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Massimo Bianchi, Francesco Fucito and Jose F. Morales

Dipartimento di Fisica, Universitá di Roma "Tor Vergata", I.N.F.N. Sezione di Roma II,
Via della Ricerca Scientifica, 00133 Roma, Italy
E-mail: bianchi@roma2.infn.it, fucitoroma2.infn.it,
moralesroma2.infn.it

ABSTRACT: We study the local dynamics of an unoriented $\mathbb{Z}_5$-quiver theory of GUT kind with gauge group U(5) and chiral matter. At strong coupling the non-perturbative dynamics is described in terms of baryon and meson variables satisfying a quantum deformed constraint. We compute the effective superpotential of the theory and show that it admits a line of supersymmetric vacua for vanishing FI terms and a phase where supersymmetry is dynamically broken via gaugino condensation for a non-vanishing FI term.

Contents

1 Introduction

The search for mechanisms of supersymmetry breaking in string and field theories has a long history. Out of the many alternatives, the idea that supersymmetry be broken dynamically remains one of the most attractive choices. In this framework, supersymmetry remains a symmetry of the effective action even after non-perturbative effects are included and is only spontaneously broken by the vacuum choice [1–5]. This is in contrast with scenarios based on explicit supersymmetry breaking where the loss of supersymmetry has more dramatic effects and can be hardly followed in a controllable way. Still the conditions under which supersymmetry is broken are very restrictive and the search for string realizations of such scenarios remains a far from obvious task (see [6, 17] for recent contributions in the subject).

In this paper we study in details an example of a $\mathcal{N} = 1$ chiral model, based on D3-branes at a $\mathbb{C}^3/\mathbb{Z}_5$ singularity, exhibiting dynamical supersymmetry breaking. At low energy the D-brane dynamics is described by a quiver gauge theory with gauge groups and matter representations specified by the choice of fractional D3-branes. We consider a $\mathbb{Z}_5$ quiver leading to a gauge theory with gauge group $U(5) \times U(1)$, two flavours of chiral matter in the **10** dimensional representation of the gauge group, three flavours of $\bar{\mathbf{5}}$, one flavour of **5** and two flavours of charged singlets. We present a complete description of the non-perturbative dynamics of the gauge theory. The strong coupling dynamics of the theory is described in terms of a set of gauge invariant composites, 'baryons' and 'mesons', satisfying a quantum deformed constraint. The vacuum structure of the theory depends

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very much on whether field dependent Fayet-Iliopoulos (FI) terms are turned on or not. In the absence of FI terms a supersymmetric line of vacua is found. Turning on FI terms, the theory undergoes a Higgs mechanism, whereby a $U(1)_1$ charged $SU(5)$ singlet gets a VEV and supersymmetry is dynamically broken by instanton effects.[1]

Our analysis will be local and we will address neither the issue of how to stabilize the VEV's of the closed-string moduli, that determine the FI terms, nor the issue of how to embed the quiver gauge theory in a global vacuum configuration of (unoriented) D-branes. At low energies the physics near the singularity is efficiently captured by the local description of the singularity and constraints such as global tadpole cancellation can be neglected at this stage. We assume that a global description exist, whose bulk dynamics allows to stabilize the twisted closed-string moduli and thus FI terms as desired, and that the other (unoriented) D-branes or fluxes needed for global R-R tadpole cancellation are located far from the singularity. We remark that although T^6 tori do not admit $\mathbb{Z}_5$ actions compatible with residual supersymmetry, 'supersymmetric' $\mathbb{C}^3/\mathbb{Z}_5$ singularites may show up in Calabi Yau compactifications as for instance in the $\mathbb{Z}_5$ quotient of the quintic.[2]

The paper is organized as follows. In section 2, we briefly review the construction of unoriented $\mathbb{Z}_n$ quiver gauge theories and identify a grand unification (GUT) like model with gauge group $U(5) \times U(1)$ and chiral matter content. In section 3 we describe the strong coupling dynamics of the theory. In section 4 we test our strong coupling description by matching the anomalies of the effective theory described in terms of baryon/meson composites with those computed in terms of elementary fields. In section 5 we identify the relevant instanton moduli space and interactions and sketch the D-instanton derivation of the non-perturbative superpotential. In section 6 we summarize our results.

2 Unoriented $\mathbb{Z}_n$ quivers

The low-energy dynamics of N D3-branes in flat spacetime is described by a $\mathcal{N} = 4$ supersymmetric gauge theory in four dimensions with gauge group $U(N)$. In the $\mathcal{N} = 1$ language this theory comprises a vector superfield V and three chiral fields Φ^I, $I = 1, 2, 3$ transforming in the adjoint representation of $U(N)$. In addition one has a cubic superpotential of the form

$$W_{\text{tree}} = \text{Tr}\, \Phi^1 [\Phi^2, \Phi^3] \tag{2.1}$$

If one puts D3-branes at a singularity $\mathbb{C}^3/\Gamma$, with Γ a discrete subgroup of $SU(3)$, the low-energy dynamics is described by an $\mathcal{N} = 1$ quiver gauge theory with bi-fundamental matter specified by the intersection matrix of cycles on $\mathbb{C}^3/\Gamma$ and a cubic superpotential descending from (2.1).

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[1]Although both $U(1)$'s are anomalous and massive after eating closed-string axions in the twisted sectors, we would like to stress the difference between these 'Stückelberg fields', whose global shift symmetries are gauged by the open-string vectors, from standard 'Higgs fields' such as the $U(1)_1$ charged $SU(5)$ singlet taking a non-zero VEV.

[2]We thank E. Witten for suggesting us this embedding.

We take $\Gamma = \mathbb{Z}_n$ acting on the $\mathbb{C}^3$ coordinates z_I as

$$\mathbb{Z}_n: \qquad z^I \to \omega^{q_I} z^I \qquad \omega = e^{2\pi i/n} \qquad \sum_{I=1}^{3} q_I = 0 \bmod n \qquad (2.2)$$

The orbifold group action on Chan-Paton indices can be represented by the γ_g block diagonal matrix

$$\gamma_g = \begin{pmatrix} \mathbb{1}_{N_0 \times N_0} & 0 & ... & ... \\ 0 & \omega \, \mathbb{1}_{N_1 \times N_1} & 0 & ... \\ ... & 0 & ... & 0 \\ ... & ... & 0 & \omega^{n-1} \, \mathbb{1}_{N_{n-1} \times N_{n-1}} \end{pmatrix} \qquad (2.3)$$

with $N = \sum_a N_a$ the original number of branes and N_a the number of fractional D3-branes of type "a". The quiver gauge theory at the singularity is defined by restricting the $\mathcal{N} = 4$ fields V, Φ^I to $N \times N$ matrices satisfying the $\mathbb{Z}_n$-invariant conditions

$$\mathbb{Z}_n: \qquad V = \gamma_g V \gamma_g^\dagger \qquad \Phi^I = \omega^{q_I} \gamma_g \Phi^I \gamma_g^\dagger \qquad (2.4)$$

This results into block matrices with non-trivial components

$$V: \sum_{a=0}^{n-1} N_a \times \bar{N}_a$$

$$\Phi^I: \sum_{a=0}^{n-1} N_a \times \bar{N}_{a+q_I} \qquad (2.5)$$

The resulting quiver gauge theory has then gauge group $G = \prod_{a=0}^{n-1} \mathrm{U}(N_a)$ and bi-fundamental matter.

