Instanton Calculus In R-R Background And The Topological String

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This talk is based on


It of course builds over a vast literature. The few references scattered on the slides are by no means intended to be exhaustive. I apologize for the many relevant ones which will be missing.
Plan of the talk

1. Introduction
2. Microscopic string description of $\mathcal{N} = 2$ SYM
3. Instanton calculus by mixed string diagrams
4. Deformation from a graviphoton background
5. Relation to topological strings on CY
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Introduction
$\mathcal{N} = 2$ SYM and Seiberg-Witten solution

$\mathcal{N} = 2$ SYM theories in $d = 4$: an important test-bed for non-perturbative physics

- Seiberg-Witten: exact expression of the prepotential $\mathcal{F}(a)$ governing the low energy dynamics on the Coulomb branch using duality and monodromy properties; this involves an auxiliary Riemann surface

- “Geometrical engineering” construction [Kachru et al 1995, Katz et al 1996]: SW solution $\leftrightarrow$ Type IIB string theory on a “local” CY manifold $\mathcal{M}$ whose geometric moduli are suitably related to the gauge theory quantities ($\Lambda, a, \ldots$)
The quest for the multi-instanton contributions

Semi-classical limit: 1-loop plus instanton contributions

\[ \mathcal{F}(a) = \frac{i}{2\pi} a^2 \log \frac{a^2}{\Lambda^2} + \sum_{k=1}^{\infty} \mathcal{F}^{(k)}(a) \]

- Important task: compute the multi-instanton contributions \( \mathcal{F}^{(k)}(a) \) within the “microscopic” description of the non-abelian gauge theory to check them against the SW solution

- Only recently fully accomplished using localization techniques to perform the integration over the moduli space of the ADHM construction of the super-instantons

The localizing deformation

Introduce a deformation of the ADHM measure on the moduli spaces exploiting the 4d chiral rotations symmetry of ADHM constraints.

- The deformed instanton partition function

\[
Z(a, \varepsilon) = \sum_k Z^{(k)}(a, \varepsilon) = \sum_k \int d\hat{M}(k)e^{-S_{\text{mod}}(a, \varepsilon; M(k))}
\]

can then be computed using localization techniques using the topological twist of its supersymmetries. One has

\[
Z(a, \varepsilon) = \exp \left( \frac{\mathcal{F}_{\text{n.p.}}(a; \varepsilon)}{\varepsilon^2} \right)
\]

\[
\lim_{\varepsilon \to 0} \mathcal{F}_{\text{n.p.}}(a; \varepsilon) = \mathcal{F}_{\text{n.p.}}(a) = \text{non-pert. part of SW prepotential}
\]
Multi-instanton calculus and topological strings

What about higher orders in the deformation parameter $\varepsilon$?

- **Nekrasov’s proposal:** terms of order $\varepsilon^{2h} \leftrightarrow$ gravitational $F$-terms in the $\mathcal{N} = 2$ eff. action involving metric and graviphoton curvatures


  $$
  \int d^4x (R^+)^2 (F^+)^{2h-2}
  $$

- When the effective $\mathcal{N} = 2$ theory is obtained from type II strings on CY via geometrical engineering, such terms
  - arise from world-sheets of genus $h$
  - are computed by the *topological string* [Bershadsky et al 1993, Antoniadis et al 1993]

- For the local CY describing the SU(2) theory the proposal has been tested [Klemm et al, 2002]
“Microscopic” string description of the gauge theory

The “semiclassical” approach to the low energy $\mathcal{N} = 2$ effective action can be given a simple (i.e., calculable) string theory realization.

- The non-abelian gauge theory d.o.f. are realized by open strings attached to (fractional) D3 branes. In the $\mathcal{N} = 2$ case, the perturbative 1-loop contributions to the prepotential are easily retrieved. [Douglas-Li 1996, Lawrence et al 1998, ...]

- The instantonic sectors of gauge theories can be realized by including D(-1) branes (a.k.a. D-instantons).
The description of gauge instantons via D(-1)-branes is more than a book-keeping device for the ADHM construction.

- The D(-1)'s act as sources that produce the actual profile of the gauge instanton solution [Billò et al. 2002]

One can include closed string backgrounds producing interesting deformations of the gauge theory. For instance

- non-commutative theories from NSNS background $B_{\mu\nu}$
- non-anticommutative theories from specific RR backgrounds [Billò et al. 2004-2005,...]
The description of gauge \textit{instantons} via D(-1)-branes is more than a book-keeping device for the ADHM construction.

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One can include \textit{closed string backgrounds} producing interesting \textit{deformations} of the gauge theory. For instance

- \textit{non-commutative} theories from NSNS background $B_{\mu\nu}$

- \textit{non-anticommutative} theories from specific RR backgrounds [Billò et al 2004-2005,...]
The aim of this work

- Reproduce the semiclassical instanton expansion of the low energy effective action for the $\mathcal{N} = 2$ SYM theory in the microscopic string realization via (fractional) D3/D(-1) branes.

- Show that the inclusion of the graviphoton of the $\mathcal{N} = 2$ bulk sugra, which comes from the RR sector,
  - leads exactly to the localization deformation on the instanton moduli space which allows to perform the integration
  - produces in the effective action the gravitational F-terms which are computed by the topological string on local CY.
The aim of this work

- Reproduce the semiclassical instanton expansion of the low energy effective action for the $\mathcal{N} = 2$ SYM theory in the microscopic string realization via (fractional) D3/D(-1) branes.

The situation is therefore as follows:

- Microscopic string description
  - deformed multi-instanton computations
  - Gravitational F-term interactions

- Geometrically engineered string description of l.e.e.t on local CY
  - topological string amplitudes at genus $h$

- The two ways to compute the same F-terms must coincide if the two descriptions are equivalent.
Microscopic string description of $\mathcal{N} = 2$ SYM
SYM from fractional branes

Consider pure SU($N$) Yang-Mills in 4 dimensions with $\mathcal{N} = 2$ susy.

