Stringy instantons and dualities

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We discuss non-perturbative corrections to the gauge kinetic functions in a four-dimensional \(N=2\) gauge theory realized with a system of D7/D3-branes in a compactification of type I\(^\prime\) theory on \(T_4/\mathbb{Z}_2 \times \mathbb{Z}_2\). The non-perturbative contributions arise when D(–1) branes, corresponding to stringy instantons, are added to the system; such contributions can be explicitly evaluated using localization techniques and precisely match the results predicted by the heterotic/type I\(^\prime\) duality. This agreement represents a very non-trivial test of the stringy multi-instanton calculus.

1 Introduction

It has been recently found \cite{1} that certain classes of D-brane instantons arising in intersecting brane models can generate effective interactions at energies that are not linked to the gauge theory scale, and for this reason they are usually called “stringy” or “exotic” instantons. This feature is very welcome in the search of semi-realistic string scenarios for the physics beyond the Standard Model. It is therefore of the greatest importance to devise techniques to determine quantitatively such exotic non-perturbative corrections through their explicit realization at the string level. In this context both the usual gauge instantons and the exotic ones can be obtained from Euclidean branes entirely wrapping some cycle of the internal space. Depending on whether this cycle coincides or not with the one wrapped by the space-filling D-branes on which the gauge theory is defined, the Euclidean branes correspond to gauge or exotic instantons, respectively.

In the simplest cases, 4d gauge instantons can be realized with bound states of space-filling D3-branes and point-like D(-1)-branes \cite{2}. In these systems the massless sector of open strings having at least one endpoint on the D(-1)’s is in one-to-one correspondence with the moduli of the gauge instanton solution. Actually, also the effective action on the moduli space, the rules of the instanton calculus and the profile of the classical solution can be explicitly obtained in this way \cite{3,4}.

In the exotic cases, the gauge and instantonic branes intersect non-trivially in the internal space and thus the open strings stretching between them have extra “twisted” directions and some instanton moduli (specifically those related to sizes and gauge orientations) disappear from the spectrum. Their supersymmetric fermionic partners remain massless, and when integrated out, they can lead to the effective interactions we alluded to above. A very simple example of this phenomenon occurs in the D(-1)/D7 brane system, which exhibits the world-sheet features of exotic instantons since mixed open strings have eight twisted directions. By adding O7-planes, this system can be embedded in type I\' string theory, a setup which possesses a computable perturbative heterotic dual. The non-perturbative contributions of D-instantons to the effective action on the D7-branes can be explicitly computed as integrals over the moduli space via localization techniques, in strict analogy with what is done for usual gauge instantons \cite{5}. One finds that all D-instanton numbers correct the quartic gauge couplings of the 8d gauge theory of the D7-branes \cite{6-8}, and this whole series of terms matches the perturbative result obtained in the dual heterotic string theory.

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In this contribution we briefly describe one example of exotic instanton calculus in a 4d setup that has been developed in [9]. We consider a perturbatively conformal $\mathcal{N} = 2$ gauge theory that admits a brane realization where exotic instantons generate a whole series of corrections to the quadratic gauge couplings and possesses a calculable heterotic dual against which these corrections can be checked. This provides a very non-trivial check of the correctness of this approach to the exotic instanton calculus.

2 A $\mathcal{N} = 2$ conformal model from an orbifold of type I$'$

We consider a $\mathcal{N} = 2$ orientifold compactification of type IIB string theory on $T_4 \times T_2$. The action of the orientifold generators selects 4 O7-planes located at the invariant points of $T_2$ and 64 O3-planes located at the fixed points of $T_4 \times T_2$. The global cancellation of the RR tadpoles requires the presence of 16 dynamical D7-branes transverse to $T_2$ and of 16 dynamical “half” D3-branes transverse to the internal 6-torus. We choose to cancel locally the RR charges in $T_2$ by placing exactly 4 D7-branes and 4 half D3-branes on top of each O7-plane. The D3’s could then be distributed over the 16 orbifold fixed points that are common to a given O7-plane. For sake of simplicity we place them at distinct points of $T_4$ and we focus on the gauge theory leaving in one of the O7 fixed plane.

