

# Notes on Quiver Gauge Theories

Louan Mol

*Université Libre de Bruxelles*

## Abstract

In these notes, we present some basic ideas around the large topic of quiver gauge theories. The goal is to reproduce and regroup the basics of quiver gauge theories. Note that this is a draft, it may contain a lot of typos, errors and imprecisions. It is only meant as a work support.

# Contents

<b>1</b>	<b>Physical setup</b>	<b>3</b>
<b>2</b>	<b>The simplest case: smooth transverse space</b>	<b>4</b>
2.1	Generalities . . . . .	4
2.2	Matter content . . . . .	4
<b>I</b>	<b>Orbifold singularities</b>	<b>5</b>
<b>3</b>	<b>Singular transverse space: orbifold singularities</b>	<b>5</b>
3.1	Generalities . . . . .	6
3.2	Representation theory realization . . . . .	6
3.3	Field content, quivers and McKay graphs . . . . .	8
3.4	A simple example: $S = \mathbb{C}^3/\mathbb{Z}_3$ . . . . .	8
3.4.1	Projection . . . . .	9
3.4.2	Quiver . . . . .	10
3.4.3	Gauge anomaly cancellation . . . . .	11
3.5	$(p+1)$ -dimensional quiver gauge theories . . . . .	11
<b>4</b>	<b>Orbifold constructions of <math>\mathcal{N} = 2</math> daughter theories</b>	<b>12</b>
4.1	$S = \mathbb{C} \times \mathbb{C}^2/\mathbb{Z}_n$ . . . . .	12
4.2	$S = \mathbb{C} \times \mathbb{C}^2/\mathcal{D}_n$ . . . . .	14
4.3	$S = \mathbb{C} \times \mathbb{C}^2/\mathcal{T}, \mathcal{O}, \mathcal{I}$ . . . . .	15
<b>5</b>	<b>Orbifold constructions of <math>\mathcal{N} = 1</math> daughter theories</b>	<b>15</b>
5.1	$S = \mathbb{C}^3/\mathbb{Z}_n$ . . . . .	15
5.1.1	Actions of $\mathbb{Z}_n$ on $\mathbb{C}^3$ . . . . .	15
5.1.2	Example: $\mathbb{Z}_5$ . . . . .	16
5.1.3	General $\mathbb{Z}_n$ . . . . .	18
5.2	$S = \mathbb{C}^3/\Delta(3n^2), \Delta(6n^2)$ . . . . .	19
5.3	$S = \mathbb{C}^3/\Sigma_{36 \times 3}, \Sigma_{60 \times 3}, \Sigma_{168 \times 3}, \Sigma_{216 \times 3}, \Sigma_{360 \times 3}$ . . . . .	19
5.4	$S = \mathbb{C}^3/(\mathbb{Z}_m \times \mathbb{Z}_n)$ . . . . .	19
5.4.1	Representations of $\mathbb{Z}_m \times \mathbb{Z}_n$ . . . . .	19
5.4.2	Projection . . . . .	20
<b>6</b>	<b>Correspondence between gauge theory and singularity</b>	<b>21</b>
6.1	From gauge theory to singularity: forward algorithm . . . . .	21
6.1.1	$S = \mathbb{C} \times \mathbb{C}^2/\mathbb{Z}_n$ . . . . .	21
6.1.2	The conifold . . . . .	21
6.2	From singularity to gauge theory: inverse algorithm . . . . .	21
<b>7</b>	<b>A note about projective representations, discrete torsion and deformations</b>	<b>21</b>
7.1	Projective representations and discrete torsion . . . . .	21
7.2	Quiver gauge theories deformations and conifold . . . . .	22
<b>II</b>	<b>Beyond orbifold singularities</b>	<b>22</b>
<b>8</b>	<b>Non-commutative resolutions</b>	<b>22</b>
<b>9</b>	<b>Determinantal varieties as transverse spaces</b>	<b>22</b>
9.1	Basic properties of determinantal varieties . . . . .	22
9.1.1	Computing the dimension . . . . .	22
9.1.2	Singularity . . . . .	23
9.1.3	Action and syzygies . . . . .	23

9.2	Young's lattice . . . . .	23
<b>A</b>	<b>Supersymmetric Yang-Mills theory from D-branes</b>	<b>25</b>
<b>B</b>	<b>Reminder on <math>\mathcal{N} = 4</math> super Yang-Mills theory in <math>D = 4</math></b>	<b>26</b>
B.1	Superconformal group $SU(2, 2 4)$ and its representations . . . . .	26
B.2	Matter content . . . . .	26
B.3	Moduli space and dynamical phases . . . . .	27
<b>C</b>	<b>Gauge anomaly</b>	<b>28</b>
<b>D</b>	<b>Properties of D-branes in type II theories</b>	<b>28</b>
<b>E</b>	<b>Some finite subgroups</b>	<b>29</b>
E.1	Finite subgroups of $SU(2)$ and $SL(2, \mathbb{C})$ . . . . .	29
E.1.1	Finite subgroups . . . . .	29
E.1.2	Irreducible representations . . . . .	30
E.1.3	Character tables . . . . .	30
E.2	Finite subgroups of $SU(3)$ . . . . .	32
<b>F</b>	<b>The McKay correspondence</b>	<b>32</b>
F.1	Classical correspondence . . . . .	32
F.2	Geometrical McKay correspondence . . . . .	34
<b>G</b>	<b>Algebraic geometry</b>	<b>34</b>
G.1	Singularities and resolutions . . . . .	34
G.2	Path algebras . . . . .	34
<b>H</b>	<b>Calabi-Yau manifolds, orbifolds and crepant resolutions</b>	<b>34</b>
<b>I</b>	<b>Graphs</b>	<b>35</b>
<b>J</b>	<b>Toric and non-toric singularities</b>	<b>35</b>
<b>K</b>	<b>Spacetime geometry: ALE space and orbifolds</b>	<b>35</b>
<b>L</b>	<b>Some derivations</b>	<b>35</b>
L.1	. . . . .	35
<b>M</b>	<b>References guide</b>	<b>38</b>
	<b>List of Markers</b>	<b>39</b>

# 1 | Physical setup

## Brane-world paradigm

We consider our world to be a slice in the ten-dimensional spacetime of type II superstring theory, i.e. the worldvolume of a D3-brane. More precisely, we consider a stack of  $N$  D3-branes in order to have  $U(N)$  Chan-Paton factors resulting in a  $U(N)$  gauge group. To have D3-branes, we need to consider type IIB superstring theory. The spacetime is therefore not necessarily  $\mathbb{R}^{1,9}$  but of the more general form

$$M = \mathbb{R}^{1,3} \times M^{(6)}.$$

This is the so-called *brane-world paradigm*.

## Supersymmetry on the worldvolume: holonomy and Calabi-Yau compactification

Independently from string theory, we can ask for the worldvolume theory to be supersymmetric. Preserving any degree of supersymmetry constrains the transverse space  $M^{(6)}$  to be compact, complex, Kähler and to have  $G \subset SU(3)$  holonomy. Namely,  $M^{(6)}$  must be a Calabi-Yau threefold.

Let us discuss the the holonomy part of the constrain in more details as it will play an important role in the following. We start from type IIB superstring theory which is 10-dimensional and has  $\mathcal{N} = 2$  supersymmetry so it possesses 32 supercharges. As usual, they transform under the minimal spinor representation of the bulk Lorentz group, here  $SO(1,9)$ . In ten dimensions this representation is 8-dimensional (complex) which is why there are  $2(8+8) = 32$  supercharges: 8 transforming in the 8-dimensional minimal spinor representation and 8 transforming in the 8-dimensional conjugate minimal representation and whole thing times two since  $\mathcal{N} = 2$ . Let us now consider a four-dimensional field theory resulting from compactification of the transverse six-dimensional space. The number of supercharges that generate supersymmetries for this theory is the number of Killing spinor (covariantly constant spinor): each Killing spinor contracted with the local supersymmetry current generates a residual supersymmetry(\*). Now the link with holonomy: a generic curved six-dimensional manifold has  $O(6)$  holonomy and  $SO(6)$  if it is orientable. Since  $SO(6) \cong SU(4)$ , minimal spinors can be viewed as having four complex component and as transforming under  $SU(4)$ . Indeed, minimal spinors in six dimensions have four complex components. In order to have one covariantly conserved spinor, we look for the biggest subgroup of  $SU(4)$  that leaves a component of the spinor invariant. This is clearly  $\{e\} \times SU(3) \subset SU(4)$  that acts trivially on the first component. The spinor  $(1, 0, 0, 0)$  is then covariantly conserved. Our transverse space must therefore have  $SU(3)$  holonomy such that the parallel transport of the spinor  $(1, 0, 0, 0)$  under any closed loop is a lower  $SU(3)$  rotation. We conclude that if the transverse Calabi-Yau has  $SU(3)$  holonomy, the worldvolume theory has  $\mathcal{N} = 1$  supersymmetry. The same reasoning can be used to obtain that  $SU(2)$  holonomy implies  $\mathcal{N} = 2$  supersymmetry.

## Non-compact transverse space

If we let the worldvolume of the D3-branes carry the requisite gauge theory while the bulk contains gravity, we can relax the compactness condition and study non-compact threefolds. In other words,  $M^{(6)}$  is taken to be Calabi-Yau variety, instead of a Calabi-Yau manifold. A Calabi-Yau varitey is an affine variety that locally models a Calabi-Yau manifold, therefore allowing for singularities.

Using a non-compact transverse space can intuitively be understood as a Kaluza-Klein compactification where we take the size of the compact dimensions to infinity. The four-dimensional gravity coupling constant being inversely proportional to this quantity, there is no gravity in this limit. This makes the analysis much simpler and therefore also serves as an argument to ignore gravity in the worldvolume theory. Consequently, we will mostly ignore gravity and not care about the metric of the spacetime, see appendix K for more details.

## Singular transverse space

The only smooth Calabi-Yau threefold is  $\mathbb{C}^3$  so we are lead to consider singular Calabi-Yau varieties. We usually denote  $S \equiv M^{(6)}$  to remind us of the singular aspect. String theory being a theory of

extended objects, it is well-defined on such singularities. In a sense that will be clarified later, this singular structure of the geometry requires to “project” the theory. As a result, the gauge group  $U(N)$  will be broken down into products of smaller gauge groups. The simplest examples of singular Calabi-Yau varieties are Calabi-Yau orbifolds. We will mainly be interested these examples.

From the point of view of the orbifold, the D3-brane is a point. Consequently, the D3-branes parametrize the transverse space. This is the first clue of the crucial relationship between the worldvolume theory and the Calabi-Yau singularity. Eventually, we will see that the classical vacuum of the gauge theory should be, in explicit coordinates, the defining equation of  $S$ . This is precisely the opposite of the projection manipulation we mentioned above: recovering the transverse space from the gauge theory.

Projecting and computing the classical vacua are therefore inverse operations with respect to each other. This suggest a bijection between the singular transverse space and the gauge theory: the former can be computed from the latter and vice-versa. This is called “forward algorithm” and “inverse algorithm”. We will of course discuss this in more details.

### Mathematical formulation

Mathematically, this brane-world paradigm is the realization of branes as supports of vector bundles (sheaf). Gauge theories on branes are intimately related to algebraic constructions of stable bundles, i.e. holomorphic or algebraic vector bundles that are stable in the sense of geometric invariant theory. In particular, D-brane gauge theories manifest as a natural description of symplectic quotients and their resolutions in geometric invariant theory.

To summarize in more mathematical terms, our D-branes, together with the stable vector bundle (sheaf) supported thereupon, resolves the transverse Calabi-Yau orbifold which is the vacuum for the gauge theory on the worldvolume as a GIT quotient (\*).

### Summary

We consider  $N$  D3-branes in type IIB superstring theory carrying a  $U(N)$  gauge group. The transverse space  $S$  is taken to be a non-compact singular Calabi-Yau variety.

## 2 | The simplest case: smooth transverse space

### 2.1 | Generalities

Let us start by considering the simplest configuration were the transverse Calabi-Yau space is non-singular, i.e. it is a proper smooth Calabi-Yau threefold. As mentioned above, the only smooth Calabi-Yau threefold is  $S = \mathbb{C}^3$ . In this case, the spacetime is simply flat space  $\mathbb{R}^{1,9} = \mathbb{R}^{1,3} \times \mathbb{R}^6$  with a choice a complex structure on  $\mathbb{R}^6$ . From the  $U(N)$  Chan-Paton factors, the worldvolume theory inherits from a  $U(N)$  gauge group. Type IIB superstring theory is a ten-dimensional  $\mathcal{N} = 2$  theory so it has 32 supercharges. The presence of the branes breaks the Lorentz symmetry of  $\mathbb{R}^{1,9}$  as

$$SO(1,9) \rightarrow SO(1,3) \times SO(6), \quad (2.1)$$

whereby breaking half of the supersymmetries, as we explained in the previous section. We are thus left with 16 supercharges. In four dimensions, this corresponds to  $\mathcal{N} = 4$ . The worldvolume theory for  $S = \mathbb{C}^3$  is therefore  $D = 4, \mathcal{N} = 4$   $U(N)$  SCFT gauge theory. This worldvolume theory, obtained in the non-singular case  $S = \mathbb{C}^3$ , is called the *parent theory*.

Note that the D3-brane will warp the flat space metric to that of  $AdS_5 \times S^5$  and the bulk geometry is not strictly  $\mathbb{C}^3$ . However, as stated above, we are only concerned with the local gauge theory and not with gravitational back-reaction, therefore it suffices to consider  $S$  as  $\mathbb{C}^3$ .

### 2.2 | Matter content

As discussed in appendix B, there is only one  $D = 4, \mathcal{N} = 4$  SCFT theory, up to a choice of gauge group  $G$ . In our case,  $G = U(N)$ . The isometry group of the transverse space  $\mathbb{R}^6$  is  $SO(6) \cong SU(4)$ .

Since the scalar fields living on the branes are interpreted as its transverse oscillations,  $SO(6)$  is a global symmetry of the field theory. These global symmetries of worldvolume theory lead to the R-symmetry group  $SU(4)_R$ . The only  $\mathcal{N} = 4$  supermultiplet can be rewritten in terms of  $\mathcal{N} = 1$  supermultiplets as follows:

$$[\mathcal{N} = 4 \text{ vector multiplet}] : V = (\lambda_\alpha, A_\mu, D) \oplus \Phi_A = (\phi^A, \psi_\alpha^A, F^A). \quad (2.2)$$

with  $A = 1, 2, 3$ . In other words, after removing the auxiliary fields  $D$  and  $F^A$ , the matter content is

- a  $U(N)$  gauge field  $A_\mu$  which transforms as a singlets under  $SU(4)_R$ :

$$\text{Gauge transformation} : A_\mu \mapsto UA_\mu U^{-1} + U\partial_\mu U^{-1}, \quad U \in U(N) \quad (2.3)$$

$$\text{R-symmetry} : A_\mu \mapsto A_\mu. \quad (2.4)$$

Note that the usual term

- 4 Weyl fermions  $\psi_\alpha^a \equiv (\lambda_\alpha, \psi_\alpha^1, \psi_\alpha^2, \psi_\alpha^3)$  that transform under the adjoint of  $U(N)$  and are mixed together under the representation  $\mathbf{4}$  of  $SU(4)_R$ . This means that each fermion  $\psi^a$  takes values in  $\mathfrak{u}(N)$ . We denote the components by  $\psi_{IJ}^a$  ( $I, J = 1, \dots, N$ ). Explicitly:

$$\text{Gauge transformation} : \psi^a \mapsto U\psi^a U^\dagger, \quad U \in U(N), \quad (2.5)$$

$$\text{R-symmetry} : \psi^a \mapsto R^a_b \psi^b, \quad R \in SU(4)_R. \quad (2.6)$$

Note that this gives us  $4N^2$  Weyl fermions in total.

- 3 complex scalar fields  $\phi^A$  transforming under the adjoint representation of  $U(N)$  and under the two-times anti-symmetric representation of  $SU(4)_R$ . This means that each  $\phi^A$  takes values in  $\mathfrak{u}(N)$  and we denote the components by  $\phi_{IJ}^A$ . Recall that  $SU(4) \cong SO(6)$  so the action of the R-symmetry can be seen as the  $\mathbf{3}$  of  $SU(3) \subset SU(4)_R$  acting on three complex scalars  $\phi^A$  or equivalently as the  $\mathbf{6}$  of  $SO(6)_R$  acting on 6 real scalars  $X^m$ , the real and imaginary parts of the  $\phi^A$ . They are interpreted as the oscillations of the branes in the transverse space. Explicitly:

$$\text{Gauge transformation} : X^m \mapsto UX^m U^\dagger, \quad U \in U(N), \quad (2.7)$$

$$\text{R-symmetry} : X^m \mapsto R^m_n X^n, \quad R \in SO(6)_R. \quad (2.8)$$

Note that this gives  $6N^2$  real scalars in total. They are the superpartners of the fermions(\*).

Note that the gauge group  $U(N)$  can also be seen as the group of isometries of the metric space  $\mathbb{C}^N$ , i.e.  $\text{Hom}(\mathbb{C}^N, \mathbb{C}^N)$ . From this point of view, the transformations (2.3)-(2.8) can be summarized as

$$A_\mu \in \text{Hom}(\mathbb{C}^N, \mathbb{C}^N), \quad (2.9)$$

$$\psi \in \mathbf{4} \otimes \text{Hom}(\mathbb{C}^N, \mathbb{C}^N), \quad (2.10)$$

$$X \in \mathbf{6} \otimes \text{Hom}(\mathbb{C}^N, \mathbb{C}^N). \quad (2.11)$$

If the transverse space is non-singular, the only possibility is  $S = \mathbb{C}^3$ . In this case, the worldvolume theory is therefore  $D = 4, \mathcal{N} = 4$   $U(N)$  SCFT gauge theory. It is called the *parent theory*.

## Part I

# Orbifold singularities

## 3 | Singular transverse space: orbifold singularities

When the transverse space is singular, the worldvolume theory corresponds to a specific projection of the parent theory that we found in the smooth case  $S = \mathbb{C}^3$ . We call it the *daughter theory*. This projections depends on the type of singularity that one considers. The simplest case is the the case of orbifold singularity, i.e. when the transverse space is a quotient space with a non-free action.

### 3.1 | Generalities

We now wish to pick a discrete group  $\Gamma$  and which acts non-trivially on  $\mathbb{R}^6$ . There are several possibilities:

- $\Gamma \subset \text{SU}(4)$  naturally acts on  $\mathbb{R}^6$ . This does not require a choice of complex structure. We get an  $\mathcal{N} = 0$  theory.
- $\Gamma \subset \text{SU}(3)$  naturally acts on  $\mathbb{C}^3$ , this also requires a choice of complex structure on  $\mathbb{R}^6$ . We get an  $\mathcal{N} = 1$  theory.
- $\Gamma \subset \text{SU}(2)$  naturally acts on the second factor of  $\mathbb{C} \times \mathbb{C}^2$ , so this requires a choice of complex structure on  $\mathbb{R}^6$ . We get an  $\mathcal{N} = 2$  theory.

