Open string boundary conditions dependence on the background fields*

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Abstract

For some particular relations between background fields the local gauge symmetries appear in the theory. They turn some Neumann boundary conditions into the Dirichlet ones, and decrease the number of the Dp-brane dimensions. In these cases, in order to keep the conformal symmetry of the open string theory on the quantum level, in addition to the space-time field equations we must impose additional conditions on the background fields.

1. Introduction

The presence of the antisymmetric background field $B_{\mu\nu}$ makes Dp-brane a noncommutative manifold [1]. Adding linear dilaton field, $\Phi(x) = \Phi_0 + a_\mu x^\mu$, [2] the Dp-brane coordinate $x = a_\mu x^\mu$ becomes commutative and the conformal part of the world-sheet metric F is a new noncommutative variable [3].

The classical string action in the presence of the dilaton field $\Phi(x)$ explicitly depends on the field F, and break conformal invariance. On the quantum level the field F decouples on the space-time field equations in closed string theory. But there is a possibility that it contributes to the open string theory through the boundary conditions [4]. The main purpose of this article is to find conditions under which open string theory is F-independent.

We consider some particular cases when a_i is a light-like vector, with respect to the closed string metric G_{ij}, $a^2 = G^{ij}a_ia_j = 0$, as well as to the effective one, G^{eff}_{ij}, $\tilde{a}^2 = (G^{-1}_{eff})^{ij}a_ia_j = 0$. As a consequence of these conditions, the first class constraints and local gauge symmetries appear. Then some of the initial Dp-brane coordinates change the corresponding boundary conditions from Neumann to Dirichlet and decrease the number of the Dp-brane dimensions.

* Work supported in part by the Serbian Ministry of Science and Environmental Protection, under contract No. 141036.
PACS: 11.10.Nx, 04.20.Fy, 11.25.-w .
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In all cases, solving the second class constraints, we find the effective theory. It is again the string theory defined in terms of the effective string variables, symmetric under transformation $\sigma \to -\sigma$, and propagating in new open string background.

2. Action and its quantum consistency condition

The action [5]

$$S = \kappa \int_\Sigma d^2\xi \sqrt{-g} \left\{ \left[\frac{1}{2} g^{\alpha\beta} G_{\mu\nu}(x) + \frac{\varepsilon^{\alpha\beta}}{\sqrt{-g}} B_{\mu\nu}(x) \right] \partial_\alpha x^\mu \partial_\beta x^\nu + \Phi(x) R^{(2)} \right\}, \tag{2.1}$$

describes the evolution of the open string in the presence of the space-time metric $G_{\mu\nu}(x)$, Kalb-Ramond antisymmetric field $B_{\mu\nu}(x)$ and the dilaton scalar field $\Phi(x)$. Further, $R^{(2)}$ denotes the scalar curvature related to the intrinsic world-sheet metric $g_{\alpha\beta}$.

Conditions for quantum conformal invariance have the role of space-time field equations [6]

$$\beta^G_{\mu\nu} \equiv R_{\mu\nu} - \frac{1}{4} B_{\mu\rho\sigma} B_\nu{}^{\rho\sigma} + 2 D_\mu a_\nu = 0\,, \tag{2.2}$$

$$\beta^B_{\mu\nu} \equiv D_\rho B^\rho{}_{\mu\nu} - 2 a_\rho B^\rho{}_{\mu\nu} = 0\,, \tag{2.3}$$

$$\beta^\Phi \equiv 2\pi\kappa \frac{D-26}{6} - \frac{1}{24} B_{\mu\rho\sigma} B^{\mu\rho\sigma} - D_\mu a^\mu + 4a^2 = 0\,, \tag{2.4}$$

where $B_{\mu\rho\sigma} = \partial_\mu B_{\nu\rho} + \partial_\nu B_{\rho\mu} + \partial_\rho B_{\mu\nu}$ and $a_\mu = \partial_\mu \Phi$. When $G_{\mu\nu}(x)$ and $B_{\mu\nu}(x)$ are constant fields and dilaton field is linear $\Phi(x) = \Phi_0 + a_\mu x^\mu$, there is an exact solution for $a^2 = \kappa\pi \frac{26-D}{12}$, [5].

