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Metastable vacua and D-branes at the conifold

Riccardo Argurio

Physique Théorique et Mathématique and International Solvay Institutes, Université Libre de Bruxelles, C.P. 231, 1050 Bruxelles, Belgium
E-mail: rargurio@ulb.ac.be

Matteo Bertolini

SISSA/ISAS and INFN - Sezione di Trieste
Via Beirut 2, I 34014 Trieste, Italy
E-mail: bertmat@sissa.it

Sebastián Franco

Joseph Henry Laboratories, Princeton University, Princeton, NJ 08544, U.S.A.
E-mail: sfranco@feynman.princeton.edu

Shamit Kachru

Department of Physics and SLAC, Stanford University, Stanford, CA 94305 U.S.A.
E-mail: skachru@stanford.edu

Abstract: We consider quiver gauge theories arising on D-branes at simple Calabi-Yau singularities (quotients of the conifold). These theories have metastable supersymmetry breaking vacua. The field theoretic mechanism is basically the one exhibited by the examples of Intriligator, Seiberg and Shih in SUSY QCD. In a dual description, the SUSY breaking is captured by the presence of anti-branes. In comparison to our earlier related work, the main improvements of the present construction are that we can reach the free magnetic range of the SUSY QCD theory where the existence of the metastable vacua is on firm footing, and we can see explicitly how the small masses for the quark flavors (necessary to the existence of the SUSY breaking vacua) are dynamically stabilized. One crucial mass term is generated by a stringy instanton. Finally, our models naturally incorporate R-symmetry breaking in the non-supersymmetric vacuum, in a way similar to the examples of Kitano, Ooguri and Ookouchi.

Keywords: Brane Dynamics in Gauge Theories, Supersymmetry Breaking, Gauge-gravity correspondence.

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1. Introduction

There has been a renaissance in the study of metastable supersymmetry breaking vacua in string theory and field theory in the past several years. Several of the examples are stringy constructions which involve brane dynamics and/or fluxes in a nontrivial gravitational background [1 – 7]. There are also purely field-theoretic constructions, like the Intriligator-Seiberg-Shih (ISS) models [8] where strong/weak coupling dualities [9] allow one to find such states, or the retrofitted models [10] where canonical theories like O'Raifeartaigh or Polonyi models can be coupled to additional dynamics in a way that naturally produces an exponentially small SUSY breaking scale. These states have potential applications both in understanding the existence and properties of stable non-supersymmetric string compactifications [11 – 13] (for recent reviews see [14, 15]), and in building realistic models of gauge-mediation [16] or sequestered, high-scale SUSY breaking [17]. It is natural to think that via AdS/CFT duality [18] or brane engineering [19], one can sometimes relate the stringy and field theoretic constructions. Indeed, many groups have recently engineered various D-brane field theories which exhibit dynamical SUSY breaking (DSB) and reduce to known DSB field theories in the decoupling limit [20 – 29].

Here, we continue our investigation [25] into the possible relations between ISS-like states in field theory, and SUSY-breaking states where SUSY is broken at the end of a "warped throat," as in [3] (where SUSY was broken by anti-D3 brane probes at the tip of a warped, deformed conifold [30]). At a qualitative level, it is natural to think that SUSY breaking at the end of a warped throat is AdS/CFT dual to dynamical supersymmetry breaking. Then one should try to interpret the SUSY breaking states involving warped antibranes, which can tunnel to supersymmetric states of the same gravitational system [3], as metastable states in the dual SUSY field theory. While it is not necessary that such a correspondence should hold (since non-supersymmetric vacua are not protected upon extrapolation in the 't Hooft coupling $g_s N$), it would be interesting and suggestive to find examples where such states can be argued to exist both in the gravitational system and the field theory dual.

In [25], we argued that the gauge/gravity duals derived by studying fractional branes in simple quotients of the conifold are a natural place to find such a correspondence. In a particular $\mathbb{Z}_2$ orbifold of the conifold, we were able to realize a close relative of SUSY QCD where the small mass parameter of the ISS models is dynamically generated, and where the dual gravitational system plausibly admits metastable anti-brane states. However, that work left several open questions. Firstly, the SQCD model that could be realized had $N_f = N_c$, while it is in the free magnetic range $N_c + 1 \leq N_f < \frac{3}{2}N_c$ that one is really confident of the existence of the metastable states in the field theory. Secondly, because we wish to dynamically generate the small mass parameter of the ISS models, we must take care to ensure that the dynamical masses are *stabilized* against relaxation to zero or infinity. In the models of [25], this delicate question rested entirely on (largely unknown) properties of the Kähler potential.

In this paper, we argue that a simple variant of the model of [25] solves both of these problems. By considering other orbifolds of the conifold, we are able to find models where we can reach the free magnetic regime necessary to prove existence of the ISS vacua, and where we can argue that the superpotential itself helps to stabilize the dynamical quark masses in the range where SUSY breaking occurs. More precisely, the models we construct have $N_f = N_c + 1$. As an added bonus, our model shares a nice feature with the model of Kitano, Ooguri and Ookouchi [27]: the quiver superpotential naturally comes with terms that break the R-symmetry of the original ISS models, which is problematic in model building applications (as it forbids a gaugino mass).

In the rest of this section, we introduce our model. Its field theory dynamics is analyzed in section 2, where we show how an effective massive SQCD arises in a given corner of its moduli space. An important role is played by a stringy instanton generated contribution to the effective superpotential, whose origin we discuss in section 3. In section 4 we prove that the quark masses can be dynamically stabilized, and in section 5 we estimate the lifetime of the metastable vacua. In section 6, we briefly discuss the gravity dual IIB description (involving fluxes and branes in the deformed geometry). We also present a IIA T-dual of the IIB picture, where the gauge theory arises from a configuration of NS 5 branes and D4 branes. These descriptions allow one to visualize many (though not all) aspects of the field theory dynamics, and in particular, make it obvious that the metastable SUSY-breaking

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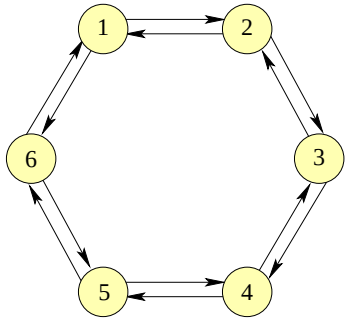

Figure 1: The quiver describing the field content of the gauge theory at the tip of the non-chiral $\mathbb{Z}_3$ orbifold. The ranks of the gauge groups can be chosen arbitrarily.

vacua are dual to models which contain anti-branes. We conclude in section 7.

1.1 The structure of the model

The models we would like to analyze are obtained by considering (fractional) D3 branes at the tip of a non-chiral $\mathbb{Z}_N$ orbifold of the conifold, which is nothing but a straightforward generalization of the system considered in [25] for the case $N = 2$.

The corresponding quiver gauge theory admits $2N$ gauge factors and $4N$ bifundamental chiral superfields X_{ij} interacting via the following quartic superpotential

$$W = h \sum_{i=1}^{2N} (-1)^{i+1} X_{i,i+1} X_{i+1,i+2} X_{i+2,i+1} X_{i+1,i} \,, \tag{1.1}$$

where the index i is understood $mod(2N)$.