We are interested in supersymmetric theories so we take  $\Gamma \subset \text{SU}(3)$  with the action

$$\cdot : \begin{pmatrix} \Gamma \times \mathbb{C}^3 & \longrightarrow & \mathbb{C}^3 \\ (\gamma, z) & \longmapsto & \gamma \cdot z \end{pmatrix} \quad (3.1)$$

is the representation of  $\Gamma$  coming from the fundamental representation of  $\text{SU}(3)$ , so  $\cdot$  is just the matrix product. We can see that the origin is always a fixed point so this action is never free. Since  $\mathbb{C}^3$  is a smooth manifold, this makes  $\mathbb{C}^3/\Gamma$  an orbifold. Note this case naturally includes the case  $\Gamma \subset \text{SU}(2)$  too (as  $\text{SU}(2) \subset \text{SU}(3)$ ), just not the case  $\Gamma \in \text{SO}(6)$ . When  $\Gamma \subset \text{SU}(2) \subset \text{SU}(3)$ , it acts trivially on one component so we write  $S = \mathbb{C} \times \mathbb{C}^2/\Gamma$ .

If  $\Gamma$  is a general finite group the condition that  $\mathbb{C}^3/\Gamma$  is an Calabi-Yau orbifold means that there must exist a resolution of this orbifold such that the corresponding smooth space is Calabi-Yau, i.e. a crepant resolution. Existence of such a resolution constrains  $\Gamma$  (\*), see appendix H.

If the transverse space is  $S = \mathbb{C}^3/\Gamma$ , the field theory must be projected to a theory which also invariant under  $\Gamma$ , seen as a subgroup of the R-symmetry group. The prescription is straihg-t-forward: we can use the elements  $\gamma \in \Gamma$  to project out that states that are not  $\Gamma$ -invariant. That is, if  $\rho$  is an embedding of  $\Gamma$  in the gauge group  $\text{U}(N)$ , only the the fields such that

$$\rho(\gamma)A_\mu\rho(\gamma)^{-1} = A_\mu, \quad (3.2)$$

$$R(\gamma)\rho(\gamma)\psi_{IJ}\rho(\gamma)^{-1} = \psi_{IJ}, \quad (3.3)$$

$$R(\gamma)\rho(\gamma)X_{IJ}\rho(\gamma)^{-1} = X_{IJ} \quad (3.4)$$

are kept in the spectrum, where  $\rho$  is a unitary representation of  $\Gamma$  on  $\mathbb{C}^N$  and  $R = \mathbf{4}, \mathbf{6}$ . Let us make two remarks:

- the term  $U\partial_\mu U^{-1}$  is absent from 3.2. Indeed,  $\Gamma$  is a finite group and the only smooth functions  $x \mapsto \Gamma$  are the constant ones. Consequently, transformations of the gauge field under a finite subgroup of its gauge group cannot depend on  $x$ .
- the fields that transform non-trivially under R-symmetry also have an extra induced action of  $\Gamma$ , in agreement with (2.9)-(2.11). The R-symmetry untouched by  $\Gamma$  will be the resulting R-symmetry of daughter theory.

### 3.2 | Representation theory realization

Let  $\{(\rho_i, V_i)\}_{i \in I}$  be a complete set of irreducible representations of  $\Gamma$ . Since  $\Gamma$  is finite, it is particular compact and those representation can be taken to be unitary. Moreover,  $i$  takes a finite number of values. Let us consider a representation of  $\Gamma$  on  $\mathbb{C}^N$ , we denote it  $(\rho, \mathbb{C}^N)$  and also take it to be unitary. In that case,  $\rho(\gamma) \in \text{U}(N)$ . This is what we mean by “the embedding of  $\Gamma$  in  $\text{U}(N)$ ”. The adjoint representation of  $\text{U}(N)$  defined as<sup>1</sup>

$$\text{Ad} : \begin{pmatrix} \text{U}(n) & \longrightarrow & \text{GL}(\mathfrak{u}(N)) \\ U & \longmapsto & \text{Ad}_U \end{pmatrix}, \quad \text{Ad}_U : \begin{pmatrix} \mathfrak{u}(N) & \longrightarrow & \mathfrak{u}(N) \\ \omega & \longmapsto & \text{Ad}_U(\omega) \equiv U\omega U^{-1} \end{pmatrix}, \quad (3.5)$$

<sup>1</sup>this is well-defined since for all  $\omega \in \mathfrak{u}(N)$  and  $U \in \text{U}(N)$ ,  $U\omega U^{-1} \in \mathfrak{u}(N)$ .

now allows us to act with  $\Gamma$  on  $\mathfrak{u}(N)$ . We use this representation in the expression (3.2)-(3.4). More formally, these relations can be rewritten as

$$\text{Ad}_{\rho(\gamma)} A_\mu = A_\mu, \quad (3.6)$$

$$R(\gamma) \text{Ad}_{\rho(\gamma)} \psi_{IJ} = \psi_{IJ}, \quad (3.7)$$

$$R(\gamma) \text{Ad}_{\rho(\gamma)} X_{IJ} = X_{IJ} \quad (3.8)$$

We can decompose  $(\rho, \mathbb{C}^N)$  as follows:

$$(\rho, \mathbb{C}^N) = \bigoplus_{i \in I} (\rho_i, V_i)^{N_i} \quad (3.9)$$

$$= \bigoplus_{i \in I} (\mathbf{1}^{N_i} \otimes \rho_i, \mathbb{C}^{N_i} \otimes V_i) \quad (3.10)$$

where  $N_i$  are integer multiplicities  $((\rho_i, V_i)^{N_i} \equiv (\rho_i, V_i)^{\oplus N_i})$  and  $\mathbf{1}$  is the trivial representation, so  $\Gamma$  acts trivially on the  $\mathbb{C}^{N_i}$ . We have  $\sum_i N_i \dim(\rho_i) = N$ . The rewriting (3.10) will be useful later on.

After the projection, the resulting gauge group is given by the  $\Gamma$ -invariant part of the gauge group, that is  $\text{Hom}(\mathbb{C}^N, \mathbb{C}^N)^\Gamma$ . We use the superscript  $\Gamma$  to indicate that we only keep the trivial representations in the decomposition, that is, we only keep those subspaces that transform trivially. First, we can see that by Schur's lemma (\*)

$$(V_i \otimes V_j^*)^\Gamma = \delta_{ij} \quad (3.11)$$

Now since  $\text{Hom}(\mathbb{C}^N, \mathbb{C}^N) \cong \mathbb{C}^N \otimes (\mathbb{C}^N)^*$ , we get

$$\text{Hom}(\mathbb{C}^N, \mathbb{C}^N)^\Gamma = (\mathbb{C}^N \otimes (\mathbb{C}^N)^*)^\Gamma \quad (3.12)$$

$$= \bigoplus_{i,j \in I} ((\mathbb{C}^{N_i} \otimes V_i) \otimes (\mathbb{C}^{N_j} \otimes V_j^*)^\Gamma)^\Gamma \quad (3.13)$$

$$= \bigoplus_{i,j \in I} (\mathbb{C}^{N_i} \otimes (\mathbb{C}^{N_j})^* \otimes V_i \otimes V_j^*)^\Gamma \quad (3.14)$$

$$= \bigoplus_{i,j \in I} (\mathbb{C}^{N_i} \otimes (\mathbb{C}^{N_j})^*)^\Gamma \otimes (V_i \otimes V_j^*)^\Gamma \quad (3.15)$$

$$= \bigoplus_{i \in I} \mathbb{C}^{N_i} \otimes (\mathbb{C}^{N_i})^* \quad (3.16)$$

so the daughter gauge group is

$$G_{\text{proj}} = \bigotimes_{i \in I} \text{U}(N_i). \quad (3.17)$$

Now it turns out that in the low energy effective field theory the  $\text{U}(1)$  factor of every  $\text{U}(N_i)$  decouples (\*) so the resulting gauge group is in fact

$$G_{\text{proj}} = \bigotimes_{i \in I} \text{SU}(N_i). \quad (3.18)$$

For the matter fields, the reasoning is similar but we now have to take into account the  $\text{R}$ -symmetry. Let  $\mathbf{4} \equiv (\rho_4, V_4)$  be the fundamental representation of  $\text{SU}(4)_R$  and  $\mathbf{6} \equiv (\rho_6, V_6)$  be the fundamental representation of  $\text{SO}(6)_R$ . Wish to compute  $(V_{\mathcal{R}} \otimes \text{Hom}(\mathbb{C}^N, \mathbb{C}^N))^\Gamma$  with  $\mathcal{R} = \mathbf{4}, \mathbf{6}$ :

$$(V_{\mathcal{R}} \otimes \text{Hom}(\mathbb{C}^N, \mathbb{C}^N))^\Gamma = \bigoplus_{i,j \in I} (V_{\mathcal{R}} \otimes (\mathbb{C}^{N_i} \otimes V_i) \otimes (\mathbb{C}^{N_j} \otimes V_j^*)^\Gamma)^\Gamma \quad (3.19)$$

$$= \bigoplus_{i,j \in I} (V_{\mathcal{R}} \otimes \mathbb{C}^{N_i} \otimes (\mathbb{C}^{N_j})^*)^\Gamma \otimes (V_i \otimes V_j^*)^\Gamma \quad (3.20)$$

$$= \bigoplus_{i,j \in I} a_{ij}^{\mathcal{R}} (\mathbb{C}^{N_i} \otimes (\mathbb{C}^{N_j})^*) \quad (3.21)$$



with

$$(\rho_{\mathcal{R}}, V_{\mathcal{R}}) \otimes (\rho_i, V_i) = \bigoplus_{j \in I} a_{ij}^{\mathcal{R}} (\rho_j, V_j). \quad (3.22)$$

This expression makes sense because  $(\rho_{\mathcal{R}}, V_{\mathcal{R}})$  is a representation of  $SU(4)$  so it is in particular a representation of  $SU(3)$  and therefore also in particular a representation of  $\Gamma$ .

Using the orthogonality of characters of irreducible non-equivalent representations, we obtain the explicit expression of the coefficient  $a_{ij}^{\mathcal{R}}$ :

$$a_{ij}^{\mathcal{R}} = \frac{1}{|\Gamma|} \sum_{\gamma=1}^r r_{\gamma} \chi^{\mathcal{R}}(\gamma) \chi^i(\gamma) \chi^j(\gamma)^* \quad (3.23)$$

where  $r_{\gamma}$  is the order of the conjugacy class containing  $\gamma$  and  $\chi^i$  is the character of  $\rho_i$ .

In the end, we can see that:

In the daughter theory, the matter fields become a total of  $a_{ij}^4$  bi-fundamental fermions and  $a_{ij}^6$  bi-fundamental bosons transforming as the  $(\mathbf{N}_i, \bar{\mathbf{N}}_j)$  of  $SU(N_i) \times SU(N_j)$  under the products of gauge groups.

### 3.3 | Field content, quivers and McKay graphs

A convenient way to represent the matter content of a daughter theory is to use *quiver diagrams*. A quiver is a finite oriented graph such that each node  $i$  represents a gauge factor  $SU(N_i)$  and each arrow  $i \rightarrow j$  represents a bi-fundamental field transforming under  $(\mathbf{N}_i, \bar{\mathbf{N}}_j)$ . The *adjacency matrix*  $A$  of the graph is a  $k \times k$  with  $k$  being the number of nodes (gauge factors) whose elements  $A_{ij}$  are the number of arrows (bi-fundamental fields) from  $i$  to  $j$ . In other words, from (3.21), the adjacency matrix of the fermions has elements  $A_{ij} = a_{ij}^4$  and the one of the scalars has elements  $A_{ij} = a_{ij}^6$ .

On the other hand, given finite group  $\Gamma$ , a representation  $(\rho_W, W)$  and a complete set of irreducible representations  $\{(\rho_i, V_i)\}_{i \in I}$  of the latter, one can construct a McKay graph (or quiver) as follows:

1. Draw a vertex for every representation  $(\rho_i, V_i)$ .
2. For every representation  $(\rho_i, V_i)$ , compute the decomposition

$$(\rho_W, V_W) \otimes (\rho_i, V_i) = \bigoplus_j (\rho_j, V_j)^{\oplus n_{ij}}$$

where  $n_{ij}$  is the multiplicity of  $(\rho_j, V_j)$  in the decomposition of  $(\rho_W, V_W) \otimes (\rho_i, V_i)$ .

3. For every  $n_{ij} > 0$ , draw  $n_{ij}$  arrows from the vertex of  $(\rho_i, V_i)$  to the one of  $(\rho_j, V_j)$ .

When  $\Gamma \subset SU(2)$  and that  $(\rho_W, W)$  is its defining representation, the McKay graphs are in one-to-one correspondence with the extended Dynkin diagrams of the simply laced Lie algebras. This is the classical McKay correspondence, see appendix F.

From (3.22) we see that  $n_{ij} = a_{ij}^{\mathcal{R}}$ , i.e. the matter quivers that we defined previously are exactly the McKay graph associated to the matter representation in question. Put differently, the matter content of the daughter theory is encapsulated in the McKay graphs of  $\Gamma$  and with respect to  $\mathcal{R}$  with  $\mathcal{R} = 4$  for spinors and  $\mathcal{R} = 6$  for scalars. This very important point allows us to use known results on McKay graphs such as the McKay correspondence for example.

### 3.4 | A simple example: $S = \mathbb{C}^3/\mathbb{Z}_3$

We illustrate the previous discussion by the simple case where  $\Gamma = \mathbb{Z}_3$  acts on  $\mathbb{C}^2$  as

$$\begin{bmatrix} z_1 \\ z_2 \\ z_3 \end{bmatrix} \mapsto \begin{bmatrix} \zeta_3 & 0 & 0 \\ 0 & \zeta_3 & 0 \\ 0 & 0 & \zeta_3 \end{bmatrix} \begin{bmatrix} z_1 \\ z_2 \\ z_3 \end{bmatrix} \quad (3.24)$$

i.e. the transverse space is the orbifold  $\mathbb{C}^3/\mathbb{Z}_3$ . This simple example is a good first approach in which we will to explain in details each step so that we can go faster afterwards.

### 3.4.1 | Projection

Let us consider a representation  $(\rho, \mathbb{C}^N)$  of  $\mathbb{Z}_3$ . A complete set of irreducible representations of  $\mathbb{Z}_3$  is given by  $\{(\rho_k, V_k)\}_{k=0,1,2}$  with  $V_k = \mathbb{C}$  and

$$\rho_i(g) = \zeta_3^k \quad (3.25)$$

where  $g$  is the generator of  $\mathbb{Z}_3$ . The representation  $(\rho, V)$  can be decomposed as

$$(\rho, V) = \bigoplus_{i=0}^2 N_i(\rho_i, V_i). \quad (3.26)$$

In other words, it is equivalent to the representation

$$\rho(g) = \begin{bmatrix} 1 & & & & & & 0 \\ & \ddots & & & & & \\ & & 1 & & & & \\ & & & \zeta_3 & & & \\ \vdots & & & & \ddots & & \vdots \\ & & & & & \zeta_3 & \\ & & & & & & \zeta_3^2 \\ & & & & & & \ddots \\ 0 & & & & & & \zeta_3^2 \end{bmatrix} \begin{matrix} \left. \vphantom{\begin{matrix} 1 \\ \vdots \\ 0 \end{matrix}} \right\} N_0 \\ \left. \vphantom{\begin{matrix} \zeta_3 \\ \zeta_3^2 \end{matrix}} \right\} N_1 \\ \left. \vphantom{\begin{matrix} \zeta_3^2 \end{matrix}} \right\} N_2 \end{matrix} \quad (3.27)$$

Since  $\dim \rho_i = 1$ , we have

$$N_0 + N_1 + N_2 = N. \quad (3.28)$$

The gauge field configurations that are left invariant under the action of  $\mathbb{Z}_n$  are therefore the ones that satisfy

$$\rho(g)A_\mu\rho(g)^{-1} = A_\mu. \quad (3.29)$$

We actually want this relation to be true for any element of  $\mathbb{Z}_n$  but in this case it is invariant under any element of  $\mathbb{Z}_n$  if and only if it is invariant under the generator  $g$  of  $\mathbb{Z}_3$ , so we only need to impose (3.29). The constrained is easily solved by using the bi-index notation  $A_{\mu;i\alpha_i,j\beta_j}$  ( $i, j = 0, 1, 2, \alpha_i, \beta_i = 1, \dots, N_i$ ) for the component of the gauge fields. From (4.2), we can see that

$$A_{\mu;i\alpha_i,i\beta_j} \mapsto \rho_i(g)A_{\mu;i\alpha_i,j\beta_j}\rho_j(g)^{-1} = \zeta_n^{i-j}A_{\mu;i\alpha_i,j\beta_j}. \quad (3.30)$$

thus only the configurations with  $A_{\mu;i\alpha_i,j\beta_j} = 0$  if  $i \neq j$  are invariant. The gauge field has therefore a block diagonal form:

$$A_\mu = \begin{bmatrix} A_{\mu;00} & & \\ & A_{\mu;11} & \\ & & A_{\mu;2,2} \end{bmatrix} \quad (3.31)$$

with  $A_{\mu;ij} \equiv (A_{\mu;i\alpha_i,j\beta_j})_{\alpha_i=0,\dots,N_i,\beta_j=0,\dots,N_j}$ . The block  $A_{ii}$  transforms under  $\mathbb{Z}_3$  under  $(\rho_i, V_i)^{N_i}$ . Consequently, the gauge group is now broken to

$$G_{\text{proj}} = \text{U}(N_0) \times \text{U}(N_1) \times \text{U}(N_2), \quad (3.32)$$

the biggest subgroup of  $\text{U}(N)$  that preserves those form of configurations.

