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SUSY breaking in monopole quivers

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ABSTRACT: We study 3d monopole quivers, theories with product gauge groups interacting through monopole superpotentials and we observe that they are a natural setup for the study of the spontaneous breaking of 3d $\mathcal{N} = 2$ supersymmetry. We give evidence of this statement by studying various examples of increasing complexity, considering quivers that are mirror dual to Wess-Zumino models. These Wess-Zumino models, in opportune regimes of parameters, break supersymmetry in perturbatively controllable (meta)stable vacua.

KEYWORDS: Supersymmetry Breaking, Solitons Monopoles and Instantons, Supersymmetry and Duality

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1 Introduction

In the recent past the three dimensional version of 2d bosonization attracted a large interest (see for example [1–5]). This is just the tip of a larger web of 3d dualities, sharing many similarities with the supersymmetric case. The possibility of deriving non-supersymmetric dualities from supersymmetry has been explored in [6–10], by explicit but controllable SUSY breaking.

An intriguing, but so far unexplored possibility, consists of deriving 3d non-supersymmetric dualities from spontaneous or dynamical supersymmetry breaking. Examples of $\mathcal{N} = 2 \rightarrow \mathcal{N} = 0$ breaking in 3d gauge theories have been discussed in [11–14], mostly

using arguments from the brane dynamics, even if a complete field theoretical analysis is missing.

Another interesting case was discussed in [15], generalizing the ISS mechanism [16] to 3d $\mathcal{N} = 2$ $U(N_c)_k$ SQCD, where k is the Chern-Simons (CS) level. In the IR the theory reduces to a Wess-Zumino (WZ) model, that, in a proper regime of couplings, gives perturbatively accessible, long lived and non-supersymmetric metastable vacua. Furthermore 3d supersymmetry breaking for WZ models has been explored in the literature in [17]. It has been observed that supersymmetry breaking can occur for large classes of models, with both classically marginal and relevant deformations, and that these vacua are often (long lived) metastable ones.

In this paper we observe that many of the models discussed in [15, 17] can be derived from new types of quiver gauge theories, defined in [18] as monopole quivers. Differently from common quivers, here the gauge nodes are connected through superpotential interactions involving monopole operators.

Here we will consider monopole quivers with $U(1)_k$ gauge factors with charged matter fields and deform some of the nodes by linear or quadratic monopole superpotentials, of the type discussed in [19, 20]. By studying a large variety of examples we show that the $\mathcal{N} = 2 \rightarrow \mathcal{N} = 0$ breaking in monopole quivers is quite generic. The analysis is possible after performing mirror symmetry at the various nodes, dualizing the monopole quivers to WZ models. Many of these WZ models reduce in the IR to the ones discussed in [15, 17], that have been shown to break supersymmetry in either stable or long lived metastable non-supersymmetric vacua.

Before starting our analysis a crucial comment is in order: in the whole paper we will assume that the WZ models, obtained after the dualization of the monopole quivers, have a canonical Kähler potential. This is a very strong assumption, because often these models correspond to strongly coupled interacting quantum field theories. A detailed analysis of the Kähler potential is beyond the scopes of the paper, but it will be compulsory in order to establish the stability of the supersymmetry breaking states associated to the monopole quivers.

The paper is organized as follows. In section 2 we collect some review material, useful for our analysis. In section 3 we discuss various examples of monopole quivers leading to (meta)stable supersymmetry breaking. In section 4 we discuss open questions and propose possible generalizations of our analysis. In appendix A we discuss the one loop effective potential in presence of spontaneous supersymmetry breaking and the bounce action for a 3d triangular barrier, necessary for the estimation of the lifetime of the supersymmetry breaking vacuum. In appendix B and in appendix C we give a detailed analysis of supersymmetry breaking for two of the WZ models obtained in section 3 that, to our knowledge, have not been analyzed yet in the literature.

2 Review

We will not review here generic aspects of 3d $\mathcal{N} = 2$ gauge theories. We refer the reader to [21] for definitions and general discussions. Here we restrict our attention to super-

symmetry breaking in 3d WZ models, to 3d dualities and to the definition of monopole quivers.

2.1 Supersymmetry breaking in 3d WZ models

In this section we overview some of the results on supersymmetry breaking in 3d WZ models obtained in [15, 17]. We refer the reader to the original references for more complete discussions. We collect in the appendix A the various tools for the analysis of the 1-loop effective potential and for the estimation of the lifetime of the metastable state through the evaluation of the bounce action for a triangular barrier.

In our analysis we consider WZ models of the form

$$W_{\text{WZ}} = fX + \sum_{i=0}^2 \sum_{a,b=1}^n X^i M_{ab}^{(i)} \phi_a \phi_b \quad (2.1)$$

where there are $n+1$ chiral superfields X and ϕ_i and the couplings are encoded in the three matrices $M^{(i)}$. A classification scheme for supersymmetry breaking can be constructed by separating the cases with either $M^{(1)} = 0$ (marginal couplings) or $M^{(2)} = 0$ (relevant couplings).

The first cases has been fully classified in [17] and we will report here the main results. The second case can be further divided in sub-classes, depending on the possible R -charge assignments. We will discuss this classifications distinguishing the various possibilities. There is a third case, where neither $M^{(1)} = 0$ nor $M^{(2)} = 0$. We will comment about this case as well.