- It is realized by the massless d.o.f. of open strings attached to fractional D3-branes in the orbifold background

$$\mathbb{R}^4 \times \mathbb{R}^2 \times \mathbb{R}^4 / \mathbb{Z}_2$$

- The orbifold breaks 1/2 SUSY in the bulk, the D3 branes a further 1/2:

$$32 \times \frac{1}{2} \times \frac{1}{2} = 8$$

real supercharges
Fields and string vertices

- **Field content:** $\mathcal{N} = 2$ chiral superfield

$$\Phi(x, \theta) = \phi(x) + \theta \Lambda(x) + \frac{1}{2} \theta \sigma^{\mu \nu} \theta F^+_{\mu \nu}(x) + \cdots$$

- **String vertices:**

$$V_A(z) = \frac{A_\mu(p)}{\sqrt{2}} \psi^\mu(z) e^{ip \cdot X(z)} e^{-\varphi(z)}$$

$$V_\Lambda(z) = \Lambda^{\alpha A}(p) S_\alpha(z) S_A(z) e^{ip \cdot X(z)} e^{-\frac{1}{2} \varphi(z)}$$

$$V_\phi(z) = \frac{\phi(p)}{\sqrt{2}} \bar{\psi}(z) e^{ip \cdot X(z)} e^{-\varphi(z)}$$

with all polariz.s in the adjoint of $U(N)$
Fields and string vertices

- **Field content:** $\mathcal{N} = 2$ chiral superfield

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- **String vertices:**

\[
V_A(z) = \frac{A_\mu(p)}{\sqrt{2}} \psi^\mu(z) e^{ip \cdot X(z)} e^{-\varphi(z)}
\]

\[
V_\Lambda(z) = \Lambda^\alpha A(p) S_\alpha(z) S_A(z) e^{ip \cdot X(z)} e^{-\frac{1}{2} \varphi(z)}
\]

\[
V_\phi(z) = \frac{\phi(p)}{\sqrt{2}} \overline{\psi}(z) e^{ip \cdot X(z)} e^{-\varphi(z)}
\]

with all polariz. s in the **adjoint** of $U(N)$
Gauge action from disks on fD3's

\[ \alpha' \to 0 \quad \text{with gauge quantities fixed.} \]

String amplitudes on disks attached to the D3 branes in the limit

\[ \alpha' \to 0 \] give rise to the tree level (microscopic) \( \mathcal{N} = 2 \) action

\[
S_{\text{SYM}} = \int d^4x \ Tr \left\{ \frac{1}{2} F_{\mu\nu}^2 + 2 D_\mu \bar{\phi} D^\mu \phi - 2 \bar{\Lambda} A \bar{D}^{\dot{\alpha} \beta} \Lambda^A_{\beta} + \sqrt{2} g \bar{\Lambda} A \epsilon^{AB} [\phi, \bar{\Lambda} B^A] + \sqrt{2} g \Lambda^A \epsilon_{AB} [\bar{\phi}, \Lambda^B] + g^2 [\phi, \bar{\phi}]^2 \right\}
\]
We are interested in the l.e.e.a. on the Coulomb branch parametrized by the v.e.v.’s of the adjoint chiral superfields:

\[ \langle \Phi_{uv} \rangle \equiv \langle \phi_{uv} \rangle = a_{uv} = a_u \delta_{uv}, \quad u, v = 1, \ldots, N, \quad \sum_u a_u = 0 \]

breaking \( SU(N) \rightarrow U(1)^{N-1} \) [we focus for simplicity on \( SU(2) \)]

Up to two-derivatives, \( \mathcal{N} = 2 \) susy forces the effective action for the chiral multiplet \( \Phi \) in the Cartan direction to be of the form

\[ S_{\text{eff}}[\Phi] = \int d^4 x d^4 \theta \mathcal{F}(\Phi) + \text{c.c} \]

We want to discuss the instanton corrections to the prepotential \( \mathcal{F} \)

Recall in our string set-up
Instanton calculus by mixed string diagrams
Consider the Wess-Zumino term of the effective action for a stack of D3 branes:

\[ \text{D. B. I. } + \int_{D3} \left[ C_3 + \frac{1}{2} C_0 \text{Tr}(F \wedge F) \right] \]

The topological density of an instantonic configuration corresponds to a localized source for the RR scalar \( C_0 \), i.e., to a distribution of D-instantons on the D3’s.

Instanton-charge \( k \) solutions of 3+1 dims. SU(\( N \)) gauge theories correspond to \( k \) D-instantons inside \( N \) D3-branes.

Stringy description of gauge instantons

\[
\begin{array}{c|cccc|cccc}
 & 1 & 2 & 3 & 4 & 5 & 6 & 7 & 8 & 9 & 10 \\
\hline
D3 & - & - & - & - & * & * & * & * & * \\
D(-1) & * & * & * & * & * & * & * & * & * \\
\end{array}
\]

\(N\) D3 branes
\(u = 1, \ldots, N\)

\(k\) D\((-1)\) branes
\(i = 1, \ldots, k\)

D3/D3, C-P: \(uv\)

D(-1)/D(-1), C-P: \(ij\)

D(-1)/D3, C-P: \(iu\)
Open strings with at least one end on a $D(-1)$ carry no momentum: they are moduli, rather than fields $\leftrightarrow$ parameters of the instanton.

<table>
<thead>
<tr>
<th>ADHM</th>
<th>Meaning</th>
<th>Vertex</th>
<th>Chan-Paton</th>
</tr>
</thead>
<tbody>
<tr>
<td>$a'_\mu$</td>
<td>centers</td>
<td>$\psi^\mu(z)e^{-\varphi(z)}$</td>
<td>adj. $U(k)$</td>
</tr>
<tr>
<td>$\chi$</td>
<td>aux.</td>
<td>$\overline{\Psi}(z)e^{-\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$D_c$</td>
<td>Lagrange mult.</td>
<td>$\overline{\eta}_{\mu\nu}^{c}(z)\psi^\nu(z)\psi^\mu(z)$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$M^{\alpha A}$</td>
<td>partners</td>
<td>$S_\alpha(z)S_A(z)e^{-\frac{1}{2}\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$\lambda_{\dot{\alpha} A}$</td>
<td>Lagrange mult.</td>
<td>$S^{\dot{\alpha}}(z)S^A(z)e^{-\frac{1}{2}\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$w_{\dot{\alpha}}$</td>
<td>sizes</td>
<td>$\Delta(z)S^{\dot{\alpha}}(z)e^{-\varphi(z)}$</td>
<td>$k \times \overline{N}$</td>
</tr>
<tr>
<td>$\bar{w}_{\dot{\alpha}}$</td>
<td>$\vdots$</td>
<td>$\overline{\Delta}(z)S^{\dot{\alpha}}(z)e^{-\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$\mu^A$</td>
<td>partners</td>
<td>$\Delta(z)S_A(z)e^{-\frac{1}{2}\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
<tr>
<td>$\bar{\mu}^A$</td>
<td>$\vdots$</td>
<td>$\overline{\Delta}(z)S_A(z)e^{-\frac{1}{2}\varphi(z)}$</td>
<td>$\vdots$</td>
</tr>
</tbody>
</table>
Among the \(-1/-1\) moduli we can single out the center \(x^\mu_0\) and its super-partners \(\theta^{\alpha A}\):

\[
a'^\mu = x^\mu_0 1_{k \times k} + y^\mu_c T^c \quad (T^c = \text{gen.s of SU}(k))
\]

\[
M^{\alpha A} = \theta^{\alpha A} 1_{k \times k} + \zeta^\alpha_c T^c
\]

The moduli \(x^\mu_0\) and \(\theta^{\alpha A}\) decouple from many interactions and play the rôle of superspace coords.

