but we choose one particular method to actually demonstrate them. Another idea which can be implemented in such

scenarios was studied recently in [42]. It would be nice to perform a more exhaustive analysis in order to find the most appealing possibility.

Since obtaining sizeable gaugino masses requires having lower energy states already in the renormalizable approximation (at least in calculable examples), it would be interesting to study the cosmological applications of it and the thermal history of such theories, analogous to the studies in [43–45]. This question is clearly very general and deserves a separate study.

Acknowledgments

A Wess-Zumino models at large fields

In this appendix we discuss a useful property of Wess-Zumino models at large fields. The simple result we present here is a useful guideline in model building and applies directly to our construction.

Consider a chiral superfield ϕ with some superpotential

$$\int d^2\theta \left(f\phi + \epsilon W(\phi) \right) , \tag{A.1}$$

where ϵ is some small parameter and the Kähler potential is canonical. A setup of this type is often encountered at low energies in some dual description. e.g. irrelevant operators in the UV can become renormalizable at low energies but they are suppressed by small parameters. The classical potential is

$$V = \left| f + \epsilon \frac{\partial W}{\partial \phi} \right|^2 . \tag{A.2}$$

Since an absolute value of a holomorphic function cannot have a local minimum with non-vanishing value for the function, there cannot be proper SUSY-breaking minima in (A.1). One could try to couple the field ϕ to some chiral superfields $P, \bar{P}$ such that the superpotential becomes

$$\int d^2\theta \left(f\phi + h\phi P\bar{P} + \epsilon W(\phi) \right) . \tag{A.3}$$

Now, for large ϕ the $P, \bar{P}$ fields are integrated out and generate a nontrivial Kähler metric at one loop. For $\phi \gg \sqrt{f}$ the contribution of the $P, \bar{P}$ fields is to modify the Kähler potential to

$$K = \left(1 + c\frac{\alpha_h}{4\pi}\log(\phi\phi^\dagger)\right)\phi\phi^\dagger \,, \tag{A.4}$$

where c is a number which can be extracted by matching with the anomalous dimension (for a more general treatment of such systems see [46]) and $\alpha_h = \frac{h^2}{4\pi}$. This means that now the scalar potential is of the form

$$V = \left(1 + c\frac{\alpha_h}{4\pi}\log(\phi\phi^\dagger)\right)^{-1}\left|f + \epsilon\frac{\partial W}{\partial \phi}\right|^2 \,, \tag{A.5}$$

where we have dropped some unimportant constant correction in the prefactor. At leading order in the loop expansion parameter this can also be written as

$$V = \left(1 - c\frac{\alpha_h}{4\pi}\log(\phi) - c\frac{\alpha_h}{4\pi}\log(\phi^\dagger)\right)\left|f + \epsilon\frac{\partial W}{\partial \phi}\right|^2 = \left|\left(1 - c\frac{\alpha_h}{4\pi}\log(\phi)\right)\left(f + \epsilon\frac{\partial W}{\partial \phi}\right)\right|^2 \,. \tag{A.6}$$

We see that again this is an absolute value of a holomorphic function and therefore has no SUSY-breaking minima (i.e. any classical solution will have either tachyons or flat directions). We could have also got this result by performing a holomorphic change of variables in (A.4).

In spite of our discussion above being in the simplest possible setup, this result is a very general property of sigma models with canonical Kähler potential: when the VEV of some fields is above all the mass parameters in the Lagrangian it is generally impossible to find SUSY-breaking minima at one-loop. This is of course consistent with the known examples. (At two-loops this can be circumvented, e.g. [47–50].)

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