Because the quiver is non-chiral, we can consistently assign arbitrary ranks to the quiver nodes, which suggests that there should be $2N - 1$ independent fractional branes one can define. This is indeed mirrored in the geometric structure of the singularity, which admits $2N-1$ independent shrinking 2-cycles the branes can wrap. In a given basis, which will be relevant later, one obtains a natural classification into N deformation fractional branes and $N - 1$ $\mathcal{N} = 2$ fractional branes, following the definition proposed in [31]. We remind the reader that deformation branes correspond to isolated nodes in the quiver (and hence gauge groups with no matter) and lead to confinement, while $\mathcal{N} = 2$ branes arise from occupying two connected nodes, which yields a product of two SQCD theories with $N_f = N_c$, and hence have a moduli space of vacua.

As we are going to show, for our present purposes it is enough to take $N = 3$ (any $\mathbb{Z}_N$ with $N > 3$ naturally works in the same way). Therefore, from now on we stick to this case, for simplicity. This specific $\mathbb{Z}_3$ quotient admits five shrinking 2-cycles. Two of them see, locally, a $\mathbb{C}^2/\mathbb{Z}_2$ singularity. The other three are dual (via geometric transitions) to compact 3-cycles A_i $(i = 1, 2, 3)$ which can be made finite by a complex deformation. The corresponding quiver is shown in figure 1.

We would like to consider the gauge theory arising from the following assignment of ranks in the quiver

$$(N_c \,,\ N_c \,,\ N_c \,,\ 1,\ 0 \,,\ 0) \,. \tag{1.2}$$

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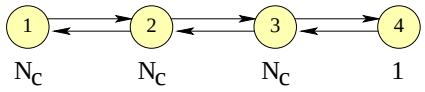

Figure 2: The quiver of the theory (1.2).

In terms of fractional branes, this may be viewed as N_c $\mathcal{N}=2$ branes at nodes 1 and 2, N_c deformation branes at node 3 and another single deformation brane at node 4 (this definition is basis-dependent, of course). The"SU(1)" fourth node is not actually a gauge group. Its interpretation is that X_{34} and X_{43} transform purely as fundamental and anti-fundamental representations of node 3, respectively. The corresponding quiver is depicted in figure 2.

As we shall show, this system, whose dynamics we are going to study in detail, reduces exactly (in a region of the moduli space to be specified below) to massive SU(N_c) SQCD with $N_f = N_c+1$ massive (but light) flavors, and therefore admits both the supersymmetric and the metastable non-supersymmetric vacua of that theory. Hence, this system provides a string embedding of an ISS model. Moreover, it has some additional virtues: the small flavor masses are dynamically generated (and stabilized), it is possible to give a simple gravity dual interpretation of the metastable non-supersymmetric vacua, and R-symmetry is explicitly broken (which could be useful in any model-building applications).

2. The dynamics of the model

The quiver gauge theory we are going to analyze is the one depicted in figure 2. This theory has a superpotential of the form

$$W = h(X_{12}X_{23}X_{32}X_{21} - X_{23}X_{34}X_{43}X_{32}) + mX_{43}X_{34} \ . \tag{2.1}$$

As already discussed, the two quartic terms follow from the conifold by standard orbifold techniques. The mass term for X_{34} and X_{43} is generated by a stringy instanton. We postpone discussion of the relevant instanton to section 3, and we proceed to analyze the above superpotential. The quartic coupling h has the dimensions of an inverse mass, and is inversely proportional to the UV scale generating the non-renormalizable interaction. In this context, it is natural to take $h \sim 1/M_s^*$. Here, M_s^* indicates the string mass scale effectively warped down to a lower value due to the RG flow, which manifests itself as a duality cascade. For the field theory interpretation to be valid, we need M_s^* to be bigger than any of the dynamical scales of the gauge groups involved in the quiver.

To start analyzing the gauge theory, we will make some assumptions about the scales of the gauge groups on every node. Node 3 is the main node where the ISS-like SQCD dynamics takes place. Node 2 acts as a subgroup of the flavor symmetry, broken as SU(N_c+ 1) $\times$ SU(N_c+1) $\supset$ SU(N_c+1) $\supset$ SU(N_c). Accordingly, we will take its dynamical scale Λ_2 to be (much) smaller than the others, so that this gauge group can be effectively considered as classical.

Node 1, which has a number of flavors which equals the number of colors, undergoes confinement so that its effective dynamics is described in terms of mesons and baryons.

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The mesons are going to supply mass terms for some of the flavors of node 3, through the superpotential couplings in (2.1), much as in [25]. These masses are subject to the constraint on the deformed moduli space of node 1. In particular, we will need the scale Λ_1 to be such that those masses are still lower than the scale of the SQCD node, Λ_3. The additional mass m will be larger than these N_c masses, but still smaller than Λ_3. We will explain later how this parameter range can be obtained.

This model is quite similar to the one in [27]. One difference is that all the small parameters are generated dynamically. In addition, this model arises naturally at a fairly simple Calabi-Yau singularity.

Assuming now that node 1 is confining, the tree level superpotential reads

$$W = h(M_{22}X_{23}X_{32} - X_{23}X_{34}X_{43}X_{32}) + mX_{43}X_{34}\,, \tag{2.2}$$

where we have defined $M_{22} = X_{21}X_{12}$. This superpotential is not quite complete: we should really implement the constraint relating the mesons and the baryons of node 1 through the introduction of a Lagrange multiplier. We delay that to later on. Let us now assume that the interactions are such that the mesonic and baryonic branches of node 1 decouple, and in particular that when the meson matrix has maximal rank the baryons are required to vanish – we will argue that this is the case in section 4. For the time being, we assume that node 1 is on the mesonic branch, where the constraint describing the quantum-deformed moduli space reads

$$\det M_{22} = \Lambda_1^{2N_c}\ . \tag{2.3}$$

This constraint is necessary for the generation of dynamical masses but does not fully fix the eigenvalues of M_{22}. In the non-supersymmetric vacua, their stabilization occurs at tree-level as we explain shortly. At the stable point, the VEV of M_{22} is proportional to the identity matrix. Hence we see that in the effective SQCD theory at node 3, we have N_c flavors of mass $\sim h\Lambda_1^2$ and 1 flavor of mass m. We will take the latter to be the heavier one, so that $h\Lambda_1^2 < m$.

Therefore, along this branch the theory on node 3 with superpotential (2.2) is nothing but SQCD with $N_f = N_c + 1$ massive flavors, together with a quartic coupling (which is irrelevant in the IR). Integrating out the flavors, we obtain pure $\mathrm{SU}(N_c)$ SYM characterized by a dynamical scale

$$\Lambda_L^{3N_c} = \Lambda_3^{2N_c-1} h^{N_c} \det M_{22}\ m = \Lambda_3^{2N_c-1}(h\Lambda_1^2)^{N_c} m\ . \tag{2.4}$$

Implementing the constraint on the deformed moduli space of node 1 with a Lagrange multiplier in an effective superpotential, it is easy to see that we indeed have a moduli space of supersymmetric vacua where M_{22} has non zero VEV, while the baryons are vanishing. When taking into account that node 2 is actually gauged, we see that at low-energies the moduli space will be described by $\mathbf{C}^{N_c-1}$ together with a residual $\mathrm{U}(1)^{N_c-1}$ gauge symmetry.

