Scale Invariance of Dirac Condition $g_e g_m = 1$ in Type 0 String Approach to Gauge Theory

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Abstract

In this letter we shall discuss a description of non-supersymmetric four-dimensional Yang-Mills theory based on Type 0 strings recently proposed by Klebanov and Tseytlin. The three brane near-horizon geometry allows one to study the UV behaviour of the gauge theory. Following Minahan and Klebanov and Tseytlin we shall discuss how the gravity solution reproduces logarithmic renormalization of coupling constant $g_e$ extracted from quark-antiquark potential and then show that effective coupling constant $g_m$ describing monopole-antimonopole interactions is of zero-charge type and Dirac condition $g_e g_m = 1$ is scale invariant in logarithmic approximation.

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1 Introduction

Non-perturbative description of Yang-Mills fields in four dimensions is still one of the most challenging problems in modern theoretical physics. One of the most popular views is that in the large $N$ limit this description must be based on some kind of string theory. In spite of the huge amount of work on the subject in the last two decades the proper description of $SU(N)$ gauge theories in the large $N$ limit is still an open problem. Nevertheless, it was a remarkable progress in the last two years which gave us a deeper understanding of non-perturbative aspects of gauge theory using new ideas in string theory. The duality of $N = 4$ supersymmetric Yang-Mills to ten dimensional supergravity on $AdS_5 \times S^5$ has been used to obtain exact results in the large $N$ limit of the strongly coupled superconformal gauge theory. This theory has vanishing beta function and is realized as the world-volume theory of $N$ coincident D3-branes of type IIB string theory. This set of branes causes the near horizon geometry of $AdS_5 \times S^5$ and the classical type IIB theory can be approximated by the compactified supergravity in the limit $\lambda_{IIB} \to 0$, $g_{eff}$ large but fixed.

The interaction between charged particles has been analyzed in this context. The massless sources might be delicate to deal with, since their long-range fields are exponentially suppressed due to conformal invariance. Hence, the efforts have been concentrated on the study of massive electric and magnetic particles. The computation of the Nambu-Goto action in an $AdS$ background for a static string configuration allowed Rey and Yee and Maldacena to find the coulomb potential between quark and antiquarks at zero temperature. Some results have been also obtained for finite temperature where a new confining branch appears. In Minahan extended the results of finite temperature to the cases of monopole-antimonopole and monopole-quark interactions. Here, he finds an attractive coulomb force for the monopole-antimonopole pair as a function of the coupling manifestly dual under the transformation $g_e \to 1/g_m$. Therefore, the Dirac condition corresponding to the gauge theory has arised as a result of stringy computations.

The number of interesting predictions due to the AdS/CFT conjecture has led to do research into a possible extension to the non-supersymmetric case. By heating up a maximally supersymmetric gauge theory all the supersymmetries get broken. In this approach the gauge theory is dual to near-extremal branes whose near-horizon geometry corresponds to a black hole in AdS space. Recently, a different approach was suggested by Polyakov: the Type 0 string theory in $d \leq 10$ could be used to extend the AdS/CFT duality to non-supersymmetric non-conformal field theories. This idea inspired the conjecture of Klebanov and Tseytlin: the existence of a duality between the non-supersymmetric four dimensional $SU(N)$ gauge theory coupled to 6 adjoint scalars fields and a background of Type 0 string theory involving a non-vanishing tachyon field.

The presence of such a tachyon instability could seem fatal. However it disappears due to a non-perturbative mechanism and the Type 0 theory can be used to describe a
tachyon-free gauge theory. In this non-supersymmetric theory the coupling constant is expected to depend on scale. Nevertheless, the linear logarithmically dependence with the scale which should appear for the effective coupling constant in accordance with the short-distance behaviour of the gauge theory [14] does not agree with the squared log dependence found for the effective string coupling [15], [16]. The problem of the running coupling constant in this theory was also discussed in [17] and [18].

In order to shed some light to this puzzle we will study in this letter the leading order for the effective electric and magnetic coupling constants. In section 2 we describe a near-horizon geometry in Type 0B theory caused by a set of N electric D3-branes. In section 3 we will obtain following calculations of Minachan [15] (see also [16]) the running electric coupling constant \( g_e(L) \) from the computation of Wilson loop describing quark-antiquark pair. In section 4 we will use a similar method based on the DBI action for a D-string to study the magnetic coupling constant \( g_m(L) \) describing monopole-antimonopole interaction and obtain the result that Dirac condition \( g_e(L)g_m(L) = 1 \) is scale invariant as it is expected from \( S \)-duality.