The 4d field theory leaving on the D7 world-volume at the selected fixed plane is a conformal $\mathcal{N} = 2$ U(4) SYM theory containing one adjoint vector multiplet, two antisymmetric hypermultiplets and four fundamental hypermultiplets which are charged under a $U(1)^4$ flavour group. The quadratic effective action for the gauge fields can be described by holomorphic Wilsonian couplings [10] that have the following structure:

$$ f = f(0) + f(1) + f_{a.p.} $$

where the subscripts $(0)$ and $(1)$ refer to the tree-level and 1-loop contributions, while the last term accounts for possible non-perturbative corrections. Writing the effective action in terms of the $\mathcal{N} = 2$ multiplet encoding the U(4) gauge degrees of freedom, $\Phi(x, \theta) = \phi(x) + \theta^\alpha \Lambda_\alpha(x) + (\theta_\gamma \theta^\mu \theta^\nu) F_{\mu \nu}(x)$, we see that there are two possible colour structures, each with its own coupling:

$$ S = \int d^4 x d^4 \theta \{ f \text{ Tr } \Phi^2 + f' (\text{Tr } \Phi)^2 \} + c.c. $$

The tree-level value for the single trace coupling $f$ can be deduced from the Born-Infeld action and is

$$ f(0) = -i t \quad (\text{Re } t = \frac{\theta_{YM}}{2\pi}, \text{ Im } t = \frac{4\pi}{g_Y^2} \sim \frac{Vol(T_4)}{g_s}) ; $$

on the other hand the tree level value for the double trace coupling $f'$ is vanishing.

The only perturbative corrections to $f$ and $f'$ come from the 1-loop threshold corrections, which in turn are related in an universal way to the string 1-loop two point-functions. The correction to the single trace coupling $f(1)$ is expected to vanish, since the gauge theory is conformal, and in fact the 1-loop string diagrams that contribute to the single trace coupling add up to zero. On the contrary, the 1-loop diagrams that contribute to the double trace structure give a non vanishing result, due to the massless states winding on $T_2$, and we have

$$ f'(1) = -8 \log \eta(U) , $$

where $U$ is complex structure of $T_2$.

In the next sections we will study the non-perturbative corrections $f_{a.p.}$ and $f'_{a.p.}$ induced by D-instantons.
3 D-instantons and their moduli spectrum

The orientifold projection that defines our model is compatible both with Euclidean E3-branes wrapped on \( T^4 \) that represent ordinary gauge instantons for the field theory living on the D7-branes, and D-instantons which describe truly stringy instanton configurations for the D7-brane gauge theory [6–8]. Here we only discuss the contributions produced by the D(-1)-branes, and show that they correct non-perturbatively the gauge kinetic functions of the \( \mathcal{N} = 2 \) U(4) theory discussed in Sect. [3].

Again we focus on the four D7-branes located at one of the orientifold fixed points, and place on them a number of fractional D-instantons. However, since there are also four D3-branes distributed in four different orbifold fixed points, we have to distinguish between two possibilities, depending on whether the D-instantons are at an empty fixed point or occupy the same position of one of the D3-branes. In the first case (case \( a \)) only the \((-1)/(-1)\) and \((-1)/7\) open strings support massless moduli, because the \((-1)/3\) strings have always a non-vanishing stretching energy due to the separation between their endpoints. In the second case (case \( b \)) we can find massless excitations also in the spectrum of the \((-1)/3\) strings. Consistently with the orientifold projections, when we set \( k \) “half” D-instantons at a given fixed point, the instanton moduli organize in representation of \( U(k) \) and in representations of the Lorentz symmetry group, which in our local system is broken to \( SO(4) \times SO(4) = SU(2)_+ \times SU(2)_- \times \hat{SU}(2)_+ \times \hat{SU}(2)_- \times SO(2) \).