We will distinguish the moduli \(\mathcal{M}(k)\) into

\[
\mathcal{M}(k) \rightarrow (x_0, \theta ; \hat{\mathcal{M}}(k))
\]

\(\hat{\mathcal{M}}(k)\) are the centred moduli.
Disk amplitudes and effective actions

Usual disks: 
- $D_3$
- $D(-1)$

Mixed disks: 
- $D_3$
- $D(-1)$

Disk amplitudes

$\alpha' \to 0$ field theory limit

Effective actions

- $\mathcal{N} = 2$ SYM action
- $D(-1)/D(-1)$ and mixed

ADHM measure

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The action for the moduli

From disk diagrams with insertion of moduli vertices, in the field theory limit we extract the ADHM moduli action (at fixed $k$)

$$S_{\text{mod}} = S^{(k)}_{\text{bos}} + S^{(k)}_{\text{fer}} + S^{(k)}_{\text{c}}$$

with

$$S^{(k)}_{\text{bos}} = \text{tr}_k \left\{ -2 [\chi^\dagger, a'_\mu] [\chi, a'^{\mu}] + \chi^\dagger \bar{w}_\alpha w^{\dot{\alpha}} \chi + \chi \bar{w}_\dot{\alpha} w^\alpha \chi^\dagger \right\}$$

$$S^{(k)}_{\text{fer}} = \text{tr}_k \left\{ i \frac{\sqrt{2}}{2} \bar{\mu}^A \epsilon_{AB} \mu^B \chi^\dagger - i \frac{\sqrt{2}}{4} M^{\alpha A} \epsilon_{AB} [\chi^\dagger, M^{B\alpha}] \right\}$$

$$S^{(k)}_{\text{c}} = \text{tr}_k \left\{ -i D_c (W^c + i \bar{\eta}^c_{\mu\nu} [a'^{\mu}, a'^{\nu}]) \right. \right.$$  

$$\left. -i \lambda^{\dot{\alpha}} A (\bar{\mu}^A w_{\dot{\alpha}} + \bar{w}_{\dot{\alpha}} \mu^A + [a'_{\alpha \dot{\alpha}}, M^{\alpha A}]) \right\}$$

- In $S^{(k)}_{\text{c}}$ the bosonic and fermionic ADHM constraints appear (string moduli span the unconstrained parameter space)
Auxiliary moduli

The quartic interactions in $S_{bos}^{(k)}$ can be disentangled using auxiliary moduli $Y_\mu$, $X_\dot{\alpha}$ and $\bar{X}_\dot{\alpha}$:

$$S_{bos}'^{(k)} = \text{tr}_k \left\{ 2 \, Y_\mu^\dagger Y_\mu + 2 \, Y_\mu^\dagger [a'^\mu, \chi] + 2 \, Y_\mu [a'^\mu, \chi^\dagger] 
+ \bar{X}_\dot{\alpha}^\dagger X_\dot{\alpha} + \bar{X}_\dot{\alpha} X_\dot{\alpha}^\dagger + \bar{X}_\dot{\alpha}^\dagger w_\dot{\alpha} \chi + \bar{X}_\dot{\alpha} w_\dot{\alpha} \chi^\dagger - \chi \bar{w}_\dot{\alpha} X_\dot{\alpha}^\dagger - \chi^\dagger \bar{w}_\dot{\alpha} X_\dot{\alpha} \right\}$$

The corresponding auxiliary vertices are

$$V_Y(z) = \sqrt{2}g_0 \, Y_\mu \, \bar{\Psi}(z) \, \psi_\mu(z)$$
$$V_X(z) = g_0 \, X_\dot{\alpha} \, \Delta(z) S_\dot{\alpha}(z) \, \bar{\Psi}(z)$$
$$V_{\bar{X}}(z) = g_0 \, \bar{X}_\dot{\alpha} \, \Delta(z) S_\dot{\alpha}(z) \, \bar{\Psi}(z)$$
An example

One of the terms in $S_{\text{bos}}'(k)$ (involving the auxiliary moduli):

$$\left\langle V\bar{X}^\dagger V_w V_\chi \right\rangle$$

$$\equiv C_0 \int \prod_i \frac{dz_i}{dV_{\text{CKG}}} \times \left\langle V\bar{X}^\dagger(z_1) V_w(z_2) V_\chi(z_3) \right\rangle$$

$$= \cdots = \text{tr}_k \left\{ \bar{X}_{\dot{\alpha}}^\dagger w^{\dot{\alpha}} \chi \right\}$$

▶ Here $C_0 = 8\pi^2/g^2$ is the normalization of $D(-1)$ disks.
Introducing scalar v.e.v.’s

To evaluate the effect of a v.e.v. $\langle \Phi_{uv} \rangle = a_{uv} = a_u \delta_{uv}$ compute systematically mixed disk diagrams with a constant scalar $\phi$ emitted from the D3 boundary. Example:

\[ \langle V \bar{X}^\dagger V_{\phi=a} V_w \rangle = \ldots = \text{tr}_k \{ \bar{X}^\dagger_{\alpha} a w^{\dot{\alpha}} \} \]
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$$\langle V \bar{X}^\dagger V \phi = a V w \rangle = \ldots = \text{tr}_k \left\{ \bar{X}^\dagger_\alpha a w^{\dot{\alpha}} \right\}$$

The resulting moduli action is obtained with the shifts

$$\chi_{ij} \delta_{uv} \rightarrow \chi_{ij} \delta_{uv} - \delta_{ij} a_{uv}, \quad \chi^\dagger_{ij} \delta_{uv} \rightarrow \chi^\dagger_{ij} \delta_{uv} - \delta_{ij} \bar{a}_{uv}$$
Introducing scalar v.e.v.'s

To evaluate the effect of a v.e.v. \( \langle \Phi_{uv} \rangle = a_{uv} = a_u \delta_{uv} \) compute systematically mixed disk diagrams with a constant scalar \( \phi \) emitted from the D3 boundary. Example:

\[
\langle V \bar{X}^\dagger V \phi=a \ V_w \rangle = ... = \text{tr}_k \left\{ \bar{X}^\dagger \alpha \ a \ w^\alpha \right\}
\]

- The action does not depend on the center super-coordinates \( x_0^\mu \) and \( \theta^\alpha A \).
Holomorphicity, \(Q\)-exactness

In the action \(S_{\text{mod}}(a, \bar{a}; M_{(k)})\) the v.e.v.’s \(a\) and \(\bar{a}\) are not on the same footing: \(a\) does not appear in the fermionic action.