2.1.1 Marginal couplings

This family has been analyzed in [17]. The main results of the analysis are

- The perturbative analysis is reliable if the marginal coupling constants encoded in $M^{(2)}$ are small numbers.
- The origin is a local maximum. This implies that the non-supersymmetric vacuum spontaneously breaks $U(1)_R$.
- The scalar potential has a classical runaway behavior.

A case study. The simplest example of this family is given by the superpotential

$$W = fX + hX^2\phi_1^2 + \mu\phi_1\phi_2 \quad (2.2)$$

where the mass dimensions of the couplings are $[h] = 0$, $[\mu] = 1$ and $[f] = \frac{3}{2}$. Next we summarize the detailed analysis of [17] for this case, because it will play a crucial role here. First it is necessary to compute the F-terms

$$F_X = f + 2hX\phi_1^2 = 0, \quad F_{\phi_1} = 2hX^2\phi_1 + \mu\phi_2 = 0, \quad F_{\phi_2} = \mu\phi_1 = 0 \quad (2.3)$$

and observe that they cannot be solved simultaneously, signaling that supersymmetry is broken at tree level.

In absence of a global supersymmetric vacuum we may look for tree level non-supersymmetric field configurations that minimize the tree level scalar potential¹

$$V_{\text{tree}} = |F_X|^2 + |F_{\phi_1}|^2 + |F_{\phi_2}|^2 \quad (2.4)$$

We obtain that supersymmetry is broken at tree level for $\langle \phi_1 \rangle = \langle \phi_2 \rangle = 0$, while there is a classical flat direction associated to the vev $\langle X \rangle$. This supersymmetry breaking locus has to be stable at tree level, which means that the classical real bosonic components of the superfields ϕ_i must be tachyon free around $\langle \phi_i \rangle = 0$. This boils down to compute the four squared masses of the real bosonic components of the superfields ϕ_i :

$$m_B^2 = \mu^2 + 2hX(hX^3 + \eta_1 f + \eta_2 \sqrt{f^2 + 2\eta f h X^3 + h^2 X^6 + X^2 \mu^2}) \quad (2.5)$$

with both $\eta_i = \pm 1$. The tachyon free condition for the candidate supersymmetry breaking vacuum at tree level is $m_B^2 > 0$ and it corresponds to

$$4fX < \mu^2 \quad (2.6)$$

The fermionic masses are obtained by setting $f = 0$ in m_B . At classical level the mass of the scalar component of the superfield X is vanishing. This scalar corresponds to a pseudo-modulus, i.e. an accidental flat direction of the tree level scalar potential that can be lifted at quantum level. The two real fermionic components of the superfield X correspond to the two goldstinos of the $\mathcal{N} = 2 \rightarrow \mathcal{N} = 0$ supersymmetry breaking.

In order to have a trustable supersymmetry breaking model one then needs to study the quantum corrections around the this tree level supersymmetry breaking locus. The knowledge of the massive tree level spectrum of the supersymmetry breaking locus can be used to study the quantum correction, through the Coleman-Weinberg (CW) potential [22]. This is the effective potential that can lift the flat direction X . In 3d the one loop contribution to the CW effective potential is given by (A.1). By performing the calculation in our case we observe that the origin is tachyonic, $m_{X=0}^2 = -(fh)^2/m$, while there is a minimum at $X \simeq 2^{1/4} \sqrt{\mu/h}$, if $f^2 \ll \mu^3$. Furthermore condition (2.6) imposes

$$h \ll \frac{1}{16\sqrt{2}} \frac{\mu^3}{f^2} \quad (2.7)$$

This signals the fact that the perturbative expansion is reliable because we are free to chose the dimensionless coupling h to be small enough to suppress higher loops.

The stability of this non-supersymmetric vacuum has to be still checked by studying the behavior of the tree level scalar potential at large vev. By parameterizing the fields as

$$X = -\frac{f}{2h\mu} e^{2\alpha}, \quad \phi_1 = \sqrt{\mu} e^{-\alpha}, \quad \phi_2 = -\frac{f^2}{2h\sqrt{\mu}^5} e^{3\alpha} \quad (2.8)$$

we have that $V_{\text{tree}} \rightarrow 0$ in the limit $\alpha \rightarrow \infty$. The presence of this runaway signals that the non-supersymmetric state is only metastable. The lifetime of this metastable state is

¹As already discussed in the introduction we assume the Kähler potential to be canonical.

studied by estimating the bounce action S_B for a triangular barrier. The decay rate of the state is given by e^{-S_B} . We have (see appendix A for details)

$$S_B \propto \sqrt{\frac{(\Delta\Phi)^6}{\Delta V}} \propto \left(\frac{\mu^3}{f^2}\right)^2 \sqrt{\frac{1}{h}} \gg 1 \quad (2.9)$$

where, because of the presence of the runaway, $\Delta\Phi$ is estimated starting at $X_{\min}$ and ending at a value $X = X^* \neq X_{\min}$ such that $V(X^*) = V_{\min}$. The denominator in (2.9) is $\Delta V = V_{\max} - V_{\min}$.

