Non-Tachyonic Type 0B Orientifolds, Non-Supersymmetric Gauge Theories and Cosmological RG Flow

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Abstract

We discuss gauge theories on D3 branes embedded in special non-tachyonic orientifolds of the 0B string theory. In general, they correspond to non-supersymmetric $SU(N)$ gauge theories with scalars in the adjoint representation and spinors in the (anti-)symmetric representation. We study these theories via the AdS/CFT correspondence and present evidence of their relation to $\mathcal{N}=4$ SYM in the planar limit. We also discuss finite $N$ properties, focusing in particular on the renormalization group flow. Up to two loops, the logarithmic running of the gauge coupling is described by the dilaton tadpole and the cosmological constant that naturally emerge on the string theory side.
1 Introduction

The study of the strong coupling regime of gauge theories is one of the most difficult problems in high energy physics. In recent years there was much progress in this direction - mostly in the study of supersymmetric theories. Part of this knowledge was gained with the help of string theory.

Lately, it was conjectured that the natural description of strongly coupled large $N$ $\mathcal{N} = 4$ SYM theory is in terms of classical supergravity. This is known as the AdS/CFT correspondence \[1\] (for a recent review and references see \[2\]). Although this conjecture is very powerful, its original version is restricted to theories with maximal supersymmetry. It seems very difficult to study non-supersymmetric and non-conformal gauge theories in this framework, since it is not known how to construct their dual gravity backgrounds.

The first step in the direction of studying non-supersymmetric conformal theories using the AdS/CFT duality was made in \[3\], where it was shown that one can break supersymmetry completely, still within a large $N$ CFT. The idea was to study the type IIB string on $AdS_5 \times X_5$, where $X_5$ is a five-dimensional space of the type $S^5/\Gamma$. This construction yields conformal theories with a reduced number of supersymmetries or no supersymmetry at all. The conformal invariance of the resulting large $N$ theories was then proved both by string theory \[4, 5\] and by direct field theory techniques \[6\].

A similar idea motivated the study of type 0B string theory \[7\], a $(-1)^F$ orbifold of type IIB. The orbifold breaks supersymmetry and the resulting string theory, purely bosonic, contains the NS-NS sector of the IIB theory and doubled sets of R-R fields. In addition, it contains a tachyon from the twisted sector, while the NS-R and R-NS fermionic sectors are projected out. More recently this theory was also related to an orbifold of M-theory \[8\].

The theory that lives on the untwisted (“dyonic”) D3 branes is an $SU(N) \times SU(N)$ gauge theory with 6 adjoint scalars for each of the gauge groups and 4 $\bar{\psi} \psi$ + 4 $\bar{\psi} \psi$ Weyl spinors. Recently, the construction of \[8\] was used to study this gauge theory \[11\], that is conformal at large $N$ \[9, 10\]. Other type 0 conformal models were studied in \[12, 13, 14\].

Although this gauge theory seems to be perfectly sensible, one might suspect that, since the dual string theory contains a tachyon, some instabilities should appear on the gauge theory side, and indeed, using the AdS/CFT correspondence, it was shown that at strong coupling the dimension of the
twisted operator $T \sim \text{tr} F_1^2 - \text{tr} F_2^2$ becomes complex \cite{15}. The presence of the tachyon makes the study of the dual gravity difficult, since the detailed structure of its coupling to the R-R fields is not known. Another (somewhat) disturbing feature is the doubling of the gauge factors in all models that live on “dyonic” D-branes \cite{13}. This is due to the doubling of R-R fields, and is seemingly unrelated to the tachyon issue.

Is there any way to project out the tachyon from the type 0 spectrum, so that one has a stable dual gauge theory?

The construction of tachyon-free models was first discussed in \cite{16}, where sign ambiguities in the Klein bottle amplitude \cite{17} were used to generalize the standard $\Omega$ projection. The non-tachyonic projection required a $U(32)$ gauge group on the D9-brane, together with chiral fermions in the $\overline{6}$ and $\overline{5}$ representations. An exhaustive analysis of tachyon-free orbifold compactifications was initiated in \cite{18}. In general, for the (would be) supersymmetric orbifolds new tachyons appear in the twisted sector of the closed string theory and, for a geometric action of the orbifold group, can not be projected out by any (generalized) $\Omega$ projection. The $T^4/Z_2$ and $T^6/Z_3$ cases, however, are exceptional, since no tachyons appear in the twisted sectors and thus one can find appropriate $\Omega$ projections that lead to non-tachyonic open descendants. The analysis was then extended to $Z_N \times Z_M$ orbifolds \cite{19} and to the non-supersymmetric $T^6/Z_2$ one \cite{20}. A geometric interpretation of the signs in the Klein bottle amplitude, corresponding to the generalized $\Omega$ projections, has been given in \cite{21}: the non-tachyonic projection involves the combined action of the world-sheet parity with the right world-sheet fermion number, $(-1)^{f_R}$.

In the simplest version of the construction, the orientifold of 0B leads to a string theory that is very similar to the bosonic part of type IIB \cite{15}. From now on, we shall refer to it as the type 0' string theory. The theory on the D3 branes close to O'3 planes of this sort was first studied in \cite{14}. It is an $SU(N)$ gauge theory with the field content in table 1.
Most of the present paper is devoted to the study of this gauge theory and its gravity dual. We will show that the theory is conformal at large $N$, using both string theory and field theory arguments. Since the bulk theory does not contain any tachyon, the gauge theory should be stable. Moreover, its study via gravity will be rather simple, and we will find an $AdS_5 \times S^5$ solution of the tree level gravity action.

We are also interested in finite $N$ non-supersymmetric theories. In particular, we consider D3 branes of non-tachyonic orientifolds of type 0B string theory, i.e. D3 branes in a background of $O'7$ planes and $32$ D7 branes (together with their images). At finite $N$, the theory is no longer conformal. On the string theory side, the breakdown of conformal invariance is due to a dilaton tadpole and to a cosmological constant that appear once $g_{st}$ is turned on. We use these stringy corrections to exhibit a qualitative logarithmic running of the coupling constant, up to two loops on the field theory side. The dilaton tadpole affects the one loop beta function, while the cosmological constant affects the two loop beta function. We also comment on the relation of the dilaton equation of motion to the all orders beta function equation.

The organization of the paper is as follows. In section 2 we describe the orientifold projections allowed in 0B models. In particular, we focus on the projections that remove the tachyon from the 0B string spectrum, and construct the gauge theory on two different systems: $O'3$-D3 in type 0B and D3 branes on type 0' backgrounds. We also discuss more involved models originating from D3 branes at orbifold singularities, that lead to other non-supersymmetric field theories on non-tachyonic string backgrounds. Sections 3 and 4 are devoted to the study of the gauge theories introduced in section 2 via the AdS/CFT correspondence. We also study the field theory renormalization group flow via gravity (see \cite{22, 23, 24, 25} for a recent discussion and

| SU(N) | 
|-------|
| vector | adj. |
| 6 scalars | adj. |
| 4 fermions | $\mathbb{1} + \mathbb{5}$ |

Table 1: The field content of the non-supersymmetric gauge theory on $N$ D3 branes close to $O'3$ planes of 0B string theory.
Another example of RG flow). In section 5 we give an independent argument for the conformal invariance of the large $N$ theory, based on the work of [27]. In section 6 we study the relation between $\mathcal{N} = 4$ SYM and the theory on the D3 branes. We suggest that these theories become equivalent in the planar limit. In section 7 we discuss briefly other field theories constructed in section 2.