In general, the matter content (2.5) is either non-chiral or anomalous and therefore of very limited interest from the phenomenological point of view. More appealing gauge theories can be realized by introducing Ω-planes or equivalently by quotienting the quiver by an anti-holomorphic involution Ω that includes world-sheet parity. Ω identifies ingoing $\{N_a\}$ and outgoing $\{\bar{N}_a\}$ boundaries and invariant fields are defined by the restriction

$$\Omega: \qquad V = \epsilon \, \gamma_\Omega V \gamma_\Omega^\dagger \qquad \Phi^I = -\epsilon \, \gamma_\Omega \Phi^I \gamma_\Omega^\dagger \qquad (2.6)$$

with ϵ a sign and γ_Ω specifying the identification of boundaries. A canonical choice is given by taking[3]

$$\bar{N}_a = N_{n-a} \qquad (2.7)$$

and

$$\Omega: \qquad \gamma_\Omega = \begin{pmatrix} \mathbb{1}_{N_0 \times N_0} & 0 & ... & ... & ... & 0 \\ 0 & 0 & ... & ... & 0 & \mathbb{1}_{N_1 \times N_1} \\ ... & ... & ... & ... & \mathbb{1}_{N_2 \times N_2} & 0 \\ ... & ... & ... & ... & ... & ... \\ ... & 0 & \mathbb{1}_{N_2 \times N_2} & 0 & ... & ... \\ 0 & \mathbb{1}_{N_1 \times N_1} & 0 & ... & ... & ... \end{pmatrix} \qquad (2.8)$$

---

[3]For n even an extra choice is possible $\bar{N}_a = N_{n-a-1}$ and γ_Ω a block matrix with non-trivial blocks $\gamma_\Omega = \oplus_{a=0}^{n-1} \mathbb{1}_{N_a \times N_{n-a-1}}$.

	$\mathbb{Z}_3$
V	$Sp/O(N_0) \times U(N_1)$
Φ^I	$3\left[N_0\bar{N}_1 + \frac{1}{2}N_1(N_1 \pm 1)\right]$
$I_a = 0$	$N_0 = N_1 \pm 4$
	$\mathbb{Z}_4$
V	$Sp/O(N_0) \times Sp/O(N_2) \times U(N_1)$
Φ^I	$2(N_1N_2 + N_0\bar{N}_1) + N_0N_2 + \frac{1}{2}N_1(N_1 \pm 1) + \frac{1}{2}\bar{N}_1(\bar{N}_1 \pm 1)$
$I_a = 0$	$N_0 = N_2$
	$\mathbb{Z}_5$
V	$Sp/O(N_0) \times U(N_1) \times U(N_2)$
Φ^I	$2\left[N_0\,\bar{N}_1 + N_1\,\bar{N}_2 + \frac{1}{2}N_2(N_2 \pm 1)\right] + N_2\,N_0 + \bar{N}_1\,\bar{N}_2 + \frac{1}{2}N_1(N_1 \pm 1)$
$I_a = 0$	$N_0 = N_1 = N_2 \pm 4$

Table 1. Low energy theories arising from the quivers described in the main text.

Moreover depending on the choice of R-R charge $\epsilon = \pm$ of the $\Omega^{\pm}$-plane one has to (anti)symmetrize the open strings with ends on image stacks [18–20]. This results into a quiver gauge theory with gauge group a product of unitary (for N_a complex) and orthogonal (symplectic) gauge group components for N_a real and $\epsilon = -(+)$. In addition chiral matter appear in bi-fundamental, symmetric or antisymmetric representations of the gauge group according to (2.6). Finally (local) twisted tadpole cancellation, or equivalently cancellation of gauge anomalies, constrains the choices of N_a [21, 22]. More precisely, denoting by

$$n_{f,a\pm} = n_{f,a} \pm \bar{n}_{f,a} \qquad n_{S,a\pm} = n_{S,a} \pm \bar{n}_{S,a} \qquad n_{A,a\pm} = n_{A,a} \pm \bar{n}_{A,a} \tag{2.9}$$

the chiral-antichiral (sum) difference of multiplets transforming in the fundamental, symmetric or antisymmetric representation of a $U(N_a)$ gauge group, the cancellation of gauge anomalies requires[4]

$$\begin{aligned} I_a &= \sum_{\mathbf{R_a}} A(\mathbf{R_a}) \\ &= \frac{1}{2}n_{f,a-} + \frac{1}{2}n_{S,a-}(N_a + 4) + \frac{1}{2}n_{A,a-}(N_a - 4) = 0 \end{aligned} \tag{2.10}$$

where the sum is over all the representations $\mathbf{R_a}$ of chiral multiplets. Solving these conditions at each node, one finds for the first few choices of $\mathbb{Z}_n$ the series of anomaly-free quiver gauge theories listed in table 1. We are interested in instanton generated superpotentials in these quiver gauge theories. A superpotential is generated by gauge instantons if and only if the number of fermionic zero modes in the instanton background $\dim\mathfrak{M}_{a,k}^{\text{ferm}}$ and the β function of the gauge theory β_a satisfy the relation [23, 24]

$$\dim\mathfrak{M}_{k_a,N_a}^{\text{ferm}} = 2k_a\beta_a - 4 \tag{2.11}$$

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[4]In our conventions: $\text{tr}_{\mathbf{R}}T^a\{T^b,T^c\} = A(\mathbf{R})d_{abc}$ and $A(\mathbf{F}) = \frac{1}{2}$.

