T-duality constraint on R-R couplings

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Abstract

It has been speculated that the metric, B-field and dilaton couplings in the low energy effective action of string theory on manifold with boundary at any order of $\alpha'$ may be found by imposing the gauge symmetries and by imposing the T-duality constraint on the effective action. We speculate that the Ramond-Ramond (R-R) couplings may also be found in this approach. In this paper, we perform the calculations explicitly at the supergravity level and found the democratic form of the R-R couplings. We also show explicitly that the presence of the Gibbons-Hawking-York term that is dictated by the T-duality constraint, is also necessary to have S-duality at the supergravity level.
One of the most exciting discoveries in perturbative string theory is T-duality [1, 2] which appears when one compactifies theory on a torus. It has been speculated that the invariance of the effective action of string theory under the standard gauge transformations and under non-standard T-duality transformations which receive \( \alpha' \)-corrections, may be used as constraints to construct the low energy effective action of the string theory on the manifolds with/without boundary [3, 4]. In this approach, using the field redefinitions freedom [5], one first constructs the most general gauge invariant couplings on the bulk and on the boundary. Then one reduces them on a circle. The reduced actions must be invariant under the standard Buscher rules [6, 7] plus their \( \alpha' \)-corrections [8, 9, 10, 11]. Using this approach, the effective action of the bosonic string theory up to order \( \alpha^3 \) for the spacetimes manifold without boundary have been found in [11, 12], and the effective action of the bosonic theory on the manifold with boundary at the leading order of \( \alpha' \) has been found in [4]. In particular, the Gibbons-Hawking-York term [13, 14] has been reproduced by the T-duality constraint [4], i.e.,

\[
S_0 + \partial S_0 = -\frac{2}{\kappa^2} \int d^D x \sqrt{-G} e^{-2\phi} \left( R + 4 \nabla_a \Phi \nabla^a \Phi - \frac{1}{12} H^2 \right) \mp \frac{4}{\kappa^2} \int d^{D-1} y \sqrt{\pm g} e^{-2\phi} K \tag{1}
\]

where \( K \) is the trace of the extrinsic curvature of the boundary in the string frame and \( g_{\alpha\beta} \) is induced metric. The plus (minus) sign in square root is when the boundary is spacelike (timelike).

Another T-duality based approach for constructing the \( D \)-dimensional effective action is the Double Field Theory [15, 16, 17, 18, 19] in which the effective action in 2\( D \)-space is constrained to be invariant under T-duality and under gauge transformations. The T-duality in this case, however, is the standard \( O(D,D) \) transformation without \( \alpha' \)-corrections whereas the gauge transformation is non-standard which receives \( \alpha' \)-corrections [20, 19, 21, 22]. This approach has been extended in [23] to type II superstring theories.

In this paper we would like to extend the first approach to the couplings in type IIA and type IIB superstring theories. The type II theories have the NS-NS fields which are the same as the fields in the bosonic string theory, as well as the R-R potentials which are bosonic fields. The \( D_p \)-branes in type II string theories carry the R-R charges [24]. These theories have also NS-R and R-NS fermionic fields in which we are not interested. The odd-form R-R potentials appear in type IIA and even-forms appear in type IIB. It is known that the compactification of type IIA theory on a circle transforms to the compactification of type IIB theory on another circle under the T-duality transformations. To study the effective action of the bosonic fields in these theories, it is convenient to collect the two theories to one theory which is called type II theory. It has both odd- and even-form R-R potentials. When compactifying this theory on a circle, the effective action then is expected to be invariant under the T-duality transformations, as in the bosonic theory.

When compactifying the theory on a circle with the killing coordinate \( y \), the T-duality transformations for the NS-NS fields are the Buscher rules [6, 7], i.e.,

\[
e^{2\phi'} = \frac{e^{2\phi}}{G_{yy}} \; ; \; \; G'_{yy} = \frac{1}{G_{yy}}
\]
\[ G'_{\mu y} = \frac{B_{\mu y}}{G_{yy}} \quad ; \quad C'_{\mu \nu} = G_{\mu \nu} - \frac{G_{\mu y} G_{\nu y} - B_{\mu y} B_{\nu y}}{G_{yy}} \]

\[ B'_{\mu y} = \frac{G_{\mu y}}{G_{yy}} \quad ; \quad B'_{\mu \nu} = B_{\mu \nu} - \frac{B_{\mu y} G_{\nu y} - G_{\mu y} B_{\nu y}}{G_{yy}} \]