2 Near-horizon geometry for the Type 0B electrically charged D3-brane

In the near-horizon limit the IIB string theory has the geometry of \( AdS_5 \times S^5 \) with a self-dual 5-form whose flux through the \( S^5 \) sphere fixes the charge of the solution. This theory has been proved to be dual to the N=4 supersymmetric \( SU(N) \) gauge theory on the boundary \( S^4 \times S^5 \) whose conformal invariance is due to the non-running dilaton.

However, in type 0B string theory this conformal invariance is broken. This model has a closed string tachyon, all the fermionic partners have been removed and the R-R sector doubled. This doubling seems to be crucial for the theory to describe a tachyon-free gauge theory. This theory also have D-branes [19]. The coupling of the tachyon field \( T \) to the \( R - R \) form gauge field strength shifts the effective mass of \( T \) stabilizing in such a way the supergravity background [13].

The ansatz for the metric corresponding to the type 0B electrically charged D3-brane background proposed in [13] is the following

\[
ds^2 = e^{\frac{1}{2}\Phi} \left( e^{\frac{1}{2}\xi - 5\eta} d\rho^2 + e^{-\frac{1}{2}\xi} dx_{\perp}^2 + e^{\frac{1}{2} - \eta} d\Omega_5^2 \right),
\]

where \( \Phi, \xi, \eta \) are functions of \( \rho \), parameter related to the radial direction \( u \) transverse to the 3-brane world-volume \( x_{\perp} \). Then the action for the model can be described by a

\[^1\text{Unfortunately we became aware of the paper [18] only after our paper was submitted. In the paper [18] besides other things it is shown that there is a magnetic screening.}\]
Toda mechanical system

\[ S = \int d\rho \left[ \frac{1}{2} \dot{\Phi}^2 + \frac{1}{2} \dot{\xi}^2 - 5 \dot{\eta}^2 - V(\phi, \xi, \eta) \right] \quad (2.2) \]

\[ V = M^2 e^{\frac{1}{2} \phi + \frac{1}{2} \xi - 5 \eta} + 20 e^{- 4 \eta} - Q^2 f^{-1}(T)e^{-2\xi}. \quad (2.3) \]

In this potential the first term comes from the tachyon mass term \( M^2 = \frac{1}{2} T^2\text{vac} \), the second one represents the curvature of \( S^5 \) and the third one is due to the electric \( R - R \) charge. There exists also a zero-energy constraint

\[ \frac{1}{2} \dot{\Phi}^2 + \frac{1}{2} \dot{\xi}^2 - 5 \dot{\eta}^2 + V(\phi, \xi, \eta) = 0 \quad (2.4) \]

with \( Q \) the charge of the brane system proportional to the number of branes \( N \), \( T \) the tachyon field and \( f(T) \) a function given by

\[ f(T) = 1 + T + \frac{1}{2} T^2 + O(T^3). \quad (2.5) \]

For a large number of branes, the mass of the tachyon can be neglected and the tachyon can be shown to be a maximum of the function \((2.3)\), \( f(T)' = 0 \). Then,

\[ T = T_{\text{vac}} = -1, \quad f(T) = \frac{1}{2}. \quad (2.6) \]

It is clear that if \( T = 0 \) we get the constant dilaton and decoupled \( \xi \) and \( \eta \) fields

\[ \Phi = \Phi_0, \quad e^\xi = 2Q\rho, \quad e^\eta = 2\rho^{1/2}, \quad \rho = \frac{1}{u^4} \quad (2.7) \]

which leads to the standard \( R - R \) D3-brane solution [20]. However, when \( T \) is nonzero we can just find approximate solutions. In the UV the dilaton is expected to be slowly varying, at least compared to \( \xi \) and \( \eta \). Under this assumption, Minahan [15] solved the equations of motion of \((2.2)\) for \( \xi \) and \( \eta \) and found

\[ e^\xi = C_1 \rho, \quad e^\eta = C_2 \rho^{1/2}, \quad (2.8) \]

\( C_1 = 2Q \) and \( C_2 = 2 \) being two constants in a first order approximation. Then, using \((2.8)\) as inputs he obtained that

\[ \exp\left( \frac{\Phi}{2} \right) = \frac{C_0}{\log \rho/\rho_0}, \quad (2.9) \]

with \( C_0 = -8C_5^2/(T^2\sqrt{C_1}) \), is a leading order solution of the equation of motion for the dilaton. The integration constant \( \rho_0 \) is assumed to be \( \rho_0 >> 1 \) to assure that the gauge
theory length scale is much greater than the string scale. From (2.9) we can see that the
gauge theory coupling constant $1/g_{YM}$ is no longer constant and depends on the energy.
Thus, if we take the gauge coupling constant as $1/g = 4\pi/g_{YM}^2$, then its leading order
behaviour is

$$
\frac{1}{g_{YM}^2} = e^{-\Phi} = 2^{-12} Q (\log u/\epsilon)^2
$$

(2.10)

where it has been set $\rho = u^{-4}$ and $\epsilon << u$ is a lower limit for the energy. Then, the
coupling constant runs but with the unexpected power two for $\log u/\epsilon$ instead of the
linear dependence which appears for the effective coupling constant in QCD.