We now move on and show that our theory also has meta-stable, SUSY breaking vacua. Since node 3 has $N_f = N_c + 1$, its low-energy dynamics is governed by a theory of mesons and baryons. This case can actually be seen as a limiting case of Seiberg duality, where

the dual magnetic gauge group is a trivial SU(1), and the dual quarks are nothing but the baryons of the electric theory. In the following, we will adopt this terminology. The bifundamentals are combined into effective mesons as $X_{i3}X_{3j} = \Lambda_3\phi_{ij}$, and the dual quarks are labeled Y_{i3} and Y_{3i}. The superpotential is

$$W = (h\Lambda_3)(M_{22}\phi_{22} - \Lambda_3\phi_{24}\phi_{42}) + m\Lambda_3\phi_{44} - \phi_{22}Y_{23}Y_{32} - \phi_{44}Y_{43}Y_{34} + \phi_{24}Y_{43}Y_{32} + \phi_{42}Y_{23}Y_{34} \,. \tag{2.5}$$

Note that we have rescaled the mesons to canonical dimension using the scale Λ_3. Accordingly, the cubic terms generated by the duality have a coupling of $\mathcal{O}(1)$, which we set to one (shifting the undetermined constant to the normalization of the canonical Kähler potential). Strictly speaking, we should also add a term linear in the determinant of the meson matrix, but it is highly irrelevant in the IR (and to our considerations).[1] It is easy to see that the above model of mesons and dual quarks would have, in the absence of the $\phi_{24}\phi_{42}$ coupling, an accidental IR $\mathrm{U}(1)_R$ symmetry. The R charges would be 2 for the mesons and 0 for the dual quarks, as in an O'Raifeartaigh model. The quadratic coupling in the mesons which arises naturally in the present model provides an explicit breaking of this R-symmetry.

This theory is now amenable to an analysis very similar to [8, 27]. There is supersymmetry breaking by the rank condition. The F auxiliary field that vanishes is the one related to the more massive flavor, i.e. the F-term of ϕ_{44}. On the other hand, the F-components of ϕ_{22} are non-vanishing. As a consequence, there is a tree level vacuum energy given by

$$V_{\text{tree}} = |h\Lambda_3|^2 \sum_{i=1}^{N_c} |M_i|^2 = N_c |h\Lambda_3\Lambda_1^2|^2 \,, \tag{2.6}$$

where we obtain the final result by extremizing on the eigenvalues of the matrix M_{22} given the constraint on its determinant. We are going to show later that indeed the constraint on the determinant is not destabilized by baryonic VEVs.

A standard analysis of this model shows that it is a sum of O'Raifeartaigh models with an additional coupling $m_{24} = h\Lambda_3^2$, which is the one quadratic in the two mesons ϕ_{24} and ϕ_{42}. As we will discuss in section 4, ϕ_{22} gets a non-zero VEV due to the one-loop potential (see also [27]). This VEV is directly related to the presence of the quadratic meson coupling m_{24}, so that we can actually estimate it as

$$|\phi_{22}| \sim |h\Lambda_3^2| \,. \tag{2.7}$$

As noted in [27], the vacua analyzed here are unstable if the size of m_{24} exceeds the larger flavor mass. Here this bound reads $|h\Lambda_3^2|^2 < |m\Lambda_3|$, or in other words (recall that $h = 1/M_s^*$)

$$\left(\frac{\Lambda_3}{M_s^*}\right)^2 < \frac{m}{\Lambda_3} \,. \tag{2.8}$$

Note that all these relation must be taken with a grain of salt, since there are factors of $\mathcal{O}(1)$ that we are not retaining (most of which are non-calculable anyway).

[1] It does play a role if we want to recover the SUSY vacua in the low-energy picture.

All other pseudomoduli are lifted by the one-loop potential, and acquire a non-tachyonic mass.

This shows that the present model has metastable supersymmetry breaking vacua, provided we can show that there is no instability towards turning on baryonic VEVs at node 1. We demonstrate in section 4 that this is the case, in an appropriate regime of parameters.[2]

3. A mass term generated by a stringy instanton

Before analyzing in detail the stability of the model, we have to explain how the mass term for the additional flavor is generated.

It turns out that somewhat novel stringy instanton effects which have recently been investigated in several other contexts [32 – 38] contribute corrections to W which depend on gauge invariants that usually do not appear in the superconformal quiver superpotential. Recall that these effects arise when Euclidean D branes wrap cycles corresponding to quiver nodes which are *not* occupied by space-filling branes. In this respect, they are specific to set ups with fractional branes. We now show that such an instanton generates the mass term m.

To understand the instanton contributions, consider a D1 instanton (a Euclidean D1 or ED1 brane) wrapping node 5 of the quiver. It is BPS and preserves precisely 1/2 of the $\mathcal{N} = 1$ supersymmetry; acting on the instanton solution with the broken supercharges then produces two fermion zero modes in the ED1 - ED1 open string sector. These are the two fermion zero modes that are necessary to give rise to a contribution to the space-time superpotential (we discuss the possibility of extra "accidental" zero modes at the end of this section). Considering the extended quiver diagram including a node for the instanton, there are also fermionic strings α, β stretching to node 4, in the $(+, \overline{N_4})$ and $(-, N_4)$ representations of the $\mathrm{U}(1) \times \mathrm{SU}(N_4)$ gauge group (in the present case we will have $N_4 = 1$). As follows from the computations in [32, 34], the fermionic spectrum in the extended quiver is the same as it would be if the instanton were actually a space-filling brane, except the fermions live in a different dimension. By the simple argument of [34], we also expect that there are no bosonic zero modes in this sector. Bosons would arise from NS sector strings, but the NS sector ground state energy receives a contribution from the number of ND boundary conditions, which pushes the ground state energy above zero in this configuration. The relevant part of the extended quiver is reported in figure 3.

In the instanton action, we expect a gauge invariant coupling

$$L = \alpha X_{43} X_{34} \beta \ . \tag{3.1}$$

In evaluating the instanton contribution to the 4d effective action, we should integrate over the only two charged fermionic zero modes α, β. This yields a simple contribution to the

[2] In the model that we presented in [25] there is potentially such an instability. In that model the SQCD node is in a confining rather than IR free regime, so that (non-calculable) corrections to the Kähler potential are present, and make it difficult to determine what happens. The gravity dual description provides another source of information; such an instability is not readily apparent there, but it is a complicated system which would benefit from further study.

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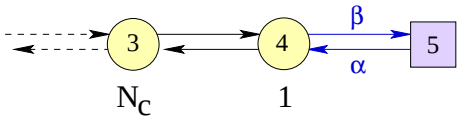

Figure 3: The extended quiver describing the interaction between the fractional brane system and a Euclidean D1 brane on node 5, represented by a square. The relevant coupling is a quartic one, involving α, β, X_{34} and X_{43}, (3.1).

superpotential

$$c\ X_{43}X_{34}\ e^{-\text{Area}}. \tag{3.2}$$

where c is a dimensionful constant and the relevant area is the area of the curve corresponding to node 5. We thus identify the mass term as $m = c\ e^{-\text{Area}}$. If we reasonably take $c \sim M_s^*$, we see that it is not difficult to assume that the area of the cycle wrapped by the instanton is such that $m < \Lambda_3$. (Roughly that would amount to assume that $\Lambda_5 \ll \Lambda_3$ if there was a gauge group on node 5.)