2 Type 0 orientifolds

In this section we construct some brane configurations that will be further analyzed in the following sections. In particular, we construct gauge theories corresponding to D3 branes embedded in non-tachyonic 0B orientifolds and O’3-D3 systems in 0B theories, where the orientifold planes involve the non-tachyonic projection of $[15]$. To this end, let us recall some known facts about the type 0B string and its orientifolds. The torus partition function

$$ T = \int_{\mathcal{F}} \frac{d^2 \tau}{\tau_2^6} \frac{1}{|\eta|^16} \left\{ |O_8|^2 + |V_8|^2 + |S_8|^2 + |C_8|^2 \right\} $$

(1)

can be conveniently written in terms of level-one $SO(8)$ characters, whose expressions in terms of Jacobi theta functions are

$$ O_8 = \frac{1}{2\eta^4}(\vartheta_3^4 + \vartheta_4^4), \quad V_8 = \frac{1}{2\eta^4}(\vartheta_3^4 - \vartheta_4^4), $$

$$ S_8 = \frac{1}{2\eta^4}(\vartheta_2^4 - \vartheta_4^4), \quad C_8 = \frac{1}{2\eta^4}(\vartheta_2^4 + \vartheta_4^4). $$

From $T$ one can easily extract the low-lying excitations, that comprise a tachyon, the graviton, the dilaton and the Kalb-Ramond field from the NS-NS sectors, and two scalars, two 2-forms and one 4-form from the R-R sectors. The doubling of the R-R sector with respect to type IIB implies that the 0B string has two types of D$p$-branes for odd $p$.

To this left-right symmetric theory one can associate orientifolds $[28, 29]$, that in general contain two different types of orientifold planes and D-branes, since the R-R sector is now doubled. A generic feature of the standard 0B orientifolds is the presence of the tachyon in the unoriented closed sector. Actually, using the properties of the underlying two-dimensional conformal
field theory \[17\], it was shown in \[16\] that the 0B string admits three inequivalent orientifold projections. These are related to additional (world-sheet and space-time) symmetries, and, in particular, the combination $\Omega(-1)^{f_R}$, with $\Omega$ the world-sheet parity and $f_R$ the world-sheet fermion number, removes the tachyon from the closed unoriented spectrum \[16\], thus stabilizing the vacuum. Moreover, the $\Omega(-1)^{f_R}$ projection keeps only one copy of the R-R sector, thus allowing for the presence of only one type of Dp-branes for any odd $p$ (as in the IIB theory). An exhaustive analysis of orbifold compactifications has shown that one can remove the closed tachyon only in the $T^4/Z_2$ and $T^6/Z_3$ \[18\], in the $T^6/Z_2 \times Z_2$ \[19\] and in the non-supersymmetric $T^6/Z_2$ \[20\] orbifolds.

2.1 The O’3-D3 system

Let us describe the simplest model on D3 branes at “non-tachyonic” orientifold planes. The O’3 planes project the D3 brane gauge field by $\Omega' = \Omega(-1)^{f_R} I_6$, where $I_6$ inverts the six (non-compact) coordinates transverse to the D3-brane.

On a flat Minkowski background, the annulus and Möbius amplitudes compatible with the type 0B string are

$$A = \frac{1}{2} \int_0^\infty \frac{dt}{t^3} \frac{1}{\eta^8} \left( 2N \bar{N} V_8 - (N^2 + \bar{N}^2) S_8 \right),$$

$$M = \pm \frac{1}{2} \int_0^\infty \frac{dt}{t^3} \frac{1}{\bar{\eta}^8} (N + \bar{N}) \hat{S}_8,$$

where the $\eta$-function and the $SO(8)$ characters depend on the modulus of the doubly-covering torus ($it/2$ for the annulus and $(1 + it)/2$ for the Möbius strip). One can simply extract the massless fields that live on the O’3-D3 system: for $N$ branes they are gauge bosons and six scalars in the adjoint of a $U(N)$ gauge group, as well as four Weyl spinors in the $\mathbf{1} + \mathbf{1}$ representations. Actually, since the space transverse to the D3 branes is non-compact, the R-R tadpole need not be imposed, and one has the freedom to revert the sign in the Möbius amplitude thus obtaining spinors in symmetric representations.

The $\Omega'$ projection acts also on the bulk theory. In fact, as we will see in the next section, in the spirit of the AdS/CFT correspondence, gauge invariant operators on D3 branes are related to harmonics of the bulk fields
on $AdS_5 \times X_5$, with $X_5$ some appropriate Einstein manifold respecting the symmetry of the brane configuration. In our case $X_5 = S^5/Z_2 = \mathbb{R}P^5$, where the freely-acting $Z_2$ corresponds to the $\Omega'$ projection, and thus only the harmonics of the projected (bulk) fields survive [34, 2].

Aside from the massless excitations, one can extract other interesting informations from the amplitudes (1) and (2). For instance, since the fundamental $(N)$ and anti-fundamental ($\bar{N}$) representations of unitary groups have identical dimensions, one can easily show that the annulus amplitude vanishes, as

$$A \sim N^2 (V_8 - S_8) \sim N^2 \left( \vartheta_3^4 - \vartheta_4^4 - \vartheta_2^4 \right) \equiv 0.$$  

This translates into a vanishing force between pairs of D3 branes, as a result of a perfect cancellation between the gravitational (NS-NS) attraction and the R-R repulsion.

Since gauge/string theory is non-supersymmetric, one would expect the presence of a dilaton tadpole and/or of a cosmological constant on the string theory side. Unfortunately, it is hard to perform explicit string theory computations in curved backgrounds and therefore, in the following we will present explicit models where these corrections are under control and a precise comparison between string theory and gauge theory can be worked out even at finite $N$.

### 2.2 D3 branes in type 0' backgrounds

The configuration that now we are going to construct corresponds to D3 branes in non-tachyonic orientifolds of the type 0B string. In particular the background is an eight-dimensional toroidal compactification of the $\Omega'$ projected type 0B. Before discussing the gauge theory arising on the D3 branes, let us introduce the background string theory and discuss its main features.

The starting point is the halved toroidal partition function of the 0B string on a two-dimensional Narain lattice with momenta

$$p_{i,L,R}^i = \frac{m_i}{R} \pm \frac{n_i R}{\alpha'}.$$  

(4)
The non-tachyonic Klein bottle projection (associated with $\Omega(-1)^{F_R} I_2$),

$$\mathcal{K} = \frac{1}{2} \int_0^\infty \frac{dt}{t^5} \frac{1}{\eta^8} \left(-O_8 + V_8 - S_8 + C_8\right) \sum_n e^{-\pi t R^2 n^2/\alpha'} ,$$

leads to the following massless spectrum: the eight-dimensional metric tensor, the moduli of the internal $T^2$, the dilaton, two abelian vectors related to the mixed components of the NS-NS $B$-field, and full ten-dimensional 0-form, 2-form and self-dual 4-form $R$-$R$ potentials.