- The moduli action has the form

\[
S_{\text{mod}}(a, \bar{a}) = Q \Xi
\]

where \(Q\) is the scalar twisted supercharge:

\[
Q^{\dot{\alpha}B} \xrightarrow{\text{top. twist}} Q^{\dot{\alpha}\dot{\beta}}, \quad Q \equiv \frac{1}{2} \epsilon^{\dot{\alpha}\dot{\beta}} Q^{\dot{\alpha}\dot{\beta}}
\]

- The parameter \(\bar{a}\) appears only in the gauge fermion \(\Xi\)
- The instanton partition function

\[
Z^{(k)}(a) \equiv \int dM_{(k)} e^{-S_{\text{mod}}(a, \bar{a})}
\]

is independent of \(\bar{a}\): variation w.r.t this parameter is \(Q\)-exact.
Holomorphicity, $Q$-exactness

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- The parameter $\bar{a}$ appears only in the gauge fermion $\Xi$
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$$Z^{(k)}(a) \equiv \int dM_{(k)} e^{-S_{\text{mod}}(a, \bar{a})}$$

is independent of $\bar{a}$: variation w.r.t this parameter is $Q$-exact.
Consider disk diagrams involving only moduli $\mathcal{M}_k$, and no D3/D3 state (these are “vacuum” contributions from the D3 point of view)

$$\alpha' \to 0 \quad \sim \quad - \frac{8\pi^2 k}{g^2} \quad - \quad S_{\text{mod}}$$

(the “pure” D(-1) disks yields $kC_0$ [Polchinski, 1994])
Consider disk diagrams involving only moduli $\mathcal{M}_{(k)}$, and no D3/D3 state (these are “vacuum” contributions from the D3 point of view).

\[
\alpha' \xrightarrow{\approx} 0 \quad - \quad \frac{8\pi^2 k}{g^2} \quad - \quad S_{\text{mod}}
\]

(the “pure” D(-1) disks yields $kC_0$ [Polchinski, 1994])

- The combinatorics of boundaries [Polchinski, 1994] is such that these D-instanton diagrams exponentiate
Instanton calculus from the string standpoint

Consider disk diagrams involving only moduli $\mathcal{M}(k)$, and no D3/D3 state (these are “vacuum” contributions from the D3 point of view)

\[
\alpha' \longrightarrow 0 \quad \equiv \quad 8\pi^2 \frac{k}{g^2} \quad \equiv \quad S_{\text{mod}}
\]

(the “pure” D(-1) disks yields $kC_0$ [Polchinski, 1994])

- The moduli must be integrated over (path integral $\rightarrow$ ordinary integral):

\[
Z^{(k)} = \int d\mathcal{M}(k) e^{-\frac{8\pi^2 k}{g^2} - S_{\text{mod}}}
\]
Field-dependent moduli action

Consider correlators of D3/D3 fields, e.g. of the scalar $\phi$ in the Cartan direction, in presence of $k$ D-instantons. It turns out that

[Green-Gutperle 2000, Billò et al 2002]

- the dominant contribution to $\langle \phi_1 \ldots \phi_n \rangle$ is from $n$ one-point amplitudes on disks with moduli insertions. The result can be encoded in extra moduli-dependent vertices for $\phi$’s, i.e. in extra terms in the moduli action containing such one-point functions

$$S_{\text{mod}}(\varphi; \hat{M}) = \phi(x_0)J_{\phi}(\hat{M}) + S_{\text{mod}}(\hat{M})$$

with

$$\phi(x_0)J_{\phi}(\hat{M}) = \phi$$
Moduli action with the unbroken scalar $\phi$

The relevant one-point diagrams are those already computed to describe the dependence on the v.e.v. $a$. We just insert the vertex $V_\phi$ at non-zero momentum. Recall

$$\langle V_{\bar{X}^\dagger} V_\phi V_w \rangle = \ldots = \text{tr}_k \left\{ \bar{X}^\dagger_{\alpha} \phi(x_0) w^{\dot{\alpha}} \right\}$$

We get a dependence on the instanton location $x_0$ ($x$ from now on)

- The field-dependent action $S_{\text{mod}}(\phi; \mathcal{M})$ is thus simply obtained by

$$a \to \phi(x)$$
Effective action for the unbroken multiplet

Other non-zero diagrams couple the components of the gauge supermultiplet to the moduli, related by the Ward identities of the susies broken by the D(-1).

- Example:

\[
\langle V \bar{X}^{\dagger} V_{\delta\phi} V_w \rangle = \langle V \bar{X}^{\dagger} \left[ \theta^{\alpha A} Q_{\alpha A}, V_{\Lambda} \right] V_w \rangle = - \langle V \bar{X}^{\dagger} V_{\Lambda} V_w \int V_{\theta} \rangle
\]

The contribution of the last diagram can be obtained simply by letting \( \phi \rightarrow \theta^{\Lambda} \)

- This iterates: further couplings with higher components of \( \Phi \) and more \( \theta \)-insertions
Effective action for the unbroken multiplet

Other non-zero diagrams couple the components of the gauge supermultiplet to the moduli, related by the Ward identities of the susies broken by the D(-1).

- The superfield-dependent moduli action $S_{\text{mod}}(\Phi; M)$ is obtained by simply letting

$$a \rightarrow \Phi(x, \theta)$$
Contributions to the prepotential

Integrating over the moduli the interactions described by $S_{\text{mod}}(\Phi; M(k))$ one gets the effective action for the long-range multiplet $\Phi$ induced by the $k$-th instanton sector:

$$S_{\text{eff}}^{(k)}[\Phi] = \int d^4 x \ d^4 \theta \ d \hat{M}_k e^{-\frac{8\pi k}{g^2} - S_{\text{mod}}(\Phi; M(k))}$$

Correspondingly, the prepotential for the low energy $\mathcal{N} = 2$ theory is given by the centred instanton partition function

$$\mathcal{F}^{(k)}(\Phi) = \int d \hat{M}_k e^{-\frac{8\pi k}{g^2} - S_{\text{mod}}(\Phi; M(k))}$$

The superfield $\Phi(x, \theta)$ is a constant w.r.t. $\hat{M}_k$. We can compute $\mathcal{F}^{(k)}$ fixing $\Phi(x, \theta) \rightarrow a$ and using the results of the literature [see e.g. Dorey et al, 2002]
Contributions to the prepotential