Finally the solution (2.9) can be used to computed the leading order correction to $C_1$
and $C_2$

$$
C_1 = 2Q \left( 1 + \frac{1}{\log u/\epsilon} \right) \quad C_2 = 2 \left( 1 + \frac{1}{\log u/\epsilon} \right)
$$

(2.11)

and the metric in the large $u$ limit

$$
ds^2 = \frac{2^6}{\sqrt{2^6 \log u/\epsilon}} \left[ \left( 1 - \frac{1}{8 \log u/\epsilon} \right) \frac{u^2}{Q} dx^2 + 9 \frac{1}{8 \log u/\epsilon} \frac{du^2}{u^2} + \right.

\left. \left( 1 - \frac{1}{8 \log u/\epsilon} \right) d\Omega^2_5 \right].
$$

(2.12)

In the Einstein frame, $ds^2 = e^{\Phi/2} ds^2_E$, the metric (2.12) corresponds to $AdS_5 \times S^5$ space
in the large $u$ limit but the lower the energy the more important the corrections are and
the curvature of $S^5$ becomes smaller than that of $AdS_5$ leading to a 10 dimensional space
with negative total curvature [15],[16] [3].

As we are interested in the high energy limit we will consider in the following that the
metric in the Einstein frame is still that of the $AdS_5 \times S^5$ space but in the string frame
we have a running dilaton

$$
ds^2 = \frac{2^5}{Q \log u/\epsilon} \left[ \frac{u^2}{2R_0^2} dx^2 + R_0^2 u^2 du^2 + R_0^2 d\Omega^2_5 \right],
$$

(2.13)

$R_0^2 = \sqrt{2Q}$ being the radius of $AdS_5$. Then, we are neglecting any sub-leading correction
which scales as $(\log u/\epsilon)^{-n}$.

### 3 The quark-antiquark interaction

A pair of massive quark-antiquark in the background of $N$ electric D3-branes can be realized as a string starting and ending on a D3-brane which has been separated

\[\text{In our units } \alpha' = 1\]

3 Higher order corrections for the Einstein metric can be found in [16].
an infinity distance from the set of \( N \) branes \([3][4]\). This separation breaks the group \( U(N + 1) \) to \( U(N) \times U(1) \) by giving an expectation value \( < H > \) to a Higgs field. The massive W-bosons which appear have a mass proportional to \( < H > \) and transform in the fundamental representation of \( U(N) \). They will play the role of quarks.

Denote the string coordinates by \( X^n(\tau, \sigma) \) where \( \tau, \sigma \) parameterize the string worldsheet, then the action for the string is

\[
S = \frac{1}{2\pi} \int d\tau d\sigma \sqrt{\det G_{mn} \partial_\alpha X^m \partial_\beta X^n},
\]

\( G_{mn} \) being the Euclidean metric in (2.13). Since we are interested in a static configuration independent of the \( \Omega_5 \) modes, we take \( \tau = t, \sigma = x \) where \( x \) is a direction along the three-branes. The action simplifies to

\[
S = \frac{1}{2\pi} \int dx \ e^{\Phi/2} \sqrt{\frac{U^4}{4R_0^4} X'^2 + \frac{1}{2} (\partial_x U)^2}
\]

where \( X' = 1, T \) is a constant due to the integration over \( t \) and

\[
e^{\Phi/2} = \frac{2^6}{\sqrt{Q} \log u/\epsilon}
\]

is the running dilaton we found in the previous section. Notice that, although the dilaton did not appear explicitly in the original action for the string (3.1), the running radius of the \( AdS \) space makes it enter the game. This fact will be relevant in order to provide the linear log dependence for the gauge coupling constant.