Let us study the scalars.  $\mathbb{Z}_3$  acts on the three complex scalars through

$$R(g) = \rho_1^{\oplus 3}(g) = \zeta_3 \mathbb{1}_3 = \begin{bmatrix} \zeta_3 & 0 & 0 \\ 0 & \zeta_3 & 0 \\ 0 & 0 & \zeta_3 \end{bmatrix} \quad (3.33)$$

or, equivalently, on the real scalars as  $R(g) = \zeta_3 \mathbb{1}_6$ . According to (2.7)-(2.8), the scalar field configurations that are left invariant satisfy

$$R(g)^m {}_n \rho(g) X^n \rho(g)^{-1} = X^m \quad (3.34)$$

for all  $g \in \mathbb{Z}_n$ . Using the bi-index notations, this becomes

$$X_{i\alpha_i, j\beta_j}^m \mapsto \zeta_n \delta_n^m {}_n \rho_i(g) X_{i\alpha_i, j\beta_j}^n \rho_j(g)^{-1} = \zeta_3^{i-j+1} X_{i\alpha_i, j\beta_j}^m \quad (3.35)$$

$$\bar{X}_{i\alpha_i, j\beta_j}^m \mapsto \zeta_3^{-1} \delta_n^m {}_n \rho_i(g) \bar{X}_{i\alpha_i, j\beta_j}^n \rho_j(g)^{-1} = \zeta_3^{i-j-1} \bar{X}_{i\alpha_i, j\beta_j}^m \quad (3.36)$$

thus only the configurations with  $X_{i\alpha_i, j\beta_j}^n = 0$  if  $i - j + 1 \neq 0$  are left invariant and only the configurations with  $\bar{X}_{i\alpha_i, j\beta_j}^n = 0$  if  $i - j - 1 \neq 0$  are left invariant. The scalar fields  $X$  have a block off-diagonal form:

$$X^m = \begin{bmatrix} 0 & X_{01}^m & 0 \\ 0 & 0 & X_{12}^m \\ X_{21}^m & 0 & 0 \end{bmatrix}, \quad \bar{X}^m = \begin{bmatrix} 0 & 0 & \bar{X}_{02}^m \\ \bar{X}_{10}^m & 0 & 0 \\ 0 & \bar{X}_{21}^m & 0 \end{bmatrix} \quad (3.37)$$

with the block notations

$$X_{ij}^m \equiv (X_{i\alpha_i, j\beta_j}^m)_{\alpha_i=0, \dots, N_i, \beta_j=0, \dots, N_j}$$

$$\bar{X}_{ij}^m \equiv (\bar{X}_{i\alpha_i, j\beta_j}^m)_{\alpha_i=0, \dots, N_i, \beta_j=0, \dots, N_j}.$$

The block  $X_{ij}^m$  is an  $N_i \times N_j$  block and transforms under the representation  $(\mathbf{N}_i, \bar{\mathbf{N}}_j)$  of  $U(N_i) \times U(N_j)$ :

$$X_{i, i+1}^m \in \mathbf{N}_{i+1} \otimes \bar{\mathbf{N}}_i \cong \text{Hom}(V_{i+1}, V_i), \quad (3.38)$$

$$\bar{X}_{i+1, i}^m \in \mathbf{N}_i \otimes \bar{\mathbf{N}}_{i+1} \cong \text{Hom}(V_i, V_{i+1}). \quad (3.39)$$

Let us make an important observation: the form of the scalar fields are the same for every  $m = 0, \dots, 5$ . This can be traced back to the fact that  $R(g) = \zeta_3 \mathbb{1}_6$  so the R-symmetry action of  $\mathbb{Z}_n$  is the same for all  $m$ .

Let us now study the four Weyl fermions  $\psi^a$ . (\*)

### 3.4.2 | Quiver

We can draw the quiver of this daughter theory. We have three types of bi-fundamental scalar fields:

$$X_{01}^m \in (\mathbf{N}_1, \bar{\mathbf{N}}_0), \quad X_{12}^m \in (\mathbf{N}_2, \bar{\mathbf{N}}_1), \quad X_{20}^m \in (\mathbf{N}_0, \bar{\mathbf{N}}_2). \quad (3.40)$$

In each representation bi-fundamental representation there are six real scalars, i.e. 3 complex scalars. They are each represented by an arrow between the right representations, see fig. 1.

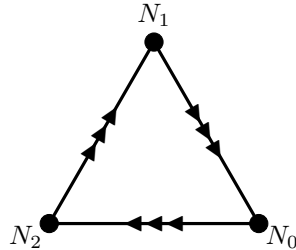


Figure 1: Quiver of the  $\mathbb{C}^3/\mathbb{Z}_3$  daughter theory.

The adjacency matrix of this quiver is

$$A = \begin{bmatrix} 0 & 3 & 0 \\ 0 & 0 & 3 \\ 3 & 0 & 0 \end{bmatrix} \quad (3.41)$$

which is coherent with the fact that the coefficients of the McKay decomposition of  $\rho_1 \oplus \rho_1 \oplus \rho_1$  are

$$\begin{aligned} n_{00} &= 0, & n_{01} &= 3 & n_{02} &= 0, \\ n_{10} &= 0, & n_{11} &= 0, & n_{12} &= 3, \\ n_{20} &= 3, & n_{21} &= 0, & n_{22} &= 0. \end{aligned} \tag{3.42}$$

### 3.4.3 | Gauge anomaly cancellation

Finally, let us discuss the gauge anomaly cancellation. Our fields transform under the adjoint representation of  $SU(N)$ , under the fundamentals of  $SU(N_i)$  and under the anti-fundamentals of  $SU(N_i)$ . The adjoint representation being real, it is self-conjugate and therefore does not contribute to the anomaly. Fundamentals of  $SU(N_i)$  have a  $+1$  contribution to the anomaly and anti-fundamentals of  $SU(N_i)$  have a  $-1$  contribution to the anomaly. Anomaly cancellation therefore imposes that the contribution of the fundamental and of the anti-fundamental of  $SU(N_i)$  cancel each other for each  $i = 0, 1, 2$ , see appendix (C). The bi-fundamental representation  $(\mathbf{N}_i, \overline{\mathbf{N}}_j)$  counts as  $N_j$  fundamentals of  $SU(N_i)$  and as  $N_i$  anti-fundamentals  $SU(N_j)$ . We get the three following conditions:

$$SU(N_0) : N_2 - N_1 = 0, \tag{3.43}$$

$$SU(N_1) : N_0 - N_2 = 0, \tag{3.44}$$

$$SU(N_2) : N_1 - N_0 = 0. \tag{3.45}$$

which immediately imply

$$N_0 = N_1 = N_2. \tag{3.46}$$

From (3.28), we get that  $N_0 = N_1 = N_2 = N/3$ , meaning that the daughter theory has quantum gauge symmetry if and only if the parent theory has gauge group  $SU(N)$  where  $N$  is a multiple of 3. Or, in other words, if the number of D-branes is a multiple of 3.

## 3.5 | $(p+1)$ -dimensional quiver gauge theories

Let us mention that we can generalize our initial brane-world paradigm and consider  $Dp$ -branes in type II string theory (type IIA if  $p$  is even and type IIB if  $p$  is odd) instead of just D3-branes. The spacetime is then of the form

$$M = \mathbb{R}^{1,p} \times \mathbb{R}^{9-p}/\Gamma \tag{3.47}$$

where  $\Gamma$  is a discrete subgroup of  $\text{Spin}(9-p)$ . If  $\Gamma$  is a subgroup of a special holonomy group, we recover a somewhat generalized version of the paradigm that we discussed above. In this case the transverse space is a Calabi-Yau orbifold and some degree of supersymmetry is preserved. Note that the fermionic and bosonic quivers that coincide. If  $\Gamma$  is not a subgroup of a special holonomy group, then the  $(p+1)$ -dimensional quiver gauge theory that we obtain in the low-energy limit is not supersymmetric. We then have different quivers for the fermions and the bosons, although with the same vertices, by definition.

Recall that the fraction of supercharges that is preserved by compactifying on a Calabi-Yau  $n$ -fold (with  $SU(n)$  holonomy) is  $2^{1-n}$ . Starting from the  $\mathcal{N} = 2$  10-dimensional type IIB string theory with 32 supercharges, this means that

- if we compactify on a 1-fold, we get 32 supercharges in 8 dimensions so  $\mathcal{N} = 2$ ,
- if we compactify on a 2-fold, we get 16 supercharges in 6 dimensions so  $\mathcal{N} = 2$ ,
- if we compactify on a 3-fold, we get 8 supercharges in 4 dimensions so  $\mathcal{N} = 2$ ,
- if we compactify on a 4-fold, we get 4 supercharges in 2 dimensions so  $\mathcal{N} = 4$ .

We will however mostly consider 4-dimensional quiver gauge theories, i.e. living on D3-branes.

## 4 | Orbifold constructions of $\mathcal{N} = 2$ daughter theories

### 4.1 | $S = \mathbb{C} \times \mathbb{C}^2/\mathbb{Z}_n$

We consider a representation  $(\rho, \mathbb{C}^N)$  of  $\mathbb{Z}_n$ . We decompose it on the set of irreducible representations of  $\mathbb{Z}_n$  as

$$(\rho, V) = \bigoplus_{i=0}^{n-1} N_i(\rho_i, V_i). \quad (4.1)$$

In other words, it is equivalent to the representation

$$\rho(g) = \left[ \begin{array}{cccccc} 1 & & \cdots & & & 0 \\ & \ddots & & & & \\ & & 1 & & & \\ \vdots & & & \ddots & & \vdots \\ & & & \zeta_n^{n-1} & & \\ & & & & \ddots & \\ 0 & & \cdots & & & \zeta_n^{n-1} \end{array} \right] \left. \begin{array}{l} \\ \\ \\ \\ \\ \end{array} \right\} \begin{array}{l} N_0 \\ \\ \\ \\ \\ N_{n-1} \end{array}. \quad (4.2)$$

Since  $\dim \rho_i = 1$ ,  $\sum_i N_i = N$ .

The gauge field configurations that are left invariant under the action of  $\mathbb{Z}_n$  are therefore the ones that satisfy

$$\rho(g)A_\mu\rho(g)^{-1} = A_\mu. \quad (4.3)$$

The constrained is easily solved by using the bi-index notation:

$$A_{\mu;i\alpha_i,j\beta_j} \mapsto \rho_i(g)A_{\mu;i\alpha_i,j\beta_j}\rho_j(g)^{-1} = \zeta_n^{i-j}A_{\mu;i\alpha_i,j\beta_j}. \quad (4.4)$$

thus only the configurations with  $A_{\mu;i\alpha_i,j\beta_j} = 0$  if  $i \neq j$  are invariant. The gauge field has therefore a block diagonal form:

$$A_\mu = \begin{bmatrix} A_{\mu;00} & & & \\ & A_{\mu;11} & & \\ & & \ddots & \\ & & & A_{\mu;n-1,n-1} \end{bmatrix} \quad (4.5)$$

with  $A_{\mu;ij} \equiv (A_{\mu;i\alpha_i,j\beta_j})_{\alpha_i=0,\dots,N_i,\beta_j=0,\dots,N_j}$ . The block  $A_{ii}$  transforms under  $\mathbb{Z}_n$  as  $(\rho_i, V_i)^{N_i}$ . For now it is only a simple generalization of the case  $\mathbb{C}^3/\mathbb{Z}_3$ . This makes sense: projection of the gauge field only depends the discrete group  $\Gamma$ , not on the way it acts on  $\mathbb{C}^3$  because it does not transform under R-symmetry.

The gauge group is now broken to

$$G_{\text{proj}} = \prod_{i=0}^{n-1} \text{U}(N_i). \quad (4.6)$$

Now for the scalar fields. The action of  $\mathbb{Z}_n$  that we consider leaves the first component of  $\mathbb{C}^3$  untouched so we take the action  $\mathbf{1} \oplus \mathbf{2}$  where  $\mathbf{2}$  is the usual action of  $\mathbb{Z}_n$  on  $\mathbb{C}^2$ . In other words,

$$R(g) = \begin{bmatrix} 1 & 0 & 0 \\ 0 & \zeta_n & 0 \\ 0 & 0 & \zeta_n^{-1} \end{bmatrix}. \quad (4.7)$$

Or, equivalently,  $R(g)^m_n = \delta^m_n A_n$  with  $A_m = (1, 1, \zeta_n, \zeta_n, \zeta_n^{-1}, \zeta_n^{-1})$ . The scalar field configurations that are left invariant satisfy

$$R(g)^m_n \rho(g) X^n \rho(g)^{-1} = X^m \quad (4.8)$$

for all  $g \in \mathbb{Z}_n$ . Using the bi-index notations, this becomes

$$X_{i\alpha_i, j\beta_j}^m \mapsto \delta_n^m A_n \rho_i(g) X_{i\alpha_i, j\beta_j}^m \rho_j(g)^{-1} = \delta_n^m A_n \zeta_n^{i-j} X_{i\alpha_i, j\beta_j}^m = \begin{cases} \zeta_n^{i-j} X_{i\alpha_i, j\beta_j}^m, & m = 0, 1 \\ \zeta_n^{i-j+1} X_{i\alpha_i, j\beta_j}^m, & m = 2, 3 \\ \zeta_n^{i-j-1} X_{i\alpha_i, j\beta_j}^m, & m = 4, 5 \end{cases} \quad (4.9)$$

$$\overline{X}_{i\alpha_i, j\beta_j}^m \mapsto \delta_n^m \overline{A}_n \rho_i(g) \overline{X}_{i\alpha_i, j\beta_j}^m \rho_j(g)^{-1} = \delta_n^m \overline{A}_n \zeta_n^{i-j} \overline{X}_{i\alpha_i, j\beta_j}^m = \begin{cases} \zeta_n^{i-j} \overline{X}_{i\alpha_i, j\beta_j}^m, & m = 0, 1 \\ \zeta_n^{i-j-1} \overline{X}_{i\alpha_i, j\beta_j}^m, & m = 2, 3 \\ \zeta_n^{i-j+1} \overline{X}_{i\alpha_i, j\beta_j}^m, & m = 4, 5 \end{cases} \quad (4.10)$$

thus only the configurations with  $X_{i\alpha_i, j\beta_j}^{0,1} = 0$  if  $i - j \neq 0$ ,  $X_{i\alpha_i, j\beta_j}^{2,3} = 0$  if  $i - j + 1 \neq 0$  and  $X_{i\alpha_i, j\beta_j}^{4,5} = 0$  if  $i - j - 1 \neq 0$  are left invariant (and similarly for the conjugated fields). The scalar fields  $X$  have a the following forms:

$$X^{0,1} = \begin{bmatrix} X_{00}^{0,1} & & 0 \\ & \ddots & \\ 0 & & X_{n-1, n-1}^{0,1} \end{bmatrix}, \quad (4.11)$$

$$X^{2,3} = \begin{bmatrix} 0 & X_{01}^{2,3} & & 0 \\ & \ddots & \ddots & \\ X_{n-1, 0}^{2,3} & & 0 & X_{n-2, n-1}^{2,3} \\ & & 0 & \end{bmatrix}, \quad X^{4,5} = \begin{bmatrix} 0 & & & X_{0, n-1}^{4,5} \\ X_{10}^{4,5} & 0 & & \\ & \ddots & \ddots & \\ 0 & & X_{n-1, n-2}^{4,5} & 0 \end{bmatrix}, \quad (4.12)$$

$$\overline{X}^{0,1} = \begin{bmatrix} \overline{X}_{00}^{0,1} & & 0 \\ & \ddots & \\ 0 & & \overline{X}_{n-1, n-1}^{0,1} \end{bmatrix}, \quad (4.13)$$

$$\overline{X}^{2,3} = \begin{bmatrix} 0 & & & \overline{X}_{0, n-1}^{2,3} \\ \overline{X}_{10}^{2,3} & 0 & & \\ & \ddots & \ddots & \\ 0 & & \overline{X}_{n-1, n-2}^{2,3} & 0 \end{bmatrix}, \quad \overline{X}^{4,5} = \begin{bmatrix} 0 & \overline{X}_{01}^{4,5} & & 0 \\ & \ddots & \ddots & \\ \overline{X}_{n-1, 0}^{4,5} & & 0 & \overline{X}_{n-2, n-1}^{4,5} \\ & & 0 & \end{bmatrix} \quad (4.14)$$

so  $X_{ij}^m$  is an  $N_i \times N_j$  block and transforms under the representation  $(\mathbf{N}_i, \overline{\mathbf{N}}_j)$  of  $U(N_i) \times U(N_j)$ :

$$X_{i,i}^{0,1} \in \mathbf{N}_i \otimes \overline{\mathbf{N}}_i \cong \text{Hom}(V_i, V_i), \quad (4.15)$$

$$X_{i, i+1}^{2,3} \in \mathbf{N}_{i+1} \otimes \overline{\mathbf{N}}_i \cong \text{Hom}(V_{i+1}, V_i), \quad (4.16)$$

$$X_{i+1, i}^{4,5} \in \mathbf{N}_i \otimes \overline{\mathbf{N}}_{i+1} \cong \text{Hom}(V_i, V_{i+1}). \quad (4.17)$$

So the scalar fields are split up in three families depending on the way they transform under R-symmetry. We now see a big difference with the case  $\mathbb{C}^3/\mathbb{Z}_3$ : since the R-symmetry does not act the same way on each directions in  $\mathbb{C}^3$ , the invariant scalar field configurations are not the same in each direction either.

Let us draw the quiver for the case  $n = 3$  so that we can compare to 1. We have  $2 \cdot 9 = 18$  real scalar fields in 9 different representations:

$$X_{00}^0, X_{00}^1 \in (\mathbf{N}_0, \overline{\mathbf{N}}_0), \quad X_{11}^0, X_{11}^1 \in (\mathbf{N}_1, \overline{\mathbf{N}}_1), \quad X_{22}^0, X_{22}^1 \in (\mathbf{N}_2, \overline{\mathbf{N}}_2), \quad (4.18)$$

$$X_{01}^2, X_{01}^3 \in (\mathbf{N}_1, \overline{\mathbf{N}}_0), \quad X_{12}^2, X_{12}^3 \in (\mathbf{N}_2, \overline{\mathbf{N}}_1), \quad X_{20}^2, X_{20}^3 \in (\mathbf{N}_0, \overline{\mathbf{N}}_2), \quad (4.19)$$

$$X_{10}^4, X_{10}^5 \in (\mathbf{N}_0, \overline{\mathbf{N}}_1), \quad X_{21}^4, X_{21}^5 \in (\mathbf{N}_1, \overline{\mathbf{N}}_2), \quad X_{02}^4, X_{02}^5 \in (\mathbf{N}_2, \overline{\mathbf{N}}_0), \quad (4.20)$$

We now only have 1 complex scalar in each representation and the quiver is given by 2.

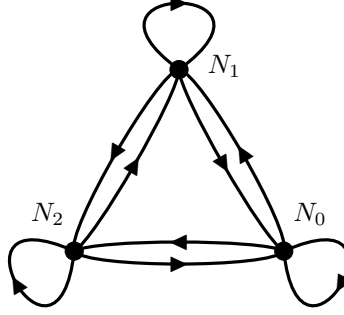


Figure 2: Quiver of the  $\mathbb{C} \times \mathbb{C}^2/\mathbb{Z}_3$  daughter theory.

It is easy to see how the construction of the the quiver generalizes for arbitrary  $n$ . The adjacency matrix is

$$A = \begin{bmatrix} 1 & \cdots & 1 \\ \vdots & \ddots & \vdots \\ 1 & \cdots & 1 \end{bmatrix} \quad (4.21)$$

which is, as it should, coincides with the McKay decomposition of  $\mathbf{1} \oplus \mathbf{2}$ :

$$\begin{aligned} n_{00} &= 1, & \dots, & & n_{0,n-1} &= 1, \\ \vdots & & & & \vdots & \\ n_{n-1,0} &= 1, & \dots, & & n_{n-1,n-1} &= 1. \end{aligned} \quad (4.22)$$

Gauge anomaly cancellation now imposes that

$$N_{i-1} - N_{i-1} + N_{i+1} + N_{i+1} = 0 \quad (4.23)$$

for  $i = 0, \dots, n-1$ . Those constrains are always satisfied so the the factors  $N_i$  are arbitrary, as long as  $\sum_i N_i = N$  of course.

This quiver is easy to scale up for any  $n$ . We therefore implement that into Mathematica and get the quivers for any  $n$ .