subgroup	matter content	gauge group	$(N_0, N_1, ..)$
$\mathbb{Z}_3$	$3\,(\mathbf{10} + \bar{\mathbf{5}})$	U(5)	(1,5)
$\mathbb{Z}_4$	$(\mathbf{10} + \bar{\mathbf{10}}) + 2(\mathbf{5} + \bar{\mathbf{5}}) + \mathbf{1}$	U(5)	(1,5,1)
$\mathbb{Z}_5$	$2(\mathbf{10} + \bar{\mathbf{5}}) + (\mathbf{5} + \bar{\mathbf{5}}) + 2 \times \mathbf{1}$	U(1) $\times$ U(5)	(1,1,5)

Table 2. Matter content of the quiver theories with a U(5) factor in the gauge group.

with k_a the number of D(-1)-branes (instanton number) sitting on top of the N_a D3-branes and[5]

$$
\begin{aligned}
\dim \mathfrak{M}_{k_a,N_a}^{\text{ferm}} &= 2k_a \left[\ell(\mathbf{Adj_a}) + \sum_{\mathbf{R_a}} \ell(\mathbf{R_a}) \right] \\
&= k_a \left[2N_a + n_{f,a+} + n_{S,a+}(N_a + 2) + n_{A,a+}(N_a - 2) \right] \\
\beta_a &= 3\,\ell(\mathbf{Adj_a}) - \sum_{\mathbf{R_a}} \ell(\mathbf{R_a}) \\
&= 3N_a - \frac{1}{2}n_{a,f+} - \frac{1}{2}n_{S,a+}(N_a + 2) - \frac{1}{2}n_{A,a+}(N_a - 2)
\end{aligned}
\tag{2.12}
$$

the number of fermionic zero modes around the instanton background. When condition (2.11) is satisfied an Affleck-Dine-Seiberg (ADS) like superpotential is generated

$$
W_{\text{non-pert}} = \frac{\Lambda^{k_a \beta_a}}{\mathcal{I}_{k_a \beta_a - 3}}
\tag{2.13}
$$

with $\mathcal{I}_{k_a \beta_a - 3}$ denoting a gauge-invariant flavor-singlet made out of $k_a \beta_a - 3$ chiral superfields.

Non-perturbative D-brane instanton effects in the $\mathbb{Z}_3$ case have been studied in [23–25]. A superpotential in this series is generated only for the U(4) gauge theory choice with Ω^--planes and Sp(6) $\times$ U(2) for Ω^+-planes. Here we are interested in quiver gauge theories with U(5) gauge groups and GUT matter. Such matter contents can be realized by taking the orthogonal series and choosing properly the N_a's in table 1. In table 2 we list the matter content of these theories. Using (2.12) one can easily check that none of these theories satisfy the condition (2.11) and therefore there is apparently no room for instanton effects. However, in presence of a tree level superpotential (2.1), this is not the end of the story, since a $\mathbf{5} - \bar{\mathbf{5}}$ pair can get mass from Yukawa interactions and decouples from the spectrum allowing for an instanton generated superpotential. The focus on this paper is on the non-perturbative dynamics of this $\mathbb{Z}_5$-quiver gauge theory.

The field content of the $\mathbb{Z}_5$ quiver gauge theory is summarized in table 3. U(5) $\sim$ SU(5) $\times$ U(1)$_5$ and U(1)$_1$ are gauge symmetries while SU(2) is a global symmetry rotating the two generations. We name the various chiral multiplets (the invariant blocks of Φ^I) as $\{A^i, B^i, C^i, C^3, E\}$. The superpotential follows from (2.1) by restricting Φ^I to their $\mathbb{Z}_5$-invariant components and can be written as

$$
W_{\text{tree}} = C^{iu} B_i E_u + C^{iu} A_{iuv} C^{3v}
\tag{2.14}
$$

---

[5]In our conventions: $\text{tr}_{\mathbf{R}} T^a T^b = \ell(\mathbf{R})\delta_{ab}$ and $\ell(\mathbf{F}) = \frac{1}{2}$.

fields	$SU(2)_F$	$SU(5)$	$U(1)_1 \times U(1)_5$
A^i_{uv}	**2**	**10**	$(0, 2)$
B^i	**2**	**1**	$(-1, 0)$
C^{iu}	**2**	$\mathbf{\bar{5}}$	$(1, -1)$
C^{3u}	**1**	$\mathbf{\bar{5}}$	$(-1, -1)$
E_u	**1**	**5**	$(0, 1)$

Table 3. Chiral matter field content. Indices $i = 1, 2$, $u = 1, ..5$ run over the fundamentals of the SU(2) flavor and SU(5) gauge groups.

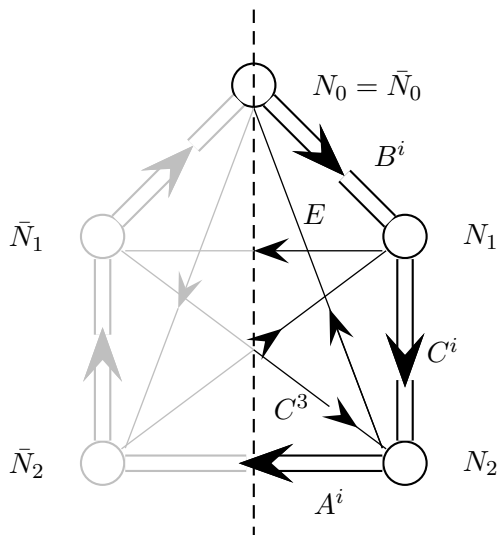**Figure 1**. The $\mathbb{Z}_5$ quiver. The gauge group $U(5) = SU(5) \times U(1)_5$ $(U(1)_1)$ is given by the stack with N_2 (N_1) branes respectively. The figure is symmetric with respect to the orientifold plane which is represented as a dashed line. Our convention is that the fields represented by entering arrows have negative quantum numbers.

with $u, v = 1, ..5$, $i = 1, 2$. The superpotential (2.14) and the quantum numbers of table 3 can also be deduced from the quiver diagram in figure 1.

3 Non-perturbative dynamics

3.1 SQCD: a short review

The studies of gauge theories at the non-perturbative level have a long history. The proto-typical case is that of SQCD, i.e. $\mathcal{N} = 1$ SU(N_c) gauge theory with N_f flavors Q^i and $\tilde{Q}_i$, $i = 1, ..N_f$, in the fundamental and anti-fundamental representations of the gauge group respectively (see [3, 26–29] for reviews). A superpotential is generated by instantons only for $N_f = N_c - 1$. For a generic N_f, the strong coupling dynamics is described in terms of composite fields: the mesons M^i_j and (for $N_f \geq N_c$) baryons $B^{[i_1...i_{N_c}]}$, $\tilde{B}^{[i_1...i_{N_c}]}$

$$M^i_j = Q^i \tilde{Q}_j \qquad B^{[i_1...i_{N_c}]} = Q^{[i_1}....Q^{i_{N_c}]} \qquad \tilde{B}^{[i_1...i_{N_c}]} = \tilde{Q}^{[i_1}....\tilde{Q}^{i_{N_c}]} \qquad (3.1)$$

The case $N_f = N_c$ is special in the sense that the $B, \tilde{B}$ are singlets and satisfy the (classical) algebraic relation

$$\det M - B\tilde{B} = 0 \tag{3.2}$$

A similar relation between composites will be found in our analysis of the $\mathbb{Z}_5$-quiver gauge theory later in this section. Before analyzing the U(5) theory, it is convenient to recall how the non-perturbative dynamics manifests in this more standard SQCD setting.