The open sector that completes the orientifold construction [28] is encoded in the annulus and Möbius amplitudes

$$\mathcal{A} = \frac{1}{2} \int_0^\infty \frac{dt}{t^5} \frac{1}{\eta^8} \sum_{j,k=1}^4 \left[V_8 \left(\bar{M}_j M_k e^{-\pi R^2 (n+a_j-a_k)^2/\alpha'} + \bar{M}_j M_k e^{-\pi R^2 (n-a_j+a_k)^2/\alpha'}\right) - S_8 \left(M_j M_k e^{-\pi R^2 (n+a_j+a_k)^2/\alpha'} + \bar{M}_j \bar{M}_k e^{-\pi R^2 (n-a_j-a_k)^2/\alpha'}\right)\right] ,$$

$$\mathcal{M} = \frac{1}{2} \int_0^\infty \frac{dt}{t^5} \frac{1}{\eta^8} \hat{S}_8 \sum_{j=1}^4 \left(M_j e^{-\pi R^2 (n+2 a_j)^2/\alpha'} + \bar{M}_j e^{-\pi R^2 (n-2 a_j)^2/\alpha'}\right) ,$$

where the $M_j$ and $\bar{M}_j$ describe Chan-Paton multiplicities. Here we have decided to equally distribute the D7 branes on the four $O'7$ planes

$$a_1 = (0,0) , \quad a_2 = (1/2,0) , \quad a_3 = (0,1/2) , \quad a_4 = (1/2,1/2) ,$$

in order to cancel tadpoles locally. After imposing (R-R) tadpole cancellation the gauge theory on each set of D7 branes has a $U(8)$ gauge group with 2 scalars in the adjoint and two spinors in the $\mathbb{H}$ and $\bar{\mathbb{H}}$.

Actually, for $\Omega'$ orientifolds one is always left with a non-vanishing dilaton tadpole coming from the $(V_8$ character in the) annulus amplitude, as can be appreciated from the massless contributions to the transverse amplitudes:

$$\tilde{\mathcal{K}} \sim -\frac{2^6}{2} S_8 ,$$

$$\tilde{\mathcal{A}} \sim \frac{2^{-6}}{2} \left[\left(\sum_{j=1}^4 M_j + \bar{M}_j\right)^2 (V_8 - S_8) - \left(\sum_{j=1}^4 M_j - \bar{M}_j\right)^2 (O_8 - C_8)\right] ,$$

$$\tilde{\mathcal{M}} \sim \frac{2}{2} \left(\sum_{j=1}^4 M_j + \bar{M}_j\right) \hat{S}_8 .$$
Taking into account R-R tadpole conditions the numerical value of the dilaton tadpole is then

\[
C^2 = 2^{-6} \left( \sum_{j=1}^{4} M_j + \bar{M}_j \right)^2 = 2^{-6} 2^{12} = 8^2 .
\]  

(7)

This has to be contrasted with the supersymmetric case, where the dilaton couples both to D-branes and O-planes and its tadpole vanishes as a result of R-R tadpole conditions. Moreover, the torus, Klein bottle and Möbius amplitudes are non-vanishing, and this translates into the emergence of a one-loop cosmological constant. Thus, dilaton tadpole and cosmological constant induce for the dilaton field a potential of the form

\[
V(\Phi) = C e^{-\Phi} + 2\Lambda ,
\]

(8)

whose properties will be further analyzed in section 4.

We are now ready to add \( N \) D3 branes in the background of this type 0’ string. The reduced \( SO(8) \rightarrow SO(4) \times SO(4) \) transverse Lorentz symmetry translates into a breaking of the \( SO(8) \) characters into products of \( SO(4) \) ones:

\[
O_8 = O_4 O_4 + V_4 V_4 , \quad S_8 = S_4 S_4 + C_4 C_4 , \quad V_8 = V_4 O_4 + O_4 V_4 , \quad C_8 = S_4 C_4 + C_4 S_4 .
\]

(9)

(10)

Putting all the D3 branes close to the orientifold plane at the origin (\( a_1 = (0,0) \)) results in the following (massless) deformation of the annulus and Möbius amplitudes:

\[
\mathcal{A} \sim (M_1 \bar{M}_1 + N \bar{N})(V_4 O_4 + O_4 V_4)
\]

\[
- \frac{1}{2} \left( M_1^2 + M_1^2 + N^2 + \bar{N}^2 \right)(C_4 C_4 + S_4 S_4)
\]

\[
+ (M_1 N + M_1 \bar{N}) C_4 O_4 - (M_1 N + \bar{M}_1 \bar{N}) C_4 O_4 ,
\]

\[
\mathcal{M} \sim (M_1 + \bar{M}_1 + N + \bar{N}) C_4 C_4 + (M_1 + \bar{M}_1 - N - \bar{N}) S_4 S_4 .
\]

(11)

(12)

The field theory on the D3 branes is thus an \( SU(N) \) gauge theory with the charged matter in table 2. (We are neglecting the \( U(1) \) factor associated with the center of mass of the brane configuration.)
Table 2: The field content of the non-supersymmetric gauge theory on $N$ D3 branes in type $0'\text{'}$ string theory.

### 2.3 D3 branes at orbifold singularities

We can now turn to more complicated models, like O'3-D3 systems on orbifold singularities. The analysis of [18] suggests to restrict oneself to the $C^2/Z_2$ and $C^3/Z_3$ singularities, since the other cases would introduce additional (twisted) tachyons in the bulk theory. For supersymmetric strings, $Z_2$ and $Z_3$ singularities would correspond to $\mathcal{N} = 2$ and $\mathcal{N} = 1$ gauge theories [3], and in our case one finds an $SU(N) \times SU(N)$ gauge group for O'3-D3 systems on a $C^2/Z_2$ singularity. The charged matter comprises four scalars in the $(\square,\square)$ and $(\square,\bar{\square})$ and Weyl spinors in the $(\bar{\square},\bar{\square})$, $(\bar{\square}+\bar{\square},1)$ and $(1,\bar{\square}+\bar{\square})$ representations.

The $C^3/Z_3$ singularities lead to a chiral spectrum containing an $SU(N-4) \times SU(N) \times SU(N)$ gauge group, as well as Weyl spinors in the $(\bar{\square}+\bar{\square},1)$, $(1,\bar{\square}+\bar{\square})$, (1,1) representations. Moreover, there are three copies of Weyl spinors in the $(\bar{\square},1)$, $(\bar{\square},1)$, $(1,\bar{\square},1)$ and $(1,1,\bar{\square})$ representations and three copies of complex scalars in the $(\bar{\square},\bar{\square},1)$, $(\bar{\square},1,\bar{\square})$, $(1,\bar{\square},\bar{\square})$ and $(1,\bar{\square},\bar{\square})$ representations.