## 4.2 | $S = \mathbb{C} \times \mathbb{C}^2/\mathcal{D}_n$

We quotient  $\mathbb{C}^3$  by the binary dihedral group  $\mathcal{D}_n$  that acts on the last two components. A set of irreducible representations of  $\mathcal{D}_n$  is given by

$$\{(\rho_i, V_i)\}_{i=0, \dots, n+2} \quad (4.24)$$

with  $V_i = \mathbb{C}$  for  $i = 0, \dots, 3$  and  $V_i = \mathbb{C}^2$  for  $i = 4, \dots, n+2$ , so there are 4 one-dimensional representations and  $n-1$  two-dimensional representations. They are explicitly given in section E.1.2. Using the bi-index notations  $A_{\mu; i\alpha_i, j\beta_j}$  with  $i, j = 0, \dots, n+2$  and  $\alpha_i, \beta_j = 0, \dots, \dim \rho_i \cdot N_i - 1$ , the invariant configurations must have  $A_{\mu; ij} = 0$  if  $i \neq j$ , i.e. it must have a diagonal block-form, wiht block of size  $N_i \times N_i$  for  $i = 0, \dots, 3$  and of size  $2N_i \times 2N_i$  for  $i = 4, \dots, n+1$ . The blocks transforming under the 2-dimensional representations must have the form

$$A_{\mu; ii} = \begin{bmatrix} A_{\mu; i0, i0} & A_{\mu; i0, i1} & \cdots & A_{\mu; i, 0, i, 2N_i-2} & A_{\mu; i, 0, i, 2N_i-1} \\ \mp A_{\mu; i0, i1} & A_{\mu; i0, i0} & \cdots & \mp A_{\mu; i, 0, i, 2N_i-1} & A_{\mu; i, 0, i, 2N_i-2} \\ \vdots & \cdots & \ddots & \vdots & \vdots \\ A_{\mu; i, 2N_i-2, i, 0} & A_{\mu; i, 2N_i-2, i, 1} & \cdots & A_{\mu; i, 2N_i-2, i, 2N_i-2} & A_{\mu; i, 2N_i-2, i, 2N_i-1} \\ \mp A_{\mu; i, 2N_i-2, i, 1} & A_{\mu; i, 2N_i-2, i, 0} & \cdots & \mp A_{\mu; i, 2N_i-2, i, 2N_i-1} & A_{\mu; i, 2N_i-2, i, 2N_i-2} \end{bmatrix} \quad (4.25)$$

so each block is composed of  $N_i^2$   $2 \times 2$  matrices of the form

$$\begin{bmatrix} A & B \\ \mp B & A \end{bmatrix}. \quad (4.26)$$

We take the “−” signs when  $i$  is even and the “+” signs when  $i$  is odd. This comes from the fact that the 2-dimensional representations involve alternating signs.

R-symmetry acts on the three complex scalar fields as

$$R(A) = A \quad (4.27)$$

(\*)

### 4.3 | $S = \mathbb{C} \times \mathbb{C}^2/\mathcal{T}, \mathcal{O}, \mathcal{I}$

(\*)

## 5 | Orbifold constructions of $\mathcal{N} = 1$ daughter theories

### 5.1 | $S = \mathbb{C}^3/\mathbb{Z}_n$

Let us now consider the general case of  $\mathbb{Z}_n$  acting on  $\mathbb{C}^3$ , i.e. we want to generalize the case that we treated in section 3.4.

#### 5.1.1 | Actions of $\mathbb{Z}_n$ on $\mathbb{C}^3$

The first thing to do is to specify a representation of  $\mathbb{Z}_n$  on  $\mathbb{C}^3$ . Any representation  $(R, \mathbb{C}^3)$  can be decomposed as

$$R = \bigoplus_{k=0}^{n-1} \rho_k^{\oplus n_k} \quad (5.1)$$

and is therefore equivalent to a block-diagonal representation. This implies that  $\sum_k n_k = 3$  so  $R$  can always be written as

$$R(g) = (\rho_a \oplus \rho_b \oplus \rho_c)(g) = \begin{bmatrix} \zeta_n^a & 0 & 0 \\ 0 & \zeta_n^b & 0 \\ 0 & 0 & \zeta_n^c \end{bmatrix} \quad (5.2)$$

with  $a, b, c$  some arbitrary exponents. We denote the representation (5.2) by  $(a, b, c)$ . Let us make a few remarks on these representations:

- Since we consider  $\mathbb{Z}_n$  as a subgroup of  $\text{SU}(3)$ , we must have

$$(a + b + c) \mod n = 0 \quad (5.3)$$

such that  $\det(R(g)) = 1$ .

- Taking  $a$  or  $a + n$  gives the same representation and the same is true for  $b$  and  $c$ , so we can actually restrict ourselves to  $0 < a, b, c < n$ . We can trade the strict inequalities and allow the values 0 or  $n$  if we want to consider trivial representations as well. In this case, we will find that at least one direction of  $\mathbb{C}^3$  is left untouched, i.e. we are in the situation where the orbifold is actually  $\mathbb{C} \times \mathbb{C}^2/\mathbb{Z}_n$ , which we will not consider here.
- Permuting  $a, b, c$  gives equivalent representations so we can take  $a, b, c$  to be ordered, so  $0 < a \leq b \leq c < n$ .

The problem has now been reformulated as follows: we are looking for all possible ordered triplets  $(a, b, c)$  such that  $0 < a \leq b \leq c < n$  and (5.3). For  $n = 1$  and  $n = 2$ , it is clear that this is not possible, the possibilities involve at least one trivial representation. For  $n = 3$ , the only possibility is  $(1, 1, 1)$ , the one



we used in 3.4. What about bigger values of  $n$ ? First, note that  $0 < a \leq b \leq c < n$  implies that the maximum value of  $a+b+c$  is  $3n-1$ . From equation (5.3), the only two possibilities are then  $a+b+c = n$  or  $a+b+c = 2n$ . What simplifies the analysis is that the second case can actually be ignored because it is already being taken care of by the first case. Let us explain that in more details: (\*)

Now that we have seen that the only possibility is  $a+b+c = n$ , we find that the maximum value that  $a$  can take is  $\lfloor n/3 \rfloor$ , so  $a = 1, \dots, \lfloor n/3 \rfloor$ .  $c$  can then be expressed in terms of  $a$  and  $b$  as  $c = n - a - b$ . The constraint  $b \leq c$  then implies  $b \leq (n-a)/2$ , i.e.  $b = a, \dots, \lfloor (n-a)/2 \rfloor$ . To summarize, the representations are all the representations of the form

$$(a, b, n - a - b) \quad (5.4)$$

with  $a = 1, \dots, \lfloor n/3 \rfloor$  and  $b = a, \dots, \lfloor (n-a)/2 \rfloor$ . For small values, we get

$$\begin{aligned} \mathbb{Z}_1 &: \text{necessarily involves a trivial representation} \\ \mathbb{Z}_2 &: \text{necessarily involves a trivial representation} \\ \mathbb{Z}_3 &: (1, 1, 1) \\ \mathbb{Z}_4 &: (1, 1, 2) \\ \mathbb{Z}_5 &: (1, 1, 3), (1, 2, 2) (*) \\ \mathbb{Z}_6 &: (1, 1, 4), (1, 2, 3), (2, 2, 2) \\ \mathbb{Z}_7 &: (1, 1, 5), (1, 2, 4), (1, 3, 3), (2, 2, 3) (*) \\ &\vdots \end{aligned}$$

For an arbitrary  $n$ , the total number of different representation is

$$\sum_{a=1}^{\lfloor n/3 \rfloor} \left\lfloor \frac{n-3a}{2} + 1 \right\rfloor = \begin{cases} 3k^2, & \text{if } n = 6k \\ 3k^2 + k, & \text{if } n = 6k + 1, \\ 3k^2 + 2k, & \text{if } n = 6k + 2, \\ 3k^2 + 3k + 1, & \text{if } n = 6k + 3, \\ 3k^2 + 4k + 1, & \text{if } n = 6k + 4, \\ 3k^2 + 5k + 2, & \text{if } n = 6k + 5 \end{cases} \quad (5.5)$$

The details are presented in appendix L.1.

### 5.1.2 | Example: $\mathbb{Z}_5$

For  $\mathbb{Z}_5$ , we saw that there are two nonequivalent ways of acting on  $\mathbb{C}^3$ , see (??). Let us first consider the action  $(1, 1, 3)$

$$R(g) = \begin{bmatrix} \zeta_5 & 0 & 0 \\ 0 & \zeta_5 & 0 \\ 0 & 0 & \zeta_5^3 \end{bmatrix} \quad (5.6)$$

i.e.  $R = \rho_1 \oplus \rho_1 \oplus \rho_3$ .

For the gauge field, the reasoning is exactly the same than for  $\mathbb{Z}_3$  and we get

$$A_\mu = \begin{bmatrix} A_{\mu;00} & & \\ & \ddots & \\ & & A_{\mu;44} \end{bmatrix} \quad (5.7)$$

where each block  $A_{\mu;ii}$  is of size  $N_i \times N_i$ . The projected gauge group is

$$G_{\text{proj}} = \text{U}(N_0) \times \text{U}(N_1) \times \text{U}(N_2) \times \text{U}(N_3) \times \text{U}(N_4). \quad (5.8)$$

The scalar fields transform as

$$X_{i\alpha_i, j\beta_j}^m \rightarrow R(g)^m_{n\rho_i}(g) X_{i\alpha_i, j\beta_j}^n \rho_j(g)^{-1} \quad (5.9)$$

so the invariant configurations must satisfy

$$X_{i\alpha_i, j\beta_j}^m = R(g)^m {}_n\zeta_4^{i-j} \rho_i(g) X_{i\alpha_i, j\beta_j}^n = \begin{cases} \zeta^{i-j+1} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 0, 1, 2, 3 \\ \zeta^{i-j+3} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 4, 5. \end{cases} \quad (5.10)$$

and are therefore of the form

$$X^m = \begin{bmatrix} 0 & X_{01}^m & 0 & 0 & 0 \\ 0 & 0 & X_{12}^m & 0 & 0 \\ 0 & 0 & 0 & X_{23}^m & 0 \\ 0 & 0 & 0 & 0 & X_{34}^m \\ X_{40}^m & 0 & 0 & 0 & 0 \end{bmatrix} \quad (5.11)$$

for  $m = 0, 1, 2, 3$  and

$$X^m = \begin{bmatrix} 0 & 0 & 0 & X_{03}^m & 0 \\ 0 & 0 & 0 & 0 & X_{14}^m \\ X_{20}^m & 0 & 0 & 0 & 0 \\ 0 & X_{31}^m & 0 & 0 & 0 \\ 0 & 0 & X_{42}^m & 0 & 0 \end{bmatrix} \quad (5.12)$$

for  $m = 4, 5$ . Gauge anomaly cancellation imposes

$$-N_{i-2} + N_{i+2} - 2N_{i+1} + 2N_{i-1} = 0 \quad (5.13)$$

for  $i = 0, 1, 2, 3, 4$  which implies that

$$N_0 = N_1 = N_2 = N_3 = N_4. \quad (5.14)$$

Now if we choose the representation  $(1, 2, 2)$ , i.e.

$$R(g) = \begin{bmatrix} \zeta_5 & 0 & 0 \\ 0 & \zeta_5^2 & 0 \\ 0 & 0 & \zeta_5^2 \end{bmatrix} \quad (5.15)$$

we get instead that the invariant field configurations must satisfy

$$X_{i\alpha_i, j\beta_j}^m = R(g)^m {}_n\zeta_4^{i-j} \rho_i(g) X_{i\alpha_i, j\beta_j}^n = \begin{cases} \zeta^{i-j+1} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 0, 1 \\ \zeta^{i-j+2} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 2, 3, 4, 5. \end{cases} \quad (5.16)$$

and are therefore of the form

$$X^m = \begin{bmatrix} 0 & X_{01}^m & 0 & 0 & 0 \\ 0 & 0 & X_{12}^m & 0 & 0 \\ 0 & 0 & 0 & X_{23}^m & 0 \\ 0 & 0 & 0 & 0 & X_{34}^m \\ X_{40}^m & 0 & 0 & 0 & 0 \end{bmatrix} \quad (5.17)$$

for  $m = 0, 1$  and

$$X^m = \begin{bmatrix} 0 & 0 & X_{02}^m & 0 & 0 \\ 0 & 0 & 0 & X_{13}^m & 0 \\ 0 & 0 & 0 & 0 & X_{24}^m \\ X_{30}^m & 0 & 0 & 0 & 0 \\ 0 & X_{41}^m & 0 & 0 & 0 \end{bmatrix} \quad (5.18)$$

for  $m = 2, 3, 4, 5$ . Gauge anomaly cancellation imposes

$$-2N_{i+2} + 2N_{i-2} - N_{i+1} + N_{i-1} = 0(*) \quad (5.19)$$

for  $i = 0, 1, 2, 3, 4$  which implies that

$$N_0 = N_1 = N_2 = N_3 = N_4. \quad (5.20)$$

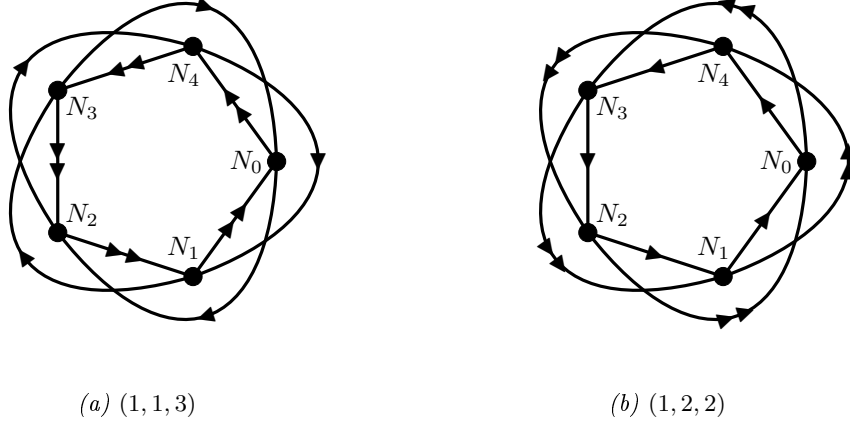


Figure 3: Quivers of the  $\mathbb{C}^3/\mathbb{Z}_5$  daughter theories.

Note that, even we thought that (1, 1, 3) and (1, 2, 2) were different representation, that are actually equivalent in the sense that  $(1, 1, 3) + (1, 1, 3) = (2, 2, 1)$ . The two quivers in fig. 3 should therefore be the same. And indeed, upon further inspection, the two are the same. To see this, we can rename the vertices in the second graph as  $N_1 \rightarrow N_3, N_2 \rightarrow N_1, N_3 \rightarrow N_4, N_4 \rightarrow N_2$ . This renaming defines the bijection between the two graphs. A more pragmatic way to see this is to look at the adjacency matrices:

$$a_{(1,1,3)} = \begin{bmatrix} 0 & 0 & 1 & 0 & 2 \\ 2 & 0 & 0 & 1 & 0 \\ 0 & 2 & 0 & 0 & 1 \\ 1 & 0 & 2 & 0 & 0 \\ 0 & 1 & 0 & 2 & 0 \end{bmatrix} \quad a_{(1,2,2)} = \begin{bmatrix} 0 & 0 & 0 & 2 & 1 \\ 1 & 0 & 0 & 0 & 2 \\ 1 & 1 & 0 & 0 & 0 \\ 0 & 2 & 1 & 0 & 0 \\ 0 & 0 & 2 & 1 & 0 \end{bmatrix}. \quad (5.21)$$

Changing the names of the vertices is equivalent to swapping line and columns. For example,

### 5.1.3 | General $\mathbb{Z}_n$

Let us now consider the general action

$$R(g) = \begin{bmatrix} \zeta_n^a & 0 & 0 \\ 0 & \zeta_n^b & 0 \\ 0 & 0 & \zeta_n^c \end{bmatrix} \quad (5.22)$$

of  $\mathbb{Z}_n$  on  $\mathbb{C}^3$ , where  $(a, b, c)$  one of the representation that we studied before. In particular, recall that  $a + b + c = n$ . Following the same reasoning than before, we get that the gauge field of of the form

$$A_\mu = \begin{bmatrix} A_{\mu;00} & & \\ & \ddots & \\ & & A_{\mu;n-1,n-1} \end{bmatrix}. \quad (5.23)$$

Invariant scalar field configurations transform as

$$X_{i\alpha_i, j\beta_j}^m = \begin{cases} \zeta_n^{i-j+a} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 0, 1 \\ \zeta_n^{i-j+b} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 2, 3 \\ \zeta_n^{i-j+c} X_{i\alpha_i, j\beta_j}^m, & \text{if } m = 4, 5 \end{cases} \quad (5.24)$$

so

$$X_{j-a, j}^{0,1} \in (\mathbb{N}_j, \overline{\mathbb{N}}_{j-a}), \quad (5.25)$$

$$X_{j-b,j}^{2,3} \in (\mathbf{N}_j, \overline{\mathbf{N}}_{j-b}), \quad (5.26)$$

$$X_{j-c,j}^{4,5} \in (\mathbf{N}_j, \overline{\mathbf{N}}_{j-c}) \quad (5.27)$$

are the only possible non-vanishing components. This allows us to quickly draw all the possible quivers for a given  $n$ . Once again, the difficulty is only computational, not conceptual. This can therefore easily be implemented into Mathematica and we get the quiver for any  $n$  and any representation  $(a, b, c)$ .

## 5.2 | $S = \mathbb{C}^3/\Delta(3n^2), \Delta(6n^2)$

### 5.3 | $S = \mathbb{C}^3/\Sigma_{36 \times 3}, \Sigma_{60 \times 3}, \Sigma_{168 \times 3}, \Sigma_{216 \times 3}, \Sigma_{360 \times 3}$

### 5.4 | $S = \mathbb{C}^3/(\mathbb{Z}_m \times \mathbb{Z}_n)$

#### 5.4.1 | Representations of $\mathbb{Z}_m \times \mathbb{Z}_n$

We consider the group  $\mathbb{Z}_m \times \mathbb{Z}_n$ . Let us denote by  $\{\mu_i\}_{i=0,\dots,m-1}$  and  $\{\sigma_j\}_{j=0,\dots,n-1}$  two complete set of irreducible representation of  $\mathbb{Z}_m$  and  $\mathbb{Z}_n$  respectively, with

$$\mu(g_1) = \zeta_m^i \quad (5.28)$$

$$\sigma_j(g_2) = \zeta_n^j \quad (5.29)$$

where  $g_1$  and  $g_2$  are the generating elements of  $\mathbb{Z}_m$  and  $\mathbb{Z}_n$  respectively.  $\mathbb{Z}_m \times \mathbb{Z}_n$  is of order  $mn$  and possesses the same number of equivalency classes (abelian). It has therefore  $mn$  irreducible representations. Since the group is abelian, they are all of dimension 1. Let us denote by  $\{T_k\}_{k=0,\dots,m+n-1}$  a complete set of irreducible representations. Since we have a product group, for any  $k$  there exists indices  $i(k)$  and  $j(k)$  such that  $T_k = \mu_{i(k)} \otimes \sigma_{j(k)}$ . We choose the indices  $i(k)$  and  $j(k)$  such that

$$\begin{aligned} T_0 &= \mu_0 \otimes \mu_0, \\ T_1 &= \mu_0 \otimes \mu_1, \\ T_2 &= \mu_0 \otimes \mu_2, \\ &\vdots \\ T_{n-1} &= \mu_0 \otimes \mu_{n-1}, \\ T_n &= \mu_1 \otimes \mu_0, \\ &\vdots \\ T_{2n-1} &= \mu_1 \otimes \mu_{n-1}, \\ &\vdots \\ T_{mn-1} &= \mu_{m-1} \otimes \mu_{n-1}. \end{aligned}$$

That is, we take

$$\begin{cases} i(k) = \lfloor k/n \rfloor \\ j(k) = k \bmod n \end{cases} \Leftrightarrow k = i(k)n + j(k). \quad (5.30)$$

Note that this is simply a dictionary between a line notation and a matrix notation and that it is indeed a bijection  $k \Leftrightarrow i, j$ . In this way, we can proceed to similar manipulations than before, were we used line notations.