We decompose the space-time coordinates $x^\mu(\xi)\,(\mu = 0,1,2,...,D-1)$ in Dp-brane coordinates $x^i(\xi)\,(i = 0,1,...,p)$ and the orthogonal ones $x^a(\xi)$ $(a = p+1, p+2, ..., D-1)$, in such a way that $G_{\mu\nu} = 0$, $(\mu = i\,, \nu = a)$. Other two background fields are nontrivial only on Dp-brane: $B_{\mu\nu} \to B_{ij}$, $a_\mu \to a_i$. The part of the action describing the string oscillation in x^a directions decouples from the rest. Choosing the conformal gauge, $g_{\alpha\beta} = e^{2F} \eta_{\alpha\beta}$, we obtain

$$S(x^i, F) = \kappa \int_\Sigma d^2\xi \left[\left(\frac{1}{2} \eta^{\alpha\beta} G_{ij} + \epsilon^{\alpha\beta} B_{ij} \right) \partial_\alpha x^i \partial_\beta x^j + 2 \eta^{\alpha\beta} a_i \partial_\alpha x^i \partial_\beta F \right]. \tag{2.5}$$

3. Noncommutativity in the case $a^2 \neq 0$ and $\tilde{a}^2 \neq 0$

In the presence of linear dilaton field, for $a^2 \neq 0$ and $\tilde{a}^2 \neq 0$, the field F becomes a new noncommutative variable, while $x_c = a_i x^i$ becomes commutative (see Refs.[3, 4]).

3.1. Canonical Hamiltonian in terms of currents

In terms of currents

$$J^i_\pm = (P^0_T)^{ij} j_{\pm j} + \frac{a^i}{2a^2} i^\Phi_\pm = j^i_\pm - \frac{a^i}{a^2} j\,, \qquad j = a^i j_{\pm i} - \frac{1}{2} i^\Phi_\pm\,, \tag{3.1}$$

$$j_{\pm i} = \pi_i + 2\kappa \Pi_{\pm ij} x'^j\,, \qquad i^\Phi_\pm = \pi \pm 2\kappa a_i x'^i\,, \qquad i^F_\pm = \frac{j}{a^2} \pm 2\kappa F'\,, \tag{3.2}$$

canonical Hamiltonian gets the form

$$H_c = \int_0^\pi d\sigma (T_- - T_+)\,, \quad T_\pm = \mp \frac{1}{4\kappa}\left(G^{ij} J_{\pm i} J_{\pm j} + \frac{j}{a^2} i^\Phi_\pm\right) + \frac{1}{2}\left(i^{\Phi\prime}_\pm - F' i^\Phi_\pm\right)\,, \tag{3.3}$$

where $(P^0_T)_i{}^j = \delta_i{}^j - \frac{a_i a^j}{a^2}$ and $\Pi_{\pm ij} = B_{ij} \pm \frac{1}{2} G_{ij}$. The canonically conjugated momenta to x^i and F are denoted by π_i and π, respectively.

3.2. Boundary conditions as canonical constraints

The form of the Neumann boundary conditions corresponding to the fields x^i and F are $\gamma^{(0)}_i\Big|_0^\pi = 0$ and $\gamma^{(0)}\Big|_0^\pi = 0$, where

$$\gamma^{(0)}_i = \Pi_{+ij} J^j_- + \Pi_{-ij} J^j_+ + \frac{a_i}{2}\left(i^F_- - i^F_+\right)\,, \quad \gamma^{(0)} = \frac{1}{2}\left(i^\Phi_- - i^\Phi_+\right)\,. \tag{3.4}$$