Integrating over the moduli the interactions described by $S_{\text{mod}}(\Phi; \mathcal{M}(k))$ one gets the effective action for the long-range multiplet $\Phi$ induced by the $k$-th instanton sector:

$$S_{\text{eff}}^{(k)}[\Phi] = \int d^4x \, d^4\theta \, d\widehat{\mathcal{M}}_k \, e^{-\frac{8\pi k}{g^2} - S_{\text{mod}}(\Phi; \mathcal{M}(k))}$$

Correspondingly, the prepotential for the low energy $\mathcal{N} = 2$ theory is given by the centred instanton partition function

$$\mathcal{F}^{(k)}(\Phi) = \int d\widehat{\mathcal{M}}_k \, e^{-\frac{8\pi k}{g^2} - S_{\text{mod}}(\Phi; \mathcal{M}(k))}$$

- The superfield $\Phi(x, \theta)$ is a constant w.r.t. $\widehat{\mathcal{M}}_k$. We can compute $\mathcal{F}^{(k)}$ fixing $\Phi(x, \theta) \rightarrow a$ and using the results of the literature [see e.g. Dorey et al, 2002]
Contributions to the prepotential (2)

One finds ($\Lambda$ is the dynamical scale)

$$F^{(k)}(\Phi) = c_k \Phi^2 \left( \frac{\Lambda}{\Phi} \right)^{4k}.$$ 

- $\Lambda^{4k}$ stems from the term $\exp(-8\pi k/g^2)$, using the $\beta$-function of the $\mathcal{N} = 2$, SU(2) theory.
- The coefficients $c_k$ (the hard part!) were finally determined using a deformation of the moduli action which localizes the integration.


We will now embed this localization deformation in our string set-up, recognizing it as the effect of a graviphoton background and deriving gravitational corrections to the non-perturbative prepotential.
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We will now embed this localization deformation in our string set-up, recognizing it as the effect of a graviphoton background and deriving gravitational corrections to the non-perturbative prepotential.
Deformation from a graviphoton background
The Weyl multiplet

- The field content of $\mathcal{N} = 2$ sugra:

  $$ h_{\mu\nu} \ (\text{metric}) , \quad \psi_\mu^{\alpha A} \ (\text{gravitini}) , \quad C_\mu \ (\text{graviphoton}) $$

  can be organized in a chiral Weyl multiplet:

  $$ W^+_{\mu\nu}(x, \theta) = F^+_{\mu\nu}(x) + \theta \chi^+_{\mu\nu}(x) + \frac{1}{2} \theta \sigma^\lambda \theta \ R^+_{\mu\nu\lambda\rho}(x) + \cdots $$

  ($\chi^{\alpha A}_{\mu\nu}$ is the gravitino field strength)

- These fields arise from massless vertices of type IIB strings on $\mathbb{R}^4 \times \mathbb{C} \times \mathbb{C}^2 / \mathbb{Z}_2$
Graviphoton vertex

The graviphoton vertex, connected by the supercharges to the other fields in the Weyl multiplet, is given by

\[
V_{\mathcal{F}}(z, \bar{z}) = \frac{1}{4\pi} \mathcal{F}^{\alpha, \beta AB}(p) \\
\times \left[ S_\alpha(z) S_A(z)e^{-\frac{1}{2}\varphi(z)} S_\beta(\bar{z}) S_B(\bar{z})e^{-\frac{1}{2}\varphi(\bar{z})} \right] e^{i p \cdot X(z, \bar{z})}
\]
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- We will insert the closed string vertices in disk diagrams bounded by the branes → suitable identifications between left- and right-movers taken into account
The graviphoton vertex, connected by the supercharges to the other fields in the Weyl multiplet, is given by

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\]

The bi-spinor graviphoton polarization is given by

\[
\mathcal{F}(\alpha\beta)[AB] = \frac{\sqrt{2}}{4} \mathcal{F}^{+}_{\mu\nu}(\sigma^{\mu\nu})^{\alpha\beta} \epsilon^{AB}
\]
Graviphoton vertex

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\[ V_F(z, \bar{z}) = \frac{1}{4\pi} \mathcal{F}^{\alpha\beta AB}(p) \times \left[ S_\alpha(z)S_A(z)e^{-\frac{1}{2}\varphi(z)} S_\beta(\bar{z})S_B(\bar{z})e^{-\frac{1}{2}\varphi(\bar{z})} \right] e^{i p \cdot X(z, \bar{z})} \]

▶ A different RR field, with a similar structure, will be useful:

\[ V_{\tilde{F}}(z, \bar{z}) = \frac{1}{4\pi} \tilde{\mathcal{F}}^{\alpha\beta \hat{A}\hat{B}}(p) \times \left[ S_\alpha(z)S_{\hat{A}}(z)e^{-\frac{1}{2}\varphi(z)} S_\beta(\bar{z})S_{\hat{B}}(\bar{z})e^{-\frac{1}{2}\varphi(\bar{z})} \right] e^{i p \cdot X(z, \bar{z})} \]

\[ \hat{A}, \hat{B} = 3, 4 \leftrightarrow \text{odd “internal” spin fields} \]
Let us investigate the effect of a graviphoton v.e.v. 

\[ \langle W_{\mu\nu}^+ \rangle \equiv \langle F_{\mu\nu}^+ \rangle \equiv f_{\mu\nu} \]

on the moduli measure.

- We have to consider disk amplitudes with open string moduli vertices on the boundary and closed string graviphoton vertices in the interior which survive in the field theory limit \( \alpha' \to 0 \).
- We will consider also insertions of vertices of type \( V_{\bar{F}} \), with constant polarization \( \bar{F}_{\mu\nu}^+ = \bar{f}_{\mu\nu} \).
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- We will consider also insertions of vertices of type $V\bar{F}$, with constant polarization $\bar{F}_{\mu\nu} = \bar{f}_{\mu\nu}$.
Non-zero diagrams

Very few diagrams contribute.

The only one involving the true graviphoton is

\[
\left\langle V_{Y^+} V_{a'} V_{\mathcal{F}, \bar{\mathcal{F}}} \right\rangle \equiv C_0 \int \frac{dz_1 \, dz_2 \, dw \, d\bar{w}}{dV_{\text{CKG}}} \left\langle V_{Y^+}(z_1) \, V_{a'}(z_2) \, V_{\mathcal{F}}(w, \bar{w}) \right\rangle
\]
Non-zero diagrams

Very few diagrams contribute.