Any representation  $R$  of  $\mathbb{Z}_m \times \mathbb{Z}_n$  can be decomposed as

$$R = \bigoplus_{i,j} N_k T_k = \bigoplus_{i,j} N_{ij} (\mu_i \otimes \sigma_j) \quad (5.31)$$

with  $N_k = N_{i(k)j(k)}$ . We must have  $\sum_k N_k = 3$ . In other words, any representation of  $\mathbb{Z}_m \times \mathbb{Z}_n$  on  $\mathbb{C}^3$  is equivalent to

$$R(g_1, g_2) = [(\mu_a \otimes \sigma_{a'}) \oplus (\mu_{b'} \otimes \sigma_{b'}) \oplus (\mu_c \otimes \sigma_{c'})](g_1, g_2) = \begin{bmatrix} \xi_m^a \xi_n^{a'} & 0 & 0 \\ 0 & \xi_m^b \xi_n^{b'} & 0 \\ 0 & 0 & \xi_m^c \xi_n^{c'} \end{bmatrix}. \quad (5.32)$$

The determinant condition is

$$\xi_m^{a+b+c} = \xi_n^{-a'-b'-c'} \quad (5.33)$$

$$\Leftrightarrow (a+b+c) \bmod m = (a'+b'+c') \bmod n \quad (5.34)$$

$$R(g_1, g_2) = \begin{bmatrix} \xi_m & 0 & 0 \\ 0 & \xi_n & 0 \\ 0 & 0 & \xi_m^{-1} \xi_n^{-1} \end{bmatrix}. \quad (5.35)$$

### 5.4.2 | Projection

Let us start by the gauge field. We consider a unitary representation  $(\rho, \mathbb{C}^N)$  of  $\mathbb{Z}_m \times \mathbb{Z}_n$  on  $\mathbb{C}^n$  and decompose it as

$$\rho = \bigoplus_{i,j} N_k T_k \quad (5.36)$$

such that

$$\rho(g) = \left[ \begin{array}{cccccc} T_0(g) & & & & & 0 \\ & \ddots & & & & \\ & & T_0(g) & & & \\ \vdots & & & \ddots & & \vdots \\ & & & & T_{mn-1}(g) & \\ & & & & & \ddots \\ 0 & & & & & T_{mn-1}(g) \end{array} \right] \left\{ \begin{array}{l} N_0 \\ \\ \\ \\ N_{mn-1} \end{array} \right\}. \quad (5.37)$$

We can now use our usual bi-index notation  $A_{\mu; k\alpha_k, l\beta_l}$  with  $k, l = 0, \dots, mn-1$  and  $\alpha_k, \beta_k = 0, \dots, N_k$  but instead it is more convenient to come back to our matrix notation by writing the block  $A_{\mu; k, l}$  as  $A_{\mu; i(k)j(k), i(l)j(l)}$  that we simply denote by  $A_{\mu; ij, i'j'}$  with  $i, i' \in \{0, \dots, n-1\}$  and  $j, j' \in \{0, \dots, m-1\}$ . So for  $m=2, n=3$  for example, the link between the two notations is

$$\begin{bmatrix} A_{\mu;00} & A_{\mu;01} & A_{\mu;02} & A_{\mu;03} & A_{\mu;04} & A_{\mu;05} \\ A_{\mu;10} & A_{\mu;11} & A_{\mu;12} & A_{\mu;13} & A_{\mu;14} & A_{\mu;15} \\ A_{\mu;20} & A_{\mu;21} & A_{\mu;22} & A_{\mu;23} & A_{\mu;24} & A_{\mu;25} \\ A_{\mu;30} & A_{\mu;31} & A_{\mu;32} & A_{\mu;33} & A_{\mu;34} & A_{\mu;35} \\ A_{\mu;40} & A_{\mu;41} & A_{\mu;42} & A_{\mu;43} & A_{\mu;44} & A_{\mu;45} \\ A_{\mu;50} & A_{\mu;51} & A_{\mu;52} & A_{\mu;53} & A_{\mu;54} & A_{\mu;55} \end{bmatrix} = \begin{bmatrix} A_{\mu;00,00} & A_{\mu;00,01} & A_{\mu;00,02} \\ A_{\mu;01,00} & A_{\mu;01,01} & A_{\mu;01,02} \\ A_{\mu;02,00} & A_{\mu;02,01} & A_{\mu;02,02} \\ A_{\mu;10,00} & A_{\mu;10,01} & A_{\mu;10,02} \\ A_{\mu;11,00} & A_{\mu;11,01} & A_{\mu;11,02} \\ A_{\mu;12,00} & A_{\mu;12,01} & A_{\mu;12,02} \end{bmatrix} \begin{bmatrix} A_{\mu;00,10} & A_{\mu;00,11} & A_{\mu;00,12} \\ A_{\mu;01,10} & A_{\mu;01,11} & A_{\mu;01,12} \\ A_{\mu;02,10} & A_{\mu;02,11} & A_{\mu;02,12} \\ A_{\mu;10,10} & A_{\mu;10,11} & A_{\mu;10,12} \\ A_{\mu;11,10} & A_{\mu;11,11} & A_{\mu;11,12} \\ A_{\mu;12,10} & A_{\mu;12,11} & A_{\mu;12,12} \end{bmatrix} \quad (5.38)$$

So instead of considering  $A_\mu$  to be a single  $mn \times mn$  matrix of element  $A_{\mu;kl}$ , where  $A_{\mu;kl}$  are  $N_k \times N_l$  matrices, we consider it as an  $m \times m$  where each element  $A_{\mu;ii'}$  (line  $i$  column  $i'$ ) is itself an  $n \times n$  matrices with elements  $A_{\mu;ij, i'j'}$  (line  $j$  column  $j'$ ), as shown above.

Using these notations, the gauge field transforms as

$$A_{\mu; ij, i'j'} \mapsto (\mu_i(g) \otimes \sigma_j(g)) A_{\mu; ij, i'j'} (\mu_{i'}(g) \otimes \sigma_{j'}(g))^{-1} = \zeta_m^{i-i'} \zeta_n^{j'-j} A_{\mu; ij, i'j'} \quad (5.39)$$

so invariant configurations can possess non-vanishing components  $A_{\mu; ij, i'j'}$  only if

$$\zeta_m^{i-i'} = \zeta_n^{j'-j} \quad (5.40)$$

$$\Leftrightarrow (i-i') \bmod m = (j'-j) \bmod n \quad (5.41)$$

$$\Leftrightarrow j' = j + |i' - i|. \quad (5.42)$$

This means that the submatrices  $A_{\mu; ii'}$  has an off-diagonal block form with offset  $|i' - i|$ . Once again, for a general  $n$ , the difficulty is only computational, not conceptual. This can therefore easily be implemented into Mathematica and we get the form of the gauge field for any  $n$ .

For the scalar fields, we have

$$X_{ij, i'j'}^m \mapsto R(g)^m_n (\mu_i(g) \otimes \sigma_j(g)) X_{ij, i'j'}^n (\mu_{i'}(g) \otimes \sigma_{j'}(g))^{-1} \quad (5.43)$$

$$= \begin{cases} \zeta_m^{i-i'+1} \zeta_n^{j-j'} X_{ij,i'j'}^m, & m = 0, 1 \\ \zeta_m^{i-i'} \zeta_n^{j-j'+1} X_{ij,i'j'}^m, & m = 2, 3 \\ \zeta_m^{i-i'-1} \zeta_n^{j-j'-1} X_{ij,i'j'}^m, & m = 4, 5. \end{cases} \quad (5.44)$$

So a configuration is invariant if and only if the only non-vanishing component satisfy

$$m = 0, 1 : (i - i' + 1) \bmod m = (j' - j) \bmod n \quad (5.45)$$

$$m = 2, 3 : (i - i') \bmod m = (j' - j - 1) \bmod n \quad (5.46)$$

$$m = 4, 5 : (i - i' - 1) \bmod m = (j' - j + 1) \bmod n. \quad (5.47)$$

## 6 | Correspondence between gauge theory and singularity

Above, we presented all the possible orbifold constructions of supersymmetric quiver gauge theories in four dimensions. We started from quotienting the transverse space and we found the corresponding (supersymmetric) gauge theory. In other words, we started from the singularity and found the gauge theory. We can therefore consider that the orbifold singularities are understood. However, not all singularities are orbifold one, such as the conifold for example. We can then ask ourselves how to obtain the gauge theory for more general singularities than the orbifold ones. Is a general approach possible? On the other hand, we can also study the converse question; is it possible to obtain the singularity from the gauge theory? And if it is, how so? In general we will see that there is a bijection between the four-dimensional supersymmetric worldvolume gauge theory and the Calabi-Yau singularity. We now detail this bijection.

### 6.1 | From gauge theory to singularity: forward algorithm

We start with the simplest question: how to recover the singularity from the gauge theory? We already mentioned that the vacuum parameter space of the scalar fields of the gauge theory is the so-called moduli space, denoted  $\mathcal{M}$ . Because our D3-brane is a point in the Calabi-Yau threefold, the vacuum moduli space  $\mathcal{M}$  is the affine coordinates of the Calabi-Yau singularity  $S$ .

$$6.1.1 \quad S = \mathbb{C} \times \mathbb{C}^2 / \mathbb{Z}_n$$

#### 6.1.2 | The conifold

### 6.2 | From singularity to gauge theory: inverse algorithm

Mathematically, a quiver gauge theory is a representation of a finite quiver with relations. The labels are  $\{N_i \in \mathbb{Z}_+\}$ , they correspond to the dimension of the vector space  $\{V_i\}$ . The gauge group is  $\prod_i \text{SU}(N_i)$ . The gauge fields are self-adjoint arrows  $\text{Hom}(V_i, V_i)$  while the matter fields are bi-fundamentals fermions/bosons and are arrows  $X_{ij} \in \text{Hom}(V_j, V_i)$ . For a quiver with adjacency matrix  $a_{ij}$ , the gauge anomaly cancellation condition can be generally expressed as

$$(a_{ij} - a_{ji})N_i = 0. \quad (6.1)$$

At last, there are some relations that arises the superpotential  $W(\{X_{ij}\})$ . The vacuum is the minima of the superpotential. In other words,

$$\frac{\partial W}{\partial X_{ij}} = 0. \quad (6.2)$$

## 7 | A note about projective representations, discrete torsion and deformations

### 7.1 | Projective representations and discrete torsion

Up until now, and in particular in all of our computations in section 3, we considered ordinary representations, i.e. linear representations. We can however use a more general representations such as projective

representations for example; representations  $\rho$  of  $\Gamma$  such that for all  $\gamma_1, \gamma_2 \in \Gamma$ ,

$$\rho(\gamma_1)\rho(\gamma_2) = A(g_1, g_2)\rho(\gamma_1\gamma_2) \quad (7.1)$$

for some factor  $A(\gamma_1, \gamma_2)$ . The  $A(\gamma_1, \gamma_2) = 1$  corresponding to linear representations. For consistency reasons, the factor  $A(g_1, g_2)$  cannot be completely arbitrary and must a cocycle condition. As a result, the possibilities for  $A(\gamma_1, \gamma_2)$  are classified by the second cohomology group  $H^2(\Gamma, \mathbb{C}^*)$ , also called *discrete torsion*. There exists a projective representation if and only if the latter does not vanish.

This new liberty, whenever admissible, provides new classes of quiver gauge theories that can be very different from the ones we considered, with no discrete torsion.

## 7.2 | Quiver gauge theories deformations and conifold

We can replace the orbifold by a conifold by deforming singular algebraic description of the orbifold with a field into a family smooth surfaces. The total space is then the conifold.

## Part II

# Beyond orbifold singularities

## 8 | Non-commutative resolutions

## 9 | Determinantal varieties as transverse spaces

### 9.1 | Basic properties of determinantal varieties

A *determinantal variety* (DV) is a space of matrices with a given upper bound on their ranks. More precisely, given  $m, n$  and  $r < \min(m, n)$ , the DV  $Y_r$  of the field  $K$  is the set of  $m \times n$  matrices over  $K$  with rank lower or equal to  $r$ :

$$Y_r \equiv \{M \in M_{m \times n}(K) | \text{rank } M \leq r\}. \quad (9.1)$$

Recall that a  $k$ -minor is the determinant of a  $k \times k$  sub-matrix and that the rank of a matrix is equal to the biggest integer such that there is a non-vanishing minor of that size. Imposing  $\text{rank } M \leq r$  is therefore equivalent to the vanishing of its  $(r+1) \times (r+1)$  minors, as it also implies the vanishing of the biggest minors. This naturally qualifies  $Y_r$  as affine varieties embedded in  $K^{mn}$ . Let us denote by  $x_{i,j}$  the independent entries of a generic matrix of  $Y_r$ , i.e.  $x_{i,j}$  are affine coordinates. The minors are therefore homogeneous polynomials of degree  $r+1$ . The ideal of  $K[x_{i,j}]$  generated by these polynomials  $(r+1) \times (r+1)$  is called the *determinantal ideal*. Homogeneity the polynomials implies that  $Y_r$  can equivalently be seen as a projective variety in  $\mathbb{A}^{mn-1}$ .

#### 9.1.1 | Computing the dimension

Let us compute the dimension of  $Y_r$  seen as an affine variety. We consider the space  $\mathbb{A}^{mn} \times \mathbf{Gr}(r, m)$ , where  $\mathbf{Gr}(r, m)$  is the Grassmannian of  $r$ -planes in an  $m$ -dimensional vector space. Let us define the subspace

$$Z_r \equiv \{(A, W) | Ax \in W \text{ for all } x \in \mathbb{A}^n\}. \quad (9.2)$$

$Y_r$  and  $Z_r$  are birationally equivalent so  $\dim Y_r = \dim Z_r$ . We want to compute  $Z_r$ . First we notice that  $Z_r$  is a vector bundle over  $\mathbf{Gr}(r, m)$  and we denote it by  $Z_r \xrightarrow{\pi_1} \mathbf{Gr}(r, m)$ . Now, over the Grassmannian  $\mathbf{Gr}(r, m)$ , there is a tautological vector bundle that we denote by  $E_{\mathbf{Gr}} \xrightarrow{\pi_2} \mathbf{Gr}(r, m)$  whose fibers are  $\pi_2^{-1}(W) = W \cong \mathbb{R}^r$ . Finally,  $K^m$  can also be seen as a vector bundle, with fibers  $\mathbb{R}^m$ . We denote it by

$E_{K^n} \xrightarrow{\pi_3} K^n$ . From  $E_{\mathbf{Gr}}$  and  $E_{K^n}$ , we can construct<sup>2</sup> the vector bundle  $\mathrm{Hom}(E_{\mathbf{Gr}}, E_{K^n}) \xrightarrow{\pi_4} \mathbf{Gr}(r, m)$ . This vector bundle has the same base space and its fibers are  $\mathrm{Hom}(\mathbb{R}^m, \mathbb{R}^r)$  which are exactly the same as the ones of  $Z_r$ . So the two vector bundles are isomorphic:

$$Z_r \cong \mathrm{Hom}(E_{\mathbf{Gr}}, E_{K^n}). \quad (9.3)$$

Finally, since the fibers of  $\mathrm{Hom}(K^n, E_{\mathbf{Gr}})$  have dimension  $nr$ , we find

$$\dim Z_r = \dim \mathrm{Hom}(K^n, E_{\mathbf{Gr}}) = \dim \mathbf{Gr}(r, m) + nr = r(m - r) + nr. \quad (9.4)$$

Finally, we conclude that  $Y_r$  is a affine variety of dimension  $r(m - r) + nr$ .

$m$	$n$	$r$	$\dim_{\mathbb{C}} Y_r$
2	2	1	3
3	2	1	4
3	3	1	5
3	3	2	8
4	2	1	5
4	3	1	6
4	3	2	10
4	4	1	7
4	4	2	12
4	4	3	15

### 9.1.2 | Singularity

The variety  $Y_r$  is singular and the singular locus is contained in the subset of matrices with rank strictly lower than  $r$ .  $Z_r$  is a resolution (over the open set of matrices with rank exactly  $r$ , this map is an isomorphism)(<sup>\*</sup>).

### 9.1.3 | Action and syzygies

$Y_r$  naturally acts on  $G = \mathrm{GL}(m, K) \times \mathrm{GL}(n, K)$

## 9.2 | Young's lattice

*Young's lattice* is a lattice  $Y$  formed by all integer partitions ordered by inclusion of their Young tableau. It is generally used to

---

<sup>2</sup>Recall that if  $E$  and  $F$  are vector bundles over  $X$ , then we can construct a new vector bundle over  $X$ , called the Hom-bundle and denoted  $\mathrm{Hom}(E, F)$ , by defining the fiber over  $x \in X$  to be  $\mathrm{Hom}(E_x, F_x)$ .



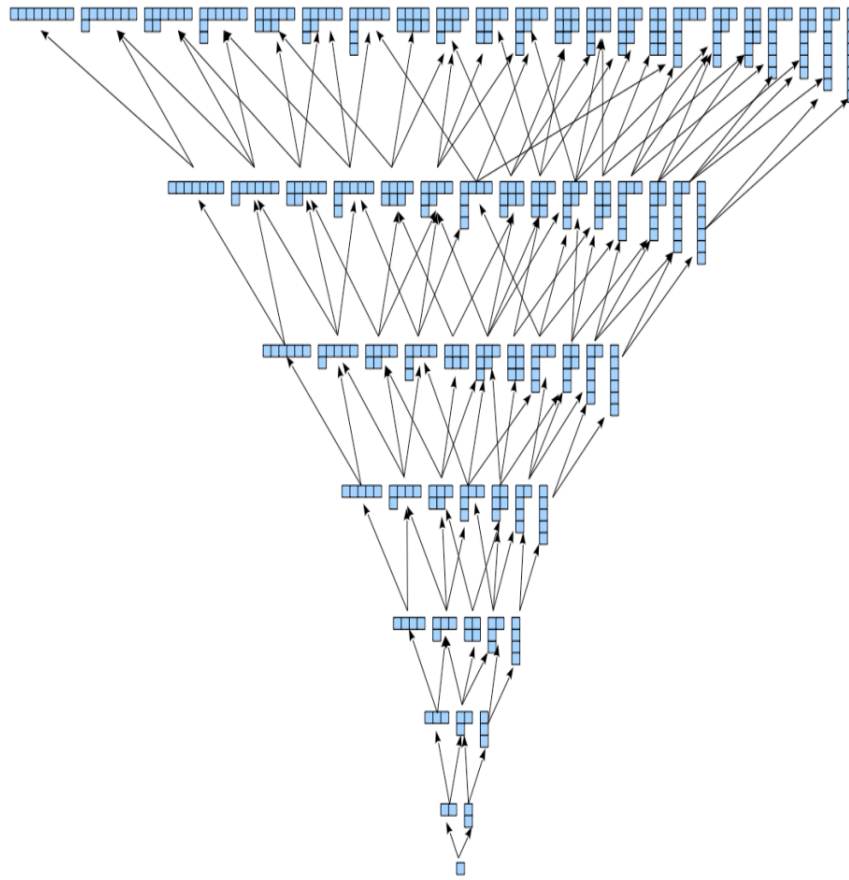


Figure 4: Young's lattice depicted as a Hasse diagram, i.e. the descendance of two tableaux is the union and the parent is the intersection.