At the quantum level the relation (3.2) is deformed to

$$\det M - B\tilde{B} = \Lambda^\beta \qquad \beta = 2N_c \tag{3.3}$$

with Λ the SQCD scale. Introducing a Lagrangian multiplier U for this constraint an effective superpotential for the theory can written as

$$W_{\text{eff}} = W_{\text{tree}} + \text{U}(\det M - B\tilde{B} - \Lambda^{2N_c}) \tag{3.4}$$

with W_{tree} the tree level superpotential. This formula can be tested by taking for the tree level superpotential a mass term for the dth quark-antiquark pair

$$W_{\text{tree}} = mQ^d\tilde{Q}_d \tag{3.5}$$

In the limit $m \to \infty$ the massive quarks decouple and one is left with SQCD with $N_f = N_c - 1$ flavors where an ADS superpotential is expected. Indeed writing

$$M = \begin{pmatrix} \hat{M}_{ij} & X_j \\ Y_i & Z \end{pmatrix} \qquad i,j = 1, ... N_c - 1 \tag{3.6}$$

and solving the F-flatness conditions for the massive components

$$\frac{\partial W_{\text{eff}}}{\partial X_j} = \frac{\partial W_{\text{eff}}}{\partial Y_i} = \frac{\partial W_{\text{eff}}}{\partial B} = \frac{\partial W_{\text{eff}}}{\partial \tilde{B}} = \frac{\partial W_{\text{eff}}}{\partial U} = 0 \tag{3.7}$$

one finds

$$X_j = Y_i = B = \tilde{B} = 0 \qquad Z = \frac{\Lambda^{2N_c}}{\det\hat{M}} \tag{3.8}$$

Plugging this into (3.4) one finds

$$W_{\text{eff}} = \frac{m\Lambda^{2N_c}}{\det\hat{M}} = \frac{\hat{\Lambda}^{2N_c+1}}{\det\hat{M}} \tag{3.9}$$

as expected for the ADS superpotential in SU(N_c) SQCD with $N_f = N_c - 1$ flavors after the identification $\hat{\Lambda}^{2N_c+1} = m\Lambda^{2N_c}$ between the scales of the two theories. Remarkable, the non-perturbative superpotential for SQCD with $N_f = N_c - 1$ origins from the quantum deformation of the classical relation between baryons and mesons in the parent theory with $N_f = N_c$! In the following section we will find similiar results for the U(5) chiral theory emerging from the unoriented $\mathbb{Z}_5$ quiver.

3.2 $\mathbb{Z}_5$-quiver superpotential

The first step in defining the gauge theory in the strong coupling regime is to identify the right gauge invariant degrees of freedom, 'mesons' and 'baryons'. The two U(1) factors are anomalous and decouple anyway in the IR. Therefore we should consider SU(5) rather than U(5) invariants. We notice that in the absence of W_{tree} the fields C^{iu} and C^{3u} can be grouped into a triplet that we will denote by C^{Iu}. SU(5) invariants will then be built out of the elementary fields A^i_{uv}, C^{Iu}, E_u and will be labelled by the indices $I = 1, 2, 3$ and $i = 1, 2$ in the fundamental of SU(3)_C and SU(2)_A flavor groups respectively. The complete list of composites reads

$$
\begin{aligned}
X^I &= C^{Iu} E_u \\
Y^i_I &= \frac{1}{2} \epsilon_{IJK} A^i_{u_1 u_2} C^{Ju_1} C^{Ku_2} \\
\tilde{Y}^{ij} &= \frac{1}{4} \epsilon^{u_1 \dots u_5} A^i_{u_1 u_2} A^j_{u_3 u_4} E_{u_5} \\
Z^{iI} &= \frac{1}{12} \epsilon^{u_1 \dots u_5} A^i_{u_1 u_2} A^j_{u_3 u_4} A_{u_5 v, j} C^{Iv}
\end{aligned}
\tag{3.10}
$$

One can check that $X, Y, \tilde{Y}, Z$ satisfy the algebraic relations

$$
Y^i_I Z^I_i = 0
$$
$$
\epsilon_{IJK} X^I Z^J_i Z^{iK} + Y_{iI} \tilde{Y}^{ij} Z^I_j = 0
\tag{3.11}
$$

Indices i, j are raised and lowered with $\epsilon^{ij}, \epsilon_{ij}$. Together with B_i one finds

$$
18 = 42_{\text{fields}} - 24_{\text{gauge}} = 20_{\text{composites}} - 2_{\text{relations}}
\tag{3.12}
$$

gauge invariant degrees of freedom which one can use as coordinates on the moduli space. As indicated, the very same number of degrees of freedom can also be obtained by subtracting to the 42 degrees of freedom coming from elementary fields the $24 = \dim[\text{SU(5)}]$ independent generators of gauge invariance.

It is important to notice that the two relations in (3.11) have engineering dimension 7 and 10 and therefore only the second one can get modified by a non-perturbative $\Lambda^\beta = \Lambda^{10}$ term. Introducing the Lagrange multipliers U, V for the two relations, the effective superpotential of the unoriented $\mathbb{Z}_5$-quiver gauge theory can then be written as

$$
W_{\text{eff}} = B_i X^i + \delta^I_i Y^i_I + V Y^i_I Z^I_i + U \left(Y^i_I \tilde{Y}_{ij} Z^{Ij} + \epsilon_{IJK} X^I Z^J_i Z^{iK} - \Lambda^{10} \right)
\tag{3.13}
$$

Notice that the first two terms coming from W_{tree} break explicitly the 'fake' $\text{SU(2)} \times \text{SU(3)}$ global symmetry to its diagonal subgroup.