The \((-1)/(1)\) moduli form the so-called neutral moduli sector, since they do not transform under the U(4) gauge group and are common to both case \( a \) and \( b \). They comprise four complex scalars transforming as vectors of the \( SO(4) \times SO(4) \) groups that rotate the coordinates of the D7 world volume, a complex scalar \( \chi \) and their fermionic partners. The charged moduli sector accounts for the \((-1)/7\) open strings. Since there are eight directions with mixed boundary conditions we only find a physical fermionic modulus \( \mu' \).

Finally, the flavored sector of the instanton moduli space arises from the \((-1)/3\) open strings. In our model, one finds two complex variables transforming as chiral spinors with respect to \( SO(4) \), and two spinors of opposite chiralities with respect to \( \hat{SO}(4) \). All physical moduli and their transformation properties are summarized in Tab. 1.

| neutral | \( SU(2)^4 \) | \( U(k) \) | charged | \( SU(2)^4 \) | \( U(k) \times U(4) \) |
|---------|----------------|----------|----------|----------------|----------------|
| \( B_{\ell} \) | \( (2, 2, 1, 1) \) | adjoint | \( \mu' \) | \( (1, 1, 1, 1) \) | \( \overline{\alpha}, \overline{\alpha} \) |
| \( M_{\alpha a} \) | \( (1, 2, 2, 1) \) | | | | |
| \( N_{\dot{\alpha} \dot{a}} \) | \( (1, 2, 1, 2) \) | \( \overline{\beta} + \overline{\dot{\beta}} \) | flavoured | \( SU(2)^4 \) | \( U(k) \times U(m) \) |
| \( B_{\ell} \) | \( (1, 1, 2, 2) \) | \( \overline{\dot{\alpha}} + \overline{\dot{\alpha}} \) | \( \overline{w}_{\dot{a}} \) | \( (2, 1, 1, 1) \) | \( \overline{\alpha}, \overline{\alpha} \) |
| \( M_{\alpha a} \) | \( (2, 1, 1, 2) \) | | | \( (1, 1, 2, 1) \) | |
| \( N_{\dot{\alpha} \dot{a}} \) | \( (2, 1, 2, 1) \) | adjoint | \( \mu_{\dot{a}} \) | \( (1, 1, 1, 2) \) | \( \overline{\alpha}, \overline{\alpha} \) |
| \( \chi \) | \( (1, 1, 1, 1) \) | | | | |

Table 1 Spectrum of physical instanton moduli.

Note that in addition to the physical moduli we have to consider extra auxiliary fields, \( d_m, D_{\alpha a}, h_\alpha \) and \( h' \), that linearize the quartic interactions among the moduli and whose equations of motion generalize the ADHM constraints on the ordinary instanton moduli space.

4 Non-perturbative corrections from localization formulæ

The corrections induced by D-instantons can be encoded in a non perturbative prepotential \( \mathcal{F}_{\text{non-p.}}(\Phi) \), which, taking into account the different instanton configurations and their multiplicity, can be written
as
\[ \mathcal{F}_{n,p}(\Phi) = 12 \mathcal{F}^{(m=0)}(\Phi) + 4 \mathcal{F}^{(m=1)}(\Phi). \] (5)