## A | Supersymmetric Yang-Mills theory from D-branes

The dynamics of D-branes is described by the Dirac-Born-Infeld action

$$S_{\text{DBI}}[X, F] = -\frac{T_p}{g_s} \int d^{p+1}\sigma \sqrt{-\det_{0 \leq a, b \leq p} (\eta_{ab} + \partial_a X^m \partial_b X_m + 2\pi\alpha' F_{ab})}. \quad (\text{A.1})$$

The latter can be expended for slowly-varying fields, which is equivalent to passing to the field theory limit  $\alpha' \rightarrow 0$ . The resulting action is the action of a  $U(1)$  gauge theory in  $p+1$  dimensions with  $9-p$  real scalar fields. This action is exactly the same than the one we would obtain by dimensionally-reducing a pure  $U(1)$  Yang-Mills gauge theory in 10 spacetime dimensions with the identification

$$g_{\text{YM}} = g_s T_p^{-1} (2\pi\alpha')^{-2} = \frac{g_s}{\sqrt{\alpha'}} (2\pi\sqrt{\alpha'})^{p-1}. \quad (\text{A.2})$$

This construction can be generalized for multiple D-branes. It now results in a non-abelian theory. The general statement is the following:

The low-energy dynamics of  $N$  parallel, coincident  $Dp$ -branes in flat space is described in static gauge by the dimensional reduction to  $p+1$  dimensions of pure 10d  $\mathcal{N} = 1$  supersymmetric Yang-Mills theory with gauge group  $U(N)$  in ten spacetime dimensions.

Recall that the 10-dimensional action is given by

$$S_{\text{YM}} = \frac{1}{4g_{\text{YM}}^2} \int d^{10}x \left[ \text{tr}(F_{\mu\nu} F^{\mu\nu}) + 2i \text{tr}(\bar{\psi} \Gamma^\mu D_\mu \psi) \right], \quad (\text{A.3})$$

where  $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu + i[A_\mu, A_\nu]$  is the non-abelian field strength of the  $U(N)$  gauge field  $A_\mu$ ,  $D_\mu = \partial_\mu - i[A_\mu, \psi]$ ,  $\Gamma^\mu$  are  $16 \times 16$  Dirac matrices (\*), and the  $N \times N$  Hermitian fermion field  $\psi$  is a 16-component Majorana-Weyl spinor of the Lorentz group  $SO(1, 9)$  which transforms under the adjoint representation of the gauge group  $U(N)$ . On-shell, there are eight on-shell bosonic, gauge field degrees of freedom, and eight fermionic degrees of freedom, after imposition of the Dirac equation  $\not{D}\psi = \Gamma^\mu D_\mu \psi = 0$ . One can verify that this action is invariant under the supersymmetry transformations

$$\begin{aligned} \delta_\epsilon A_\mu &= \frac{i}{2} \bar{\epsilon} \Gamma_\mu \psi, \\ \delta_\epsilon \psi &= \frac{1}{2} F_{\mu\nu} [\Gamma^\mu, \Gamma^\nu] \epsilon, \end{aligned}$$

where  $\epsilon$  is an Majorana-Weyl spinor.

Using (A.3), we can construct a supersymmetric Yang-Mills gauge theory in  $p+1$  dimensions with 16 independent supercharges by dimensional reduction: we take all fields to be independent of the coordinates  $X^{p+1}, \dots, X^9$ , then the ten-dimensional gauge field  $A_\mu$  splits into a  $(p+1)$ -dimensional  $U(N)$  gauge field  $A_a$  plus  $9-p$  Hermitian scalar fields  $\Phi^m = X^m/2\pi\alpha'$  in the adjoint representation of  $U(N)$ . The  $Dp$ -brane action is thereby obtained from the dimensionality reduced field theory as

$$S_{Dp} = -\frac{T_p g_s (2\pi\alpha')^2}{4} \int d^{p+1}\sigma \text{tr} \left( F_{ab} F^{ab} + 2D_a \Phi^m D^a \Phi_m + \sum_{m \neq n} [\Phi^m, \Phi^n]^2 + \text{fermions} \right) \quad (\text{A.4})$$

where  $a, b = 0, \dots, p$ ,  $m, n = p+1, \dots, 9$ . We do not explicitly display the fermionic contributions for the moment. In conclusion, the low-energy brane dynamics is described by a supersymmetric Yang-Mills theory on the  $Dp$ -brane worldvolume which is dynamically coupled to the transverse, adjoint scalar fields  $\Phi^m$ .

The scalar potential is given by

$$V(\Phi) = \sum_{m \neq n} [\Phi^m, \Phi^n]^2. \quad (\text{A.5})$$

It is negative definite because  $[\Phi^m, \Phi^n]^\dagger = [\Phi^n, \Phi^m] = -[\Phi^m, \Phi^n]$ . A classical vacuum of the field theory defined by (A.4) corresponds to a static solution of the equations of motion whereby the potential energy of the system is minimized. It is given by the field configurations which solve simultaneously the equations  $F_{ab} = D_a \Phi^m = \psi^a = 0$  and  $V(\Phi) = 0$ . Since all term in  $V(\Phi)$  have the same sign, the equation  $V(\Phi) = 0$  is equivalent to the equation  $[\Phi^m, \Phi^n] = 0$  for all  $m, n$  and at each point in the  $(p+1)$ -dimensional worldvolume of the branes. This implies that the  $N \times N$  hermitian matrix fields  $\Phi^m$  are simultaneously diagonalizable by a gauge transformation, so that we may write

$$\Phi^m = U \begin{bmatrix} X_1^m & & 0 \\ & X_2^m & \\ & & \ddots \\ 0 & & & X_N^m \end{bmatrix} U^{-1}, \quad (\text{A.6})$$

the matrix  $U$  is independent of  $m$ . The simultaneous, real eigenvalues  $X_i^m$  give the positions of the  $N$  distinct D-branes in the  $m$ -th transverse direction. It follows that the moduli space of classical vacua for the  $(p+1)$ -dimensional field theory (A.4) is the quotient space  $(\mathbb{R}^{9-p})^N / S_N$ , where the factors of  $\mathbb{R}$  correspond to the positions of the  $N$  D $p$ -branes in the  $(9-p)$ -dimensional transverse space, and  $S_N$  is the symmetric group acting by permutations of the  $N$  coordinates  $X_i$ . The group  $S_N$  corresponds to the residual Weyl symmetry of the  $U(N)$  gauge group acting in (A.6). It represents the permutation symmetry of a system of  $N$  *indistinguishable* D-branes.

From (A.4) one can easily deduce that the masses of the fields corresponding to the off-diagonal matrix elements are given precisely by the distances  $|x_i - x_j|$  between the corresponding branes. This description means that an interpretation of the D-brane configuration in terms of classical geometry is only possible in the classical ground state of the system, whereby the matrices  $\Phi^m$  are simultaneously diagonalizable and the positions of the individual D-branes may be described through their spectrum of eigenvalues. This gives a simple and natural dynamical mechanism for the appearance of “non-commutative geometry” at short distances, where the D-branes cease to have well-defined positions according to classical geometry.

The end of this section has to be rewritten (\*).

## B | Reminder on $\mathcal{N} = 4$ super Yang-Mills theory in $D = 4$

### B.1 | Superconformal group $SU(2, 2|4)$ and its representations

Conformal transformations and supersymmetries do not commute so the presence of conformal symmetry in addition to  $\mathcal{N} = 4$  supersymmetry leads to an even larger group of symmetry known as the *superconformal group*. In the  $D = 4, \mathcal{N} = 4$  case, the superconformal group is the super group<sup>3</sup>  $SU(2, 2|4)$ . The different component of the latter are

- **Conformal symmetries:** they form the 15-dimensional subgroup  $SO(2, 4)$  and are generated by  $P_\mu, M_{\mu\nu}, K_\mu$  and  $D$ .
- **R-symmetry:** they form the 15-dimensional subgroup  $SO(6)_R$  and are generated by  $T^A$  ( $A = 1, \dots, 15$ ).
- **Poincaré supersymmetries:** they form the 16-dimensional sub group (\*) and are generated by  $Q_\alpha^I$  and  $\bar{Q}_{\dot{\alpha}}^I$ .
- **Conformal supersymmetries:** they form the 16-dimensional subgroup (\*) and are generated by  $S_{\alpha I}$  and  $\bar{S}^{\dot{\alpha} I}$ .

Conformal invariance of this theory can be seen as a consequence of the non-renormalization theorems.

### B.2 | Matter content

For  $D = 4, \mathcal{N} = 4$ , there is only one kind of supermultiplet, the vector multiplet. Therefore, from an  $\mathcal{N} = 4$  perspective, the only  $\mathcal{N} = 4$  is a pure SYM. For extended supersymmetry, it is easier to express

<sup>3</sup>Supermanifold which is also a group with smooth product and inverse maps.

it in terms of  $\mathcal{N} = 1$  superfield on  $\mathcal{N} = 1$  superspace instead of looking to construct a superspace for  $\mathcal{N} = 4$ . In this case, we can see that the  $\mathcal{N} = 4$  vector superfield can be expressed in terms of  $\mathcal{N} = 1$  representations as one vector supermultiplet and three chiral scalar supermultiplets:

$$[\mathcal{N} = 4 \text{ vector multiplet}] : V = (\lambda_\alpha, A_\mu, D) \oplus \Phi^A = (\phi^A, \psi_\alpha^A, F^A). \quad (\text{B.1})$$

with  $A = 1, 2, 3$  and

$$\phi^A = \phi_a^A T^a, \quad \psi_\alpha^A = \psi_{\alpha,a}^A T^a, \quad F^A = F_a^A T^a, \quad (\text{B.2})$$

$$\lambda^A = \lambda_a^A T^a, \quad A_\mu^A = A_{\mu,a}^A T^a, \quad D^A = F_a^A T^a, \quad (\text{B.3})$$

$$V = V_a T^a, \quad \Phi^A = \Phi_a^A T^a, \quad (\text{B.4})$$

where  $T^a$  ( $a = 1, \dots, \dim G$ ) are the generators of  $\mathfrak{g}$ . The propagating degrees of freedom are therefore a vector field, three complex scalars and four gauginos. The Lagrangian is very much constrained by  $\mathcal{N} = 4$  supersymmetry. First, the chiral superfields  $\Phi^A$  should transform in the adjoint representation of the gauge group  $G$ , since internal symmetries commute with supersymmetry. This means that all fields transform in the adjoint of  $G$ .

Moreover, there is a large R-symmetry group<sup>4</sup>:  $\text{SU}(4)_R$ . The four Weyl fermions transform in the fundamental of  $\text{SU}(4)_R$ , while the six real scalars in the two times anti-symmetric representation, which is nothing but the fundamental representation of  $\text{SO}(6)$ . The auxiliary fields are singlets under the R-symmetry group. Using  $\mathcal{N} = 1$  superfield formalism the Lagrangian reads

$$\begin{aligned} \mathcal{L}_{\text{SYM}}^{\mathcal{N}=4} = & \frac{1}{32\pi} \text{Im} \left( \tau \int d^4x \text{tr}(W^\alpha W_\alpha) \right) + \int d^2\theta d^2\bar{\theta} \text{tr} \sum_{A=1}^3 \bar{\Phi}^A e^{2gV} \Phi^A \\ & - \int d^2\theta \sqrt{2g} \text{tr} \Phi_1 [\Phi_2, \Phi_3] + \text{h.c.} \end{aligned} \quad (\text{B.5})$$

where as usual  $W_\alpha = -\frac{1}{4} \overline{D\bar{D}}(e^{-V} D_\alpha e^V)$  is the gaugino superfield. This lagrangian is indeed invariant under the superPoincaré algebra and under the gauge transformations

$$e^V \rightarrow e^{i\bar{\Lambda}} e^V e^{-i\Lambda} \quad (\text{which implies that } W_\alpha \rightarrow e^{i\Lambda} W_\alpha e^{-i\Lambda}), \quad (\text{B.6})$$

$$\Phi^A \rightarrow e^{i\Lambda} \Phi^A. \quad (\text{B.7})$$

The large  $\text{SU}(4)_R$  R-symmetry group forbids of having a superpotential. The commutator in the third term of (B.5) appears for the same reason as for the  $\mathcal{N} = 2$  Lagrangian. Notice that the choice of a single  $\mathcal{N} = 1$  supersymmetry generator breaks the full  $\text{SU}(4)_R$  R-symmetry to  $\text{SU}(3) \times \text{U}(1)_R$ . The three chiral superfields transform in the  $\mathbf{3}$  of  $\text{SU}(3)$  and have R-charge  $R = 2/3$  under the  $\text{U}(1)_R$ . It is an easy but tedious exercise to perform the integration in superspace and get an explicit expression in terms of fields. Finally, one can solve for the auxiliary fields and get an expression where only propagating degrees of freedom are present, and where  $\text{SU}(4)_R$  invariance is manifest.

### B.3 | Moduli space and dynamical phases

The scalar potential in (B.5) can be written in a rather compact form in terms of the six real scalars  $X^i$  making up the three complex scalars  $\phi^A$  and reads

$$V(X_1, \dots, X_6) = \frac{1}{2} g^2 \text{tr} \sum_{i,j=1}^6 [X_i, X_j]^2. \quad (\text{B.8})$$

The positive definite behavior of the Cartan-Killing form on the compact gauge algebra  $\mathfrak{g}$  implies that each term in the sum is positive or zero. In other words,  $V = 0$  is equivalent to

$$[X^i, X^j] = 0, \quad i, j = 1, \dots, 6. \quad (\text{B.9})$$

<sup>4</sup>The fact that the scalar fields transform under the fundamental representation of  $\text{SO}(6)$ , which is real, makes the R-symmetry group of the  $\mathcal{N} = 4$  theory being at most  $\text{SU}(4)$  and not  $\text{U}(4)$ , in fact).

This means that the potential vanishes whenever the scalar fields belong to the Cartan subalgebra of the gauge group  $G$ . At a generic point of the moduli space, the gauge group is broken to  $U(1)^r$  where  $r$  is the rank of  $\mathfrak{g}$ . This equations admit two classes of solutions:

- $\langle X^i \rangle = 0$  for all  $i = 1, \dots, 6$ . This is the *superconformal phase*. Neither the gauge symmetry nor the superconformal symmetry is broken. The physical states and operators are gauge invariant and transform under unitary representations of  $SU(2, 2|4)$ .
- $\langle X^i \rangle \neq 0$  for at least one  $i$ . This is the *spontaneously broken Coulomb phase*. The gauge algebra  $\mathfrak{g}$  is going to be broken to  $U(1)^r$ , where  $r \equiv \text{rank } \mathfrak{g}$ . The low energy behavior is then the one of  $r$  copies of  $\mathcal{N} = 4$   $U(1)$  gauge theories. Superconformal symmetry is spontaneously broken since the non-zero VEV  $\langle X^i \rangle$  sets a scale.

**$\mathcal{N} = 4$  Yang-Mills theory.** There is only one  $D = 4, \mathcal{N} = 4$  Yang-Mills theory and it contains 3  $\mathcal{N} = 1$  chiral scalar supermultiplet and 1  $\mathcal{N} = 1$  vector supermultiplet (up to  $g$  and  $\tau$ ). This theory is conformal and can be recovered from dimensional reduction of  $D = 10, \mathcal{N} = 1$  Yang-Mills on  $\mathbb{T}^6$ .

## C | Gauge anomaly

The *anomaly degree*  $A(\rho)$  of a representation  $\rho$  is defined as

$$\frac{1}{2} \text{tr}(T_a \{T_b, T_c\}) = A(\rho) d_{abc} \quad (\text{C.1})$$

where  $d_{abc}$  is an invariant symmetric tensor of the Lie algebra of  $G$ , independent of the representation. One can show that  $A(\rho^*) = -A(\rho)$  so self dual representation have  $A(\rho) = 0$  in particular. The only simple Lie groups that allow for a complex non-self-conjugate representation are  $SU(n)$  with  $n \geq 3$ . We can normalize  $d_{abc}$  such that  $A(\rho) = 1$  for the fundamental  $n$ -dimensional representations.

## D | Properties of D-branes in type II theories

The minimal irreducible representation in 10 dimensions is a Majorana-Weyl representation of dimension 8. In type II theories, we have  $\mathcal{N} = (1, 1)$  for IIA and  $\mathcal{N} = (2, 0)$  for IIB. Because of the string origin of the generators, the two supersymmetry generators  $\epsilon_L$  and  $\epsilon_R$  (Majorana-Weyl spinors) satisfy

$$\epsilon_L = \Gamma_{11} \epsilon_L, \quad \epsilon_R = \eta \Gamma_{11} \epsilon_R \quad (\text{D.1})$$

with  $\eta = +1$  for IIB and  $\eta = -1$  for IIA theory. For a  $Dp$ -brane, the supersymmetry projections is the following:

$$\epsilon_L = \Gamma_0 \dots \Gamma_p \epsilon_R. \quad (\text{D.2})$$

In other words, the supersymmetries with generators of the form

$$Q_\alpha + \Gamma_0 \dots \Gamma_p \bar{Q}_{\dot{\alpha}} \quad (\text{D.3})$$

are preserved by the  $Dp$ -brane while the one with generators of the form

$$Q_\alpha - \Gamma_0 \dots \Gamma_p \bar{Q}_{\dot{\alpha}} \quad (\text{D.4})$$

are broken. They violate the boundary conditions. Since there is the same number of generators of the form (D.3) than of the form (D.4), exactly half of the supersymmetry is broken. The idea that one spacetime direction would break one supercharge could be reasonable if supersymmetries were transforming as vectors which not the case; supercharges transform as spinors. It would also be incompatible with the T-duality because two branes of different dimensions must have the same number of unbroken supercharges if there is a T-duality relating them: the number of unbroken supercharges is the same for all dual descriptions (a necessary condition for the equivalence). And indeed, in the correct theory, that's the case. Every type II D-brane breaks half of the supercharges.