The F-terms can be written as $F_s = \frac{\partial W_{\text{eff}}}{\partial \Phi_s}$ with $\Phi_s = (B, X, Y, \tilde{Y}, Z, U, V)$. Explicitly

$$
\begin{aligned}
F_{B_i} &= X^i \\
F_{X^I} &= B_i \delta_I^i + U\, \epsilon_{IJK}\, Z_i^J\, Z^{iK} \\
F_{Y_I^i} &= \delta_i^I + V\, Z_i^I + U\, \tilde{Y}_{ij}\, Z^{jI} \\
F_{\tilde{Y}^{ij}} &= U\, Y_{I(i}\, Z_{j)}^I \\
F_{Z_i^I} &= V\, Y_I^i + U\, (Y_{Ij}\, \tilde{Y}^{ij} + 2\, \epsilon_{IJK} X^K\, Z^{iJ}) \\
F_U &= Y_{Ii}\, \tilde{Y}^{ij} Z_j^I + \epsilon_{IJK}\, X^I\, Z_i^J\, Z^{iK} - \Lambda^{10} \\
F_V &= Y_I^i\, Z_i^I
\end{aligned}
\tag{3.14}
$$

We will momentarily discuss the solution of the above F-term equations and then take into account the effect of FI terms.

3.3 The supersymmetric vacuum

There is a one-parameter solution to the F-flatness conditions $F_s = 0$ with F_s given in (3.14). This is given by setting all composite fields to zero except for

$$
Z^{21} = -Z^{12} \qquad\qquad X^3 = \frac{\Lambda^{10}}{2(Z^{12})^2} \qquad\qquad V = \frac{1}{Z^{12}}
\tag{3.15}
$$

Supersymmetric vacua are parametrized by the single modulus Z^{12}. The superpotential W_{eff} vanishes along the valley of minima.

3.4 Dynamical supersymmetry breaking

Now let us consider the effect of turning on FI terms. In string theory this can be achieved by giving non-trivial VEV's to the two scalar fields (blowing up modes) belonging to the closed-string twisted sectors and thus living at the $\mathbb{Z}_5$-singularity. They are partners of the two axions acting as Stückelberg fields for the anomalous $U(1)_1$ and $U(1)_5$ gauge bosons.[6] The $U(1) \times U(5)$ D-terms can be written as

$$
\begin{aligned}
D &= -|B_i|^2 + |C^{\alpha u}|^2 - |C^{3u}|^2 + \xi_1 \\
D_v^u &= \bar{A}_i^{uw}\, A_{wv}^i - \bar{C}_{iv} C^{iu} - \bar{C}_{3v} C^{3u} + \bar{E}^u E_v + \frac{1}{5}\xi_2 \delta^u{}_v
\end{aligned}
\tag{3.16}
$$

The F-terms following from W_{tree} can be written as

$$
\begin{aligned}
F_{B_i} &= C^{iu} E_u \\
F_{A_i} &= C^{iu} C^{3v} \\
F_{C^3} &= A_{iuv} C^{iu} \\
F_{C^i} &= B_i E_u + A_{iuv} C^{3v}
\end{aligned}
\tag{3.17}
$$

---

[6]The precise combinations of twisted scalars entering the definition of the FI terms $\xi_{1,2}$ is determined by disk amplitudes with insertion of a closed-string scalar vertex in the bulk and an open-string auxiliary D-term vertex on the boundary. A similar computation determines the axion mixing with the anomalous $U(1)$ gauge bosons [22, 30, 31]

For $\xi_2 = 0$, $\xi_1 > 0$, there is a U(5) preserving solution to both D and F-flatness conditions given by taking

$$\langle B_i \rangle = (m, 0) \qquad\qquad \xi_1 = m^2 \qquad\qquad (3.18)$$

and setting all remaining fields A^i, C^I, E to zero. This defines a perturbative supersymmetric vacuum of the theory. We notice that for $m \neq 0$ the superpotential interactions result into a mass term for C^{1u} and E^u via the coupling

$$W_{\text{tree}} = m C^{1u} E_u + \dots \qquad\qquad (3.19)$$

In the limit $m \to \infty$ these fields decouple from the spectrum and we find a U(5) gauge theory with two generations of $(\mathbf{10} + \bar{\mathbf{5}})$. The condition (2.11) is now satisfied with $\beta = 11$ and $\dim\mathfrak{M}_{k=1}^{\text{ferm}} = 18$ and therefore a superpotential is generated by gauge instantons. The form of the superpotential is completely fixed by dimensional analysis and $\text{SU}(5) \times \text{SU}(2)$ invariance. Indeed there is a unique $\text{SU}(5) \times \text{SU}(2)$ invariant of dimension eight: $Z^{i\alpha} Z_{i\alpha}$. The ADS superpotential can then be written as

$$W_{\text{ADS}} = \frac{\hat{\Lambda}^{11}}{Z^{i\alpha} Z_{i\alpha}} \qquad\qquad (3.20)$$

with $\alpha = 2, 3$ labelling the massless $C^{\alpha u}$ components and $\hat{\Lambda}^{11} = m\Lambda^{10}$ the effective scale of the massless theory. The effective superpotential can be written as the sum of W_{ADS} and W_{tree} with all massive fields set to zero.

$$W_{\text{eff}} = A_{uv}^1 C^{2u} C^{3v} + \frac{m\Lambda^{10}}{Z^{i\alpha} Z_{i\alpha}} \qquad\qquad (3.21)$$

Summarizing, for $m >> \Lambda$ the unoriented $\mathbb{Z}_5$ quiver GUT-like theory is effectively described in terms of an $\text{SU}(5)$ gauge theory with two generations of matter $(\mathbf{10} + \bar{\mathbf{5}})$ and the superpotential (3.21).

Similarly to SQCD, in which the $N_f = N_c - 1$ case is recovered from that with $N_f = N_c$ by decoupling a massive flavor, (3.21) can be derived from the effective superpotential (3.13) after solving the F-flatness conditions for the massive fields in favor of the massless ones. The gauge invariant massless degrees of freedom are encoded in the following composites

$$Y_1^i = A_{u_1 u_2}^i C^{\alpha u_1} C_\alpha^{u_2}$$
$$Z^{i\alpha} = \epsilon^{u_1 \dots u_5} A_{u_1 u_2}^i A_{u_3 u_4}^j A_{u_5 j, k} C^{\alpha v} \qquad\qquad (3.22)$$

and the singlets B_i. Altogether they provide us with the expected $8 = 32 - 24$ degrees of freedom. Solving the F-flatness conditions for the remaining composites one finds

$$X^1 = \frac{\Lambda^{10}}{Z^{i\alpha} Z_{i\alpha}} \qquad X^2 = \frac{Y_1^2}{m} \qquad U = -\frac{m}{Z^{i\alpha} Z_{i\alpha}} \qquad V = \frac{Z^{23}}{Z^{i\alpha} Z_{i\alpha}}$$
$$\tilde{Y}_{22} = -\frac{2 Z^{13}}{m} \qquad \tilde{Y}_{12} = \frac{Z^{23}}{m} \qquad B_i = (m, 0) \qquad\qquad (3.23)$$

with all remaining fields set to zero. Plugging (3.23) into (3.13) one finds

$$W_{\text{eff}} = Y_1^1 + \frac{m\Lambda^{10}}{Z^{i\alpha}Z_{i\alpha}} \tag{3.24}$$

in perfect agreement with (3.21).