In section 3 we will find many quiver gauge theories that in the IR reduce to the WZ model with the superpotential given in formula (2.2). In all these cases a long lived supersymmetry breaking metastable minimum exists if the various parameters satisfy the constraints discussed here.

2.1.2 Relevant couplings

Models with relevant couplings give rise to a quite different analysis. Indeed, as observed in [15], there is a problem with the reliability of the perturbative expansion in these cases. This issue was solved in [15] by adding an explicit R-symmetry breaking deformation.² The case discussed in [15] corresponds to a WZ model with superpotential

$$W = \frac{1}{2}h\mu\epsilon X^2 - h\mu^2 X + h\mu\epsilon\phi_3\phi_4 + hX\phi_1\phi_2 + h\mu(\phi_1\phi_3 + \phi_2\phi_4) \quad (2.10)$$

where the mass dimensions of the couplings are $[h] = [\mu] = \frac{1}{2}$ and $[\epsilon] = 0$. This theory has a non-supersymmetric vacuum at $\phi_i = 0$ and $X \simeq \epsilon\mu b$, with $b \equiv \hat{b}_h^\mu \equiv \frac{4\pi\mu}{(3-2\sqrt{2})h}$, dimensionless. This vacuum is close to the origin if $\epsilon b \ll 1$. Higher loops are negligible if the relevant coupling h remains small at the mass scale set by X , $h \ll X$. This boils down to the inequality $h \ll \epsilon b\mu$. All in all there is a trustable non-supersymmetric vacuum if

$$\left(\frac{h}{\mu}\right)^2 \ll \epsilon \hat{b} \ll \frac{h}{\mu} \quad (2.11)$$

2.1.3 Marginal and relevant couplings

One can consider more generic situations, where both $M^{(1)} \neq 0$ and $M^{(2)} \neq 0$. This possibility has been discussed in [17], by considering the superpotential

$$W = fX + hX\phi_1^2 + m\phi_1\phi_2 + \lambda X^2\phi_3^2 + \mu\phi_3\phi_4 \quad (2.12)$$

that combines (2.2) and the R-symmetric part of (2.10). The analysis shows that there are regimes of masses and coupling such that this theory has an unstable origin, a long lived metastable R-symmetry and supersymmetry breaking minimum and a runaway in the large field region.

²Observe that, depending on the structure of the superpotential, there can be models with relevant couplings that lead to spontaneous R-symmetry breaking. In such cases the requirement of an explicit R-symmetry breaking deformation is not always necessary.

Summarizing. The general message that can be extracted from the analysis is that (meta)stable supersymmetry breaking in 3d WZ models is reliable at perturbative level if R-symmetry is broken, either explicitly or spontaneously.³ Spontaneous R-symmetry breaking is possible whenever there are marginal couplings involving the pseudo-modulus in the superpotential. On the other hand the explicit breaking is often necessary for models with relevant couplings. Actually there are situations with spontaneous R-symmetry breaking also for models with relevant couplings, depending on the R-charge assignments to the chiral superfields [17].

2.2 Monopole quivers

Here we introduce the notion of monopole quivers. They have been originally defined in [18], in the study of the dimensional reductions of 4d dualities in terms of the D-brane engineering.

Let's consider 4d $\mathcal{N} = 1$ $U(N_c)$ SQCD with N_f fundamental and N_f anti-fundamental chiral superfields, i.e. N_f fundamental flavors. When this theory is dimensionally reduced on a circle of radius r it gives rise to a 3d effective gauge theory with the same field content of the 4d parent, a gauge coupling g_3 and a monopole superpotential for the Kaluza-Klein (KK) monopole, $W_{\text{KK}} = \eta e^{\Sigma_{N_c} - \Sigma_1}$ [26], where $\eta \propto e^{-\frac{1}{rg_3^2}}$. At a generic point in the Coulomb branch the superpotential is described by a further contribution due to the BPS monopoles

$$W = W_{\text{BPS}} + W_{\text{KK}} = \sum_{i=1}^{N_c-1} e^{\Sigma_1 - \Sigma_{i+1}} + \eta e^{\Sigma_{N_c} - \Sigma_1} \quad (2.13)$$

One can consider more complex situations, where the gauge group is broken, by the real scalars in the vector multiplet, into a product of r $U(n_i)$ factors, with $n_1 + n_2 + \dots + n_r = N_c$. Moreover the real scalars in the background flavor symmetry can be used to set f_i massless fundamental flavors at each $U(n_i)$ factor, such that $f_1 + f_2 + \dots + f_r = N_f$. In this case a monopole superpotential, similar to (2.13) is generated

$$W = \sum_{i=1}^r \lambda_i e^{\Sigma_{n_i}^{(i)} - \Sigma_1^{(i+1)}} \quad (2.14)$$

with $n_{r+1} = n_1$. This construction is the prototypical example of the monopole quiver studied in [18]. The definition is generalized to include real gauge groups, but we will not consider this possibility here. The coupling constants λ_i have a simple interpretation in the brane setup. In this case the gauge groups live on D3 branes separated along the compact direction by D1 branes, signaling the monopole interactions. The couplings λ_i in (2.14) are related to the distance between the D3 branes in this setup.