To obtain the previous relations, we start by the ones from M-theory and compactify the 11th direction, getting type IIA theory.  $\Gamma_{11}$  then plays the role of the chiral projector in 10 dimensions; the supersymmetry parameters are related by  $\epsilon_L = \frac{1}{2}(1 + \Gamma_{11})\epsilon$  and  $\epsilon_R = \frac{1}{2}(1 - \Gamma_{11})\epsilon$ . The relations for type IIB theory are then obtained by T-duality. Under a T-duality over the  $\hat{i}$  direction, the supersymmetry parameters transform as

$$\begin{aligned}\epsilon_L &\mapsto \epsilon_L, \\ \epsilon_R &\mapsto \Gamma_i \epsilon_R.\end{aligned}$$

The tension of a D $p$ -brane is given by

$$T_p = \frac{1}{(2\pi)^p g_s l_s^{p+1}}. \quad (\text{D.5})$$

This completely fixes the Newton constant: the tension of electric-magnetic duals must satisfy:

$$T_p T_{D-p-4} = \frac{2\pi}{16\pi G_D}. \quad (\text{D.6})$$

In ten dimensions, this gives  $G_{10} = 8\pi^6 g_s^2 l_s^8$ .

The dualities are defined as follows:

$$\begin{aligned}\text{S-duality} : g_s &\mapsto \frac{1}{g_s}, & l_s^2 &\mapsto g_s l_s^2, \\ \text{T-duality} : R &\mapsto \frac{l_s^2}{R}, & g &\mapsto g_s \frac{l_s}{R}.\end{aligned}$$

## E | Some finite subgroups

### E.1 | Finite subgroups of SU(2) and SL(2, ℂ)

#### E.1.1 | Finite subgroups

The first thing to recall is that every finite subgroup of SL(2, ℂ) is isomorphic to a subgroup of SU(2, ℂ) and vice-versa, so we equivalently talk about the subgroups of SU(2). The finite subgroups of SU(2), called the *binary polyhedral groups*, are the double covers of the finite subgroups of SO(3) that are called *polyhedral groups*. They simply constitute the symmetries of the Platonic solids. The groups fall into two infinite series, associated to the regular polygons, as well as three exceptional, associated with the 5 regular polyhedra: the tetrahedron (self-dual), the cube (and its dual octahedron), the icosahedron (and its dual dodecahedron).

More precisely, the finite subgroups of SL(2, ℂ) are

- $\mathbb{Z}_n$  : cyclic group of order  $n$  ( $n \geq 2$ ) generated by

$$\begin{bmatrix} \zeta_n & 0 \\ 0 & \zeta_n^{-1} \end{bmatrix} \quad (\text{E.1})$$

- $2\mathcal{D}_n$  : *binary dihedral groups* (also known as the *dicyclic group*) of order  $4n$  ( $n \geq 1$ ) generated by

$$A \equiv \begin{bmatrix} \zeta_{2n} & 0 \\ 0 & \zeta_{2n}^{-1} \end{bmatrix} \quad \text{and} \quad B \equiv \begin{bmatrix} 0 & i \\ i & 0 \end{bmatrix} \quad (\text{E.2})$$

One can show that  $A^n = B^2$  and that  $AB = BA^{-1}$  so that  $2\mathcal{D}_n = \{B^b A^a | 0 \leq b \leq 3, 0 \leq a \leq n-1\}$ . This rewriting of the most general element of the group will be useful.

- $2\mathcal{T}$  : *binary tetrahedral group* of order 24 generated by  $D_2$  and

$$C \equiv \frac{1}{\sqrt{2}} \begin{bmatrix} \zeta_8 & \zeta_8^3 \\ \zeta_8 & \zeta_8^7 \end{bmatrix} \quad (\text{E.3})$$

- $2\mathcal{O}$  : *binary octahedral group* of order 48 generated by  $\mathcal{T}$  and

$$D \equiv \begin{bmatrix} \zeta_8^3 & 0 \\ 0 & \zeta_8^5 \end{bmatrix} \quad (\text{E.4})$$

- $2\mathcal{I}$  : *binary icosahedral group* of order 120 generated by

$$E \equiv -\frac{1}{\sqrt{5}} \begin{bmatrix} \zeta_5^4 - \zeta_5 & \zeta_5^2 - \zeta_5^3 \\ \zeta_5^2 - \zeta_5^3 & \zeta_5 - \zeta_5^4 \end{bmatrix} \quad \text{and} \quad F \equiv -\frac{1}{\sqrt{5}} \begin{bmatrix} \zeta_5^2 - \zeta_5^4 & \zeta_5^4 - 1 \\ 1 - \zeta_5 & \zeta_5^3 - \zeta_5 \end{bmatrix} \quad (\text{E.5})$$

with  $\zeta_m \equiv e^{i\frac{2\pi}{m}}$  such that  $(\zeta_m)^m = 1$ . Note that the orders are all divisible by 2. This is because the center of  $SU(2)$  is  $\mathbb{Z}_2$ .

### E.1.2 | Irreducible representations

- $\mathbb{Z}_n$  has  $n$  irreducible representations. They are all 1-dimensional (since  $\mathbb{Z}_n$  is abelian) and are given by

$$\rho_k(g) = \zeta_n^k \quad (\text{E.6})$$

with  $k = 0, \dots, n-1$ .

- $2\mathcal{D}_n$  has  $n+3$  irreducible representations: 4 of dimension 1 and  $n-1$  of dimension 2. The 1-dimensional ones are given by

$n$	$\rho(A)$	$\rho(B)$	$\rho(B^b A^a)$
even	1	1	1
		-1	$(-1)^b$
	-1	1	$(-1)^a$
		-1	$(-1)^{a+b}$
odd	1	1	1
		-1	$(-1)^b$
	-1	$i$	$(-1)^a i^b$
		$-i$	$(-1)^a (-i)^b$

and the 2-dimensional ones are given binary by

$$\rho_r(A) = \begin{bmatrix} e^{i\frac{\pi}{n}r} & 0 \\ 0 & e^{-i\frac{\pi}{n}r} \end{bmatrix}$$

$$\rho_r(B) = \begin{bmatrix} 0 & (-1)^r \\ 1 & 0 \end{bmatrix}$$

with  $r = 1, \dots, n-1$ .

### E.1.3 | Character tables

conj. class repr.	$e$	$M$	$M^2$	...	$M^{n-1}$
conj. class order	1	1	1	...	1
$V_0$	1	1	1	...	1
$V_1$	1	$\zeta_n$	$\zeta_n^2$	...	$\zeta_n^{n-1}$
$V_2$	1	$\zeta_n^2$	$\zeta_n^4$	...	$\zeta_n^{2(n-1)}$
$V_3$	1	$\zeta_n^3$	$\zeta_n^6$	...	$\zeta_n^{3(n-1)}$
$\vdots$	$\vdots$	$\vdots$	$\vdots$	$\ddots$	$\vdots$
$V_{n-1}$	1	$\zeta_n^{(n-1)}$	$\zeta_n^{2(n-1)}$	...	$\zeta_n^{(n-1)^2}$
$W$	2	$2 \cos\left(\frac{2\pi}{n}\right)$	$2 \cos\left(\frac{4\pi}{n}\right)$	...	$2 \cos\left(\frac{2\pi(n-1)}{n}\right)$

Table 1: Character table of  $\mathbb{Z}_n$ .

conj. class repr.	$e$	$B^2$	$B$	$BA$	$A$	$A^2$	$\dots$	$A^{n-1}$
conj. class order	1	1	$n$	$n$	2	2	$\dots$	2
$V_0$	1	1	1	1	1	1	$\dots$	1
$V_1$	1	1	-1	-1	1	1	$\dots$	1
$V_2$	1	1 ou -1	1 ou $i$	-1 ou $-i$	-1	1	$\dots$	$(-1)^{n-1}$
$V_3$	1	1 ou -1	-1 ou $-i$	1 ou $i$	-1	1	$\dots$	$(-1)^{n-1}$
$V_4$	2	-2	0	0	$2 \cos \frac{\pi}{n}$	$2 \cos \frac{2\pi}{n}$	$\dots$	$2 \cos \frac{(n-1)\pi}{n}$
$V_5$	2	2	0	0	$2 \cos \frac{2\pi}{n}$	$2 \cos \frac{4\pi}{n}$	$\dots$	$2 \cos \frac{2(n-1)\pi}{n}$
$\vdots$	$\vdots$	$\vdots$	$\vdots$	$\vdots$	$\vdots$	$\vdots$	$\ddots$	$\vdots$
$V_{n+2}$	2	$2(-1)^{n-1}$	0	0	$2 \cos \frac{(n-1)\pi}{n}$	$2 \cos \frac{2(n-1)\pi}{n}$	$\dots$	$2 \cos \frac{(n-1)^2\pi}{n}$
$W$	2	-2	0	0	$2 \cos \left(\frac{\pi}{n}\right)$	$2 \cos \left(\frac{2\pi}{n}\right)$	$\dots$	$2 \cos \left(\frac{\pi}{n}(n-1)\right)$

Table 2: Character table of  $2\mathcal{D}_n$ .

conj. class repr.	$e$	$B^2$	$B$	$C$	$C^2$	$C^4$	$C^5$
conj. class order	1	1	6	4	4	4	4
$V_0$	1	1	1	1	1	1	1
$V_1$	2	-2	0	1	-1	-1	1
$V_2$	3	3	-1	0	0	0	0
$V_3$	2	-2	0	$e^{i\frac{2\pi}{3}}$	$-e^{i\frac{2\pi}{3}}$	$-e^{i\frac{4\pi}{3}}$	$e^{i\frac{4\pi}{3}}$
$V_3^\vee$	2	-2	0	$e^{i\frac{4\pi}{3}}$	$-e^{i\frac{4\pi}{3}}$	$-e^{i\frac{2\pi}{3}}$	$e^{i\frac{2\pi}{3}}$
$V_4$	1	1	1	$e^{i\frac{2\pi}{3}}$	$e^{i\frac{2\pi}{3}}$	$e^{i\frac{4\pi}{3}}$	$e^{i\frac{4\pi}{3}}$
$V_4^\vee$	1	1	1	$e^{i\frac{4\pi}{3}}$	$e^{i\frac{4\pi}{3}}$	$e^{i\frac{2\pi}{3}}$	$e^{i\frac{2\pi}{3}}$
$W$	2	-2	0	1	-1	-1	1

Table 3: Character table of  $2\mathcal{T}$ .

conj. class repr.	$e$	$B^2$	$B$	$C$	$C^2$	$D$	$BD$	$D^3$
conj. class order	1	1	6	8	8	6	12	6
$V_0$	1	1	1	1	1	1	1	1
$V_1$	2	-2	0	1	-1	$-\sqrt{2}$	0	$\sqrt{2}$
$V_2$	3	3	-1	0	0	1	-1	1
$V_3$	4	-4	0	-1	1	0	0	0
$V_4$	3	3	-1	0	0	-1	1	-1
$V_5$	2	-2	0	1	-1	$\sqrt{2}$	0	$-\sqrt{2}$
$V_6$	1	1	1	1	1	-1	-1	-1
$V_7$	2	2	2	-1	-1	0	0	0
$W$	2	-2	0	1	-1	$-\sqrt{2}$	0	$\sqrt{2}$

Table 4: Character table of  $2\mathcal{O}$ .



conj. class repr.	$e$	$E^2$	$E$	$F$	$F^2$	$EF$	$(EF)^2$	$(EF)^3$	$(EF)^4$
conj. class order	1	1	30	20	20	12	12	12	12
$V_0$	1	1	1	1	1	1	1	1	1
$V_1$	2	-2	0	1	-1	$\varphi^+$	$-\varphi^-$	$\varphi^-$	$-\varphi^+$
$V_2$	3	3	-1	0	0	$\varphi^+$	$\varphi^-$	$\varphi^-$	$\varphi^+$
$V_3$	4	-4	0	-1	1	1	-1	1	-1
$V_4$	5	5	1	-1	-1	0	0	0	0
$V_5$	6	-6	0	0	0	-1	1	-1	1
$V_6$	4	4	0	1	1	-1	-1	-1	-1
$V_7$	2	-2	0	1	-1	$\varphi^-$	$-\varphi^+$	$\varphi^+$	$-\varphi^-$
$V_8$	3	3	-1	0	0	$\varphi^-$	$\varphi^+$	$\varphi^+$	$\varphi^-$
$W$	2	-2	0	1	-1	$\varphi^+$	$-\varphi^-$	$\varphi^-$	$-\varphi^+$

Table 5: Character table of  $2\mathcal{I}$ , with  $\varphi^\pm \equiv (1 \pm \sqrt{5})/2$ .

## E.2 | Finite subgroups of $SU(3)$

The finite subgroups of  $SU(3)$  are

- the finite subgroups of  $SU(2)$

so there are 2 infinite series and 5 exceptional subgroups. Note that they are all divisible by 3 because the center of  $SU(3)$  is  $\mathbb{Z}_3$ .

## F | The McKay correspondence

### F.1 | Classical correspondence

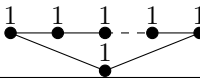
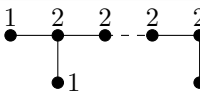
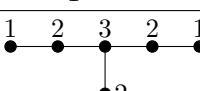
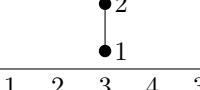
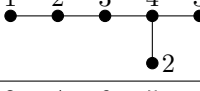
$\Gamma \subset SU(2)$	Platonic solids	McKay graph	Variety
$\mathbb{Z}_n$		 $(n \text{ nodes})$	$z^n + xy = 0$
$2\mathcal{D}_n$	$n$ -polygon	 $(n + 3 \text{ nodes})$	$x^2 + y^2z + z^{n-1} = 0$
$2\mathcal{T}$	tetrahedron	 $(7 \text{ nodes})$	$x^2 + y^3 + z^4 = 0$
$2\mathcal{O}$	cube octahedron	 $(8 \text{ nodes})$	$x^2 + y^3 + yz^3 = 0$
$2\mathcal{I}$	icosahedron dodecahedron	 $(9 \text{ nodes})$	$x^2 + y^3 + z^5 = 0$

Table 6: Binary polyhedral groups and their McKay graphs. Labels over the vertices are the dimension of the representation. We erase the arrow ends if they go in both directions and erase the label if it is equal to 1.

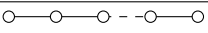
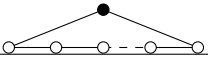
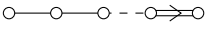
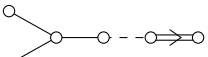
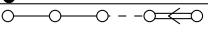
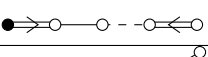
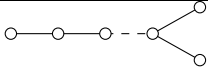
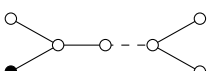
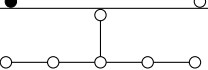
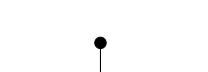
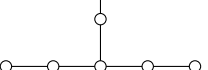
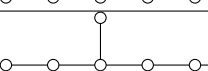

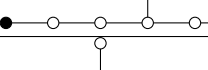
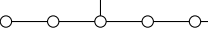
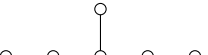
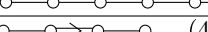
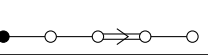
Simple Lie algebra	Simply laced	Dynkin diagram Extended Dybkin diagram
$\mathfrak{sl}(n+1, \mathbb{C}), n \geq 1$	yes	$A_n$ :  ( $n$ nodes) $\tilde{A}_n$ :  ( $n+1$ nodes)
$\mathfrak{so}(2n+1, \mathbb{R}), n \geq 2$	no	$B_n$ :  ( $n$ nodes) $\tilde{B}_n$ :  ( $n+1$ nodes)
$\mathfrak{sp}(2n, \mathbb{C}), n \geq 3$	no	$C_n$ :  ( $n$ nodes) $\tilde{C}_n$ :  ( $n+1$ nodes)
$\mathfrak{so}(2n, \mathbb{R}), n \geq 4$	yes	$D_n$ :  ( $n$ nodes) $\tilde{D}_n$ :  ( $n+1$ nodes)
$\mathfrak{e}_6$	yes	$E_6$ :  (6 nodes) $\tilde{E}_6$ :  (7 nodes)
$\mathfrak{e}_7$	yes	$E_7$ :  (7 nodes) $\tilde{E}_7$ :  (8 nodes)
$\mathfrak{e}_8$	yes	$E_8$ :  (8 nodes) $\tilde{E}_8$ :  (9 nodes)
$\mathfrak{f}_4$	no	$F_4$ :  (4 nodes) $\tilde{F}_4$ :  (5 nodes)
$\mathfrak{g}_2$	no	$G_2$ :  (2 nodes) $\tilde{G}_2$ :  (3 nodes)

Figure 5: Simple Lie algebras and their (extended) Dynkin diagrams. The first four algebras are the classical simple Lie algebras and the last five are the exceptional simple Lie algebras.

Finally, we can see the following correspondence between the extended Dynkin diagrams and the McKay graphs.

Simply Lie group	Simply laced Lie algebra	Extended Dynkin diagram	Finite subgroup of $\mathrm{SO}(3)$	Finite subgroup of $\mathrm{SU}(2)$
$\mathrm{SU}(n+1)$	$\mathfrak{sl}(n+1, \mathbb{C})$	$\widetilde{A}_n$	$\mathbb{Z}_{n+1}$	$\mathbb{Z}_{n+1}$
$\mathrm{SO}(2n), \mathrm{Spin}(2n)$	$\mathfrak{so}(2n, \mathbb{R})$	$\widetilde{D}_n$	$\mathcal{D}_{2(n-2)}$	$2\mathcal{D}_{2(n-2)}$
$E_6$	$\mathfrak{e}_6$	$\widetilde{E}_6$	$\mathcal{T}$	$2\mathcal{T}$
$E_7$	$\mathfrak{e}_7$	$\widetilde{E}_7$	$\mathcal{O}$	$2\mathcal{O}$
$E_8$	$\mathfrak{e}_8$	$\widetilde{E}_8$	$\mathcal{I}$	$2\mathcal{I}$

Figure 6: Classical McKay correspondence.

## F.2 | Geometrical McKay correspondence

## G | Algebraic geometry

### G.1 | Singularities and resolutions

A *rational map* from a variety  $X$  to another  $Y$  is a morphism from a non-empty subset  $U \subset X$  to  $Y$ . Recall that, by definition of the Zariski topology, a non-empty open subset is always dense. Concretely, a rational map can be written in coordinates using ration functions (quotient of polynomials). A *birational map* is an invertible rational map. It induces an isomorphism between two non-empty open subsets. In this case,  $X$  and  $Y$  are said to be *birationally equivalent*.

The *resolution of a singularity* of an algebraic variety  $V$  is a non-singular variety  $W$  with a proper birational map  $W \rightarrow V$ . For varieties over fields of characteristic 0, it was proven (Hironaka, 1964) that (\*).

### G.2 | Path algebras

## H | Calabi-Yau manifolds, orbifolds and crepant resolutions

Simply put, a *Calabi-Yau manifold* is a Kähler manifold with trivial canonical bundle or, equivalently, with a Kähler metric whose global holonomy is contained in  $\mathrm{SU}(n)$ . This is equivalent to having a trivial canonical bundle.