All these models have the same main features of the O'3-D3 system in flat Minkowski space-time: there is no brane-brane interaction, and one expects that a (half-loop) dilaton tadpole and a (one-loop) cosmological constant be
3 The low energy action and the AdS/CFT correspondence

The orientifold projection introduced in the previous section leads to a non-supersymmetric string theory with a rather simple low energy action. The action of $\Omega' = \Omega(-1)^fR$ on the oscillators of the underlying CFT is [21]

\[
\begin{align*}
\Omega' \alpha^\mu_n \Omega' &= \tilde{\alpha}^\mu_n \\
\Omega' \Psi^\mu_I \Omega' &= \tilde{\Psi}^\mu_I \\
\Omega' \tilde{\Psi}^\mu_I \Omega' &= \Psi^\mu_I \\
\Omega' |0\rangle_{NS-NS} &= - |0\rangle_{NS-NS}
\end{align*}
\]

Therefore, whereas type 0B string theory contains a tachyon and two sets of R-R fields, the orientifold projection removes the tachyon and one set of R-R fields from the bulk. It picks the R-R two-form from the $(R^+, R^+)$ sector and the zero-form and four-form from the $(R^-, R^-)$ sector. Thus the resulting ten-dimensional $0'$ theory is similar to the bosonic part of type IIB

\[
S^{(\Omega)}_0 = \int d^{10}x \sqrt{-G} \left( e^{-2\Phi} (R + 4\partial_\mu \Phi \partial^\mu \Phi) - \frac{1}{2} \left( |F_1|^2 + |F_3|^2 + \frac{1}{2} |F_5|^2 \right) \right)
\]

Note that the tachyon is projected out and hence we expect the dual theory on the D3 branes to be more stable than for other type 0 models. We will make this statement more precise in the following. Note also that the NS-NS 2-form $B_{\mu\nu}$ is projected out. A final remark concerning the action (17) is that, in contrast to the unprojected type 0 action, we have only one set of R-R fields. The doubling of the R-R fields in the type 0B action leads to a semi-simple $SU(N) \times SU(N)$ gauge theory on the dyonic D3 branes. Here we expect that the dual theory on the brane have only one gauge factor. Indeed, the large $N$ theory has a simple $SU(N)$ gauge group.

In order to study the type $0'$ theory in the presence of D3 branes we first perform 2 T-dualities along the 8, 9 directions and than place $N$ D3 branes. The resulting background consists of $O'7$ and 8 D7 branes and localized $N$ D3 branes. (Here we are interested in the degrees of freedom living close to
the orientifold plane at the origin.) The near horizon background, $AdS_5 \times S^3$ (where the $S^3$ sits inside the $S^5$) respects the $SU(2)$ flavor symmetry of the model (see table 2). At infinite $N$ the 8 D7 branes can be neglected and the flavor symmetry is enhanced to $SU(4)$. Accordingly, in this limit one finds the standard $AdS_5 \times S^5$ background metric

$$ds^2 = d\tau^2 + e^{2\tau} dx_i^2 + d\Omega_5^2$$ (18)

(with a radius $R = 1$). The $AdS_5$ part of the metric corresponds to an $SO(2,4)$ symmetry and suggests that the large $N$ non-supersymmetric dual theory is conformal. The $S^5$ part corresponds, in the present case, to an $SU(4)$ global flavor symmetry.

The situation for the O$^3$-D3 system is apparently quite different. In this case the bulk theory is the full 0B string and (in the large $N$ limit where the charge of the orientifold 3-plane can be neglected and the configuration is indeed extremal) the near horizon geometry is $AdS_5 \times RP^5$. However, the $Z_2$ identification in $RP^5$ keeps only the invariant harmonics of the bulk fields, and the resulting spectrum coincides with the ($\Omega'$) projected bulk theory. Further evidence for the conformal invariance of these models will be given in sections 5 and 6.

The AdS/CFT correspondence can be used to study the stability of these non-supersymmetric models[15]. In the type 0 (unprojected) theory, the bulk theory contains a tachyon. The squared-mass of the tachyon is shifted to a positive value at small radius of the $AdS$ metric [9]. However, at large values of the radius, when the bulk theory is flat, we really have a tachyon. It is related, by holography, to the twisted operator

$$T \sim tr F_1^2 - tr F_2^2,$$ (19)

and therefore it was argued [15] that at large ‘t Hooft coupling $\lambda$ (large $AdS$ radius) the dimension of $tr F_1^2 - tr F_2^2$ will become complex, $\Delta(\lambda) = 2 + \sqrt{4 - 2\sqrt{2}}$. This indicates an instability in the gauge theory at strong coupling.

The present bulk theory contains no tachyon. Therefore, the corresponding gauge theory should be stable even at strong coupling. This conclusion is not surprising since, as we shall see in section 6, this theory is related to $\mathcal{N} = 4$ SYM.
4 Cosmological renormalization group flow

In the previous section we started a study of the gauge theory using the bulk theory. The discussion was restricted to large $N$, where the gauge theory is expected to be conformal.

Here we would like to study $\frac{1}{N}$ corrections to the gauge coupling of the theory on the D3 branes in the type 0' background (the field content is described in table 2). Our model is the non-supersymmetric analogue of the one studied in [31, 32]. At finite $N$ the coupling is expected to run, and indeed the value of the one-loop beta-function coefficient is

$$b_1 = -8.$$  \hspace{1cm} (20)

In calculating (20) we took into account the contribution of the 8 flavors of fermions and scalars (arising from the open strings that connect the D3-D7 branes). Note that the value (20) is $\frac{1}{N}$ suppressed with respect to the general expected behavior of beta function coefficients and that the theory is IR free. The higher order beta function coefficients are also expected to behave as

$$b_n \sim N^{n-1},$$  \hspace{1cm} (21)

as befits with the conformal invariance of the model at large $N$.

The beta function equation (up to two loops) reads

$$\frac{dg}{d \log u} = -\frac{g^3}{(4\pi)^2} b_1 - \frac{g^5}{(4\pi)^4} b_2 + ...,$$ \hspace{1cm} (22)

and its solution is

$$g^2(u) = \frac{8\pi^2}{b_1 (\log \frac{u}{u_0} + \frac{b_2}{b_1^2} \log \log \frac{u}{u_0})} + ... .$$ \hspace{1cm} (23)

In order to study these $\frac{1}{N}$ corrections on the gauge theory side, we should go beyond the tree-level classical gravity on the string theory side. Before entering into details, we would like to mention that several attempts to study the running of the coupling were already made in various type 0 models [33]. Thus, the coupling of the tachyon to the dilaton was used to extract the logarithmic running of the gauge coupling in several non conformal models, but it was difficult to find precise results, since the detailed structure of
the tachyon coupling to the dilaton is not known. Moreover, these results depend on the minimum of the tachyon potential (the function that describes the dilaton-tachyon couplings), and it is not clear whether such a minimum really exists.

In the present case, the spectrum does not contain a tachyon, and therefore the mechanism that we present seems more controlled. An important remark is that at finite \( N \) the background will no longer be an \( AdS_5 \times S^3 \), due to the back-reaction of the dilaton on the metric. However, we shall see that the leading \( \frac{1}{N} \) correction is determined by the tree level \( AdS_5 \times S^3 \) background.