A further generalization, relevant for our analysis, includes an asymmetric content of fundamentals/anti-fundamentals and CS terms. Moreover these monopole quivers can be defined also for pure 3d models in $\mathbb{R}^3$, without requiring the presence of any extra S^1 . In

³This analysis is reminiscent of the 4d results of [23–25]. It should be interesting to further study the relation between the 3d and the 4d results for a full understanding of the role of the R-symmetry in 3d $\mathcal{N} = 2 \rightarrow \mathcal{N} = 0$ supersymmetry breaking.

Figure 1. Graphical representation of a monopole superpotential between two gauge nodes. This represents one of the building blocks of the monopole quivers.

general they give rise to linear rather than circular monopole quivers with unitary gauge groups.

In order to provide a graphical representation of the monopole quiver we refer to the interaction $e^{\Sigma_{n_i}^{(i)} - \Sigma_1^{(j)}} \equiv \tilde{T}_{n_i} T_{n_j}$ as an oriented sequence of triangles connecting the two gauge nodes $U(N_i)$ and $U(N_j)$. The orientation of the triangles is fixed as in figure 1.

2.3 A survey of 3d $\mathcal{N} = 2$ dualities

Here we survey the 3d $\mathcal{N} = 2$ dualities relevant for our discussion. We restrict ourself to $U(N_c)_k$ gauge factors with N_f fundamentals and N_a anti-fundamentals. We will be interested in two broad families of 3d $\mathcal{N} = 2$ dualities: mirror symmetry and Seiberg like dualities.

Mirror symmetry. Here we review mirror symmetry for a $U(1)$ gauge theory with a pair of charged fields, Q and $\tilde{Q}$ with opposite charge, $+1$ and -1 respectively. This model, usually referred to as 3d $\mathcal{N} = 2$ SQED can be equivalently described in terms of three chiral fields X , Y and Z interacting through the superpotential $W = XYZ$. The duality maps the field X to the gauge invariant combination $Q\tilde{Q}$ in SQED, while the other two singlets Y and Z correspond to the monopole T and the anti-monopole $\tilde{T}$ in SQED.

By assigning a large real mass to one of the fields in the electric theory a new duality emerges, relating $U(1)_{\frac{1}{2}}$ with a field Q to a singlet. Such a singlet corresponds to the monopole in the electric phase. This is the minimal version of mirror symmetry, originally discussed by [27, 28].

In the rest of the paper we will make a large use of these abelian mirror symmetries. The nature of the interactions in the monopole quivers allows to use abelian mirror symmetry separately at the various nodes.⁴ For example let us consider a pair of SQEDs connected by a monopole superpotentials⁵

$$W_{\text{mon}} = T_1 \tilde{T}_2 \quad (2.15)$$

⁴Observe that a similar procedure has been used in [26], by considering the mirror of a single abelian sector in a more complicated model. This was possible because the different gauge sectors interacted only through monopole superpotentials.

⁵From now on we suppress the couplings λ_i in (2.14) by absorbing them in the definition of the monopoles, but we turn the coupling on in the WZ dual models, treating the fields as free.

This is a simple example of monopole quiver with two gauge groups. In this case we can apply mirror symmetry separately on each node and obtain the following duality map

$$Q_1 \tilde{Q}_1 \leftrightarrow X_1, \quad Q_2 \tilde{Q}_2 \leftrightarrow X_2, \quad T_1 \leftrightarrow Y_1, \quad T_1 \leftrightarrow Y_2, \quad T_2 \leftrightarrow Z_1, \quad \tilde{T}_2 \leftrightarrow Z_2 \quad (2.16)$$

where Q_1 (Q_2) and $\tilde{Q}_1$ ($\tilde{Q}_2$) are the flavors of the first (second) abelian gauge group. The dual phase then corresponds to a pair of WZ models of the XYZ type discussed above, coupled by a massive interaction $Y_1 Z_2$. The final superpotential in this case is

$$W = h_1 X_1 Y_1 Z_1 + h_2 X_2 Y_2 Z_2 + m Y_1 Z_2 \quad (2.17)$$

Observe that we could have modified the superpotential (2.15) as

$$W_{\text{mon}} = T_1 \tilde{T}_2 + T_2 \tilde{T}_1 \quad (2.18)$$

This modification implies a circular shape for the monopole quiver, with the net effect to add a new reversed arrow to the one connecting the two gauge nodes in figure 1. In this case the superpotential of the dual WZ model is

$$W = h X_1 Y_1 Z_1 + h X_2 Y_2 Z_2 + m(Y_1 Z_2 + Y_1 Z_2) \quad (2.19)$$

Seiberg dualities and its generalizations. To complete this section we give a brief overview of three dimensional non-abelian dualities generalizing 4d Seiberg duality. Three dimensional non-abelian dualities assume different forms, depending on the presence of CS terms in the action and of possible superpotentials involving monopole operators. Here we list the possible dualities by specifying the superpotential, the matter content and the possible presence of CS terms in the electric theory. Moreover we refer to the papers in which they have been first derived.