A *Calabi-Yau orbifold* is the quotient of a smooth Calabi-Yau manifold by a discrete group action which generically has fixed points. From an algebraic geometry perspective we can try to resolve the orbifold singularity. A resolution  $(X, \pi)$  of  $\mathbb{C}^n/\Gamma$  is a non-singular complex manifold  $X$  of dimension  $n$  with a proper biholomorphic map

$$\pi : X \rightarrow \mathbb{C}^n/\Gamma \quad (\text{H.1})$$

that induces a biholomorphism between dense open sets. A resolution  $(X, \pi)$  of  $\mathbb{C}^n/\Gamma$  is called a *crepant resolution*<sup>5</sup> if the canonical bundles of  $X$  and  $\mathbb{C}^n/\Gamma$  are isomorphic, i.e.

$$K_X \cong \pi^*(K_{\mathbb{C}^n/\Gamma}).$$

Since Calabi-Yau manifolds have trivial canonical bundle, to obtain a Calabi-Yau structure on  $X$  one must choose a crepant resolution of singularities.

It turns out that the amount of information we know about a crepant resolution of singularities of  $\mathbb{C}^n/\Gamma$  depends dramatically on the dimension  $n$  of the orbifold:

- $n = 2$ : a crepant always exists and is unique. Its topology is entirely described in terms of the finite group  $\Gamma$  (via the McKay correspondence).

<sup>5</sup>For a resolution of singularities we can define a notion of discrepancy. A crepant resolution is a resolution without discrepancy.

- $n = 3$ : a crepant resolution always exists but it is not unique; they are related by flops. However all the crepant resolutions have the same Euler and Betti numbers: the *stringy* Betti and Hodge numbers of the orbifold.
- $n \geq 4$ : very little is known; crepant resolution exists in rather special cases. Many singularity are terminal, which implies that they admit no crepant resolution.

## I | Graphs

A *quiver*  $Q = (E, V, t, h)$  is a piece of combinatorial data where  $E$  is a set of edges,  $V$  is a set of vertices, and  $h, t : V \rightarrow E$  are two maps that respectively assign a head and a tail to every arrow. A *morphism*  $m = (m_V, m_E)$  between two quivers  $(V, E, t, h)$  and  $(V', E', t', h')$  consists in two maps

$$m_V : V \rightarrow V' \quad (\text{I.1})$$

$$m_E : E \rightarrow E' \quad (\text{I.2})$$

such that  $m_V \circ t = t' \circ m_E$  and  $m_V \circ h = h' \circ m_E$ .

The *dimer diagram* of a quiver gauge theory is a graph whose faces represent the gauge groups, the edges represent the bi-fundamental fields and the vertices represent the superpotentials.

## J | Toric and non-toric singularities

All abelian orbifolds are toric varieties.

## K | Spacetime geometry: ALE space and orbifolds

Asymptotically locally euclidean (ALE) spaces are a particularly interesting choice of string background to probe with branes for mainly four reasons

- they are the resolution (blow-ups) of orbifolds
- they are completely classified: they fall in the ADE classification
- they only break half of the supersymmetry
- they are non-compact therefore we can study them for self-dual type II theory<sup>(\*)</sup>.

Mathematically, an ALE space is complete riemannian  $n$ -manifold  $M$  such that there exists a compact set  $K \subset M$  such that  $M \setminus K$  is diffeomorphic to  $(\mathbb{R}^n \setminus B_0(R))/G$ , where  $R \in \mathbb{R}_0^+$  is a radius and  $G \subset O(n)$  a subgroup. Additionally, it is asked that the pulled back metric on  $\mathbb{R}^n \setminus B_0(R)$  tends to the euclidean flat metric at infinity.

If one considers string theory on the orbifold  $\mathbb{R}^4/\Gamma$  where  $\Gamma$  is a finite sub group of  $SU(2)$ , massless states appear from the twisted sector. They are precisely the moduli needed to deform the theory to the one with smooth spacetime, i.e. the resolution of the orbifold. In that sense, it is said that the strings know about the metric ALE space and that it is said that strings resolve the singularity. The metric of the ALE space can be recovered if the lagrangian of the resulting field theory is explicitly known, such as for the Wess-Zumino-Witten model. However, it is often not the case.

## L | Some derivations

### L.1 |

We want to compute the sum

$$\sum_{a=1}^{\lfloor n/3 \rfloor} \left\lfloor \frac{n-3a}{2} + 1 \right\rfloor = \left\lfloor \frac{n}{3} \right\rfloor + \sum_{a=1}^{\lfloor n/3 \rfloor} \left\lfloor \frac{n-3a}{2} \right\rfloor. \quad (\text{L.1})$$

Let us write  $n \in \mathbb{N}$  as  $n = 3m + r$  with  $r = 0, 1$  or  $2$  and  $m \in \mathbb{N}$ . Regardless of  $r$ , we have  $\lfloor n/3 \rfloor = m$  and

$$\sum_{a=1}^{\lfloor n/3 \rfloor} \left\lfloor \frac{n-3a}{2} \right\rfloor = \sum_{a=1}^m \left\lfloor \frac{3}{2}(m-a) + \frac{r}{2} \right\rfloor = \sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a + \frac{r}{2} \right\rfloor. \quad (\text{L.2})$$

- if  $r = 0$ , then (L.2) becomes

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a \right\rfloor = \sum_{a=0}^{m-1} a + \sum_{a=0}^{m-1} \left\lfloor \frac{a}{2} \right\rfloor = \frac{(m-1)m}{2} + \sum_{a=0}^{m-1} \left\lfloor \frac{a}{2} \right\rfloor. \quad (\text{L.3})$$

Now if  $m$  is even, we have

$$\sum_{a=0}^{m-1} \left\lfloor \frac{a}{2} \right\rfloor = 2 \sum_{a=0}^{\lfloor \frac{m-1}{2} \rfloor} a = 2 \sum_{a=0}^{\frac{m}{2}-1} a = \left( \frac{m}{2} - 1 \right) \frac{m}{2} \quad (\text{L.4})$$

and if  $m$  is odd,

$$\sum_{a=0}^{m-1} \left\lfloor \frac{a}{2} \right\rfloor = 2 \sum_{a=0}^{\lfloor \frac{m-2}{2} \rfloor} a + \left\lfloor \frac{m-1}{2} \right\rfloor = 2 \sum_{a=0}^{\frac{m-3}{2}} a + \frac{m-1}{2} = \frac{(m-1)^2}{4} \quad (\text{L.5})$$

so

$$\sum_{a=0}^{m-1} \left\lfloor \frac{a}{2} \right\rfloor = \begin{cases} \left( \frac{m}{2} - 1 \right) \frac{m}{2}, & \text{if } m \text{ is even} \\ \frac{(m-1)^2}{4}, & \text{if } m \text{ is odd} \end{cases}. \quad (\text{L.6})$$

and

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3a}{2} \right\rfloor = \begin{cases} \frac{m(3m-4)}{4}, & \text{if } m \text{ is even} \\ \frac{(m-1)(3m-1)}{4}, & \text{if } m \text{ is odd} \end{cases}. \quad (\text{L.7})$$

- if  $r = 1$ , then (L.2) becomes

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a + \frac{1}{2} \right\rfloor = \sum_{a=0}^{m-1} a + \sum_{a=0}^{m-1} \left\lfloor \frac{a+1}{2} \right\rfloor = \frac{(m-1)m}{2} + \sum_{a=0}^{m-1} \left\lfloor \frac{a+1}{2} \right\rfloor \quad (\text{L.8})$$

and

$$\sum_{a=0}^{m-1} \left\lfloor \frac{a+1}{2} \right\rfloor = \sum_{a=1}^m \left\lfloor \frac{a}{2} \right\rfloor = \sum_{a=0}^m \left\lfloor \frac{a}{2} \right\rfloor = \begin{cases} \frac{m^2}{4}, & \text{if } m \text{ is even} \\ \frac{m^2-1}{4}, & \text{if } m \text{ is odd} \end{cases} \quad (\text{L.9})$$

by (L.6) so

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a + \frac{1}{2} \right\rfloor = \begin{cases} \frac{m(3m-2)}{4}, & \text{if } m \text{ is even} \\ \frac{3m^2-2m-1}{4}, & \text{if } m \text{ is odd} \end{cases} \quad (\text{L.10})$$

- if  $r = 2$ , then (L.2) becomes

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a + 1 \right\rfloor = m + \sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a \right\rfloor. \quad (\text{L.11})$$

so

$$\sum_{a=0}^{m-1} \left\lfloor \frac{3}{2}a + 1 \right\rfloor = \begin{cases} \frac{3m^2}{4}, & \text{if } m \text{ is even} \\ \frac{3m^2+1}{4}, & \text{if } m \text{ is odd} \end{cases} \quad (\text{L.12})$$

from (L.7).

Finally, we can write  $m = 2k$  if  $m$  is even and  $m = 2k + 1$  if  $m$  is odd in order to distinguish the six different cases. We get

$$a(n) \equiv \sum_{a=1}^{\lfloor n/3 \rfloor} \left\lfloor \frac{n-3a}{2} + 1 \right\rfloor = \begin{cases} 2k + \frac{2k(6k-4)}{4}, & \text{if } n = 6k \\ 2k + \frac{2k(6k-2)}{4}, & \text{if } n = 6k + 1, \\ 2k + \frac{12k^2}{4}, & \text{if } n = 6k + 2, \\ (2k+1) + \frac{2k(6k+2)}{4}, & \text{if } n = 6k + 3, \\ (2k+1) + \frac{3(2k+1)^2 - 2(2k+1) - 1}{4}, & \text{if } n = 6k + 4, \\ (2k+1) + \frac{3(2k+1)^2 + 1}{4}, & \text{if } n = 6k + 5 \end{cases} \quad (\text{L.13})$$

$$= \begin{cases} 3k^2, & \text{if } n = 6k \\ 3k^2 + k, & \text{if } n = 6k + 1, \\ 3k^2 + 2k, & \text{if } n = 6k + 2, \\ 3k^2 + 3k + 1, & \text{if } n = 6k + 3, \\ 3k^2 + 4k + 1, & \text{if } n = 6k + 4, \\ 3k^2 + 5k + 2, & \text{if } n = 6k + 5 \end{cases} \quad (\text{L.14})$$

Starting from  $n = 1$ , the first value of this sequence is : 0, 0, 1, 1, 2, 3, 4, 5, 7, 8, 10, 12, ... Upon further analysis, this corresponds to the sequence [A001399](#), that have several interpretations:

- the number of partitions of  $n$  into at most 3 parts. This makes sense with our initial problem: finding all the  $a, b, c$ 's such that  $a + b + c = n$ ,
- the number of connected graphs with 3 nodes and  $n$  edges (where multiple edges between the same nodes are allowed),
- the number of non-negative solutions to  $b + 2c + 3d = n$ ,

as well as many others. Finally, we note that we can simply write

$$a(n) = \text{round} \left( \frac{n^2}{12} \right). \quad (\text{L.15})$$

## M | References guide

- He: review: [He04], thesis: [He02]
- Orbifold construction for  $\Gamma \subset \mathrm{SU}(2)$ : type  $A$ : [DM96], type  $D$  and  $E$ : [JM97]
- Orbifold construction for  $\Gamma \subset \mathrm{SU}(3)$ : [HH99]
- Formalization of projection to daughter theories: [LNV98], [KS98]
- On toric singularities: [Ful93], [Oda88]
- Forward algorithm for toric singularities: developments: [Asp94], [DGM97], [MP98], [Bea+00], [DD00], formalization: [FHH01a], [FHH01b], [Fen+02]
- Toric diagrams, dimer diagrams and Higgsing: [AFP08]
- Fractional branes: [Arg+08]
- Quiver representations and quiver varieties: [Kir16]

# List of Markers

1	.....	3
2	.....	4
3	.....	5
4	.....	6
5	.....	7
6	.....	7
7	.....	10
8	.....	15
9	.....	15
10	.....	16
11	.....	16
12	.....	16
13	.....	17
14	.....	23
15	.....	25
16	.....	26
17	.....	26
18	.....	26
19	.....	34
20	.....	35



## References

- [He04] Yang-Hui He. *Lectures on D-branes, Gauge Theories and Calabi-Yau Singularities*. 2004. arXiv: hep-th/0408142 [hep-th].
- [He02] Yang-Hui He. *On Algebraic Singularities, Finite Graphs and D-Brane Gauge Theories: A String Theoretic Perspective*. 2002. DOI: 10.48550/ARXIV.HEP-TH/0209230. URL: <https://arxiv.org/abs/hep-th/0209230>.
- [DM96] Michael R. Douglas and Gregory Moore. *D-branes, Quivers, and ALE Instantons*. 1996. arXiv: hep-th/9603167 [hep-th].
- [JM97] Clifford V. Johnson and Robert C. Myers. “Aspects of type IIB theory on asymptotically locally Euclidean spaces”. In: *Phys. Rev. D* 55 (10 May 1997), pp. 6382–6393. DOI: 10.1103/PhysRevD.55.6382.
- [HH99] Amihay Hanany and Yang-Hui He. “Non-abelian finite gauge theories”. In: *Journal of High Energy Physics* 1999.02 (Feb. 1999), pp. 013–013. DOI: 10.1088/1126-6708/1999/02/013. URL: <https://doi.org/10.1088/1126-6708/1999/02/013>.
- [LNV98] Albion Lawrence, Nikita Nekrasov, and Cumrun Vafa. “On conformal field theories in four dimensions”. In: *Nuclear Physics B* 533.1-3 (Nov. 1998), pp. 199–209. ISSN: 0550-3213. DOI: 10.1016/s0550-3213(98)00495-7.
- [KS98] Shamit Kachru and Eva Silverstein. “4D Conformal Field Theories and Strings on Orbifolds”. In: *Physical Review Letters* 80.22 (June 1998), pp. 4855–4858. ISSN: 1079-7114. DOI: 10.1103/physrevlett.80.4855.
- [Ful93] William Fulton. *Introduction to toric varieties*. Princeton, N.J: Princeton University Press, 1993. ISBN: 9780691000497.
- [Oda88] Tadao Oda. *Convex bodies and algebraic geometry : an introduction to the theory of toric varieties*. Berlin: Springer, 1988. ISBN: 978-3-642-72549-4.
- [Asp94] P. Aspinwall. *Resolution of Orbifold Singularities in String Theory*. 1994. DOI: 10.48550/ARXIV.HEP-TH/9403123. URL: <https://arxiv.org/abs/hep-th/9403123>.
- [DGM97] Michael R. Douglas, Brian R. Greene, and David R. Morrison. “Orbifold resolution by D-branes”. In: *Nuclear Physics B* 506.1-2 (Nov. 1997), pp. 84–106. DOI: 10.1016/s0550-3213(97)00517-8. URL: [https://doi.org/10.1016/s0550-3213\(97\)00517-8](https://doi.org/10.1016/s0550-3213(97)00517-8).
- [MP98] David R. Morrison and M. Ronen Plesser. “Non-Spherical Horizons, I”. In: (1998). DOI: 10.48550/ARXIV.HEP-TH/9810201. URL: <https://arxiv.org/abs/hep-th/9810201>.
- [Bea+00] Chris Beasley et al. “D3-branes on partial resolutions of abelian quotient singularities of Calabi-Yau threefolds”. In: *Nuclear Physics B* 566.3 (Feb. 2000), pp. 599–641. DOI: 10.1016/s0550-3213(99)00646-x. URL: [https://doi.org/10.1016/s0550-3213\(99\)00646-x](https://doi.org/10.1016/s0550-3213(99)00646-x).
- [DD00] Duiliu-Emanuel Diaconescu and Michael R. Douglas. *D-branes on Stringy Calabi-Yau Manifolds*. 2000. DOI: 10.48550/ARXIV.HEP-TH/0006224. URL: <https://arxiv.org/abs/hep-th/0006224>.
- [FHH01a] Bo Feng, Amihay Hanany, and Yang-Hui He. “D-brane gauge theories from toric singularities and toric duality”. In: *Nuclear Physics B* 595.1-2 (Feb. 2001), pp. 165–200. DOI: 10.1016/s0550-3213(00)00699-4. URL: [https://doi.org/10.1016/s0550-3213\(00\)00699-4](https://doi.org/10.1016/s0550-3213(00)00699-4).
- [FHH01b] Bo Feng, Amihay Hanany, and Yang-Hui He. “Phase structure of D-brane gauge theories and toric duality”. In: *Journal of High Energy Physics* 2001.08 (Aug. 2001), pp. 040–040. DOI: 10.1088/1126-6708/2001/08/040. URL: <https://doi.org/10.1088/1126-6708/2001/08/040>.
- [Fen+02] Bo Feng et al. “Symmetries of Toric Duality”. In: *Journal of High Energy Physics* 2002.12 (Dec. 2002), pp. 076–076. DOI: 10.1088/1126-6708/2002/12/076. URL: <https://doi.org/10.1088/1126-6708/2002/12/076>.

- [AFP08] Riccardo Argurio, Gabriele Ferretti, and Christoffer Petersson. “Instantons and toric quiver gauge theories”. In: *Journal of High Energy Physics* 2008.07 (July 2008), pp. 123–123. DOI: 10.1088/1126-6708/2008/07/123. URL: <https://doi.org/10.1088/1126-6708/2008/07/123>.
- [Arg+08] Riccardo Argurio et al. “Gauge/gravity duality and the interplay of various fractional branes”. In: *Physical Review D* 78.4 (Aug. 2008). DOI: 10.1103/physrevd.78.046008. URL: <https://doi.org/10.1103/physrevd.78.046008>.
- [Kir16] Alexander Kirillov. *Quiver representations and quiver varieties*. Providence, Rhode Island: American Mathematical Society, 2016. ISBN: 978-1-4704-2307-0.
- [CV20] Andrés Collinucci and Roberto Valandro. “A string theory realization of special unitary quivers in 3 dimensions”. In: *Journal of High Energy Physics* 2020.11 (Nov. 2020). ISSN: 1029-8479. DOI: 10.1007/jhep11(2020)157.
- [Lin12] Ulf Lindström. *Supersymmetric Sigma Model geometry*. 2012. arXiv: 1207.1241 [hep-th].
- [Hit+87] Nigel J. Hitchin et al. “Hyperkahler Metrics and Supersymmetry”. In: *Commun. Math. Phys.* 108 (1987), p. 535. DOI: 10.1007/BF01214418.
- [NST03] Lei Ni, Yuguang Shi, and Luen-Fai Tam. “Ricci Flatness of Asymptotically Locally Euclidean Metrics”. In: *Transactions of the American Mathematical Society* 355.5 (2003), pp. 1933–1959. ISSN: 00029947.
- [Ber21] Matteo Bertolini. *Introduction to Supersymmetry*. 2021. URL: <https://people.sissa.it/~bertmat/teaching.htm>.
- [HW97] Amihay Hanany and Edward Witten. “Type IIB superstrings, BPS monopoles, and three-dimensional gauge dynamics”. In: *Nuclear Physics B* 492.1–2 (May 1997), pp. 152–190. ISSN: 0550-3213. DOI: 10.1016/s0550-3213(97)80030-2.
- [Arg98] Riccardo Argurio. “Brane physics in M theory”. PhD thesis. Brussels U., 1998. arXiv: hep-th/9807171.
- [DF02] Eric D’Hoker and Daniel Z. Freedman. *Supersymmetric Gauge Theories and the AdS/CFT Correspondence*. 2002. arXiv: hep-th/0201253 [hep-th].