We consider them as canonical constraints. Using Dirac procedure we find two infinite sets of conditions. With the help of Taylor expansion, we rewrite all the constraints at $\sigma = 0$ in a more compact σ-dependent form

$$\Gamma_i(\sigma) = \Pi_{+ij} J^j_-(\sigma) + \Pi_{-ij} J^j_+(-\sigma) + \frac{a_i}{2}\left[i^F_-(\sigma) - i^F_+(-\sigma)\right]\,, \tag{3.5}$$

$$\Gamma(\sigma) = \frac{1}{2}\left[i^\Phi_-(\sigma) - i^\Phi_+(-\sigma)\right]\,. \tag{3.6}$$

In the same fashion we obtain similar expressions from the constraints at $\sigma = \pi$, which is solved by periodicity condition, with period 2π [3]. From $\{H_c\,, \Gamma_i\} = \Gamma'_i$ and $\{H_c\,, \Gamma\} = \Gamma'$, we conclude that there are no more constraints in the theory.

The algebra of the complete set of the constraint is

$$\{\chi_A(\sigma), \chi_B(\overline{\sigma})\} = -\kappa M_{AB} \delta'\,, \quad \chi_A = (\Gamma_i, \Gamma)\,, \quad \left[\delta' = \partial_\sigma \delta(\sigma - \overline{\sigma})\right]\,, \tag{3.7}$$

where

$$M_{AB} = \begin{pmatrix} \tilde{G}_{ij} & 2a_i \\ 2a_j & 0 \end{pmatrix}\,. \tag{3.8}$$

To the space-time component

$$\tilde{G}_{ij} = G^{eff}_{ij} + \frac{4}{a^2}(Ba)_i (aB)_j\,, \quad \left(G^{eff}_{ij} = G_{ij} - 4B_{ik}G^{kl}B_{lj}\right) \tag{3.9}$$

we will refer as effective or open string metric. The determinant

$$\det M_{AB} = -4\tilde{a}^2 \det \tilde{G}_{ij} = -4\frac{(\tilde{a}^2)^2}{a^2} \det G^{eff}_{ij} \,, \tag{3.10}$$

shows that for $\tilde{a}^2 \neq 0$ and $a^2 \neq 0$ all constraints are of the second class.

3.3. Solution of the second class constraints and noncommutativity

The variables x^i and F and the corresponding momenta π_i and π can be expressed in terms of their σ-symmetric and antisymmetric parts as $x^i = q^i + \bar{q}^i$, $F = f + \bar{f}$, $\pi_i = p_i + \bar{p}_i$ and $\pi = p + \bar{p}$. From the constraint equations $\Gamma_i(\sigma) = 0$ and $\Gamma(\sigma) = 0$, we can solve σ-symmetric variables (effective ones) in terms of antisymmetric (bar) variables and express the string variables in the form

$$\pi_i = p_i \,, \qquad x^i(\sigma) = q^i(\sigma) - 2\int_0^\sigma d\sigma_1 \Big[\Theta^{ij} p_j(\sigma_1) + \Theta^i p(\sigma_1)\Big] \,, \tag{3.11}$$

$$\pi = p \,, \quad F(\sigma) = f(\sigma) + 2\Theta^i \int_0^\sigma d\sigma_1 p_i(\sigma_1) \,, \tag{3.12}$$

where the parameters are defined as

$$\Theta^{ij} = -\frac{1}{\kappa}(\tilde{P}_T B P^0_T)^{ij} \,, \qquad \Theta^i = \frac{(aB\tilde{G}^{-1})^i}{2\kappa a^2} = \frac{(\tilde{a}BG^{-1})^i}{2\kappa \tilde{a}^2} \,. \tag{3.13}$$

In analogy with the string currents (3.1) and (3.2), we introduce effective currents denoted by $\tilde{i}^\Phi_\pm$, $\tilde{j}_{\pm i}$, $\tilde{j}$ and $\tilde{J}^i_\pm$, substituting the string variables and metric tensor with effective ones and omitting the antisymmetric field. Correlating the string currents with the effective currents, we obtain effective energy-momentum tensor components

$$T_\pm = \tilde{T}_\pm \,, \quad \tilde{T}_\pm = \mp\frac{1}{4\kappa}\left[(\tilde{G}^{-1})^{ij}\tilde{J}_{\pm i}\tilde{J}_{\pm j} + \frac{\tilde{j}}{\tilde{a}^2}\tilde{i}^\Phi_\pm\right] + \frac{1}{2}\left(\tilde{i}'^\Phi_\pm - f'\tilde{i}^\Phi_\pm\right) \,. \tag{3.14}$$