W_{ele}	k	$N_a - N_f$	Reference
$W = 0$	$k = 0$	$= 0$	[29]
$W = 0$	$k \neq 0$	$= 0$	[30]
$W = 0$	$k \neq 0$	$\neq 0$	[31]
$W = T \tilde{T}$	$k = 0$	$= 0$	[26]
$W = T + \tilde{T}$	$k = 0$	$= 0$	[19]
$W = T^2 + \tilde{T}^2$	$k = 0$	$= 0$	[19]
$W = T$	$k \neq 0$	$\neq 0$	[19]
$W = T$	$k = 0$	$= 0$	[19]
$W = T^2$	$k = 0$	$= 0$	[20]
$W = T^2$	$k \neq 0$	$\neq 0$	[20]

The monopole operators T and $\tilde{T}$ in the table refer here to the ones with flux $(\pm 1, 0, \dots, 0)$. The CS level can be integer or semi-integer depending on the parity of $N_a - N_f$. We do not provide further explanations of these dualities here and refer the reader to the original papers for definitions and details.

3 Supersymmetry breaking from monopole quivers

This is the main section of the paper, where we study various monopole quivers leading to spontaneous supersymmetry breaking in the IR. We organize the section by discussing examples of increasing complexity, increasing the number of gauge groups and of charged fields. We restrict to models with only abelian gauge factors and with either one flavor (with $k = 0$) or one fundamental (or anti-fundamental) (with $k = \frac{1}{2}$). Mirror symmetry can be performed at each node separately and this allows us to construct WZ models starting from monopole quivers. Massive or more general superpotential deformations of the original quivers can be added to the discussion without spoiling the duality. By considering complex masses for the favors and linear or quadratic monopole deformations we construct WZ models with the structure discussed in [17] that lead to (meta)stable supersymmetry breaking in the IR.

3.1 Model I: two abelian gauge groups

This is the simplest example that we consider. There are two $U(1)$ gauge groups. The $U(1)$ factor associated to the first node has a CS term at level $k = \frac{1}{2}$ and there is just one charged field at charge -1 . The second $U(1)$ factor has one field at charge 1 and one at charge -1 . There is a monopole superpotential, that connects the two gauge nodes as represented in figure 2, of the form

$$W_{\text{mon}} = T_1 \tilde{T}_2 \quad (3.1)$$

We deform the model by two superpotential terms, one consists of a quadratic term for the monopole, $W_{T_2} = T_2^2$ while the other one is a mass term, $W_{\text{mass}} = m Q \tilde{Q}$. The final form of the superpotential is

$$W = W_{\text{mon}} + W_{T_2} + W_{\text{mass}} = T_1 \tilde{T}_2 + T_2^2 + m Q \tilde{Q} \quad (3.2)$$

This model can be dualized by applying mirror symmetry separately on both gauge nodes. Here mirror symmetry maps the monopoles and the singlets of the electric theory to singlets in the dual WZ model. This mapping is

$$T_1 \leftrightarrow W, \quad T_2 \leftrightarrow Z, \quad \tilde{T}_2 \leftrightarrow Y, \quad Q \tilde{Q} \leftrightarrow X \quad (3.3)$$

The mirror theory has superpotential

$$W = fX + hXYZ + m_1 Y^2 + m_2 ZW \quad (3.4)$$

In the regime $m_1 \gg m_2 \gg f^{\frac{2}{3}}$ we integrate out the field Y and arrive to the superpotential

$$W = fX + \hat{h}X^2Z^2 + m_2ZW \quad (3.5)$$

where $h \propto \frac{h^2}{m}$ and $[\hat{h}] = 0$. This superpotential corresponds to the one in formula (2.2) and it breaks supersymmetry if the parameters are chosen in the regime discussed in subsection 2.1.

Figure 2. Summary of the monopole quivers studied in the paper.

3.2 Model II: adding a singlet

A second model is obtained by adding a massive singlet to the previous quiver and by considering the superpotential

$$W = T_1 \tilde{T}_2 + T_2 + \lambda Q M \tilde{Q} + m M^2 \quad (3.6)$$

The mirror dual is a WZ model with superpotential

$$W = h X Y Z + f Y + m_X X M + m M^2 + \mu Z W \quad (3.7)$$

with the duality map still given by (3.3). In the regime $m_X, m \gg \mu$ we can integrate out the massive field X and obtain the superpotential

$$W = \hat{h} Y^2 Z^2 + f Y + \mu Z W \quad (3.8)$$

that corresponds to (2.2) and breaks supersymmetry if the parameters are chosen as discussed in sub-section 2.1.