(2)

where \( \mu, \nu \) denote any direction other than \( y \). In above transformations the metric is in the string frame. If one assumes fields transforms covariantly under the coordinate transformations, then the above transformations receive \( \alpha' \)-corrections \([8, 9, 10, 11]\). The T-duality transformations of the R-R fields at the leading order of \( \alpha' \) have been found in \([25]\), i.e.,

\[ C'^{(n)}_{\mu \cdots \nu y} = C^{(n-1)}_{\mu \cdots \nu} - (n - 1) \frac{G^{(n-1)}_{[\mu \cdots \nu y]} G_{[\alpha]y}}{G_{yy}} \]

\[ C'^{(n)}_{\mu \cdots \nu \alpha \beta y} = C^{(n+1)}_{\mu \cdots \nu \alpha \beta} + n C^{(n-1)}_{[\mu \cdots \nu \alpha} B_{|\beta]y} + n(n - 1) \frac{C^{(n-1)}_{[\mu \cdots \nu y]} B_{[\alpha]y} G_{[\beta]y}}{G_{yy}} \]

(3)

They may have also \( \alpha' \) corrections in which are not interested in this paper. The T-duality transformations \((2)\) and \((3)\) are such that they are consistent with the fact that \( D_p \)-brane in type II theory transform to \( D_{p-1} \)-brane or \( D_{p+1} \)-brane depending on whether the brane is along or orthogonal to the circle on which the T-duality is imposed. In fact the R-R fields couple to the \( D_p \)-brane as

\[ \int_{M^{p+1}} e^B C \]

(4)

where \( C = \sum_{n=0}^{8} C^{(n)} \). It is invariant under the R-R gauge transformation \( \delta C = d\Lambda + H\Lambda \) where \( \Lambda = \sum_{n=0}^{7} \Lambda^{(n)} \). The T-duality transformations \((2)\) and \((3)\) produce the following transformations:

\[ (e^B C)'_{\cdots y} = (e^B C)_{\cdots} ; \quad (e^B C)'_{\cdots} = (e^B C)_{\cdots y} \]

(5)

where dots represent some world-volume indices. In other words, the coupling \((4)\) is covariant under the T-duality transformations.

The effective action of type II string theory on the closed manifolds at the leading order of \( \alpha' \) is the well-known type II supergravity (see e.g., \([30]\)). The first higher derivative corrections to this action is at order \( \alpha'^3 \). Only the Riemann curvature couplings are known in the literature \([26, 27, 28]\). There are many other couplings at this order that are not known. There are also boundary terms at this order when manifolds have boundary which are not known either. We expect all these couplings might be calculated by imposing the T-duality constraint on the effective action. In fact, the known Riemann curvature couplings are reproduced by this method in \([29]\). In this paper we present the technical details for imposing the T-duality constraint to find the R-R couplings at the supergravity level and leave the calculations at order \( \alpha'^3 \) for the future works. The NS-NS couplings at the leading order of \( \alpha' \) are the same as the corresponding couplings in the bosonic theory, i.e., \([1]\).
The R-R couplings, as in (1), should be invariant under the R-R gauge transformations. For non-constant R-R field, however, the couplings should be in terms of the R-R field strength, i.e.,

\[ F^{(n)} = dC^{(n-1)} + H \wedge C^{(n-3)} \]  

which is invariant under the R-R gauge transformations. At the two-derivative level, the gauge invariance requires the following couplings in the string frame:

\[ S_0^{RR} = -\frac{2}{\kappa^2} \int d^{10}x \sqrt{-G} \sum_{n=1}^{9} \frac{a_n}{n!} F^{(n)} \wedge G^{a_1 \cdots a_n} F^{(n)}_{b_1 \cdots b_n} \]

where \( a_1, a_2, \ldots, a_9 \) are 9 parameters that the R-R gauge symmetry can not fix them. We are going to show that they can be fixed by the T-duality constraint. We did not include \( F^{(0)} \) and \( F^{(10)} \) terms in above couplings because they do not include dynamical fields. In writing the above couplings we assume the R-R fields are all independent. The on-shell physics, however, requires not all components of \( C^{(4)} \) to be independent. Moreover, the fields \( C^{(5)}, C^{(6)}, C^{(7)}, C^{(8)} \) are not independent. That means in the equations of motion one has to impose some extra constraints on the R-R field strengths to reproduce the standard equations of motion.