► More explicitly,

\[
\left\langle \mathcal{V}_{Y^\dagger} \mathcal{V}_{a'} \mathcal{V}_{\mathcal{F}} \right\rangle = \frac{1}{4\pi} \text{tr}_k \left\{ Y_{\mu}^\dagger a'_\nu f_{\lambda\rho} \right\} (\sigma^{\lambda\rho})^{\alpha\beta} \epsilon^{AB} \int \frac{dz_1 dz_2 dwd\bar{w}}{dV_{\text{CKG}}} \times
\]

\[
\langle e^{-\varphi(z_2)} e^{-\frac{1}{2} \varphi(w)} e^{-\frac{1}{2} \varphi(\bar{w})} \rangle \langle \Psi(z_1) S_A(w) S_B(\bar{w}) \rangle \\
\langle \psi^\mu(z_1) \psi^\nu(z_2) S_\alpha(w) S_\beta(\bar{w}) \rangle
\]
Non-zero diagrams

Very few diagrams contribute.

► Result: (same also with $\bar{f}^{\mu\nu}$)

$$\langle V_{Y^\dagger} V_{a'} V_{\mathcal{F}} \rangle = -4i \text{tr}_k \left\{ Y_{\mu}^\dagger a_{\nu}' f^{\mu\nu} \right\}$$

► Moreover, term with fermionic moduli and a $V_{\bar{\mathcal{F}}}$:

$$\langle V_{M} V_{M} V_{\bar{\mathcal{F}}} \rangle = \frac{1}{4\sqrt{2}} \text{tr}_k \left\{ M^\alpha A M^\beta B \bar{f}_{\mu\nu} \right\} (\sigma^{\mu\nu})_{\alpha\beta} \epsilon_{AB}$$
The deformed moduli action

Including the backgrounds $f, \bar{f}$ besides the chiral v.e.v.'s $a, \bar{a}$:

$$S_{\text{mod}}(a, \bar{a}; f, \bar{f}) =$$

$$- \text{tr}_k \left\{ \left( [\chi^\dagger, a'_{\alpha\dot{\beta}}] + 2\bar{f}_c(\tau^c a')_{\alpha\dot{\beta}} \right) \left( [\chi, a'^{\dot{\beta}\alpha}] + 2f_c(a'\tau^c)^{\dot{\beta}\alpha} \right) \right.$$

$$- (\chi^\dagger \bar{w}_{\dot{\alpha}} - \bar{w}_{\dot{\alpha}} \bar{a})(w^{\dot{\alpha}} \chi - a w^{\dot{\alpha}}) - \left( \chi \bar{w}_{\dot{\alpha}} - \bar{w}_{\dot{\alpha}} a \right) (w^{\dot{\alpha}} \chi^\dagger - \bar{a} w^{\dot{\alpha}}) \left. \right\}$$

$$+ i \frac{\sqrt{2}}{2} \text{tr}_k \left\{ \tilde{\mu}^A \epsilon_{AB} (\mu^B \chi^\dagger - \bar{a} \mu^B) \right.$$

$$- \frac{1}{2} M^{\alpha A} \epsilon_{AB} \left( [\chi^\dagger, M^B_{\alpha}] + 2\bar{f}_c (\tau^c)^{\alpha\dot{\beta}} M^{\beta B} \right) \right\} + S^{(k)}_c$$

▶ The constraint part of the action, $S^{(k)}_c$, is not modified
The deformed moduli action

Including the backgrounds $f, \bar{f}$ besides the chiral v.e.v.'s $a, \bar{a}$:

$$S_{\text{mod}}(a, \bar{a}; f, \bar{f}) =$$

$$- \text{tr}_k \left\{ ([\chi^\dagger, a'_{\alpha\dot{\beta}}] + 2\bar{f}_c (\tau^c a')_{\alpha\beta}) ([\chi, a'_{\beta\dot{\alpha}}] + 2f_c (a'_{\tau^c})_{\beta\dot{\alpha}}) ight\}$$

$$- (\chi^\dagger \bar{w}_{\dot{\alpha}} - \bar{w}_{\dot{\alpha}} a) (w^{\dot{\alpha}} \chi - a w^{\dot{\alpha}}) - (\chi \bar{w}_{\dot{\alpha}} - \bar{w}_{\dot{\alpha}} a) (w^{\dot{\alpha}} \chi^\dagger - \bar{a} w^{\dot{\alpha}})$$

$$+ i \frac{\sqrt{2}}{2} \text{tr}_k \left\{ \bar{\mu}^A \epsilon_{AB} (\mu^B \chi^\dagger - \bar{a} \mu^B) ight\}$$

$$+ \frac{1}{2} M^{\alpha A} \epsilon_{AB} ([\chi^\dagger, M^B_{\alpha}] + 2\bar{f}_c (\tau^c)_{\alpha\beta} M^{\beta B}) \right\} + S_c^{(k)}$$

The effect of the $f, \bar{f}$ background amounts to the shift

$$[\chi, (\bullet)_\alpha] \rightarrow [\chi, (\bullet)_\alpha] + 2f_c (\tau^c \bullet)_{\alpha} , \ [\chi^\dagger, (\bullet)_\alpha] \rightarrow [\chi^\dagger, (\bullet)_\alpha] + 2\bar{f}_c (\tau^c \bullet)_{\alpha}$$
Holomorphicity, $Q$-exactness

Also the deformed moduli action has the form

$$S_{\text{mod}}(a, \bar{a}; f, \bar{f}) = Q \Xi$$

where $Q$ is the scalar twisted supercharge.

- The parameters $\bar{a}, \bar{f}_c$ appear only in the gauge fermion $\Xi$
- The instanton partition function

$$Z^{(k)} \equiv \int dM_{(k)} \, e^{-S_{\text{mod}}(a, \bar{a}; f, \bar{f})}$$

is independent of $\bar{a}, \bar{f}_c$: variation w.r.t these parameters is $Q$-exact.
Graviphoton and localization

The moduli action obtained inserting the graviphoton background coincides exactly with the “deformed” action considered in the literature to localize the moduli space integration if we set

\[ f_c = \frac{\varepsilon}{2} \delta_3 c, \quad \bar{f}_c = \frac{\bar{\varepsilon}}{2} \delta_3 c, \]

and moreover (referring to the notations in the above ref.s)

\[ \varepsilon = \bar{\varepsilon}, \quad \varepsilon = \epsilon_1 = -\epsilon_2 \]

The “shift” rule which yields the deformed action was interpreted as “gauging” the chiral rotations in 4d Euclidean space which are symmetries of the ADHM constraints
The moduli action obtained inserting the *graviphoton* background coincides *exactly* with the “deformed” action considered in the literature to localize the moduli space integration if we set

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and moreover (referring to the notations in the above ref.s)

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The *localization deformation* of the \( \mathcal{N} = 2 \) ADHM construction is produced, in the type IIB string realization, by a *graviphoton* background.
Deformation of the gauge action

- The graviphoton background can be inserted also in D3 disks, producing extra terms in the gauge theory action on the D3 branes.
- Collecting all diagrams, one gets

\[ S_{YM} + \int d^4x \text{Tr}\left\{ -2ig F_{\mu\nu} \bar{\phi} f^{\mu\nu} - g^2 (\bar{\phi} f^{\mu\nu})^2 \right\} \]