3.3 Model III: a circular monopole quiver

Here we consider a monopole quiver with two $U(1)$ gauge groups, each one with a pair of charged fields, with opposite charge ± 1 . The associated quiver is given in figure 2. The monopole interactions is

$$W_{\text{mon}} = T_1 \tilde{T}_2 + T_2 \tilde{T}_1 \quad (3.9)$$

and we deform the model by the superpotential interactions

$$W_{\text{def}} = m Q \tilde{Q} + h (P \tilde{P})^2 \quad (3.10)$$

We can apply mirror symmetry on both nodes and obtain a WZ model. After defining the mapping

$$Q\tilde{Q} \leftrightarrow X_1, \quad T_1 \leftrightarrow Y_1, \quad \tilde{T}_1 \leftrightarrow Z_1, \quad P\tilde{P} \leftrightarrow X_2, \quad T_2 \leftrightarrow Y_1, \quad \tilde{T}_2 \leftrightarrow Y_1 \quad (3.11)$$

the superpotential of the dual WX model is

$$W = fX_1 + h_1X_1Y_1Z_1 + m(Y_1Z_2 + Y_2Z_1) + h_2X_2Y_2Z_2 + \mu X_2^2 \quad (3.12)$$

Following the general discussion of [15, 17] here a perturbative regime is possible if we add a term $h(Q\tilde{Q})^2$ to the electric superpotential. This boils down to a term ϵX^2 in the dual phase, that breaks the R-symmetry and allows the existence of a perturbative regime.

In appendix B we study the behavior of the non-supersymmetric state. We find that this model has a perturbatively accessible metastable non-supersymmetric vacuum, with a parametrically long lifetime, if the parameters are chosen properly.

3.4 Model IV: a circular quiver with an extra singlet

Here we add a further singlet M to the model discussed above and consider the superpotential

$$W_{\text{def}} = mQ\tilde{Q} + h(Q\tilde{Q})^2 + MP\tilde{P} + m_M^2 M \quad (3.13)$$

This superpotential has to be added to the monopole superpotential (3.9). Mirror symmetry on both nodes gives

$$W = fX_1 + \epsilon X_1^2 + X_1Y_1Z_1 + m(Y_1Z_2 + Y_2Z_1) + X_2Y_2Z_2 + \mu X_2M + m_M^2 M \quad (3.14)$$

where the duality map is again given by (3.11). We consider the regime $m_M, \mu \gg m$ such that, after integrating out the massive fields M and X_2 , we arrive at

$$W = fX_1 + \epsilon X_1^2 + X_1Y_1Z_1 + m(Y_1Z_2 + Y_2Z_1) + \frac{m_M^2}{\mu} Y_2Z_2 \quad (3.15)$$

This superpotential corresponds to the one in formula (2.10) if the parameters are chosen such that $\epsilon = \frac{m_M^2}{\mu}$, and it breaks supersymmetry in the regime (2.11).

3.5 Model V: three abelian factors

Here we discuss a monopole quiver with three gauge groups $U(1)_{1/2} \times U(1)_0 \times U(1)_{1/2}$. The $U(1)_{1/2}$ nodes have $N_f^{(1)} = 1$ and $N_a^{(1)} = 0$ each while the node $U(1)_0$ has $N_f^{(2)} = 1$ and $N_a^{(2)} = 1$. The quiver is shown in figure 2. The monopole superpotential is

$$W_{\text{mon}} = T_1\tilde{T}_2 + T_2\tilde{T}_3 \quad (3.16)$$

and we deform the model by the interaction

$$W_{\text{def}} = mQ\tilde{Q} \quad (3.17)$$

We use mirror symmetry on each node obtaining a WZ model with superpotential

$$W = m_T(T_1\tilde{T}_2 + T_2\tilde{T}_3) + mX + hXT_2\tilde{T}_2 \quad (3.18)$$

We could also have added a monopole deformation T_1^2 that would have led to the superpotential

$$W = m_T T_2\tilde{T}_3 + mX + \lambda X^2 T_2^2 \quad (3.19)$$

corresponding to the one in formula (2.2).

3.6 Model VI: four abelian factors

Here we consider a monopole quiver with gauge group $U(1)_{\frac{1}{2}} \times U(1) \times U(1) \times U(1)_{\frac{1}{2}}$ as in figure 2. The monopole superpotential is

$$W_{\text{mon}} = T_+^{(1)}T_+^{(2)} + T_-^{(2)}T_+^{(3)} + T_-^{(3)}T_-^{(4)} \quad (3.20)$$

and we add the deformation

$$W_{\text{def}} = \mu Q\tilde{Q} + h(P\tilde{P})^2 \quad (3.21)$$

Applying mirror symmetry on each node we arrive at a WZ model with superpotential

$$W = m(\phi_1\phi_3 + \phi_2\phi_4 + \phi_5\phi_6) + h_X X\phi_1\phi_2 + fX + h_Y Y\phi_4\phi_5 + m_Y Y^2 \quad (3.22)$$

In this case a perturbative regime for the tree level non-supersymmetric state can be realized by adding a term $h(Q\tilde{Q})^2$ to the electric superpotential. This boils down to a term ϵX^2 in the dual phase, that breaks the R-symmetry and allows a perturbative study. In appendix C we show that the WZ model with superpotential (3.22) has a perturbatively accessible metastable non-supersymmetric vacuum, with a parametrically large lifetime, in the opportune regime of parameters.

4 Conclusions

In this paper we have proposed a mechanism of supersymmetry breaking for 3d $\mathcal{N} = 2$ monopole quivers. We restricted our attention to quivers with only abelian nodes because in such cases the analysis is simplified by the use of mirror symmetry. The presence of monopole interactions between the various gauge groups justifies the action of abelian mirror symmetry separately at each node. This is an example of local mirror symmetry, similar to the one proposed in [32–34]. The low energy descriptions consist of WZ models.