The effective theory has the same form as the initial one, but in terms of the effective variables, q^i, f, p_i, and p, in new background $G_{ij} \to \tilde{G}_{ij}$, $B_{ij} \to 0$, and $\Phi \to \Phi_0 + a_i q^i$.

Introducing the center of mass variables

$$x^i_{cm} = \frac{1}{\pi}\int_0^\pi d\sigma x^i(\sigma) \,, \quad x^i(\sigma) = x^i_{cm} + X^i(\sigma) \,, \tag{3.15}$$

(similar for F) and using the solution of the constraints (3.11)-(3.12), it is easy to calculate Poisson brackets

$$\{X^i(\sigma), X^j(\overline{\sigma})\} = \Theta^{ij}\Delta(\sigma + \overline{\sigma}) \,, \tag{3.16}$$

$$\{X^i(\sigma), \mathcal{F}(\overline{\sigma})\} = \Theta^i\Delta(\sigma + \overline{\sigma}) \,, \qquad \{\mathcal{F}(\sigma), \mathcal{F}(\overline{\sigma})\} = 0 \,, \tag{3.17}$$

where $\Delta(x)$ is different from zero only at string endpoints, $\Delta(0)=-1$ and $\Delta(2\pi)=1$.

From the relations $a_i\Theta^{ij}=0$ and $a_i\Theta^i=0$, we get that coordinate $x_c\equiv a_i x^i$ becomes a commutative one.

4. Light-like dilaton gradient and noncommutativity

It turns out that for some particular choices of the background fields the first class constraints appear. The first class constraints generate gauge symmetries in the theory, which existence enables us to fix unphysical degrees of freedom.

4.1. Case $a^2=0$

For $a^2=0$, the current j (3.1) is a constraint, because it does not depend on velocities. The canonical Hamiltonian has the form

$$\mathcal{H}_c=T_- - T_+ = \frac{1}{4\kappa}G^{ij}(j^*_{-i}j^*_{-j}+j^*_{+i}j^*_{+j})\,,\quad j^*_{\pm i}=j_{\pm i}\pm 2\kappa a_i F'\,. \tag{4.1}$$

From

$$\{H_T\,,j\}=0\,,\quad H_T=\int d\sigma(\mathcal{H}_c+\lambda j)\,, \tag{4.2}$$

where λ is Lagrange multiplier, we obtain that j is a first class constraint. It generates local gauge symmetry transformations

$$\delta X=\{X\,,G\}\;,\quad G=\int d\sigma\eta(\sigma)j(\sigma)\,. \tag{4.3}$$

Applying these expressions to the variables x^i, F, π_i, and π, we get

$$\delta_\eta x^i=a^i\eta\,,\quad \delta_\eta F=-\frac{1}{2}\eta\,,\quad \delta_\eta\pi_i=2\kappa a^j B_{ji}\eta'\,,\quad \delta_\eta\pi=0\,. \tag{4.4}$$

After gauge fixing $F=0$, we treat F and j as second class constraints. The boundary conditions Γ_i and Γ (3.5)-(3.6) change the form

$$\check\Gamma_i(\sigma)=\Pi_{+i}{}^j j_{-j}(\sigma)+\Pi_{-i}{}^j j_{+j}(-\sigma)\,,\quad \check\Gamma(\sigma)=a^i\left[j_{-i}(\sigma)-j_{+i}(-\sigma)\right]\,, \tag{4.5}$$

and satisfy the algebra

$$\{\check\chi_A(\sigma),\check\chi_B(\overline{\sigma})\}=-\kappa\check M_{AB}\delta'\,,\quad \check\chi_A=(\check\Gamma_i\,,\check\Gamma)\,, \tag{4.6}$$

where

$$\check M_{AB}=\begin{pmatrix} G^{eff}_{ij} & 2a_i \\ 2a_j & 0\end{pmatrix}\,. \tag{4.7}$$

With the help of the relation

$$\det\check M_{AB}=-4\tilde a^2\det G^{eff}_{ij}\,, \tag{4.8}$$

we conclude that for $\tilde a^2\neq 0$ all constraints $\check\chi_A$ are of the second class.