Now, a similar instanton on node 6, with fermionic strings stretching to node 1, produces another term in the superpotential. The gauge invariant coupling in the instanton action is

$$L = \alpha X_{12}X_{21}\beta\ . \tag{3.3}$$

If we let a, b denote gauge indices at node 1 and f, g denote gauge indices at node 2, then the gauge contractions in (3.1) give rise to $\alpha_{\bar{a}} X_{12}^{af} X_{21,f\bar{b}} \beta^b$. So performing the integral over the α, β fermions, which are now a set of $2N_c$ zero modes, gives the contribution

$$c' B\tilde{B}\ e^{-\text{Area}'} \tag{3.4}$$

to the 4d effective theory, where c' is similarly a dimensionful constant and the area is now the one of the curve corresponding to node 6. The coupling (3.4) provides a mass term for the baryons $B = \det X_{12}$ and $\tilde{B} = \det X_{21}$ of node 1. We will see in the next section that this term does not however play an important role in the stabilization of the baryons.

Let us end this section with a comment regarding a subtle point. With the above reasoning, the coefficients c and c' have been determined up to a dimensionless number whose precise value we cannot directly compute in our geometric set-up. A crucial ingredient for such coefficients not to vanish involves the number of uncharged fermionic zero modes on the ED1 brane. Before taking into account the quiver branes back-reaction, there are four, since the ED1 is a 1/2 BPS state in the Calabi-Yau. While, as already discussed, two zero modes are necessary to provide the chiral superspace integral for the superpotential term (3.2), the other two would provide a dangerous vanishing contribution. Still, one should take into account the full back-reacted closed string background, which includes non-trivial fluxes. This background preserves only 4 supercharges out of the 8 preserved by the CY, so that the instanton has only 2 zero modes associated to broken supersymmetries. Then, at least in many backgrounds, it is reasonable to expect that the extra zero modes get lifted by the interactions with other background fields. This is an interesting problem in itself, but we leave it for further research. Instead, in order to provide a background

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where we can explicitly identify an object responsible for lifting the additional zero modes, we can introduce orientifold planes in such a way that the instanton wraps a cycle that is mapped to itself. In this way, half of the zero modes are projected out from the start. One concrete embedding of our model in an orientifold that accomplishes this task, while not spoiling all other nice features of our model, is described in appendix A. We therefore conclude that the coefficient c is non-vanishing in many suitable models.[3]

4. Stabilizing dynamical masses

We have explained how our theory has metastable supersymmetry breaking vacua under certain assumptions regarding the stability of the dynamically generated masses. An important question is whether the dynamical masses relax to zero by turning on expectation values for the baryonic fields at node 1. We show in this section that these a priori dangerous directions are lifted because $\langle \phi_{22} \rangle \neq 0$. New superpotential interactions generated by the D1 instanton wrapping node 6 of the quiver also contribute to stabilization, although they are not the dominant effect.

We thus first sketch how the one-loop potential gives the crucial VEV to the field ϕ_{22}. We start from (2.5) and expand around the metastable vacuum. This is characterized by VEVs for $Y_{34}, Y_{43} = \sqrt{m\Lambda_3}$ and, at tree level, by an arbitrary ϕ_{22}. The latter is also the field with non-vanishing F-terms. The superpotential for the fluctuations of the fields, expanded about the vacuum, takes the form

$$W = m_{22}^2 \phi_{22} - Y_{32}\phi_{22}Y_{23} + m_{44}\phi_{24}Y_{32} + m_{44}\phi_{42}Y_{23} - m_{24}\phi_{24}\phi_{42} \,, \tag{4.1}$$

In writing this down, we have dropped several fields: ϕ_{44}, δY_{34} and δY_{43} do not feel SUSY breaking at this order so will cancel out of the one-loop energy. The masses appearing above are given by $m_{22}^2 = h\Lambda_3\Lambda_1^2$, $m_{44}^2 = m\Lambda_3$ and $m_{24} = h\Lambda_3^2$. All fields have a canonical Kähler potential.

It is straightforward to realize that the F-terms set ϕ_{22} to a diagonal form. Then, we obtain just a superpotential for N_c decoupled O'Raifeartaigh like models. Each such model has, besides the field with the non-zero F-term, 4 other fields. An important parameter is the coupling m_{24} of the meson bilinear $\phi_{24}\phi_{42}$.

The tree level vacuum energy is just:

$$E_{\rm vac} = N_c \, |m_{22}|^4 \ . \tag{4.2}$$

We had already noted that the eigenvalues of M_{22} are all trivially stabilized at their common values [25].

To compute the one-loop vacuum energy of this model, we simply compute the boson/fermion masses as a function of pseudo-moduli using

$$m_0^2 = \begin{pmatrix} W^{\dagger ac}W_{cb} & W^{\dagger abc}W_c \\ W_{abc}W^{\dagger c} & W_{ac}W^{\dagger cb} \end{pmatrix} \quad , \quad m_{1/2}^2 = \begin{pmatrix} W^{\dagger ac}W_{cb} & 0 \\ 0 & W_{ac}W^{\dagger cb} \end{pmatrix} \tag{4.3}$$

[3] A more complete discussion of these issues, for backgrounds where a simpler worldsheet CFT description is available, will appear in [38]. A general discussion on the introduction of O-planes in generic Calabi-Yau geometries will appear in [39].

and plug them into the famous 1-loop Coleman-Weinberg result. The eigenvalues of both the fermionic and the bosonic mass square matrices can be computed analytically. In this model ϕ_{22} remains massless at tree level and its Fermi partners do too. The other 8 eigenvalues split in pairs as follows. The bosonic ones are given by

$$|m_{44}|^2+\frac{1}{2}(|\phi_{22}|^2+|m_{24}|^2\pm|m_{22}|^2)\pm'\frac{1}{2}\sqrt{(|\phi_{22}|^2-|m_{24}|^2\pm|m_{22}|^2)^2+4|m_{44}\phi_{22}^\dagger+m_{44}^\dagger m_{24}|^2}\,, \tag{4.4}$$

with the $\pm$ and $\pm'$ standing for two independent sign choices. The fermionic ones are given by the same expression as the bosonic ones, save that we formally set $m_{22}=0$ (the fermionic sector does not talk directly to the F-term).

If we set $m_{24}=0$ we obtain the classic O'Raifeartaigh result for the one-loop energy, i.e.

$$E_1=\frac{N_c}{32\pi^2}|m_{22}|^4\left(y^{-2}(1+y)^2\log(1+y)+y^{-2}(1-y)^2\log(1-y)+2\log\left(\frac{|m_{44}|^2}{\Lambda^2}\right)\right)\,, \tag{4.5}$$

where we have defined $y=|\frac{m_{22}}{m_{44}}|^2$ and Λ is the UV cutoff. In this case, ϕ_{22} stabilizes around zero.

When $m_{24}\neq 0$, the analytical form of the one-loop energy is not very illuminating. However it is reasonable to expect that ϕ_{22} will pick up a tadpole around zero, so that its vacuum expectation value is displaced to a non-zero value. Moreover, the size of the VEV is directly controlled by m_{24}, as they enter almost symmetrically in the expressions for the eigenvalues. This is confirmed by a numerical analysis, which also shows the existence of tachyons when m_{24} is too close to or larger than m_{44}. Indeed, one might have guessed that in this range some dangerous mixings can occur.

Let us remark at this stage on a possibility which could have been considered. We could have actually tried to generate the higher masses like m_{44} dynamically in the same way as the lower ones, m_{22}. That would be simply implemented in a 5-node quiver with ranks N_f-N_c at the 4th and 5th node. The model would be very similar to the above, except that every O'Raifeartaigh model would have now $1+4(N_f-N_c)$ fields. If the masses were dynamical, one would have a sum of N_f-N_c contributions like (the generalization to $m_{24}\neq 0$ of) eq. (4.5). The latter potential attracts all of the higher masses to smaller values. However one can see that the deformed moduli space constraint is not sufficient in this case to stabilize them. Indeed, in the dominant contribution (the $\log|m_{44}|^2$ piece), the constraint gives trivially a constant. The rest of the potential asymptotes to a constant for very large m_{44}. Hence, it will always be favorable to bring down some masses while sending the other(s) to infinity. This is the reason why we cannot really access the full IR free region $N_c+1\leq N_f<\frac{3}{2}N_c$ of SQCD in this class of models.