We studied these models in an opportune regime of couplings, such that the SUSY broken phase is perturbatively accessible. In this analysis we assumed a canonical Kähler potential for the WZ model. In general deviations from the canonicity are expected and we believe that the brane picture can be a helpful tool for further studies in this direction. For example a possible strategy to attack the problem stays in the discussion of [35] in terms of the brane realization of the model. Such an analysis can be relevant because all the models and the dualities studied here can be engineered in terms of D-branes.

Moreover here we did not find any dynamical mechanism that fixes the hierarchy of couplings that leads to the existence of stable and calculable non-supersymmetric states. This is a crucial issue that requires the study of the UV completions of the models treated here. We hope to come back to these problem in the future.

It should be also interesting to generalize our analysis to non-abelian gauge groups. In some cases supersymmetry breaking is expected because the Seiberg-like dual phases, in presence of monopole superpotential deformations, are of the type discussed in subsection 2.3, and there is a negative rank for the dual gauge group. This is a signal of a supersymmetry breaking state in the spectrum. It is not possible to use abelian mirror symmetry in such cases, but it may be interesting to study these models along the analysis of [36].

Another interesting case that may deserve a deeper analysis corresponds to the circle compactification of SQCD. In presence of a light mass deformations in the 4d IR free window the ISS [16] mechanism is at work. It should be interesting to study its generalization in 3d in presence of a KK monopole superpotential. Furthermore it should be possible to find new SUSY breaking models by increasing the number of gauge nodes, both for cases of circular quivers and for cases with linear ones. We leave such an analysis to the future.

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A Tools in 3d WZ

A.1 One loop effective potential

In the body of the paper we found a large set of WZ models obtained as dual descriptions of abelian monopole quivers. At tree level we observed that supersymmetry breaking takes place in these WZ models because the F-terms cannot be solved simultaneously. The configuration with minimal energy breaks supersymmetry at tree level, if the scalars are considered at finite vevs. Often this is not just a vacuum state, but there are scalar flat directions. Such moduli are often just pseudo-moduli and they can acquire corrections at quantum level. If the model can be treated as perturbative around the tree level supersymmetry breaking vacuum, the quantum effects can be computed by studying the one or higher loop corrections. These corrections can either lift the moduli, isolating a single vacuum state or provide a negative squared mass, destabilizing the vacuum, along the run-away or along possible supersymmetric vacua, due to the strong dynamics that can arise at large vev. The one loop effective potential is given by the CW formula

$$V_{\text{CW}} = -\frac{1}{12\pi} \text{STr} |\mathcal{M}|^3 \equiv -\frac{1}{12\pi} \text{Tr} (|\mathcal{M}_B|^2 - |\mathcal{M}_F|^2) \quad (\text{A.1})$$

where $\mathcal{M}_B$ and $\mathcal{M}_F$ are the tree level bosonic and fermionic mass matrices respectively.

A.2 Bounce action

The lifetime of a metastable non-supersymmetric state in 3d can be estimated through the calculation of the bounce action S_B for a triangular barrier. This is a sensible estimation if the two vacua are very far in field space (or in presence of a runaway potential, as discussed in [37]). The result has been obtained in [17], along the lines of [38], computing S_B as the difference between the tunneling configuration and the metastable vacuum in the euclidean action. The analysis has been performed in terms of a single field ϕ decaying from the false vacuum ϕ_F to the true vacuum ϕ_T . The result is expressed in terms of a dimensionless parameter $c \equiv -\lambda_T/\lambda_F$, where

$$\lambda_F = \frac{V_{\max} - V_F}{\phi_{\max} - \phi_F} \equiv \frac{\Delta V_F}{\Delta \phi_F}, \quad \lambda_T = \frac{V_{\max} - V_T}{\phi_{\max} - \phi_T} \equiv -\frac{\Delta V_T}{\Delta \phi_T} \quad (\text{A.2})$$

and F and T stay for true and false. By using these definitions the bounce action is

$$S_B = \frac{16\sqrt{6}\pi}{5} \frac{1+c}{(3+2c-3(1+c)^{2/3})^{3/2}} \sqrt{\frac{\Delta \phi_F^6}{\Delta V_F}} \quad (\text{A.3})$$

B Susy breaking for model III

In this section we study the stability of the non-supersymmetric state discussed in section 3.3. For simplicity we re-organize the superpotential as

$$W = hX\phi_1\phi_2 - fX + hY\phi_3\phi_4 + h\mu(\phi_1\phi_3 + \phi_2\phi_4) + \frac{1}{2}h\mu(\epsilon_X X^2 + \epsilon_Y Y^2) \quad (\text{B.1})$$

where the parameters h and μ have mass dimensions $\frac{1}{2}$, the supersymmetry breaking scale has mass dimensions $\frac{3}{2}$ and the parameters ϵ_X and ϵ_Y are dimensionless.