It is easy to check that the F-flatness conditions for the massless fields cannot all be simultaneously satisfied and therefore *supersymmetry is dynamically broken* by instanton effects [3–5]. The U(5) gauge theory under consideration here with superpotential (3.21) is indeed one of the classic examples of dynamically supersymmetry breaking. Nicely string theory provides a precise realization of this theory where the superpotential is not an *ad hoc* choice but rather a result of string dynamics. The fact that supersymmetry is broken can be seen alternatively by considering the Konishi anomaly equation

$$\frac{1}{4}\bar{D}^2(\Phi^r e^{gV}\Phi_s^\dagger) = \Phi^r \frac{\partial W_{tree}}{\partial \Phi^s} + \frac{g^2}{32\pi^2}\delta^r{}_s tr_{R_s}(W^\alpha W_\alpha) \tag{3.25}$$

valid for any r and s. Using the fact that $\bar{D}$ annihilate supersymmetric states the VEV of the leftt hand side of this equation on a supersymmetric vacuum is zero. Taking the diagonal term in (3.25) (no sum over s) and the lowest component of the superfields, we get

$$\left\langle \phi^s \frac{\partial W_{tree}}{\partial \phi^s} \right\rangle + \frac{g^2}{32\pi^2} \text{Tr}_{R_s}\langle \lambda\lambda \rangle = 0 \tag{3.26}$$

The first term cancels for $\phi_s = A^2$ since this field is absent from the tree level super-potential (3.21). The presence of a gaugino condensate $\text{Tr}_{R_s}\langle \lambda\lambda \rangle \neq 0$ then implies that supersymmetry is broken.

3.5 The weak coupling description

The results presented in the previous subsection can be further supported by a weak coupling analysis. To this purpose we notice that the SU(5) β-function and the number of fermionic zero modes[7] (i.e. the dimension of the fermionic moduli space) for the $k = 1$ gauge instanton, are given by (2.12)

$$\dim\mathfrak{M}_{k=1}^{\text{ferm}} = 10_\lambda + 3_{\psi_C} + 1_{\psi_E} + 6_{\psi_A} = 20$$
$$\beta = 10 \tag{3.27}$$

where we have indicated in the first line the origin of the fermionic zero modes. $\dim\mathfrak{M}_{k=1}^{\text{ferm}}$ and β do not satisfy the relation (2.11) and therefore a superpotential of the ADS type is not generated by instantons in this case. Nonetheless an instanton dominated correlator is easily identified to be

$$\langle X^I \, Z^{iJ} \, Z^{jK} \rangle = \epsilon^{IJK}\,\epsilon^{ij}\,\Lambda^{10} \tag{3.28}$$

Indeed each scalar soaks up a gaugino and a matter fermion zero mode, $XZ^2 \sim EC^3A^6$ is precisely what is needed in order to soak up the fermionic zero modes in (2.12). Here

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[7]This number is $2\ell(\mathbf{R})$ for each fermion in the $\mathbf{R}$ representation.

fields	A_{uv}^i	C^{iu}	C^{3u}	E_u	X^3	X^i	Y_3^i	Y_j^i	$\tilde{Y}^{(ij)}$	Z^{i3}	Z^{ij}	U	V
SU(2)	**2**	**2**	**1**	**1**	**1**	**2**	**2**	**3, 1**	**3**	**2**	**3, 1**	**1**	**1**
$\mathrm{U(1)}_M$	1	-3	0	0	0	-3	-5	-2	2	3	0	0	2
$\mathrm{U(1)}_N$	0	0	1	-1	0	-1	0	1	-1	1	0	0	-1
SU(5)	**10**	**$\bar{5}$**	**$\bar{5}$**	**5**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**

Table 4. $\mathrm{SU(2)} \times \mathrm{U(1)}^2$ charges of elementary and composites fields. The last row displays the multiplicities of the given field

we use the gauge invariant combinations X^I, Z^{iJ}, defined in (3.10), only for the sake of book-keeping. The correlator (3.28) can be computed in the weak coupling regime, where the relevant degrees of freedom are the elementary fields A, C, E and X, Z are quadratic and quartic gauge invariant combinations of these variables.

Using the Konishi identity (3.26) with $\Phi^s = E$ and taking $B_i = (m, 0)$ one finds

$$\frac{g^2}{32\pi^2} \langle \mathrm{tr}\, \lambda\lambda\, Z^{i\alpha} Z_{i\alpha} \rangle = m\, \Lambda^{10} \tag{3.29}$$

in agreement with (3.23) and (3.24). Alternatively this result can be justified by noticing that an insertion of $\lambda\lambda$ soaks up two gaugino zero modes, Z^2 16 fermionic zero modes, and the Yukawa interaction $B\psi_C\psi_E$ the last 2 fermionic zero modes.

4 Anomaly matching conditions

In this section we test the strong coupling description of the $\mathbb{Z}_5$-quiver gauge theory by showing that the anomalies of the global symmetries computed in the 'microscopic' theory match those of the 'macroscopic' theory described in terms of baryons and mesons. Omitting the SU(5) singlet fields B_i fields, which contribute to the strong and weak coupling phases in an identical way, and neglecting the Yukawa couplings for the time being, the global symmetry of the theory turns out to be $\mathrm{SU(3)}_C \times \mathrm{SU(2)}_A \times \mathrm{U(1)}_A \times \mathrm{U(1)}_C \times \mathrm{U(1)}_E$ with the subscripts indicating the fields with unit charge which the corresponding U(1) acts on. We now consider a particular solution of (3.11) preserving the largest possible symmetry

$$X^3 = \Lambda^2 \qquad Z^{iI} = \frac{1}{\sqrt{2}}\Lambda^4 \epsilon^{iI} \qquad I = 1, 2 \tag{4.1}$$