4.2. Case $\tilde{a}^2 = 0$

From the expression (3.10) we conclude that, for $\tilde{a}^2 = 0$ and $a^2 \neq 0$, $\det M_{AB}$ has two zeros at $\tilde{a}^2 = 0$. In that case according to the algebra (3.7) the constraints

$$\Gamma_1 = 2(\tilde{a}BG^{-1})^i\Gamma_i\,, \quad \Gamma_2 = \tilde{a}^i\Gamma_i - \frac{1}{2}\Gamma\,, \tag{4.9}$$

are of the first class, while the rest ones

$$(\Gamma_T)_i = (\hat{P}_T)_i{}^j\Gamma_j\,, \quad \Gamma\,, \quad \left((\hat{P}_T)_i{}^j = \delta_i{}^j - \frac{a_i a^j}{a^2} + \frac{4}{a^2}(Ba)_i(\tilde{a}B)^j\right) \tag{4.10}$$

are of the second class.

The expression for gauge transformation is the same as in Eq.(4.3) with the generator

$$G = \int d\sigma[\eta_1(\sigma)\Gamma_1(\sigma) + \eta_2(\sigma)\Gamma_2(\sigma)]\,. \tag{4.11}$$

Applying this to the effective coordinates we obtain

$$\delta q^i = \tilde{a}^i(\eta_1)_s + 2(\tilde{a}BG^{-1})^i(\eta_2)_s\,, \qquad \delta f = -\frac{1}{2}\,(\eta_1)_s\,, \tag{4.12}$$

$$\delta \overline{q}^i = \tilde{a}^i(\eta_2)_a + 2(\tilde{a}BG^{-1})^i(\eta_1)_a\,, \qquad \delta \overline{f} = -\frac{1}{2}\,(\eta_2)_a\,, \tag{4.13}$$

while the momenta are invariant. Here the indices "s" and "a" denote σ symmetric and antisymmetric parts of the parameters η_1 and η_2. We choose gauge fixing

$$(aB)_i q^i = 0\,, \quad (aB)_i\overline{q}^i = 0\,, \quad f = 0\,, \quad \overline{f} = 0\,. \tag{4.14}$$

4.3. Solution of the constraints and noncommutativity

For $a^2 = 0$ we solve gauge condition $F = 0$, first class constraint $j = 0$, and, the second class constraints $\check{\Gamma}_i = 0$ and $\check{\Gamma} = 0$ (4.5). In the case $\tilde{a}^2 = 0$, solution of the the first class constraints, $\Gamma_1 = 0$ and $\Gamma_2 = 0$, and the second class ones, $(\Gamma_T)_i = 0$, is equivalent to the solution of all constraints, $\Gamma_i = 0$. Then the complete set of equation consists of $\Gamma_i = 0$, $\Gamma = 0$, and the gauge conditions (4.14).

Solving the corresponding set of equations, we obtain that the solution in both considered cases has the same form

$$x^i_{Dp}(\sigma) = Q^i(\sigma) - 2\Theta^{ij}\int_0^\sigma d\sigma_1 P_j(\sigma_1)\,, \quad \pi_i^{Dp}(\sigma) = P_i(\sigma)\,, \tag{4.15}$$

$$F = 0\,, \quad \pi = 2\tilde{a}^i P_i\,, \quad x_1 \equiv (aB)_i x^i|_0^\pi = 0\,, \quad \pi_1 \equiv (\tilde{a}B)^i\pi_i = 0\,, \tag{4.16}$$

where we introduced the notation

$$Q^i = (P_{Dp})^i{}_j q^j \; , \; P_i = (P_{Dp})_i{}^j p_j \; , \; \Theta^{ij} = -\frac{1}{\kappa} \, (G_{eff}^{-1} P_{Dp} B G^{-1} P_{Dp})^{ij} \; . \quad (4.17)$$