When \( g_{st} \) is turned on, the classical action (17) acquires corrections, and we find two important contributions: a cosmological constant due to the loss of supersymmetry and a dilaton tadpole arising from the annulus diagram (see section 2). Generally, the cosmological constant is expected to be proportional to \( N \) (the number of branes), but is also suppressed by \( g_{st}^2 \sim \frac{1}{N^2} \) with respect to the tree level action. The dilaton tadpole is suppressed by \( g_{st} \sim \frac{1}{N} \). The corrected action is

\[
S = S_0^{(\Omega)} + \int d^8x \sqrt{-G} (Ce^{-\Phi} + 2\Lambda),
\]

and the eight-dimensional equations of motion in the string frame are

\[
e^{-2\Phi} \left( 8(\nabla\Phi)^2 - 8\nabla^2\Phi - 2R \right) - Ce^{-\Phi} = 0,
\]

\[
e^{-2\Phi} (R_{\mu\nu} + 2\nabla_{\mu}\nabla_{\nu}\Phi) - \frac{1}{2}G_{\mu\nu} \left( \frac{C}{2}e^{-\Phi} + 2\Lambda \right) + f_{\mu\nu}(RR) = 0,
\]

where the traceless tensor \( f_{\mu\nu} \) depends on the R-R fields, and in particular on the R-R 5-form field strength. Substituting the trace of (26) in (25), we obtain for the dilaton the following equation:

\[
\nabla^2 e^{-2\Phi} = \frac{5C}{2}e^{-\Phi} + 8\Lambda.
\]

Clearly, \( \Phi = \Phi_0 \) is no longer a solution. In principle, one should also study the equation of the metric. However, as an approximation, we will assume that the background is still \( AdS_5 \times S^3 \). This assumption is valid (in the string frame, where the metric does not depend on the dilaton) because the
left hand side of (27) is dominated by the tree level background. The back-
reaction of the dilaton on the metric will be of higher order in $\frac{1}{N}$. Assuming
$\Phi = \Phi(\tau)$ and substituting (18) in (27), we obtain
\[ \partial^2_\tau e^{-2\Phi} + 4\partial_\tau e^{-2\Phi} = \frac{5C}{2} e^{-\Phi} + 8\Lambda. \] (28)
The Ansatz
\[ e^{-\Phi} = A\tau + B \log \tau, \] (29)
gives an approximate solution at large values of $\tau$, if
\[ A = \frac{5C}{16}, \] (30)
\[ B = \frac{16\Lambda}{5C} - \frac{5C}{64}. \] (31)
Generally, the $\frac{5C}{64}$ in (31) is expected to be negligible with respect to $\frac{16\Lambda}{5C}$ in
the large $N$ limit, since $\Lambda \sim O(N)$.
Recall that the coordinate $\tau$ is related to the energy scale $u$, through the
standard $AdS_5$ metric parameterization, by
\[ \tau = \log \frac{u}{u_0}, \] (32)
while the dilaton is related to the gauge coupling by
\[ g^2YM \sim g^{st} = e^\Phi. \] (33)
Thus we obtain
\[ \frac{1}{g^2YM} \sim A \log \frac{u}{u_0} + B \log \log \frac{u}{u_0}, \] (34)
with a nice qualitative agreement between gravity (34) and field theory cal-
culations (23) up to two loops. The dilaton tadpole determines the one-loop
beta function, while the cosmological constant determines the two-loop beta
function. Remarkably, in our model both the dilaton tadpole (4) and $b_1$ (24)
are proportional to the number of D7 branes. This result is encouraging,
since it links the AdS/CFT correspondence to more realistic models where
conformal invariance and supersymmetry no longer exist.
It might look surprising that we find an agreement between gravity and
perturbative Yang-Mills. After all, we do not expect that dual descriptions
in terms of gluons and gravitons apply at the same time. The reason for
the qualitative agreement is the following: As we shall see in section 6, the
non-supersymmetric models (table 1 and table 2) converge at infinite \(N\) to \(\mathcal{N} = 4\) SYM. It is believed that in this case, the background is unaffected by
\(\alpha'\) corrections. In the present case, we add a perturbation of \(\frac{1}{N}\). Our hope is
that for large enough \(N\) and fixed energy \(u\), our results are unchanged.

Finally, note that the dilaton equation of motion reproduces qualitatively,
at least to leading order in \(\frac{1}{N}\), all orders of the running coupling equation.
Generally, in the presence of an orientifold, the string theory effective action
should receive corrections of any power of \(g_{\text{st}}\) (even and odd), and therefore
the dilaton equation (27), should read

\[
\nabla^2 e^{-2\Phi} = \sum_{n=-1}^{\infty} C_n e^{n\Phi}. \tag{35}
\]

With the identification \(g_{\text{st}} \sim g^2\) and an \(AdS_5 \times X_5\) background, we find

\[
\partial^2 \frac{1}{g^4} + 4\partial \frac{1}{g^4} = \sum_{n=-1}^{\infty} C_n g^{2n}. \tag{36}
\]

The two derivative term (the first term) in (36) gives a negligible contribution
at large \(\tau\) and large \(N\), since the additional derivative with respect to \(\tau\) yields
terms which are \(\frac{1}{\tau}\) with respect to one derivative contributions. An example is
the \(\frac{5C}{64}\) in eq. (31), that comes from the two derivative term. Upon neglecting
the two derivative term, one obtains

\[
\frac{dg}{d\log u} = -\sum_{n=-1}^{\infty} \frac{C_n}{16} g^{2n+5}, \tag{37}
\]

the all orders beta function equation for the Yang-Mills coupling !

5 String theory amplitudes and the planar limit

In this section we want to use the results of [27] to give plausible arguments
about the finiteness of the gauge theories previously introduced. For simplicity
we will focus on the O'3-D3 system in the large \(N\) limit, although similar
arguments apply also to the other model.
In a theory with D-branes and orientifold planes string amplitudes involve world-sheets with $b$ boundaries, $c$ crosscaps and $h$ handles, and are weighted by

$$ (Ng_{st})^b g_{st}^{c2h-2} = \lambda^{2h-2+b+c} N^{c-2h+2}. \tag{38} $$

In particular, for the large $N$ behavior of $M$-point correlators of fields charged under the gauge group there are two classes of diagrams that have to be considered: diagrams without handles and crosscaps, and diagrams with handles and/or crosscaps, that from (38) are subleading in the large $N$ limit. In the same way, non-planar diagrams with only boundaries are suppressed with respect to planar ones.

For planar diagrams with $b$ boundaries and $M$ external lines, the boundary conditions must satisfy (in the notation of [27])

$$ \gamma^{\mu_2}_{\alpha_1} = \prod_{s=2}^b \gamma^{\mu_s}_{\alpha_s}, \tag{39} $$

where the $\gamma$’s define the action of a (generic) orbifold group on the Chan-Paton matrices and $s = 1$ refers to the outer boundary, while $s = 2, \ldots, b$ refer to the inner boundaries.