The prepotentials \( \mathcal{F}^{(m)}(\Phi) \) can be expressed as integrals over the “centered” moduli space (containing all moduli except the “center of mass” coordinates \( x \) and \( \theta \)) of the instanton branes. To compute \( \mathcal{F}_{n,p}(\Phi) \) we exploit the fact [5] that, after suitable deformations of the instanton action, the modular integrals localize around isolated points in the instanton moduli space. To obtain explicit formulas we first take \( \Phi = \text{diag}(a_1, a_4, -a_1, -a_4) \), where \( a_i \) are constant expectation values along the Cartan directions of \( U(4) \), and then consider the \( \epsilon \)-deformed instanton partition function
\[ Z^{(m)}(a, \epsilon) = \sum_k q_k Z_k^{(m)}(a, \epsilon) = \sum_k q_k \int dM_{k,m} e^{-S^{\text{mod}}_{\epsilon}(M_{k,m}, a)}. \] (6)

where \( S^{\text{mod}}_{\epsilon} \) is obtained by deforming the moduli action with Lorentz breaking terms parameterized by four parameters \( \epsilon_i \) describing rotations along the four Cartan directions of \( \text{SO}(4) \times \text{SO}(4) \). From the string perspective, these deformations can be obtained by switching on suitable RR background fluxes on the D7-branes, as shown in [7, 11, 12]. Notice that integrals in (6) run over all moduli, including \( x \) and \( \theta \). In presence of the \( \epsilon \)-deformations it is rather easy to see that the integration over the super-space yields a volume factor growing as \( 1/(\epsilon_{1,2}) \) in the limit of small \( \epsilon_{1,2} \). Therefore, to obtain the integral over the centered moduli this factor has to be removed. In addition, we have to notice that the \( k \)-th order in the \( \epsilon \)-expansion receives contributions not only from genuine \( k \)-instanton configurations but also from disconnected ones. Thus, we are led to consider
\[ \mathcal{F}^{(m)}(a, \epsilon) = \epsilon_1 \epsilon_2 \log Z^{(m)}(a, \epsilon). \] (7)

The prepotential will be extracted from \( \mathcal{F}^{(m)}(a, \epsilon) \) by sending \( \epsilon_1 \to 0 \) and \( a \to \Phi \).

The localization procedure is based on the co-homological structure of the instanton moduli action which is exact with respect to a suitable BRST charge \( Q \), namely
\[ S_{\text{mod}} = Q \Xi. \] (8)

We can choose as \( Q \) any component of the supersymmetry charges preserved on the brane system. Of course, since these charges transform as spinors of \( \text{SO}(4) \times \text{SO}(4) \), the choice of \( Q \) breaks this symmetry to the \( \text{SU}(2)^3 \equiv \text{SU}(2)_- \times \tilde{\text{SU}}(2)_- \times \text{diag} \left[ \text{SU}(2)_+ \times \tilde{\text{SU}}(2)_+ \right] \) subgroup which preserves this spinor. After this identification is made we can see that all the moduli but \( \chi \) form BRST doublets, which we will schematically denote as \( (\phi, \psi \equiv Q\phi) \), and the moduli action can indeed be written in the form [5].

To localize the integral over moduli space, it is necessary to make the charge \( Q \) equivariant with respect to all symmetries, which in our case are the gauge symmetry \( U(k) \times U(4) \times U(4) \) and the residual Lorentz symmetry \( \text{SU}(2)^3 \). After the equivariant deformation, the charge \( Q \) becomes nilpotent up to an element of the symmetry group. In the basis provided by the weights \( \tilde{q} \equiv (\tilde{q}_{U(k)}, \tilde{q}_{U(4)}, \tilde{q}_{U(m)}, \tilde{q}_{\text{SU}(2)}) \), \( Q \) acts diagonally
\[ Q\phi_\eta = \psi_\eta, \quad Q\psi_\eta = \Omega_\eta \phi_\eta, \] (9)

where \( \Omega_\eta = \tilde{\chi} \cdot \tilde{q}_{U(k)} + \tilde{a} \cdot \tilde{q}_{U(4)} + \tilde{b} \cdot \tilde{q}_{U(m)} + \tilde{c} \cdot \tilde{q}_{\text{SU}(2)} \), parametrize the equivariant deformation in terms of the Cartan components of the group parameters \( \tilde{\chi}, \tilde{a}, \tilde{b}, \tilde{c} \).