(in agreement with couplings between gauge and Weyl multiplets in \( \mathcal{N} = 2 \) sugra)

- At linear order in \( g \), field eq.s for \( \phi \)

\[ D^2 \phi = -i \sqrt{2} g \epsilon_{AB} \Lambda^\alpha_A \Lambda^B_\alpha - 2i g f_{\mu\nu} F^{\mu\nu} \]

agree with the one implied by the deformed ADHM construction.
Weyl multiplet dependence of the effective prepotential

- Just as for the case of the scalar v.e.v’s only, we can compute

\[ S_{\text{mod}}(\Phi, W^{+}; \mathcal{M}(k)) \]

containing the one-point couplings of the fields in the gauge and Weyl multiplets to the moduli by simply promoting

\[ a \rightarrow \Phi(x; \theta), \quad f_{\mu\nu} \rightarrow W_{\mu\nu}(x, \theta) \]

- The l.e.e.a. for \( \Phi \) and \( W^{+} \) in the instanton \( \# k \) sector is

\[ S_{\text{eff}}^{(k)}[\Phi, W^{+}] = \int d^4x \ d^4\theta \ d\hat{\mathcal{M}}(k) e^{-\frac{8\pi k}{g^2}} - S_{\text{mod}}(\Phi, W^{+}; \mathcal{M}(k)) \]

and the prepotential reads thus

\[ \mathcal{F}^{(k)}(\Phi, W^{+}) = \int d\hat{\mathcal{M}}(k) e^{-\frac{8\pi k}{g^2}} - S_{\text{mod}}(\Phi, W^{+}; \mathcal{M}(k)) \]
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$$\mathcal{F}^{(k)}(\Phi, W^+) = \int d\widehat{\mathcal{M}}(k) \, e^{-\frac{8\pi k}{g^2}} - S_{\text{mod}}(\Phi, W^+; \mathcal{M}(k))$$
Expansion of the prepotential

\( \phi(x, \theta) \) and \( W^+_{\mu\nu}(x, \theta) \) are constant w.r.t. the integration variables \( \hat{M}_k \).

We can compute \( \mathcal{F}^{(k)}(a; f) \) and reinstate the full multiplets in the result.

- From the explicit form of \( S_{\text{mod}}(a, 0; f, 0) \) it follows that partition function \( \mathcal{F}^{(k)}(a; f) \) is invariant under

\[
a, f_{\mu\nu} \rightarrow -a, -f_{\mu\nu}
\]

- We need a regular expansion for \( f \rightarrow 0 \), and no odd powers of \( af_{\mu\nu} \), Altogether, reinstating the superfields,

\[
\mathcal{F}^{(k)}(\phi, W^+) = \sum_{h=0}^{\infty} c_{k,h} \phi^2 \left( \frac{\Lambda}{\phi} \right)^{4k} \left( \frac{W^+}{\phi} \right)^{2h}
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The non-perturbative prepotential

Sum over the instanton sectors:

\[ \mathcal{F}_{n.p.}(\Phi, W^+) = \sum_{k=1}^{\infty} \mathcal{F}^{(k)}(\Phi, W^+) = \sum_{h=0}^{\infty} C_h(\Lambda, \Phi)(W^+)^{2h} \]

with

\[ C_h(\Lambda, \Phi) = \sum_{k=1}^{\infty} C_{k,h} \frac{\Lambda^{4k}}{\Phi^{4k+2h-2}} \]

- Many different terms in the eff. action connected by susy.
  Saturating the \( \theta \) integration with four \( \theta \)'s all from \( W^+ \)

\[ \int d^4x \, C_h(\Lambda, \phi) (R^+)^2 (\mathcal{F}^+)^{2h-2} \]

Freezing \( \phi \rightarrow a \), this is a purely gravitational \( F \)-term
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Evaluation via localization

- To determine the coefficients $c_{k,h}$, constant bkg values $\Phi \rightarrow a$ and $W_{\mu\nu}^+ \rightarrow f_{\mu\nu}$ are enough.

- The localization deformation of the multi-instanton partition function $Z^{(k)}(a, \varepsilon)$ is obtained for

\[
f_{\mu\nu} = \frac{1}{2} \varepsilon \eta^{3}_{\mu\nu} , \quad \bar{f}_{\mu\nu} = \frac{1}{2} \bar{\varepsilon} \eta^{3}_{\mu\nu}
\]

- (Holomorphicity:) $Z^{(k)}(a, \varepsilon)$ does not smoothly depend on $\bar{\varepsilon}$.
- However, $\bar{\varepsilon} = 0$ is a limiting case: some care is needed.

- $\mathcal{F}^{(k)}(a; \varepsilon)$ is well-defined. $S^{(k)}[a; \varepsilon]$ diverges because of the (super)volume integral $\int d^4x \, d^4\theta$

- $\bar{\varepsilon}$ regularizes the superspace integration by a Gaussian term. One can then work with the effective action, i.e., the full instanton partition function.
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Example: the case \( k = 1 \)

The moduli action in the \( k = 1 \) sector (at \( \bar{a} = 0 \)) is

\[
S_{\text{mod}}^{(k=1)} = -2\bar{\epsilon}\epsilon x^2 - \frac{\bar{\epsilon}}{2} \theta^\alpha A_{\epsilon AB}(\tau_3) \alpha\beta \theta^\beta B + \hat{S}_{\text{mod}}^{(k=1)}(a)
\]

The last term does not depend on \( \epsilon, \bar{\epsilon} \)

- Using the \( \epsilon, \bar{\epsilon} \)-independence of the centred partition function

\[
\mathcal{F}^{(k=1)}(a) = \int d\hat{M}^{(k=1)} e^{-8\pi^2 g^2 - \hat{S}_{\text{mod}}^{(k=1)}(a)}
\]

we have for the \( k = 1 \) instanton partition function

\[
Z^{(k=1)}(a, \epsilon) = \int d^4 x d^4 \theta e^{-2\bar{\epsilon}\epsilon x^2 - \frac{1}{2} \bar{\epsilon} \theta \cdot \theta} \mathcal{F}^{(k=1)}(a) = \frac{1}{\epsilon^2} \mathcal{F}^{(k=1)}(a)
\]

- Effectively, with the full deformation, we have the rule

\[
\int d^4 x d^4 \theta \to \frac{1}{\epsilon^2}
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- Effectively, with the full deformation, we have the rule

$$\int d^4 x \, d^4 \theta \rightarrow \frac{1}{\epsilon^2}$$
The deformed partition function vs the prepotential