If the orbifold group is trivial, in the large $N$ limit the $M$-point amplitude behaves as in the $\mathcal{N} = 4$ one without orientifold planes (with gauge group $U(N)$). For non-trivial (space-time) orbifold groups, it has been shown in [27] that diagrams with twisted Chan-Paton matrices are vanishing or suppressed in the large $N$ limit, thus leading to the conjecture that supersymmetric O’3-D3 systems on orbifold singularities have the same leading behavior as the $\mathcal{N} = 4$ $U(N)$ gauge theory.

Actually, the model that we are considering can be thought of as an orbifold of the supersymmetric O’3-D3 system, where the orbifold group corresponds to the internal $(-1)^F$ symmetry. This projection leads to two different $\gamma$ matrices such that:

$$ \text{tr}(\gamma_1) = (N + \bar{N}), \quad \text{tr}(\gamma_2) = i(N - \bar{N}). \tag{40} $$

It is then evident how to apply the previous arguments to the present case. Non-planar and unoriented amplitudes are suppressed in the large $N$ limit, and among the planar diagrams, the ones with $\gamma_1$ boundary conditions on both external and internal boundaries give the same behavior of $\mathcal{N} = 4$ $U(N)$
gauge theory. On the other hand, diagrams with at least one inner boundary with a $\gamma_2$ insertion are vanishing since they are proportional to

$$\sum \text{tr}(\gamma_{a_1}^{1} \lambda_1 \ldots \lambda_M) \prod_{s=2}^{b} \text{tr}(\gamma_{a_s}^{\mu})$$  \hspace{1cm} (41)

and at least one of the inner $\gamma$'s is

$$\gamma_{a_s}^{\mu} = \gamma_2 \Rightarrow \text{tr}(\gamma_{a_s}^{\mu}) = i(N - \bar{N}) = 0.$$  \hspace{1cm} (42)

We thus conclude that in the large $N$ limit our model indeed behaves like the $N = 4$ supersymmetric $U(N)$ gauge theory. Similar arguments can also be applied to $O'3$-D3 systems of type 0B on orbifold singularities.

6 Field-theory analysis - The relation to $\mathcal{N} = 4$ SYM

Arguments in favor of conformal invariance can also be presented in a field theory context. Our gauge theory corresponds to an $SU(N)$ gauge group with 6 scalars in the adjoint representation and 4 spinors in the $\mathbb{1}$ representation and its conjugate. The string theory analysis of the previous sections proved that this theory is conformal in the large $N$ limit. In fact, as we shall see, the large $N$ limit of this theory is exactly $\mathcal{N} = 4$ SYM.

The similarity to $\mathcal{N} = 4$ SYM is rather clear. The only difference between the two theories is that the spinors transform in a different representation, and therefore supersymmetry is broken explicitly in the type 0 case at finite $N$. However, we will argue that at large $N$ adjoint spinors and spinors in $\mathbb{1} + \overline{\mathbb{1}}$ give the same contribution to any amplitude.

Our arguments are similar to those given in [6] for the case of orbifold field theories, where it was shown that the planar graphs of the $\mathcal{N} = 4$ SYM and of the daughter theory, an orbifold truncation of $\mathcal{N} = 4$, are the same up to a rescaling of the gauge coupling in the latter.

Let us therefore examine the relation between the present non-supersymmetric gauge theory and the $\mathcal{N} = 4$ theory. Since we are interested in the planar limit, it will be convenient to use 't Hooft’s double index notation, where fields that transform in the adjoint representation carry two lines with
arrows in opposite directions. This is the case, since the adjoint representation is obtained by a multiplication of fundamental and anti-fundamental representations. Similarly, a field transforming in the anti-symmetric representation carries two lines with arrows in the *same direction*, since the anti-symmetric representation is obtained by multiplication of two fundamentals. Thus the fermion propagator in the non-supersymmetric theory will be represented by two parallel lines with arrows pointing in the same direction, while bosonic propagators are represented by two parallel lines with arrows pointing in opposite directions.

The Feynman rules of the non-supersymmetric theory and of the supersymmetric theory are very similar. The only differences are in the vertices which couple fermions and bosons, such as the coupling of the gluon to the fermions or the Yukawa coupling. In these graphs, one of the arrows of the fermion lines is reversed. The fermionic propagator and the fermion-boson coupling are described in figure 1 below.

![Figure 1: Feynman rules for $\mathcal{N} = 4$ SYM and for the non-supersymmetric theory. a) Fermion propagator and fermion-boson vertex. b) Feynman rules for the supersymmetric theory: both fermions and bosons are in the adjoint representation. c) In the non-supersymmetric orientifold theory bosons transform in the adjoint representation whereas fermions transform in the $\mathfrak{u}$ representation (and in the $\mathfrak{u}$ rep.).](image)

Although we will not present a rigorous proof, we will argue that this change of the Feynman rules will not affect any calculation that only depend on planar diagrams: every planar diagram in $\mathcal{N} = 4$ SYM can be deformed
so that it represents the non-supersymmetric theory and vice versa. As a result, there is a one-to-one correspondence between the planar graphs of both theories.

The rule is simple: one should follow the fermionic double lines and reverse the arrow of one of the lines, so that the two arrows along the fermionic lines point in the same direction. It is easy to see that this rule is consistent with the Feynman rules, and in particular the interaction of bosons and fermions is unaffected by this deformation: the bosons will remain in the adjoint representation.

Let us give a simple example. Consider the fermionic loop contribution to the one-loop vacuum polarization in both theories (figure 2). It is clear that the planar diagram of the supersymmetric theory can be deformed into the planar diagram of the non-supersymmetric theory.

![Figure 2](image)

Figure 2: a) The fermionic loop contribution to the vacuum polarization. b) A planar contribution in the supersymmetric theory. Fermions are in the adjoint representation. c) The arrows in the internal loop were reversed such that the fermions are now in the \[ \text{adjoint} \] (and \[ \text{anti-adjoint} \]) representations.

One might be tempted to believe that all diagrams (not only the planar ones) of the two theories are equivalent. This, however, is not true. There are non-planar diagrams that exist only in one of the theories. For instance, the one loop beta function is \( b_1 = -\frac{16}{3} \) in the non-supersymmetric theory, but \( b_1 = 0 \) in the supersymmetric case.

The reason why the deformation works in the planar case is that planar graphs form a collection of non-intersecting lines. One should follow the fermionic line and simply reverse one of the arrows, and this procedure does not lead to inconsistencies. The fermionic lines can be closed loops, or lines with an end (in the case of external legs). Since lines do not intersect and fermion number is conserved, the deformation can not change the representations of the bosons, that remain in the adjoint.
The conclusion of this discussion is that in the infinite $N$ limit the non-supersymmetric theory becomes supersymmetric, since there is no difference between the anti-symmetric representation and the adjoint representation at large $N$. In this regime the theory described by the orientifold of type 0 becomes the $\mathcal{N} = 4$ SYM.