4.1 Stabilization of baryonic directions

At this stage, we have seen that as long as we can be on the mesonic branch at node 1, we will successfully obtain a model of metastable supersymmetry breaking with no small parameters added by hand.

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An important question now arises, however. Assuming we are working in the regime $h\Lambda_1^2 < m$, the energy of the SUSY breaking vacuum is given by

$$V \sim N_c\, h^2 \Lambda_3^2\, \Lambda_1^4\,, \tag{4.6}$$

since it arises by summing the masses squared of the N_c lightest flavors of the SU(N_c) gauge group at node 3 [8]. The quantum moduli space constraint for the mesons M_{22} is really

$$\det M_{22} - B\tilde{B} = \Lambda_1^{2N_c}\ . \tag{4.7}$$

So, at least naively, it appears that by relaxing the mesonic VEVs and turning on $B, \tilde{B}$, one can lower the vacuum energy to zero, destabilizing the SUSY breaking vacuum. It is conceivable that the Kähler potential (which is not computable) introduces a barrier that prevents such relaxation, but confidence in the construction would be considerably enhanced if additional superpotential terms were present which prevent the baryons from 'turning on' when one expands around the point (2.3).

We now begin estimating the leading $B, \tilde{B}$ mass matrix by expanding the potential around the would-be non-supersymmetric vacuum. To do so, we assume a canonical Kähler potential. We consider this is reasonable, since both potential instabilities and stabilizing effects arise under this assumption. The leading off-diagonal term is

$$V_{,B\tilde{B}} = V_{,\tilde{B}B} = -\,h^2\, \Lambda_3^2/\Lambda_1^{2N_c-4}\ . \tag{4.8}$$

This contribution appears at tree-level and favors the condensation of baryons as discussed above.

However, there are several further terms in the (super)potential which impart diagonal terms in the mass matrix, and overwhelm the tachyonic contribution (4.8) for reasonable choices of parameters. One source of such a term is the tree-level coupling of ϕ_{22} in (2.5). This will in fact turn out to be the dominant effect stabilizing $B, \tilde{B}$ at zero. For completeness, we also include the sub-dominant effect caused by the stringy instantons discussed in the previous section.

Putting (3.4) together with (2.5), we can check for stability of the mesonic branch VEVs (2.3) as follows. Assume the M_{22} matrix is diagonal, with equal eigenvalues given by x. x is then determined by the deformed quantum moduli space constraint of the node 1 gauge theory to be

$$x^{2N_c} = \Lambda_1^{2N_c} - B\tilde{B}\ . \tag{4.9}$$

We could impose this constraint by adding a Lagrange multiplier λ_1 to the superpotential, multiplying the constraint equation. Then, subject to the constraint, we should minimize the potential

$$V = \Lambda_1^2\, |h\, \Lambda_3\, \phi_{22} + \lambda_1 x^{N_c-1}|^2 + \Lambda_1^{2N_c-2}\, |\lambda_1 B + c_1 B|^2 + \Lambda_1^{2N_c-2}\, |\lambda_1 \tilde{B} + c_1 \tilde{B}|^2\,, \tag{4.10}$$

where the first term arises from $|F_x|^2$, the second term from $|F_{\tilde{B}}|^2$, and the third from $|F_B|^2$, and we have redefined $c_1 = c'\, e^{-\mathrm{Area}'}$ with respect to (3.4).

Equation (4.10) only contributes to the diagonal

$$V_{,BB} = V_{,\tilde{B}\tilde{B}} = \frac{2}{\Lambda_1^{2N_c+2}}(c_1\Lambda_1^{2N_c} - h\Lambda_1^2\Lambda_3\phi_{22})^2 \ . \tag{4.11}$$

From (4.11) we can obtain the leading diagonal terms in the matrix of second derivatives of the potential. The leading contribution is a non-zero expectation value $\phi_{22} \sim m_{24} = h\Lambda_3^2$. The net result is

$$V_{,BB} = V_{,\tilde{B}\tilde{B}} = h^4\Lambda_3^6/\Lambda_1^{2N_c-2}. \tag{4.12}$$

Both h and the instanton coefficient c_1 are suppressed by M_s^*, as $h \sim M_s^{*-1}$ and $c_1 \sim M_s^{*3-2N_c}$. (There is also a suppression by the volume of the curve representing node 6 for the instanton, but since this effect plays no role in our theory as the ϕ_{22} VEV already stabilizes the baryons of node 1, we can take that volume to be anything $\geq \mathcal{O}(1)$. For simplicity we've chosen $\mathcal{O}(1)$ here). The only consistency requirement on the relevant scales is then that $M_s^* > \Lambda_3$.

The eigenvalues of the matrix of second derivatives of the potential are $V_{BB} \pm V_{B\tilde{B}}$. Then, we are free of tachyons provided that $V_{BB} \geq |V_{B\tilde{B}}|$. From (4.8) and (4.12) we conclude the conditions for stability of the baryonic directions are

$$(\Lambda_1/\Lambda_3) < (\Lambda_3/M_s^*) \tag{4.13}$$

We see that we can always satisfy the above inequality, as well as the ones coming from the hierarchy of mass scales in the low-energy model (as in e.g. [27]), by imposing the following hierarchy

$$\Lambda_1 \ll \Lambda_3 < M_s^* \ \text{ and } \ m < \Lambda_3 \ . \tag{4.14}$$

As discussed previously, we also need m to satisfy the bound (2.8).

To summarize, we have checked that the potential baryonic instability is cured. To do this, the one-loop generated VEV of ϕ_{22} is enough. On the other hand, to generate m, a crucial role was played by an additional term in the superpotential, generated by a string instanton. This makes our model a bona-fide version of SQCD with dynamically generated exponentially small quark masses, and allows it to stably display the related metastable vacua.

5. Lifetime of the meta-stable vacuum

In this section we study the possible decay channels for our non-supersymmetric ISS-like vacuum.

As reviewed in section 2, there is a SUSY vacuum where the mesons of node 3 acquire VEVs, which in turn are fixed by the VEVs of the mesons of node 1. This is what we refer to as the mesonic branch, and is the usual SUSY vacuum of SQCD, as considered e.g. in [8]. In addition to this SUSY vacuum, there is also another direction of possible decay, which is precisely the one discussed in the previous section. Along this direction, the baryons of node 1 acquire VEVs, and we are essentially led to an SQCD at node 3

with N_c massless and one massive flavors. By standard arguments used when discussing the deep IR of cascading quivers, one can see that after a Seiberg duality on node 3, the quiver reduces to SQCD with one flavor at node 2. The flavor acquires a mass which is directly related to m.

We thus want to estimate the decay rate towards these two (classes of) vacua. We do this by estimating the bounce action in the following form, using the triangular approximation [40]

$$S \sim \frac{(\Delta\Phi)^4}{\Delta V}, \tag{5.1}$$

where $\Delta\Phi$ is roughly the width of the barrier while ΔV is its height. We will see below that we do not need to be more precise, since we are really interested in lower bounds for the bounce action anyway. If we can tune these lower bounds to be large enough, we can be confident that decay through tunneling is suppressed and meta-stability is not affected.