After the complete localization the integral is given by the (super)-determinant of \( Q^2 \) evaluated at the fixed points of \( Q \) [5,13], and its explicit expression can be deduced by considering, for each modulus \( \phi \) in

\[ ^1 \text{The Cartan directions of the residual Lorentz group } \text{SU}(2)^3 \text{ are parametrized by } \epsilon_J (J = 1, \ldots, 4) \text{ subject to the constraint } \epsilon_1 + \epsilon_2 + \epsilon_3 + \epsilon_4 = 0. \]
Table 2 BRST structure and symmetry properties of the D(-1)/D3/D7 moduli. The first two columns report the transformation properties under the symmetry groups. The last column collects the eigenvalues $\tilde{e} \cdot \tilde{q}_{\SU^2}$ for the positive weights $\tilde{q}$’s specified in the third column.

| $(\phi, \psi)$ | $U(k) \times U(4) \times U(m)$ | $\SU^2$ | $\tilde{e} \cdot \tilde{q}_{\SU^2}$ |
|---------------|---------------------------------|---------|-----------------------------|
| $(B_\ell, M_\ell)$ | $(\text{adj}, 1, 1)$ | $(2, 1, 2)$ | $\epsilon_1, \epsilon_2$ |
| $(B_\mu, M_\mu)$ | $(\mathbb{1}, 1, 1)$ + h.c. | $(1, 2, 2)$ | $\epsilon_3, \epsilon_4$ |
| $(N_{\alpha \alpha}, D_{\alpha \alpha})$ | $(\mathbb{1}, 1, 1)$ + h.c. | $(2, 2, 1)$ | $\epsilon_2 + \epsilon_3, \epsilon_1 + \epsilon_3$ |
| $(N_{m}, d_{m})$ | $(\text{adj}, 1, 1)$ | $(1, 1, 3)$ | $0_{R}, \epsilon_1 + \epsilon_2$ |
| $(\chi, \eta)$ | $(\text{adj}, 1, 1)$ | $(1, 1, 1)$ | $0_{R}$ |
| $(\mu', h')$ | $(\mathbb{1}, 1, 1)$ + h.c. | $(1, 1, 1)$ | $0$ |
| $(w_\alpha, \rho_\alpha)$ | $(\mathbb{1}, 1, 1)$ + h.c. | $(1, 1, 2)$ | $(\epsilon_1 + \epsilon_2)/2$ |
| $(\mu_i, h_i)$ | $(\mathbb{1}, 1, 1)$ + h.c. | $(1, 2, 1)$ | $(\epsilon_3 - \epsilon_2)/2$ |

The integral over $\chi_i$ in this expression has to be thought of as a multiple contour integral, according to the prescription introduced in Ref. [13].

In order to obtain the non-perturbative prepotential from the partition function $Z^{(m)}(a, b, \epsilon)$, we set $b_\ell = 0$, since the D3-branes are fixed at one of the orbifold fixed-points and we take the limit $\epsilon_1 \rightarrow 0$ to remove the Lorentz breaking deformations. A simple inspection of the explicit results for $\log Z^{(m)}(a, \epsilon)$ [9] shows that this expression diverges as $1/(\epsilon_1 \epsilon_2 \epsilon_3 \epsilon_4)$ in this limit. Such a divergence is typical of interactions in eight dimensions, where the $\mathcal{N} = 2$ super-space volume grows like $\int d^8 x d^8 \theta \sim 1/(\epsilon_1 \epsilon_2 \epsilon_3 \epsilon_4)$. These contributions can be thought of as coming from regular D(-1)-instantons moving in the full eight-dimensional world-volume of the D7-branes and can in fact be associated to a universal quartic prepotential $F_{IV}(a)$ [7] defined as

$$F_{IV}(a) = \lim_{\epsilon_1 \rightarrow 0} \epsilon_1 \epsilon_2 \epsilon_3 \epsilon_4 \log Z^{(m)}(a, \epsilon).$$