To impose the T-duality constraint on this action, we have to consider a background with \( U(1) \) isometry. It is convenient to use the following background for the metric, B-field and dilaton:

\[
G_{ab} = \begin{pmatrix}
\bar{g}_{\mu\nu} + e^{\varphi} g_{\mu\nu} & e^{\varphi} g_{\mu} \\
e^{\varphi} g_{\nu} & e^{\varphi}
\end{pmatrix},
B_{ab} = \begin{pmatrix}
\bar{b}_{\mu\nu} + \frac{1}{2} b_{\mu} g_{\nu} - \frac{1}{2} b_{\nu} g_{\mu} & b_{\mu} \\
-b_{\nu} & 0
\end{pmatrix}, \Phi = \bar{\phi} + \varphi/4
\]

where \( \bar{g}_{\mu\nu}, \bar{b}_{\mu\nu}, \bar{\phi} \) are the metric, the B-field and the dilaton in the base space, and \( g_{\mu}, b_{\mu} \) are two vectors in this space. Inverse of the above 10-dimensional metric is

\[
G^{ab} = \begin{pmatrix}
\bar{g}^{\mu\nu} & -g^{\mu} \\
-g^{\nu} & e^{-\varphi} + g_{\alpha} g^\alpha
\end{pmatrix}
\]

where \( \bar{g}^{\mu\nu} \) is the inverse of the base metric which raises the index of the vectors. In this parametrization the T-duality transformations of the R-R fields (3) become

\[
C_{\mu_1 \cdots \mu_y}^{(n)} = C_{\mu_1 \cdots \mu_a}^{(n-1)} - (n - 1) C_{[\mu_1 \cdots \mu_y] y}^{(n-1)}
\]

\[
C_{\mu_1 \cdots \mu_\beta}^{(n)} = C_{\mu_1 \cdots \mu_\beta}^{(n-1)} + n C_{[\mu_1 \cdots \mu_\beta] y}^{(n-1)} + n(n-1) C_{[\mu_1 \cdots \mu_y] y}^{(n-1)} b_\alpha g^\beta
\]

which are nonlinear.

To simplify the above nonlinear T-duality transformations to linear transformations, we use the following parametrizations for the R-R fields as well:

\[
C_{\mu_1 \cdots \mu_n}^{(n)} = \frac{c_{\mu_1 \cdots \mu_n}^{(n)}}{c_{\mu_1 \cdots \mu_{n-1}}^{(n-1)} g_{\mu n}} + n c_{[\mu_1 \cdots \mu_n] y}^{(n-1)}
\]

\[
C_{\mu_1 \cdots \mu_y}^{(n+1)} = \frac{c_{\mu_1 \cdots \mu_n}^{(n)}}{c_{\mu_1 \cdots \mu_{n-1}}^{(n-1)} g_{\mu y}}
\]
where $\tilde{c}^{(n)}$ are R-R potentials in the 9-dimensional base space. Our notation for making antisymmetry is such that e.g., $3\tilde{c}^{(2)}_{\mu_1 \mu_2}g_{\mu_3} = g_{\mu_1} \tilde{c}^{(2)}_{\mu_2 \mu_3} - g_{\mu_2} \tilde{c}^{(2)}_{\mu_1 \mu_3} + g_{\mu_3} \tilde{c}^{(2)}_{\mu_1 \mu_2}$. The Buscher rules (2) and (3) in the parametrizations (8) and (11) then become the following linear transformations:

$$\varphi' = -\varphi, \quad g'_\mu = b_\mu, \quad b'_\mu = g_\mu, \quad \tilde{g}'_{\alpha \beta} = \tilde{g}_{\alpha \beta}, \quad \tilde{b}'_{\alpha \beta} = \tilde{b}_{\alpha \beta}, \quad \tilde{c}'_{\mu_1 \cdots \mu_n} = \tilde{c}_{\mu_1 \cdots \mu_n}$$

They form a $\mathbb{Z}_2$-group, i.e., $(x')^n = x$ where $x$ is any field in the base space. These transformations receive higher derivative corrections in which we are not interested in this paper.