with all remaining composite fields set to zero. The symmetry preserved by this solution is $G_F = \mathrm{SU(2)} \times \mathrm{U(1)}_M \times \mathrm{U(1)}_N$ where the two U(1)'s are defined in such a way that fields appearing in the solution (4.1) are uncharged . Table 4 summarizes the charges of elementary and composite fields under this symmetry. The SU(2) part is the diagonal subgroup of the $\mathrm{SU(3)}_C \times \mathrm{SU(2)}_A$ global symmetry as follows from the invariance of (4.1). One can easily check that mixed $\mathrm{U(1)}_{M,N}\mathrm{SU(5)}^2$ anomalies cancel i.e. $\mathrm{U(1)}_{M,N}$ are truly conserved currents of the theory. Indeed using

$$I_{\mathrm{U(1)}_a\mathrm{SU(5)}^2} = \sum_s \dim(\mathbf{R}_{\mathrm{SU(2)}}^s)\, Q_a(\Phi_s)\ell(\mathbf{R}_{\mathrm{SU(5)}}^s) \qquad a = M, N \tag{4.2}$$

where s runs over all elementary fields in table 4, one finds

$$I_{\mathrm{U}(1)_M\mathrm{SU}(5)^2} = 2\frac{3}{2} + 2\frac{1}{2}(-3) = 0$$

$$I_{\mathrm{U}(1)_N\mathrm{SU}(5)^2} = \frac{1}{2}(1-1) = 0 \tag{4.3}$$

Now let us consider cubic anomalies of the global currents. Since SU(2) admits no complex representation there is no cubic anomaly made only of SU(2) currents.[8] The remaining cubic anomalies can be written as

$$I_{\mathrm{U}(1)_M^n\mathrm{U}(1)_N^{3-n}} = \sum_s \dim(\mathbf{R}_{\mathrm{SU}(2)\times\mathrm{SU}(5)}^s)\, Q_M(\Phi_s)^n\, Q_N(\Phi_s)^{3-n}$$

$$I_{\mathrm{U}(1)_a\mathrm{SU}(2)^2} = \sum_s \dim(\mathbf{R}_{\mathrm{SU}(5)}^s)\, Q_a(\Phi_s)\ell(\mathbf{R}_{\mathrm{SU}(2)}^s) \qquad a = M, N \tag{4.4}$$

The spectrum of $\mathrm{U}(1)_N$ charges is non-chiral for elementary fields and therefore all anomalies involving $\mathrm{U}(1)_N$ cancel. This is in agreement with the cubic anomalies computed in the strong coupling regime where the theory is described in terms of the baryons/mesons $X, Y, \tilde{Y}, Z$

$$I_{\mathrm{U}(1)_N^3}^{XYZ} = (-1)^3(2+3+1) + (1)^3(4+2) = 0$$

$$I_{\mathrm{U}(1)_N^2\mathrm{U}(1)_M}^{XYZ} = 2(-3) + 4(-2) + 3(2) + 2(3) + (2) = 0$$

$$I_{\mathrm{U}(1)_N\mathrm{U}(1)_M^2}^{XYZ} = 2(-9) + 4(4) + 3(-4) + 2(9) + (-4) = 0$$

$$I_{\mathrm{U}(1)_N\mathrm{SU}(2)^2}^{XYZ} = \frac{1}{2}(-1+1) + 2(1-1) = 0 \tag{4.5}$$

On the other hand the non-trivial anomalies in the microscopic theory are

$$I_{\mathrm{U}(1)_M^3}^{ACE} = 20\,(1)^3 - 10(-3)^3 = -250$$

$$I_{\mathrm{U}(1)_M\mathrm{SU}(2)^2}^{ACE} = 10\frac{1}{2} - 15\frac{1}{2} = -\frac{5}{2} \tag{4.6}$$

matching again those in the macroscopic theory

$$I_{\mathrm{U}(1)_M^3}^{XYZ} = 2(-3)^3 + 2(-5)^3 + 4(-2)^3 + 3(2)^3 + 4(3)^3 + 2 = -250$$

$$I_{\mathrm{U}(1)_M\mathrm{SU}(2)^2}^{XYZ} = \frac{1}{2}(-3-5+3) + (2-2) = -\frac{5}{2} \tag{4.7}$$

The perfect match between the anomalies in the two regimes provides a robust consistency test of the proposed non-perturbative description of the $\mathbb{Z}_5$-quiver gauge theory.

5 D-instanton description

In this section we briefly describe the D-instanton derivation of the non-perturbative effects relevant to our analysis of the $\mathbb{Z}_5$ quiver theory.

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[8]Including B^i the number of doublets is even and no global anomaly appears.

origin	fields	$\mathrm{SU}(2)_\mathrm{F}$	$\mathrm{SU}(5)$	$\mathrm{U}(1)_1 \times \mathrm{U}(1)_5 \times \mathrm{U}(1)_\mathrm{inst}$
D3D3	A^i_{uv}	**2**	**10**	$(0, 2, 0)$
	B^i	**2**	**1**	$(-1, 0, 0)$
	C^{iu}	**2**	$\bar{\mathbf{5}}$	$(1, -1, 0)$
	C^{3u}	**1**	$\bar{\mathbf{5}}$	$(-1, -1, 0)$
	E_u	**1**	**5**	$(0, 1, 0)$
D(-1)D(-1)	$(a_{\alpha\dot\alpha}, \Theta^0_\alpha, \bar\Theta^{\dot\alpha}_0)$	**1**	**1**	$(0, 0, 0)$
	$(\bar\chi_i, \bar\Theta^{\dot\alpha}_i)$	**2**	**1**	$(0, 0, -2)$
	χ^i	**2**	**1**	$(0, 0, 2)$
D(-1)D3	$(\bar{w}^u_\alpha, \bar\nu^{0u})$	**1**	**5**	$(0, 1, -1)$
	$(w^{\dot\alpha}_u, \nu^0_u)$	**1**	$\mathbf{5}^*$	$(0, -1, 1)$
	ν^{iu}	**2**	**5**	$(0, 1, 1)$
	$\bar\nu^i$	**2**	**1**	$(1, 0, -1)$
	$\bar\nu^3$	**1**	**1**	$(-1, 0, -1)$
	ν^3	**1**	**1**	$(0, 0, 1)$
	$d\mathfrak{M}$	**1**	**1**	$(-1, -10, 0)$

Table 5. Chiral fields and supermoduli

Non-perturbative effects, generated by D-brane instantons, have been the subject of a dedicated effort of several groups in the past couple of years. The emerging picture suggests that there are two kinds of D-brane instantons. The first class, termed "gauge" instantons correspond to Euclidean D-branes wrapping the same cycle as a stack of "physical" branes present in the background [24, 32, 33]. The second class, termed "exotic" or "stringy" instantons, correspond to Euclidean D-branes wrapping a cycle not wrapped by any "physical" branes realizing the gauge theory [34, 36]. Obviously, in a quiver gauge theory, what looks as a 'gauge' instanton for a given gauge group may look as an 'exotic' instanton for a different gauge group factor.