- With $\varepsilon, \tilde{\varepsilon}$, the partition function $Z(a; \varepsilon)$ can be computed:
  - (super)volume divergences $\rightarrow \varepsilon$ singularities
  - $a$ and $\varepsilon, \tilde{\varepsilon}$ deformations localize completely the integration over moduli space which can be carried out


- With $\tilde{\varepsilon} \neq 0$ (complete localization) a trivial superposition of instantons of charges $k_i$ contributes to the sector $k = \sum k_i$

- Such disconnected configurations do not contribute when $\tilde{\varepsilon} = 0$

- The partition function computed by localization corresponds to the exponential of the non-perturbative prepotential:

$$Z(a; \varepsilon) = \exp \left( \frac{\mathcal{F}_{n.p.}(a, \varepsilon)}{\varepsilon^2} \right) = \exp \left( \sum_{k=1}^{\infty} \frac{\mathcal{F}(k)(a, \varepsilon)}{\varepsilon^2} \right)$$

$$= \exp \left( \sum_{h=0}^{\infty} \sum_{k=1}^{\infty} \frac{\varepsilon^{2h-2} \left( \frac{\Lambda}{a} \right)^{4k}}{a^{2h}} \right)$$
The computation via localization techniques of the multi-instanton partition function $Z(a; \varepsilon)$ determines the coefficients $c_{k,h}$ which appear in the gravitational $F$-terms of the $\mathcal{N} = 2$ effective action

$$\int d^4 x \ C_h(\Lambda, \phi) (R^+)^2 (\mathcal{F}^+)^{2h-2}$$

via the relation

$$C_h(\Lambda, \phi) = \sum_{k=1}^{\infty} c_{k,h} \frac{\Lambda^{4k}}{\phi^{4k+2h-2}}$$

The very same gravitational $F$-terms can be extracted in a completely different way: topological string amplitudes on suitable Calabi-Yau manifolds
Summarizing

- The computation via localization techniques of the multi-instanton partition function $Z(a; \varepsilon)$ determines the coefficients $c_{k,h}$ which appear in the gravitational $F$-terms of the $\mathcal{N} = 2$ effective action

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- The very same gravitational $F$-terms can been extracted in a completely different way: topological string amplitudes on suitable Calabi-Yau manifolds
Relation to topological strings on CY
Geometrical engineering and topological strings

- **SW:** low energy $\mathcal{N} = 2 \leftrightarrow$ (auxiliary) Riemann surface
- **Geometrical engineering:** embed directly the low energy theory into string theory as type IIB on a suitable local CY manifold $\mathcal{M}$
  
  
  - geometric moduli of $\mathcal{M} \leftrightarrow$ gauge theory data ($\Lambda, a$);
  - The coefficients $C_h$ in the l.e.e.a. gravitational F-terms
  
  $$C_h (R^+)^2 (\mathcal{F}^+)^{2h-2}$$

  are given by topological string amplitudes at genus $h$
  

- For the local CY $\mathcal{M}_{\text{SU}(2)}$ the couplings $C_h$ were checked to coincide with those given by the deformed multi-instanton calculus as proposed by Nekrasov
  
  [Klemm et al 2002]
Microscopic vs effective string description

**Orbifold space**
- with D3/D(-1) system
- Moduli action depends on gauge theory data $\Lambda$, $a$
- open and closed strings

2 $h$ disks
connected by integration over moduli

\[ \chi = 2h - 2 \]

**Local CY manifold**
- Geometric moduli determined from gauge theory data $\Lambda$, $a$
- No branes - closed strings only

\[ \chi = 2h - 2 \]

Same gravitational F-term interactions

\[ C_h(\Lambda, a) \left( R^+ \right)^2 \left( F^+ \right)^{2h-2} \]

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Perspectives
Some interesting directions to go...

- Study of the instanton corrections to $\mathcal{N} = 2$ eff. theory in the gauge/gravity context: modifications of the classical solution of fD3’s
- Application of similar techniques to (euclidean) D3’s along a CY orbifold to derive BH partition functions in $\mathcal{N} = 2$ sugra (which OSV conjecture relates to $|Z_{\text{top}}|^2$)
- Non-perturbative corrections to $\mathcal{N} = 1$ superpotentials by Euclidean D3’s along orbifold directions
- ...

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Some references
Multi-instanton contributions in $\mathcal{N} = 2$


A. S. Losev, A. Marshakov and N. A. Nekrasov, [arXiv:hep-th/0302191];


String perspective on instanton calculus


Some notations
String fields in the orbifold space

- In the six directions transverse to the brane,
  \[ Z \equiv (X^5 + iX^6)/\sqrt{2}, \quad Z^1 \equiv (X^7 + iX^8)/\sqrt{2}, \quad Z^2 \equiv (X^9 + iX^{10})/\sqrt{2}, \]
  \[ \Psi \equiv (\psi^5 + i\psi^6)/\sqrt{2}, \quad \Psi^1 \equiv (\psi^7 + i\psi^8)/\sqrt{2}, \quad \Psi^2 \equiv (\psi^9 + i\psi^{10})/\sqrt{2} \]

  the \( \mathbb{Z}_2 \) orbifold generator \( h \) acts by
  \[(Z^1, Z^2) \rightarrow (-Z^1, -Z^2), \quad (\Psi^1, \Psi^2) \rightarrow (-\Psi^1, -\Psi^2)\]

- Under the \( \text{SO}(10) \rightarrow \text{SO}(4) \times \text{SO}(6) \) induced by \( \text{D3}'s \), \( S_{\hat{A}} \rightarrow (S_{\alpha} S_{A'}, S_{\hat{\alpha}} S_{A'}) \)
- Under \( \text{SO}(6) \rightarrow \text{SO}(2) \times \text{SO}(4) \) induced by the orbifold,

\[
\begin{array}{c|cccc|ccccc|c}
S^A' & \text{notat.} & \text{SO}(2) & \text{SO}(4) & S_A' & \text{notat.} & \text{SO}(2) & \text{SO}(4) & h \\
S^{+++} & S^A & \frac{1}{2} & (2, 1) & S^{--} & S_A & -\frac{1}{2} & (2, 1) & +1 \\
S^{+-} & S_{\hat{A}} & A=1, 2 & & S^{--} & S_{\hat{A}} & A=1, 2 & & \\
S^{-++} & S^{\hat{A}} & -\frac{1}{2} & (1, 2) & S^{++} & S_{\hat{A}} & \hat{A}=3, 4 & (1, 2) & -1 \\
S^{--+} & S_{\hat{\alpha}} & \hat{A}=3, 4 & & S^{+--} & S_{\hat{\alpha}} & \hat{A}=3, 4 & & \\
\end{array}
\]