The action (3.2) does not show any explicit dependence on \( X \) and that allows us to solve its Euler-Lagrange equation of motion

\[
\partial_x \left[ e^{\Phi/2} \frac{\frac{U^4}{4R_0^4}}{\sqrt{\frac{U^4}{4R_0^4} + \frac{1}{2} (\partial_x U)^2}} \right] = 0.
\]

Since we are working in the approximation of a large value for \( \log u/\epsilon \), the \( \partial_x \log u/\epsilon \) is a sub-leading correction and the factor \( e^{\Phi/2} \) can be thought of as a constant. Thus, we can approach (3.4) to

\[
e^{\Phi/2} \partial_x \left[ \frac{\frac{U^4}{4R_0^4}}{\sqrt{\frac{U^4}{4R_0^4} + \frac{1}{2} (\partial_x U)^2}} \right] = 0
\]

which leads to

\[
\left[ \frac{\frac{U^4}{4R_0^4}}{\sqrt{\frac{U^4}{4R_0^4} + \frac{1}{2} (\partial_x U)^2}} \right] = \frac{U_0^2}{2R_0^2}
\]
where the leading order of $U_0$, the minimum value of $U$, is a constant to be determined. If we consider the quark placed at $x = L/2$ and the antiquark at $x = -L/2$ the minimum value for $U$ occurs at $x = 0$ by symmetry. Therefore we can write $x$ as a function of $U$

$$x = \sqrt{2} \frac{R_0^2}{U_0} \int_1^y \frac{dy}{y^2 \sqrt{y^4 - 1}} , \quad y = U/U_0$$

(3.7)

where $U_0$ must satisfy the condition

$$\frac{L}{2} = \sqrt{2} \frac{R_0^2}{U_0} \int_1^\infty \frac{dy}{y^2 \sqrt{y^4 - 1}} = \frac{R_0^2}{U_0} 2\pi^{3/2}. \Gamma(1/4)^2.$$ \quad (3.8)

Now taking into account (3.6), (3.7) and the value for $U_0$ (3.8) we can return to (3.2) to compute the total energy of the quark-antiquark configuration

$$E_{q\bar{q}} = \sqrt{\frac{2}{2\pi}} \frac{U_0}{2} \int_1^\infty dy e^{y^2/2} \frac{y^2}{\sqrt{y^4 - 1}}$$

$$\frac{2^{11/2} U_0}{Q\pi} \int_1^\infty dy \frac{1}{\log(yU_0/\epsilon) \sqrt{y^4 - 1}}.$$ \quad (3.9)

Since $\epsilon$ is an UV cutoff for the energy $\epsilon << U_0$ we can approach

$$\log(yU_0/\epsilon) = \log y + \log U_0/\epsilon \sim \log L_0/L$$ \quad (3.10)

where we have used the relation (3.8) with $L_0 = \frac{R_0^2\pi^{3/2}}{\alpha(1/4)^2}$ the UV cutoff for the distance $L << L_0$. Therefore we have found that the leading contribution to the energy is

$$E_{q\bar{q}} = \frac{2^{11/2} U_0}{Q\pi} \frac{1}{\Gamma(1/4)^2} \frac{1}{L \log (L_0/L)} \int_1^\infty dy \frac{y^2}{\sqrt{y^4 - 1}}.$$ \quad (3.11)

This result is infinity because we are including the masses of the quarks which correspond to strings stretched from the D3-brane at $U_{max}$ to the $N$ D3-branes at $U = 0$. Then, by integrating the energy up to $U_{max}$ we will subtract the regularized mass and find that the remaining energy is finite \[E_{q\bar{q}} = \frac{2^6 2Q \pi^2}{\Gamma(1/4)^4} \frac{1}{L \log (L_0/L)} \quad (3.12)\]

and, as claimed in [15], the effective coupling between a heavy quark and its antiquark does show the expected linear dependence

$$\frac{1}{g^2} \propto \log(L_0/L). \quad (3.13)$$

This result leads to reconsider the relation $1/g = 4\pi/g_{YM}^2$ between the string and the gauge coupling constants.
4 The monopole-antimonopole interaction

The pair of massive monopole-antimonopole in the background of $N$ electric D3-branes can be simulated in a similar way as the quark-antiquark pair of the previous section. For the monopole-antimonopole configuration we will take a type 0 electric D-string with both ends on a D3-brane separated an infinity distance from the set of $N$ branes. The same approach was used in [9] in case of $N = 4$ superconformal Yang-Mills theory. The world sheet action for the D-string spread in the 01 plane is the Born-Infeld action

$$S = \frac{1}{2\pi} \int dt dx \ e^{-\Phi} \sqrt{\det G_{mn} \partial_\alpha X^m \partial_\beta X^n},$$

(4.1)