In the unoriented $\mathbb{Z}_5$ quiver theory, the relevant D(-1)-instanton is an $k = 1$ instanton sitting on the U(5) node of the quiver. Being a gauge instanton, the supermoduli space will comprise both neutral D(-1)D(-1) and charged D(-1)D3 zero-modes.

The field content of the D3D(-1) gauge theory describing the dynamics of the D-instanton is summarized in table 5. States are organized according to their charges under the $\mathrm{U}(1)_1 \times \mathrm{U}(1)_5 \in \mathrm{U}(1) \times \mathrm{U}(5)$ gauge symmetries, the SU(2) flavor group and $\mathrm{U}(1)_\mathrm{inst}$ D(-1)-symmetry. The various indices run over the following domains $\alpha, \dot\alpha = 1, 2$ (Lorentz spinors), $i = 1, 2$ (flavor) , $u = 1, ..5$ (gauge fundamentals). This content follows from that of $k_2 = k_3 = 1$ fractional D(-1) and $N_2 = N_3 = 5$, $N_0 = N_1 = N_3 = 1$ fractional D3-brane system in flat space after keeping the components invariant under $\mathbb{Z}_5$ and Ω.

The superpotential is defined by the following integral over the instanton moduli space

$$\int d^4x\, d^2\Theta\, W(\Phi_s) = \Lambda^{10} \int d\mathfrak{M}\, e^{-\mathcal{S}_B - \mathcal{S}_F} \tag{5.1}$$

with

$$\mathcal{S}_F = \bar{\Theta}_{0\dot{\alpha}}(w_u^{\dot{\alpha}}\bar{\nu}^{0u} + \nu_u^0\bar{w}^{\dot{\alpha}u}) + \chi^i\bar{\nu}^3\bar{\nu}_i + \bar{\chi}_i\nu_u^0\nu^{u,i}$$
$$+\overline{A}_{iuv}\nu^{u,i}\bar{\nu}^{0v} + \overline{E}_u\nu^3\bar{\nu}^{0u} + \overline{C}_3\nu_u^0\bar{\nu}^3 + \overline{C}_i^u\nu_u^0\bar{\nu}^i \tag{5.2}$$
$$+B_i\nu^3\bar{\nu}^i + C_u^3\nu^{ui}\bar{\nu}_i + C_u^i\nu^{ui}\bar{\nu}_3 \tag{5.3}$$
$$\mathcal{S}_B = \overline{\chi}_i\chi^i\overline{w}_{\dot{\alpha}}^u w_u^{\dot{\alpha}} + \chi^i\overline{w}_{\dot{\alpha}}^u\overline{w}^{v\dot{\alpha}}\overline{A}_{iuv} + \overline{\chi}_i w_u^{\dot{\alpha}}w_{v\dot{\alpha}}A^{iuv}$$
$$+w_u^{\dot{\alpha}}\overline{w}_{\dot{\alpha}}^v\left[\overline{A}_{ivw}A^{iwu} + \overline{E}_v E^u + C_v^3\overline{C}_3^u + C_v^i\overline{C}_i^u\right] \tag{5.4}$$

The fermionic zero-modes, except for Θ^0, are lifted by Yukawa-type interactions in (5.3). Roughly the fermionic integrals bring down 9 scalar fields in the numerator such that the $18 = \dim\mathfrak{M}_{U(5),k=1}^{\mathrm{ferm}} - 2$ fermionic zero modes besides $\Theta^{0\alpha}$ are soaked up. Bosonic gaussian integrals should bring scalar fields in the denominator in such a way that the resulting expression for $W(\Phi_f)$ is holomorphic as expected. Although the structure of the integrand supports our previous derivation of the non-perturbative superpotential, the explicit evaluation of the integral is rather involved and goes beyond the scope of this work.

6 Summary of results

In this paper we have studied the non-pertubative dynamics of an unoriented $\mathbb{Z}_5$ quiver gauge theory with GUT-like gauge group $U(5)\times U(1)$, chiral matter in the $(2\times\mathbf{10}+3\times\overline{\mathbf{5}}+\mathbf{5}+2\times\mathbf{1})$ representations and a cubic superpotential. At strong coupling, the dynamics of the gauge theory is described by an effective superpotential given in terms of a set of gauge invariant variables, the baryons and mesons, satisfying a quantum deformed constraint. The gauge theory has two distinct phases depending on whether field dependent FI terms are turned on or not. Although a local analysis such as ours does not allow to determine the VEV's of the closed-string moduli that govern the FI terms, one can assume that bulk dynamics, possibly associated to branes and fluxes needed for the global consistency of our model, can fix the FI terms to the desired values. At any rate, in the absence of FI terms the effective superpotential admits a line of supersymmetric vacua. Turning on a FI term for the $U(1)_1\not\subset U(5)$, the theory enters a Higgs phase,[9] whereby a $U(1)_1$ charged $U(5)$ singlet gets a vacuum expectation value and a $(\mathbf{5}+\overline{\mathbf{5}})$ pair gets mass from Yukawa interactions. The resulting $U(5)$ gauge theory with two generations of $(\mathbf{10}+\overline{\mathbf{5}})$ is one of the classical examples of chiral gauge theories with dynamical supersymmetry breaking via gaugino condensation [3–5]. We remark that supersymmetry is dynamically broken in this theory only for a bizarre choice of the tree level superpotential that breaks the $SU(2)$ flavor symmetry rotating the two generations. Remarkable, this is a characteristic feature of the superpotential resulting from the unoriented $\mathbb{Z}_5$ quiver in the Higgs phase. The proposed strong coupling dynamics is tested by matching the anomalies of the effective theory described in terms of baryon and meson composites with those computed in terms of elementary fields.

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[9]As already mentioned, we would like to stress the difference between 'genuine' Higgs fields (e.g.open-string charged singlets) and Stückelberg fields (e.g.closed-string axions).

Finally we have sketched how to derive the non-perturbative superpotential as a D-instanton effect. We stress that non-perturbative effects considered here origin only from gauge instantons i.e. fractional D(-1)-branes sitting at the U(5) node of the quiver. Indeed, in the quiver under consideration there is no room for exotic D-brane instantons i.e. those D(-1) branes sitting on top of nodes with no D3 branes. An alternative scenario of supersymmetry breaking based on exotic instanton interactions in presence of a $G_{(0,3)}$-flux was recently studied in [37, 38].

Acknowledgments

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