Using the reductions (3), (9) and (11), it is straightforward to reduce different terms in (7). The reduction of $\sqrt{-G}$ and the R-R couplings $|F^{(1)}|^2$, $|F^{(2)}|^2$ and $|F^{(3)}|^2$ are the following:

$$\sqrt{-G} = e^{\varphi/2} \sqrt{-\bar{g}}$$

$$|F^{(1)}|^2_G = e^{-\varphi/2} \left( e^{\varphi/2} |F^{(1)}|^2_G \right)$$

$$|F^{(2)}|^2_G = e^{\varphi/2} \left( e^{-\varphi/2} |\bar{F}^{(1)}|^2_G + \frac{e^{\varphi/2}}{2!} |\bar{F}^{(2)} + \tilde{c}^{(0)}| V |^2_G \right)$$

$$|F^{(3)}|^2_G = e^{\varphi/2} \left( e^{-\varphi/2} |\bar{F}^{(2)} + \tilde{c}^{(0)}| W |^2_G + \frac{e^{\varphi/2}}{3!} |\bar{F}^{(3)} + \tilde{H} + \tilde{c}^{(1)}| V |^2_G \right)$$

where $V$ is field strength of the $U(1)$ gauge field $g_\mu$, i.e., $V = dg$, $W$ is field strength of the $U(1)$ gauge field $b_\mu$, i.e., $W = db$, $\tilde{F}^{(n)}$ is field strength of 9-dimensional R-R potentials $\tilde{c}^{(n-1)}$, i.e., $\tilde{F}^{(n)} = d\tilde{c}^{(n-1)}$ and the three-form $\tilde{H}$ is defined as

$$\tilde{H}_{\mu \nu \alpha} = \tilde{H}_{\mu \nu \alpha} - \frac{1}{2} g_{[\mu W_{\nu \alpha}]} - \frac{1}{2} b_{[\mu V_{\nu \alpha}]}$$

where $\tilde{H}$ is field strength of the T-duality invariant two-form $\tilde{b}_{\mu \nu}$, i.e., $\tilde{H} = db$. It is evident that $\tilde{H}$ is invariant under the T-duality transformations (12).

The reduction of $|F^{(n)}|^2_G$ for $n > 3$ can be written as

$$|F^{(n)}|^2_G = e^{-\varphi/2} \left( \frac{e^{\varphi/2}}{(n-1)!} |\bar{F}^{(n-1)}| + \tilde{H} \wedge \tilde{c}^{(n-4)} + \tilde{c}^{(n-3)} \wedge W |^2_G \right)$$

$$+ \frac{e^{\varphi/2}}{n!} |\bar{F}^{(n)} + \tilde{H} \wedge \tilde{c}^{(n-3)} + \tilde{c}^{(n-2)} \wedge V |^2_G$$

Using the fact that the non-dynamical field strength $F^{(10)}$ in the 10-dimensional spacetime has been ignored, one should also ignore the non-dynamical fields in the 9-dimensional base space. Hence the reduction of $|F^{(9)}|^2_G$ becomes

$$|F^{(9)}|^2_G = e^{-\varphi/2} \left( \frac{e^{\varphi/2}}{8!} |\bar{F}^{(8)} + \tilde{H} \wedge \tilde{c}^{(5)} + \tilde{c}^{(6)} \wedge W |^2_G \right)$$

The subscript $\bar{g}$ in $\cdots |^2_{\bar{g}}$ indicates that the indices are contracted with the base metric $\bar{g}^{\mu \nu}$.

We now impose the T-duality constraint to fix parameters $a_1, \cdots, a_9$ in (17). According to this proposal, the effective action should satisfy the following relation:

$$S_{\text{eff}}(\psi) - S_{\text{eff}}(\psi') = \text{TD}$$

(17)
where $S_{\text{eff}}$ is the reduction of the 10-dimensional action on the circle, $\psi$ represents all massless fields in the base space and $\psi'$ represents their transformations under the T-duality transformations (12). On the right-hand side, TD represents some total derivative terms in the base space which become zero if the base space has no boundary. They should be reproduced by the boundary action if the base space has boundary [4].

Using the reductions (13), (15), (16) and the T-duality transformations (12), one can easily observes that the effective action (7) satisfies the following relation for $a_1 = a_2 = \cdots = a_9$:

$$S_{\text{eff}}(\psi) - S_{\text{eff}}(\psi') = 0$$

(18)