$G_{mn}$ being the Euclidean metric in (2.13). We will suppose a static configuration independent of the $\Omega_5$ modes and, in such a way, this action can be written as follows

$$S = \frac{T}{2\pi} \int dx \ e^{-\Phi/2} \sqrt{\frac{U^4}{4R_0^4}} X'^2 + \frac{1}{2}(\partial_x U)^2$$

(4.2)

where $X' = 1$ and

$$e^{-\Phi/2} = \sqrt{Q \log u/\epsilon}$$

(4.3)

It is remarkable to notice that this action differs from that of the quark-antiquark (3.2) just in the factor $e^{-\Phi}$. This factor will be crucial to obtain a different log dependence for the monopole effective gauge coupling constant. Again the running radius of the AdS space changes the dilatonic factor of the action providing, as we will see, the linear log dependence for the effective magnetic coupling.

The action (4.2) does not depend explicitly on $X$ either and the Euler-Lagrange equation of motion for this variable reads

$$\partial_x \left[ e^{-\Phi/2} \frac{U^4}{4R_0^4} \right] = 0.$$  

(4.4)

We can suppose again that the factor $e^{-\Phi/2}$ has a slow variation and use

$$\frac{U^4}{4R_0^4} \left[ \frac{1}{\sqrt{\frac{U^4}{4R_0^4} + \frac{1}{2}(\partial_x U)^2}} \right] = \frac{U_0^2}{2R_0^2}$$

(4.5)

A tachyon-dependent function $k(T)$ might appear as a multiplicative factor [13]. However, in case of a constant tachyon it is not likely to play a relevant role in the determination of the log dependence of the gauge coupling constant.
to be able to write $x$ as a function of $U$ (3.7) and to fix the value of $U_0$ (3.8). Hence the minimum of the energy for the monopole-antimonopole configuration happens to be the same constant as that of the quark-antiquark case in a first approximation. However, from (3.4) and (4.4) we expect them to differ due to higher order corrections.

Now we can plug the results (3.7) and (3.8) into the expression (4.2) to obtain the monopole potential

$$E_{q\bar{q}} = \frac{\sqrt{2}U_0}{2\pi} \int_1^\infty dy \ e^{-\Phi/2} \frac{y^2}{\sqrt{y^4 - 1}} =$$

$$\frac{2^{13/2}U_0}{Q} \int_1^\infty dy \ \log \left(yU_0/\epsilon\right) \frac{y^2}{\sqrt{y^4 - 1}}.$$  (4.6)

Here, as in the previous section, taking into account that $\epsilon$ is an UV cutoff for the energy $\epsilon << U_0$ we can use (3.10) to get the leading infinite contribution to the energy

$$E_{m\bar{m}} = \frac{2^7\sqrt{Q\pi^{1/2}}}{{\Gamma(1/4)^2}} \log \left(L_0/L\right) \int_1^\infty dy \ \frac{y^2}{\sqrt{y^4 - 1}}.$$  (4.7)

In this case, the infinity is due to the masses of the monopoles, D-strings joining the D3-brane at $U_{max}$ to the N D3-branes at $U = 0$. The integration the energy up to $U_{max}$ will allow us to subtract the regularized mass and obtain a finite result

$$E_{m\bar{m}} = \frac{2^6\sqrt{2Q\pi^2}}{\Gamma(1/4)^4} \log \left(L_0/L\right) \frac{1}{L}.$$  (4.8)

Therefore, we have found that the effective coupling between a heavy monopole and antimonopole is of zero-charge type and shows the inverse log dependence

$$\frac{1}{g_m^2} \propto \frac{1}{\log \left(L_0/L\right)}$$  (4.9)

cmpared to $\frac{1}{g_e^2}$ (3.13). Hence the electric and magnetic coupling constants verify the Dirac condition

$$g_e(L)g_m(L) = 1$$  (4.10)

which has not been destroyed by running coupling constant.

5 Conclusion and Acknowledgments

The dual gravity description of non-supersymmetric gauge theories has been used in this paper to check that Dirac relation between electric and magnetic charges is RG invariant (at least in a logarithmic approximation). It seems that this is a necessary
element for the self-consistency of the theory. It was also shown that in this theory the running coupling constant is rather proportional to $\exp(-\Phi/2)$ than to $\exp(-\Phi)$ as one could expect. The fact that there is a square of logarithm in $\exp(-\Phi)$ was crucial to obtain the correct renormalization of the magnetic coupling constant. We did not answer the question why the coupling constant is square root of what we might expect naively and this seems to be open important question. It will be nice to find running coupling constant in this theory using other approaches, for example studying instanton effects and we plan to return to this issue in the future.

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