Note that Θ^{ij} is antisymmetric and for $a^2 = 0$ $(P_{Dp})^i{}_j \equiv (P_T^1)_i{}^j = \delta_i{}^j - \frac{4}{\tilde{a}^2}(Ba)_i(\tilde{a}B)^j$, while for $\tilde{a}^2 = 0$ $(P_{Dp})^i{}_j \equiv (\hat{P}_T^1)_i{}^j = \delta_i{}^j + \frac{4}{a^2}(Ba)_i(\tilde{a}B)^j$.
Let us introduce an effective current

$$\tilde{j}_{\pm i} = P_i \pm \kappa (P_{Dp} G_{eff})_{ij} Q'^j \, , \quad (4.18)$$

and, using the solution (4.15), correlate with the currents defined in Eqs.(3.1) and (3.2). Substituting the relation which connects the initial string currents with effective ones in the corresponding expression for canonical Hamiltonian, we obtain effective Hamiltonian

$$\tilde{\mathcal{H}}_c = \tilde{T}_- - \tilde{T}_+ \, , \quad \tilde{T}_\pm = \mp \frac{1}{4\kappa} (G_{eff}^{-1} P_{Dp})^{ij} \tilde{j}_{\pm i} \tilde{j}_{\pm j} \, . \quad (4.19)$$

Effective theory is expressed in terms of the effective string variables, Q^i and P_i, and lives in the background $G_{ij} \to (P_{Dp} G^{eff})_{ij}$, $B_{ij} \to 0$ and $\Phi \to 0$.
Introducing the center of mass variables as in Eq.(3.15), we get the noncommutativity relation

$$\{X^i_{D_p}(\sigma), X^j_{D_p}(\overline{\sigma})\} = \Theta^{ij} \Delta(\sigma + \overline{\sigma}) \, , \quad (4.20)$$

where the function $\Delta(x)$ is defined after Eq.(3.17). The coordinate x_1 and the field F satisfy Dirichlet boundary conditions and decrease the number of Dp-brane dimensions. From $a_i Q^i \neq 0$ and $a_i \Theta^{ij} = 0$ we have that $x_c = a_i x^i_{D_p} = a_i Q^i$ is a commutative coordinate, because it is a momentum independent variable.

5. Concluding remarks

In this paper we have considered the conditions for quantum conformal invariance of the open string theory. In all cases we treat the conformal part of the world sheet metric F as a dynamical variable.

In the first case ($\tilde{a}^2 \neq 0 \, , a^2 \neq 0$) the coordinate in the direction of dilaton gradient is commutative, and the field F is a new noncommutative variable.

In the second case ($a^2 = 0$), there is one first class constraint, while in the third case ($\tilde{a}^2 = 0$) two second class constraints originating from boundary conditions turn to the first class ones. In both cases, the first class constraints generate local symmetries and after gauge fixing, the gauge conditions and the first class constraints are treated as second class constraints. Solving them together with the real second class constraints we obtain effective coordinates and momenta describing the string dynamics. In both cases, the Dp-brane is p-dimensional with one commutative coordinate, in the direction of the dilaton gradient a_i, and $p-1$ noncommutative coordinates.

The conformal invariance of the open string theory is preserved on the quantum level if the boundary conditions are F-independent. Particulary, it occurs for $a^2=0$ and $\tilde{a}^2=0$ where the field F vanishes. So, in addition to the space-time field equations we must impose additional conditions on the background fields in order to keep the conformal symmetry of the open string theory on the quantum level.

The conformal anomaly can be canceled also by introducing Liouville action, which is a kinetic term for the field F. The main results obtained in this framework [7] are in full correspondence with the corresponding ones in the absence of the Liouville action [4].

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