Since this theory becomes $\mathcal{N} = 4$ SYM in the planar limit, we arrive at several interesting conclusions. In this limit, the theory has BPS states, a moduli space of vacua and possesses Olive-Montonen duality. These features can be proved, in the infinite $N$ limit, by the relation to $\mathcal{N} = 4$, via Feynman diagrams or via the bulk string theory. The existence of a moduli space of vacua can be seen both calculating the force between parallel D3 branes using the cylinder diagram of the type 0$'$ string theory, or calculating the Coleman-Weinberg potential in the field theory. Both calculations yield a no force result, and thus the scalar v.e.v. is indeed a modulus. The existence of BPS states in field theory is due to the existence of such objects in IIB string theory, or can be seen by a field theory calculation in the planar limit. Finally, the fact that the large $N$ beta-function is zero is either due to S-duality of the IIB action or to S-duality in $\mathcal{N} = 4$ SYM.

A more interesting question is whether some of these features survive at finite $N$. Since the theory is not supersymmetric and even the Bose-Fermi degeneracy is destroyed at finite $N$ (we have $8(N^2 - 1)$ bosons and $8(N^2 - N)$ fermions), we do not expect a supersymmetric behavior away from the planar limit. Clearly S-duality is broken, since the $\beta$ function is no longer zero. However, some of these features might be broken more softly than naively expected.

One example is the existence of a moduli space. The cylinder diagram yields a zero force between D3 branes, as a result of the balance between the exchange of the NS-NS fields $G_{\mu\nu}$ and $\Phi$ and of the R-R four form. The calculation of the Coleman-Weinberg one-loop effective potential yields the same result. The bosonic contribution to the potential is \[ V_{\text{bos}}(r) = \frac{1}{4\pi^2}r^2\log\frac{r^2}{\Lambda^2}, \] where $r$ is the v.e.v. of the scalar (the separation between the D-branes) and $\Lambda$ is some cutoff. The fermionic contribution is $V_{\text{fer}}(r) = -V_{\text{bos}}(r)$, and therefore the net force is zero, to this order. Note that, in contrast with the $SU(N) \times SU(N)$ gauge theory on the D3 branes of type 0B, where some
directions were unstable and the gauge group was dynamically broken to its $U(1)$’s \[U(1)\], in the present case there is no $\frac{1}{N}$ potential at tree level. In particular there is no instability and there seems to be a moduli space even at finite $N$. However, as we stated above, higher order perturbation theory will probably lift the vacuum energy at generic points in the moduli space.

7 Non-conformal large $N$ field theories

The procedure described in the previous section is quite general. Any supersymmetric theory with fermions in the adjoint or bi-fundamental representations can be deformed into a non-supersymmetric theory by changing the (color) representation of the fermions, in such a way that the large $N$ theory will converge to its parent supersymmetric theory. The rules are: adjoint $\rightarrow \begin{array}{c} \begin{array}{c} \text{adj} \\ + \end{array} \\ \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array} \rightarrow \begin{array}{c} \begin{array}{c} \text{bi-fund} \\ \text{representation} \end{array} \\ \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array}$ and $\begin{array}{c} \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array} \rightarrow \begin{array}{c} \begin{array}{c} \text{bi-fund} \\ \text{representation} \end{array} \\ \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array}$. In particular, non-conformal $\mathcal{N} = 2$ and $\mathcal{N} = 1$ theories can be used as parents.

Without the background of a string theory, we will not be able to say much about the non-supersymmetric daughter theory. In fact, generally we expect a tachyonic instability at strong coupling[15], and for this reason we restrict ourself to theories that live on D-branes with a non-tachyonic bulk theory. Two such theories were constructed in section 2 (other cases were recently discussed within the type I [36, 37] and type IIB [38] frameworks).

The orientifold projection of the type 0B theory on the $Z_2$ orbifold will converge in the planar limit to the $\mathcal{N} = 2$ $SU(N) \times SU(M)$ gauge theory with hypermultiplets in the $\begin{array}{c} \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array}$ and in the $\begin{array}{c} \begin{array}{c} \text{adj} \\ \text{bi-fund} \\ \text{representation} \end{array} \end{array}$ representations. Thus, up to $\frac{1}{N}$ corrections, $\mathcal{N} = 2$ results apply to this model.

The second example is the $Z_3$ orbifold. It is interesting to take the limit when two of the gauge couplings of the theory are set to zero and the rank of these groups is equal. We also take only one copy of the matter multiplet. In this limit, we will have an $SU(N_c) \times SU(N_f) \times SU(N_f)$ theory with the non-supersymmetric field content of table 3.

21
\[ SU(N_c) \times SU(N_f) \times SU(N_f) \]

|        | adj. | 1   | 1   |
|--------|------|-----|-----|
| vector |      |     |     |
| fermion|     | 1   | 1   |
| scalar |     | 1   |     |
| scalar |     |     |     |
| fermion|     |     | 1   |
| fermion|     | 1   |     |

Table 3: The field content of an “\( \mathcal{N} = 1 \) ” theory

In the large \( N_c, N_f \) limit (with the ratio \( N_f/N_c \) fixed), this theory converges to \( \mathcal{N} = 1 \) SQCD. Therefore, in the planar limit this theory is dual in the infra-red to a ”magnetic” \( SU(N_f - N_c) \) theory\[39\]. It would be interesting if some sort of similar duality holds at finite \( N \) as well. Note that, in contrast to the type 0 unprojected theory\[13\], the one loop beta function receives a finite \( N \) correction (due to the orientifold)

\[
b_1 = 3N_c - N_f - \frac{4}{3},
\]

and that the \( U_R(1), U_R^3(1) \) anomalies match only in the planar limit. Therefore, the duality map should receive a \( \frac{1}{N} \) modifications.

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**References**
[1] J. M. Maldacena, “The Large N Limit of Superconformal Field Theories and Supergravity”, Adv. Theor. Math. Phys. 2 (1998) 231-252, hep-th/9711200.

S. S. Gubser, I. R. Klebanov and A. M. Polyakov, “Gauge Theory Correlators from Non-Critical String Theory”, Phys. Lett. B428 (1998) 105-114, hep-th/9802109.

E. Witten, “Anti De Sitter Space and Holography”, Adv. Theor. Math. Phys. 2 (1998) 253-291, hep-th/9802150.

[2] O. Aharony, S. S. Gubser, J. Maldacena, H. Ooguri and Y. Oz, ”Large N Field Theories, String Theory and Gravity”, hep-th/9905111.

[3] S. Kachru and E. Silverstein, “4d Conformal Field Theories and Strings on Orbifolds”, Phys. Rev. Lett. 80 (1998) 4855, hep-th/9802183.

[4] M. Bershadsky, Z. Kakushadze and C. Vafa, “String Expansion as Large N Expansion of Gauge Theories”, Nucl. Phys. B523 (1998) 59, hep-th/9803070.

[5] A. Lawrence, N. Nekrasov and C. Vafa, “On Conformal Theories in Four Dimensions”, Nucl. Phys. B533 (1998) 199, hep-th/9803013.

[6] M. Bershadsky and A. Johansen, “Large N limit of orbifold field theories”, Nucl. Phys. B536 (1998) 141, hep-th/9803249.

[7] I.R. Klebanov and A.A. Tseytlin, “D-Branes and Dual Gauge Theories in Type 0 Strings”, Nucl. Phys. B546 (1999) 155, hep-th/9811035.