Let us first consider decay towards the mesonic branch SUSY vacua. Here ΔV is readily evaluated to be of the order of (4.6), since the energy of the metastable state and the peak only differ by a numerical factor. To estimate $\Delta\Phi$, we first note that, as in [8], the fields which have the biggest variation are the mesons of node 3. In the SUSY vacuum, their VEVs are given by

$$\phi_{22} = \Lambda_3 \left(\frac{m}{\Lambda_3}\right)^{\frac{1}{N_c}}, \qquad \phi_{44} = \Lambda_3 \frac{h\Lambda_1^2}{m} \left(\frac{m}{\Lambda_3}\right)^{\frac{1}{N_c}}. \tag{5.2}$$

It is obvious from the above that $\phi_{22} \gg \phi_{44}$ for the range of parameters discussed in the previous sections. Moreover, recall that in the metastable state ϕ_{22} had a VEV of the order of $h\Lambda_3^2$. This is however very small with respect to the VEV it has in the SUSY vacuum, since we assume that $h\Lambda_3 = (\Lambda_3/M_s^*) \ll 1$. We can then identify $\Delta\Phi$ with the VEV of ϕ_{22} listed above.

Putting all together, we have the following estimate for the bounce action towards the mesonic branch

$$S_{\text{mesonic}} \sim \left(\frac{m}{\Lambda_3}\right)^{\frac{4}{N_c}} \left(\frac{\Lambda_3}{\Lambda_1}\right)^4 \left(\frac{M_s^*}{\Lambda_3}\right)^2. \tag{5.3}$$

Every factor in the expression above is (much) greater than one, and we thus conclude that decay towards the mesonic branch is highly suppressed.

As for the decay towards the baryonic branch of node 1, let us provide the most conservative estimate. The field which varies the most along the path is taken to be a representative baryon B. Its variation, after the field has been canonically normalized, is taken to be $\Delta\Phi \sim \Lambda_1$. Note that since Λ_1 is the smallest scale in the game, this is really the most adverse situation. As for ΔV, we can take (4.12) and plug in the maximal VEV of the baryons $B \sim \Lambda_1^{N_c}$, so that we get $\Delta V \sim h^4\Lambda_3^6\Lambda_1^2$. (Note that this ΔV is much larger than the energy of the metastable vacuum.) The bounce action is thus

$$S_{\text{baryonic}} \sim \left(\frac{\Lambda_1}{\Lambda_3}\right)^2 \left(\frac{M_s^*}{\Lambda_3}\right)^4. \tag{5.4}$$

Consistently with the inequality (4.13), the above bounce action can be made parametrically large, and thus also the decay towards the baryonic branch is suppressed.

In this crude estimate, it seems that the latter decay channel is the dominant one. To conclude that this is really so would require a more serious investigation of the potential and tunneling path. In any case we see that the simplest estimates indicate that the apparent decay channels are parametrically suppressed.

6. The string dual description

As we have seen, the quiver gauge theory we have analyzed in previous sections admits a number of supersymmetric as well as metastable non-supersymmetric vacua. A natural question is whether is it possible to provide a supergravity/string dual description of such vacua. The well defined type IIB string embedding of our model outlined in section 1.1 makes this a realistic task.

6.1 On the gravity dual

In what follows we sketch the structure of only those vacua which are most relevant to our story: the ISS-like vacua and their supersymmetric counterparts, i.e. the vacua belonging to the mesonic branch. The discussion is very similar to the one for the $\mathbb{Z}_2$ conifold quotient presented in our previous work [25], to which we refer for details.

In the present construction a crucial role is played by the presence of an additional fractional brane, that we have to treat as a probe since its backreaction cannot be captured classically (by definition, we cannot take the large N limit for a single brane). Moreover, a second equally crucial ingredient is that there is a mass term constraining the position of this probe brane. Indeed, it can be checked that for $m = 0$ both classes of supersymmetric vacua become runaway. Again, the mass term is the product of a stringy instanton which is not expected to backreact on the classical geometry in any simple manner.

Below, we will take the pragmatic point of view that, because of the mass term, we can roughly integrate out the effect of the additional probe brane. We are then left with the same gravity dual as discussed in [25], albeit embedded in a higher singularity. The effects of the additional fractional brane presumably show up as (important) $1/N$ corrections to the geometry.

Our brane system can be embedded into a weakly curved gravity dual background by adding (a large number of) regular D3 branes. One can easily show that the resulting fractional/regular brane system enjoys a duality cascade, i.e. a non-trivial RG-flow along which the effective number of regular branes diminishes (in units of N_c, in this case). Hence, choosing $N = kN_c$ regular D3 branes (with k as large as we wish), the IR field theory at the end of the cascade will be the quiver field theory we have been studying. The number of cascade steps k will just control the final warp factor in the IR region of the gravity dual.

As discussed in section 2, the region of the moduli space where an effective massive SQCD SU(N_c) theory emerges is along the mesonic branch of node 1. This corresponds to

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having the N_c $\mathcal{N}=2$ branes at a distance $(\det M_{22})^{1/N_c} \propto \Lambda_1^2$ along $\mathbb{C}$, the complex direction parametrizing the VEV's of the adjoint scalar of the corresponding effective SU(N_c) $\mathcal{N}=2$ vector superfield. We will refer to the $\mathcal{N}=2$ fractional branes as wrapped D5's, throughout this section.

The $\mathbb{Z}_3$ orbifold of the conifold[4] is described by the following equation in $\mathbb{C}^4$

$$x^3 y^3 = uv \ . \tag{6.1}$$

As already discussed in section 1.1, this geometry supports three independent complex deformations, leading to the completely smooth geometry

$$\prod_{i=1}^{3}(xy - \epsilon_i) = uv \ . \tag{6.2}$$

Consequently, there are three non trivial 3-cycles A_i whose minimal size is given by ϵ_i

$$\int_{A_i} \Omega = \epsilon_i \ . \tag{6.3}$$

In our particular case we would like to consider the case where only two of the three 3-cycles are blown-up, and moreover they have the same size [25]

$$(xy - \epsilon)^2 xy = uv \ . \tag{6.4}$$

This deformation is triggered by the N_c deformation branes we have at node 3. The geometry above has a $\mathbb{C}^2/\mathbb{Z}_2$ line of singularities (also called A_1-singularities, not to be confused with the label of the 3-cycles above) at the locus $xy = \epsilon$, $u = v = 0$. Moreover, it has an innocuous conifold singularity at $x = y = u = v = 0$.

We construct the geometrical dual to the supersymmetric vacua of our theory, which were discussed at the beginning of section 2, in the following way. After a geometric transition, we expect the brane at node 3 to transmute and turn into flux,

$$\int_A G_3 = N_c \ , \quad \int_B G_3 = \frac{i}{g_s} \, k \tag{6.5}$$

where A is the compact 3-cycle absorbing the RR flux of the original branes, B its non-compact dual, and k is the number of duality cascade steps (and can be naturally taken to be very large). As already noticed, the D5's wrapping the $\mathbb{C}^2/\mathbb{Z}_2$ singularity are instead explicitly present in the dual geometry, lying somewhere along the mesonic branch. Finally, the single deformation brane at node 4 cannot transmute, and remains as a probe at the remaining conical singularity, which we expect to be slightly deformed by the instanton discussed in section 3. It can be checked that the above set up has the same charges and supersymmetric moduli space as our theory.