At small X the theory has a non-supersymmetric state in the regime $\epsilon_X \ll 1$. If we compute the bosonic masses for the fields Y and ϕ_i in this region we have

$$m_Y^2 = h^2 m^2 \epsilon_Y^2, \quad m_{\phi_i}^2 = \left\{ \frac{h}{2} \left(a_{\pm} + \sqrt{a_{\pm}^2 - b_{\pm}} \right), \quad \frac{h}{2} \left(a_{\pm} - \sqrt{a_{\pm}^2 - b_{\pm}} \right) \right\} \quad (\text{B.2})$$

with

$$a_{\pm} = 2h\mu^2 + hX^2 \pm (f - h\mu X\epsilon_X), \quad b_{\pm} = 4h\mu^2(h\mu^2 \pm (f - h\mu X\epsilon_X)) \quad (\text{B.3})$$

The tree level scalar potential for the field X together to the one loop correction from the CW potential is

$$V(X) \simeq |8\pi\mu^2\epsilon_h|^2 |\epsilon_X X - \mu\epsilon_f|^2 - |8\pi^2\mu^4\epsilon_f\epsilon_h^3(\epsilon_f(X^2(3\epsilon_X^2 - 2) + 8\mu^2) - 16\mu X\epsilon_X)| \quad (\text{B.4})$$

where we expressed it in terms of the dimensionless parameters $\epsilon_X \ll 1$, $\epsilon_f \equiv \frac{f}{h\mu^2} \ll 1$ and $\epsilon_h \equiv \frac{h}{8\pi\mu}$. The non-supersymmetric vacuum is placed at

$$\langle X \rangle = \left| \frac{\mu(\epsilon_h - 1)}{\frac{\epsilon_f}{\epsilon_X} \frac{\epsilon_h}{4} + \frac{\epsilon_X}{\epsilon_f}} \right| \quad (\text{B.5})$$

In order to study the supersymmetry breaking scenario in such a vacuum we need to impose the following constraints on the parameters

- The vacuum has to be close to the origin. This is necessary for having a parametrically large lifetime. This corresponds to the requirement

$$|\langle X \rangle| \ll |\mu| \quad (\text{B.6})$$

On the other hand the supersymmetric vacuum, placed at $|\langle X \rangle| = \frac{|\epsilon_f|}{|\epsilon_X|} |\mu|$, must be far in the field space

- The perturbative approximation has to be valid, i.e. higher loops must be suppressed. This corresponds to the requirement

$$|h| \ll |\langle X \rangle| \quad (\text{B.7})$$

Observe that this fixes $\epsilon_h \ll 1$ as well.⁶

- There should be no tachyons at the non-supersymmetric vacuum. This corresponds to the requirement

$$|\langle X \rangle| \ll \left| \frac{\mu(\epsilon_f \pm 1)}{\epsilon_X} \right| \quad (\text{B.8})$$

The various requirements are satisfied in the regime

$$2\pi\epsilon_h^2 \ll \frac{\epsilon_X}{\epsilon_f} \ll \frac{1}{4}\epsilon_h \ll 1 \quad (\text{B.9})$$

The squared mass at the vacuum is positive, $m_X^2 = 32\pi^2\mu^4\epsilon_h^3\epsilon_f^2$, and the bounce action is

$$S_B \propto \frac{\epsilon_f^2}{16\pi\epsilon_X^3\epsilon_h} \gg \frac{1}{16\pi\epsilon_X\epsilon_h} \gg 1 \quad (\text{B.10})$$

One can then conclude that in this regime of parameters there is a non-supersymmetric metastable vacuum with a parametrically large lifetime.

C Susy breaking for model VI

The other WZ model that has not been studied in the literature yet and that we have obtained in the body of the paper is model VI in sub-section 3.6. Here we re-organize the superpotential as

$$W = hX\phi_1\phi_2 - fX + h\mu(\phi_1\phi_3 + \phi_2\phi_4 + \psi_5\psi_6) + hY\phi_4\psi_5 + \frac{1}{2}h\mu(X^2\epsilon_X + Y^2\epsilon_Y) \quad (\text{C.1})$$

where again the parameters h and μ have mass dimensions $\frac{1}{2}$, the scale f has mass dimension $\frac{3}{2}$ and the parameters ϵ_X and ϵ_Y are dimensionless.

The non-supersymmetric state is still at small X if $\epsilon_X \ll 1$. The bosonic masses for the fields Y and ψ are supersymmetric and the non-supersymmetric contribution to the

⁶We further assume $\epsilon_h > 0$, *a posteriori* one can see that the opposite regime is not compatible with the various constraints.

CW potential for the field X is due to the masses of the fields ϕ_i . The masses for the bosonic components can be summarized as

$$m_{\psi_i}^2 = h^2 \mu^2, m_Y^2 = h^2 \mu^2 \epsilon_Y^2, m_{\phi_i} = \frac{h}{2} \left\{ \left(a_{\pm} + \sqrt{a_{\pm}^2 - b_{\pm}} \right), \left(a_{\pm} - \sqrt{a_{\pm}^2 - b_{\pm}} \right) \right\} \quad (\text{C.2})$$

with $a_{\pm}$ and $b_{\pm}$ given by (B.3). This coincidence and the equivalence of the tree level contribution to the scalar potential with the one of the model studied in appendix B are enough to claim that the rest of the analysis coincide with the one done there. We can again conclude that in the regime (B.9) supersymmetry is broken in a perturbatively accessible metastable vacuum with a parametrically large lifetime.

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