(11)

We can then extract a finite quadratic prepotential by subtracting the divergence coming from $F_{IV}(a)$:

$$F_{II}^{(m)}(a) = \lim_{\epsilon_1 \rightarrow 0} \left( \epsilon_1 \epsilon_2 \log Z^{(m)}(a, \epsilon) - \frac{1}{\epsilon_3 \epsilon_4} F_{IV}(a) \right).$$

(12)

Since the moduli measure is dimensionless no dynamically generated scale may appear and the contributions at all instanton numbers must be constructed only out of the $a$’s; we find

$$F_{II}^{(m=0)}(a) = \left( - \sum_{i<j} a_i a_j \right) q + \left( \sum_{i<j} a_i a_j - \frac{1}{4} \sum_i a_i^2 \right) q^2 + \left( - \frac{4}{3} \sum_{i<j} a_i a_j \right) q^3 + \cdots,$$

$$F_{II}^{(m=1)}(a) = \left( - 3 \sum_{i<j} a_i a_j \right) q + \left( \sum_{i<j} a_i a_j + \frac{7}{4} \sum_i a_i^2 \right) q^2 + \left( 4 \sum_{i<j} a_i a_j \right) q^3 + \cdots.$$
We can now promote the vacuum expectation values $a$’s to the dynamical superfield $\Phi(x, \theta)$ and determine $S_{n,p}(\Phi)$ taking into account the contributions from the various $m = 0, 1$ configurations according to Eq. (6). Performing the $\theta$-integration, we then obtain the quadratic non-perturbative action:

$$S_{n,p} = 4 \int d^4x \left[ 2(\text{tr} F)^2 - \text{tr} F^2 \right] q^2 + O(q^4) + \text{c.c.}$$

and read the non-perturbative part of the holomorphic couplings. Considering also the perturbative contributions written above we have

$$f = -4t - 4q^2 + O(q^4), \quad f' = -8 \log \eta(U)^2 + 8 q^2 + O(q^4).$$

We would like to stress that the vanishing of the contributions at the one and three instanton level is due to the non-trivial cancellations between contributions coming from configurations $a$) and $b$).

The heterotic model dual to the Type I’ description of the previous sections can be built from the $U(16)$ compactification of the SO(32) heterotic string on $T_4/\mathbb{Z}_2$ (with standard embedding of the orbifold curvature into the gauge bundle) and further reduced on $T_2$ with Wilson lines that break $U(16)$ to $U(4)^4$. The gauge kinetic terms in this heterotic set-up are corrected at 1-loop by an infinite tower of world-sheet instantons wrapping $T_2$, which are dual to the D-instantons of the type I’ theory [15] and read [9]:

$$f = -i S + 8 \log \left( \frac{\eta(T)^2}{\eta(T)^4} \right), \quad f' = -8 \log \eta(U)^2 + 8 \log \left( \frac{\eta(T)^2}{\eta(T)^4} \right).$$

These couplings are exact and do not receive any kind of corrections beyond 1-loop. Therefore they must contain all information, both perturbative and non-perturbative, on the corresponding type I’ couplings, including the (exotic) instanton corrections computed above. Indeed, when we expand for large values of $T$ and use the duality map that relate the Kähler modulus of the heterotic theory $T$ to the axio-dilaton $\lambda$ of the type I’ model: $T/4 \leftrightarrow \lambda$, these heterotic formulas predict no instanton corrections at $k = 1$ and $k = 3$, and a relative coefficient $-2$ between the $k = 2$ corrections to $f$ and $f'$, in perfect agreement with the results obtained in the type I’ setting.

We regard these results as a nice and non-trivial confirmation of the validity of the exotic instanton calculus, which can then be applied with confidence also to four-dimensional theories and to models for which the heterotic dual is not known or does not exist.

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