[4] As already mentioned, we can embed the same 4-node quiver in a $\mathbb{Z}_N$ orbifold with $N > 3$. A reason to go to larger N might be to achieve the desired range of scales, since there would be more geometrical quantities to tune.

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We now move to the description of the metastable state. As originally discussed in [3], and recently applied in similar contexts by many authors, a natural way to construct metastable non-supersymmetric vacua is by adding anti-branes. The positive vacuum energy is proportional to the number of such branes. The fact that the vacuum energy is exponentially small is due, in the gravity dual, to the warping of the anti-brane tension.

In order for these configurations to describe states in the same gauge theory, one should check that the supergravity charges, at infinity, are unchanged. In the present context this can be achieved by adding N_c anti-D3 branes and simultaneously jumping the NS fluxes by one unit

$$\int_B G_3 = \frac{i}{g_s}\, k \longrightarrow \int_B G_3 = \frac{i}{g_s}\, (k+1)\,, \tag{6.6}$$

so as to leave the full D3 brane charge untouched.[5] It is a nice check of our proposal that it is only by adding N_c such branes (no more, no less), hence providing the correct energetics for the ISS metastable vacua, that we can leave the global charges at infinity untouched, and hence describe non-supersymmetric states in the same gauge theory. Notice that as far as fluxes are concerned, this shift corresponds to moving one step down in the cascade. That this is the case, will become apparent when we discuss the type IIA T-dual description of this system in the next subsection.

Due to the F_5 background in the gravity solution (dual to the large number of D3 branes present in the cascade), the anti-D3 branes are attracted to the tip of the geometry. The metastable configuration presumably has all anti-D3s absorbed and dissolved as gauge flux into the N_c D5 branes

$$\int_{\mathcal{C}} \mathcal{F} = -N_c\,, \tag{6.7}$$

where $\mathcal{C}$ is the 2-cycle which the N_c D5s wrap. This flux, via the Chern-Simons coupling in the DBI action of the D5 branes, accounts for the N_c units of anti-brane charge. Stability against decay through the Myers effect can be argued as in [25], but of course a more detailed study would be valuable.

A natural question is to ask whether one of the anti-branes can annihilate with the deformation brane associated to node 4, which is sitting at the conifold singularity. In the next subsection we will provide a simple argument as to why this is energetically disfavoured.

The supersymmetric vacua corresponding to the baryonic branch were discussed in [25]. It is not immediately clear how one would directly relate them to the metastable states.

6.2 Type IIA dual

In this section we study the Type IIA dual configurations describing our model. These constructions were first introduced in [41]. This approach provides a vivid picture of how the anti-branes arise in the non-supersymmetric state.

figure 4 shows the type IIA T-dual brane configuration for our 4 node quiver for $\langle M_{22}\rangle = m = 0$.

[5] Let us remind the reader that the full D3 charge is $Q_3 = \int H_3 \wedge F_3 + N - \overline{N}$, where N and $\overline{N}$ are the net number of D3 and anti-D3 branes, respectively.

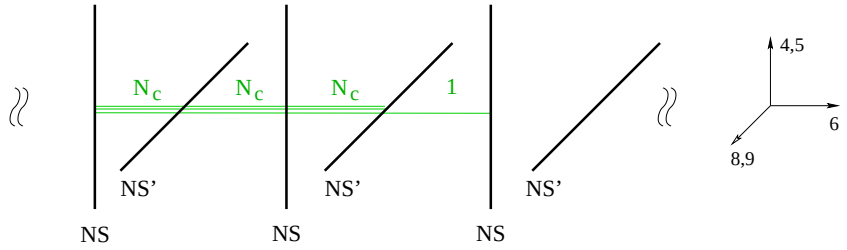

Figure 4: Type IIA configuration for the electric theory with $\langle M_{22} \rangle = m = 0$.

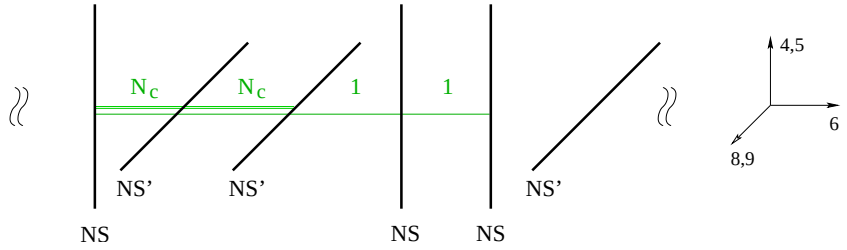

Figure 5: Type IIA configuration for the magnetic theory.

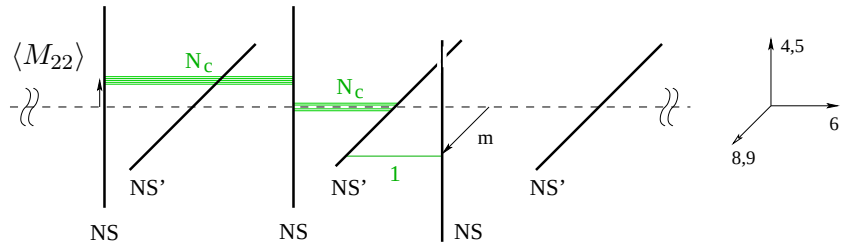

Figure 6: Type IIA configuration for the electric theory with non-zero $\langle M_{22} \rangle$ and m.

Performing a Seiberg duality on the middle node corresponds to moving the second NS brane and second NS' brane across each other [42]. In the process, some anti D4 branes are generated in the middle interval, but they are annihilated against D4 branes sitting on top of them. The result is shown in figure 5.

A non-zero VEV for $\langle M_{22} \rangle$ in the electric theory corresponds to moving the D4-branes stretched between the first and second NS-brane in figure 4 in the 45 directions. Similarly, the mass m is mapped to a displacement of the third NS brane together with the D4-branes that stretch from it to the second NS'-branes in 89. This is shown in figure 6.

If we now perform the Seiberg duality, we obtain the configuration in figure 7. The anti-D4's are not annihilated, since they are now displaced from the D4s due to the meson VEVs (following the discussion in section 3, $\langle M_{22} \rangle$ is stabilized at tree level).

figure 7 shows that the avatar of SUSY breaking in the ISS vacuum is explicit un-annihilated anti D4 branes. The failure to annihilate these branes is a direct consequence of the meson VEVs and m (the dynamically generated masses). Relative to the IIB story we described in the previous subsection, this IIA configuration is an intermediate picture between the state with explicit anti-D3 branes and the (final) state where they are dissolved

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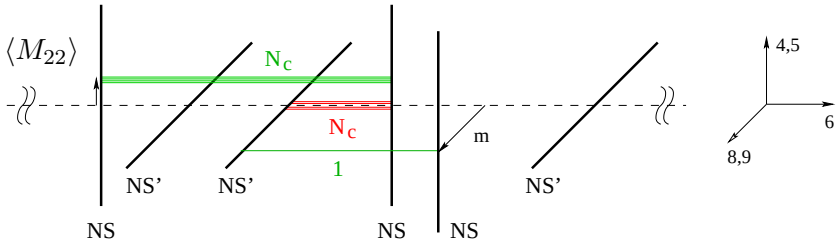

Figure 7: Type IIA configuration for the magnetic theory with non-zero $\langle M_{22} \rangle$ and m. Anti-D4 branes are indicated in red.