[8] O. Bergman and M. R. Gaberdiel, “Dualities of Type 0 Strings”, JHEP 9907 (1999) 022, hep-th/9906053.

[9] I.R. Klebanov and A.A. Tseytlin, “A non-supersymmetric large N CFT from type 0 string theory”, JHEP 9903 (1999) 015, hep-th/9901101.

[10] N. Nekrasov and S.L. Shatashvili, “On Non-Supersymmetric CFT in Four Dimensions”, hep-th/9902110.

[11] R. Blumenhagen, C. Kounnas and D. Lust, “Continuous Gauge and Supersymmetry Breaking for Open Strings on D-branes”, hep-th/9910094.
[12] M. Billo', B. Craps and F. Roose, “On D-branes in Type 0 String Theory”, Phys. Lett. B457 (1999) 61, hep-th/9902196.

[13] A. Armoni and B. Kol, “Non-Supersymmetric Large N Gauge Theories from Type 0 Brane Configurations”, JHEP 9907 (1999) 011, hep-th/990608.

[14] R. Blumenhagen, A. Font and D. Lust, “Non-Supersymmetric Gauge Theories from D-Branes in Type 0 String Theory”, hep-th/9906101.

[15] I. R. Klebanov, “Tachyon Stabilization in the AdS/CFT Correspondence”, hep-th/9906220.

[16] A. Sagnotti, “Some Properties of Open String Theories”, hep-th/9509080; A. Sagnotti, “Surprises in Open String Perturbation Theory”, hep-th/9702093.

[17] D. Fioravanti, G. Pradisi and A. Sagnotti, “Sewing Constraints and Non-Orientable Open Strings”, Phys. Lett. B321 (1994) 349, hep-th/9311183.

G. Pradisi, A. Sagnotti and Y.S. Stanev, “The Open Descendants of Non-Diagonal SU(2) WZW Models”, Phys. Lett. B356 (1995) 230, hep-th/9506014; “Planar Duality in SU(2) WZW Models”, Phys. Lett. B354 (1995) 279, hep-th/9503207; “Completeness Conditions on Boundary Operators in 2D CFT”, Phys. Lett. B381 (1996) 97.

[18] C. Angelantonj, “Non-tachyonic open descendants of the 0B string theory”, Phys. Lett. B444 (1998) 309, hep-th/9810214.

[19] K. Förger, “On Non-tachyonic $Z_N \times Z_M$ Orientifolds of Type 0B String Theory”, hep-th/9909010.

[20] R. Blumenhagen and A. Kumar, “A Note on Orientifolds and Dualities of Type 0B String Theory”, Phys. Lett. B464 (1999) 46, hep-th/9906234.

[21] R. Blumenhagen, A. Font and D. Lust, “Tachyon-free Orientifolds of Type 0B Strings in Various Dimensions”, hep-th/9904069.
[22] L. Girardello, M. Petrini, M. Porrati and A. Zaffaroni, “Novel Local CFT and Exact Results on Perturbations of N=4 Super Yang Mills from AdS Dynamics”, JHEP 9812 (1998) 022, hep-th/9810126.

[23] M. Porrati and A. Starinets, ”RG Fixed Points in Supergravity Duals of 4-d Field Theory and Asymptotically AdS Spaces”, Phys. Lett. B454 (1999) 77, hep-th/9903085.

[24] J. de Boer, E. Verlinde and H. Verlinde, “On the Holographic Renormalization Group”, hep-th/9912012.

[25] E. Verlinde and H. Verlinde, “RG-Flow, Gravity and the Cosmological Constant”, hep-th/9912018.

[26] I. R. Klebanov, N. A. Nekrasov, “Gravity Duals of Fractional Branes and Logarithmic RG Flow”, hep-th/9911096.

[27] Z. Kakushadze, “Gauge Theories from Orientifolds and Large N Limit”, Nucl. Phys. B529 (1998) 157, hep-th/9803214.

[28] A. Sagnotti, in “Non-Perturbative Quantum Field Theory”, eds G. Mack et al (Pergamon Press, 1988), p. 521.

[29] M. Bianchi and A. Sagnotti, “On the systematics of Open String Theories”, Phys. Lett. B247 (1990) 517.

[30] E. Witten, ”Baryons And Branes In Anti de Sitter Space”, JHEP 9807 (1998) 006, hep-th/9805112.

[31] O. Aharony, J. Pawelczyk, S. Theisen and S. Yankielowicz, ”A Note on Anomalies in the AdS/CFT Correspondence”, Phys. Rev. D60 (1999) 066001, hep-th/9901134.

[32] E. Gava, K. S. Narain and M. H. Sarmadi, ”Instantons in N = 2 Sp(N) Superconformal Gauge Theories and the AdS/CFT Correspondence”, hep-th/9908125.

[33] J.A. Minahan, “Glueball Mass Spectra and Other Issues for Supergravity Duals of QCD Models”, JHEP 9901 (1999) 020, hep-th/9811156.
I.R. Klebanov and A.A. Tseytlin, “Asymptotic Freedom and Infrared Behavior in the Type 0 String Approach to Gauge Theory”, Nucl. Phys. B547 (1999) 143, hep-th/9812083;

J. A. Minahan, “Asymptotic Freedom and Confinement from Type 0 String Theory”, JHEP 9904 (1999) 007, hep-th/9902074;

A. Armoni, E. Fuchs and J. Sonnenschein, ”Confinement in 4D Yang-Mills Theories from Non-Critical Type 0 String Theory”, JHEP 9906 (1999) 027, hep-th/9903096;

M. Alishahiha, A. Brandhuber and Y. Oz, “Branes at Singularities in Type 0 String Theory”, JHEP 9905 (1999) 024, hep-th/9903180.

[34] K. Zarembo, “Coleman-Weinberg Mechanism and Interaction of D3-Branes in Type 0 String Theory”, Phys. Lett. B462 (1999) 70, hep-th/9901106.

[35] A. A. Tseytlin and K. Zarembo, “Effective potential in non-supersymmetric SU(N) × SU(N) gauge theory and interactions of type 0 D3-branes”, Phys. Lett. B457 (1999) 77, hep-th/9902093.

[36] I. Antoniadis, E. Dudas and A. Sagnotti, “Brane Supersymmetry Breaking”, Phys. Lett. B464 (1999) 38, hep-th/9908023.

[37] C. Angelantonj, I. Antoniadis, G. D’Appollonio, E. Dudas and A. Sagnotti, “Type I vacua with brane supersymmetry breaking”, hep-th/9911081.

[38] G. Aldazabal and A. M. Uranga, ”Tachyon-free Non-supersymmetric Type IIB Orientifolds via Brane-Antibrane Systems”, JHEP 9910 (1999) 024, hep-th/9908072;

A. M. Uranga, ”Comments on Non-Supersymmetric Orientifolds at Strong Coupling”, hep-th/9912145.

[39] N. Seiberg, “Electric-Magnetic Duality in Supersymmetric Non-Abelian Gauge Theories”, Nucl. Phys. B435 (1995) 129, hep-th/9411149.