Hence the T-duality constraint fixes all 9 parameters in (7) in terms of one normalization parameter. Since there is no total derivative terms on the right-hand side of the above relation, the boundary action has no R-R couplings, as expected. The T-duality constraint on the NS-NS couplings, however, satisfies the relation (17) with some total derivative terms on the right-hand side which is cancelled by the the Gibbons-Hawking-York boundary term. Therefore, the T-duality constraint fixes the low energy effective action of type II string theory on the spacetime manifold with boundary as

$$S_0 + \partial S_0 = -\frac{2}{\kappa^2} \int d^{10}x \sqrt{-G} \left[ e^{-2\Phi} \left( R + 4 \nabla_a \Phi \nabla^a \Phi - \frac{1}{12} H^2 \right) + a_1 \sum_{n=1}^9 |F(n)|^2_G \right]
$$

$$\pm \frac{4}{\kappa^2} \int d^9y \sqrt{\pm g} e^{-2\Phi} K$$

(19)

Up the overall factor $a_1$, the R-R couplings are the democratic R-R couplings that have been found in [31]. If the spacetime has no boundary, then the term in the second line does not appear in the effective action.

The R-R potentials in above action are not all independent on-shell. In the equations of motion which represent the on-shell physics, they should satisfy the following constraints [31]:

$$F(9) = *F(1), \quad F(8) = -*F(2), \quad F(7) = -*F(3), \quad F(6) = *F(4), \quad F(5) = *F(5)$$

(20)

It has been shown in [31] that imposing the above relations to the equations of motion resulting from (19), one finds the same equations of motion resulting from the standard supergravity action. Hence the bulk action for type IIB theory can be rewritten in S-duality invariant form (see e.g., [30]). In performing this calculation, one first changes the string frame metric to the Einstein frame metric, i.e., $G_{ab} = e^{\Phi/2} G^{E}_{ab}$. The total derivative term resulting from this change of the frames is ignored. The Einstein frame couplings then can be written in S-duality invariant form. However, for the spacetime manifolds with boundary, the total derivative term must be cancelled with the corresponding terms in the boundary. In fact the total derivative term is produced by transforming the scalar curvature to the Einstein frame, i.e.,

$$R \rightarrow e^{-\Phi/2} \left( R - \frac{9}{2} \nabla_a \nabla^a \Phi - \frac{9}{2} \nabla_a \Phi \nabla^a \Phi \right)$$

(21)
The second term above produces the following total derivative term in the Einstein frame:

$$\frac{9}{\kappa^2} \int d^{10}x \sqrt{-G^E} \nabla_a \nabla^a \Phi = \pm \frac{9}{\kappa^2} \int d^9 y \sqrt{\pm g^E} \nabla^a \Phi n^E_a \tag{22}$$

where we have used the Stokes’s theorem as well. In the right hand side, $n^E$ is unite vector orthogonal to the surface in the Einstein frame.

We now show this term can be cancelled by the boundary action. To this end, we first transform the boundary term to the Einstein frame. Different terms in this action transform as the following:

$$\sqrt{\pm g} \rightarrow e^{\mp \Phi/4} \sqrt{\pm g^E}$$

$$K = \nabla_a n^a \rightarrow e^{-\Phi/4} (K^E + \frac{9}{4} n^E_a \nabla^a \Phi) \tag{23}$$

where $K^E = \nabla_a (n^E)^a$ is the trace of the extrinsic curvature of the boundary in the Einstein frame, and we have used the fact that the unite vector $n^a$ in the string frame and in the Einstein frame should be related as $n^a = e^{-\Phi/4} (n^E)^a$, because their lengths are one in both frames, i.e., $G_{ab} n^a n^b = 1 = G^E_{ab} (n^E)^a (n^E)^b$. The last term in the second like produces the following term:

$$\mp \frac{9}{\kappa^2} \int d^9 y \sqrt{\pm g^E} \nabla^a \Phi n^E_a \tag{24}$$

which cancels exactly the total derivative term in (22). The other term in the boundary action in the Einstein frame is

$$\mp \frac{4}{\kappa^2} \int d^9 y \sqrt{\pm g^E} K^E \tag{25}$$

which is invariant under the S-duality in which the Einstein frame metric is invariant.

It is important to note that one may study the S-duality in the string frame. Then the bulk action would be invariant under the S-duality up to a total derivative term, as in the T-duality \[4\]. The total derivative term resulting from the S-duality transformations is cancelled with the S-duality transformations of the boundary term, as in the T-duality \[4\].

Hence the low energy effective action of type II string theories on the spacetime manifolds with boundary which we have found by the T-duality constraint is also invariant under the S-duality. The boundary term found by the T-duality is necessary to have the S-duality. We expect the $\alpha'^3$ corrections to this action can also be found by the T-duality constraint and satisfies the S-duality as well. It would be interesting to perform these calculations.

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