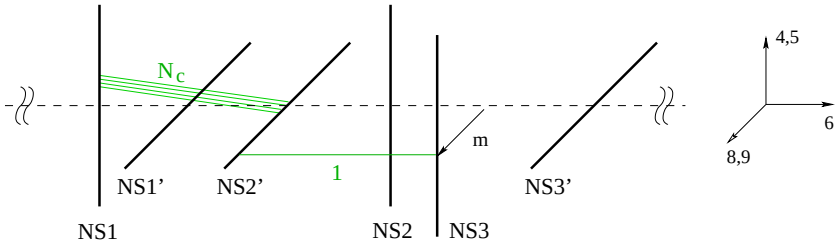

Figure 8: The type IIA configuration once the anti-D4 branes has been dissolved into the D4's as gauge flux.

within the N_c D5s at the A_1 singularity.

Collapsing D4 and anti D4-branes as much as possible, we obtain figure 8 for the IIA picture of the ISS vacuum. We have labeled the NS branes to simplify the discussion. We have annihilated the anti-D4's against D4's between NS1 and NS2 (and not NS2' and NS3) because, since $|h\langle M_{22} \rangle| < m$, this clearly produces a lower energy configuration. Notice that the objects the gauge flux (represented here by the anti-D4's) combines with are precisely D4 branes stretched between the parallel NS1 and NS2. The resulting tilted D4 branes are dual to the D5/anti-D3 bound states of the IIB configuration.

In this type IIA setting it is also easy to see how the configuration in figure 7 is related to the addition of anti D3-branes and jumping fluxes in the gravity dual, as described in section 6.1. An anti-D3 brane maps to an anti-D4 brane wrapping the entire x_6 circle in Type IIA. We can form such a complete anti-D4 by adding D4/anti-D4 pairs to all intervals with the exception of the one between the second NS' and second NS. Grouping D4 branes in each interval together, we are left with the configuration in figure 9. It corresponds to N_c full anti-D4 branes (T-dual to N_c anti-D3 branes) and the number of D4 branes in each interval corresponds to moving up one step in the cascade from the magnetic theory (the last step). Increasing the cascade by one step is exactly how increasing the NS flux by one unit manifests in this context. This matches nicely with our previous type IIB description of the metastable non-supersymmetric vacua.

7. Conclusions

We have presented a D-brane construction that engineers metastable vacua closely related to those of [8]. The construction has some interesting conceptual features and some inter-

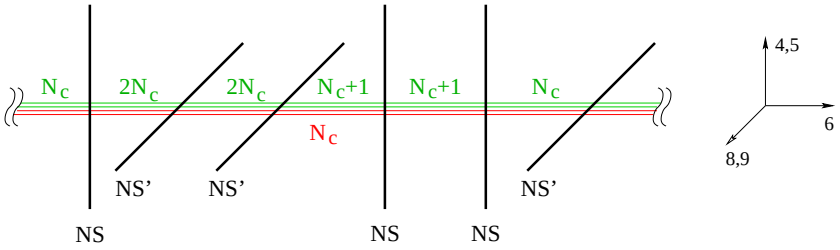

Figure 9: The ISS vacuum configuration of figure 7 can be interpreted as adding N_c anti-D3 branes and moving down one step in the cascade.

esting features for model building.

Conceptually, the most interesting thing about the construction is that it readily admits a IIB gravity dual description (in the general framework of AdS/CFT). The metastable states cannot be directly followed from weak 't Hooft coupling to strong 't Hooft coupling, but they quite plausibly match on to strong-coupling analogues where the SUSY breaking is well described by the presence of anti-D3 branes, and a picture very similar to the one in [3]. In addition, the quivers that arise in our construction are some of the simplest cases where the new string instanton effects explored in many recent works make important contributions.

One may also wish to construct pseudo-realistic models of SUSY breaking and/or direct mediation using quiver gauge theories. In this case, our model has two virtues: the small dynamical masses of the ISS model are explained naturally without any fine tuning of parameters (as could also be done by retrofitting [10]), and the problematic R-symmetry which could forbid gaugino masses is lifted by the extra terms that automatically appear in our superpotential. We note that because the R-symmetry is broken by an irrelevant operator suppressed by the high scale M_s^*, and because for metastability it is necessary to take M_s^* somewhat higher than the SUSY-breaking scale, it is likely that in real model-building applications, our theory would produce low gaugino masses. Given the lower bounds on gaugino masses, this would necessitate heavy squarks and sleptons, resulting in the need for a (mild) tune to obtain a reasonable Higgs mass.

It would be very interesting to study the gravity dual geometry in further detail. While there aren't BPS or protected quantities that are guaranteed to match between weak and strong coupling, one may find interesting patterns of qualitative agreement between the two classes of non-supersymmetric states. Conversely, some quantities (e.g. lifetimes or barrier heights) may change in a striking way upon extrapolation in $g_s N$.

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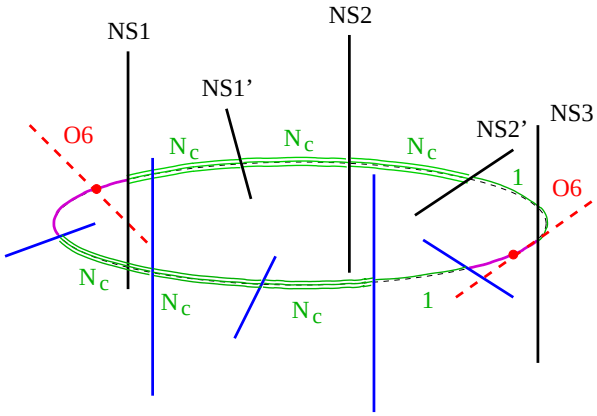

Figure 10: Type IIA T-dual configuration for an embedding of our model in an orientifold of a $\mathbb{Z}_5$ orbifold of the conifold. The image NS-branes are indicated in blue. The ED1 branes are T-dualized to ED0s, which are shown in magenta.

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A. Fermionic zero modes and orientifolds

We now briefly explain how it is possible to project out two fermionic zero modes on each instanton by embedding our model in an orientifold.

The most intuitive way of visualizing the orientifold is by means of the Type IIA T-dual setup, along the lines of [43], to which we refer the reader for further details. Everything in this construction can be mapped into a type IIB set-up.

The simplest embedding of our model is shown in figure 10. It corresponds to removing the last NS' in figure 4 and adding the orientifold images. The O-plane extends along 01237 and is at 45 degrees with respect to the 45 and 89 planes. Such an O-plane maps NS to NS' branes and vice versa. Before orientifolding, the corresponding geometry is a $\mathbb{Z}_5$ orbifold of the conifold.

If the O-plane extended along 45 or 89, the images of NS1 and NS3 would also be NS branes. We have not chosen this possibility because the instantons stretched between parallel NS-branes (in type IIB, this corresponds to the instantons wrapping $\mathbb{P}^1$'s in A_1 singularities) would have additional adjoint fermionic zero modes, some of which also survive the orientifold projection.

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We take the charge of the O-plane to be positive so that the gauge group on the instantons is $O(1)$. Two of the fermionic zero modes are projected to the symmetric representation and the other two to the antisymmetric representation of $O(1)$. The antisymmetric representation of $O(1)$ vanishes, so this orientifold projects out precisely two fermionic zero modes.

Only the cycles wrapped by the instantons are affected by the orientifold, thus our discussion of the IR applies without changes. The cascade and the corresponding supergravity solution section 6.1 require some small modifications in the presence of orthogonal gauge groups/orientifold planes, probably along lines similar to [44].

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