On Noncommutativity in String Theory and D-Branes

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Abstract

String theory in a constant B-field exhibits noncommutative structure of space-time. By considering the B-field dynamical and studying its interaction with Ramond-Ramond (RR) background we observe the breaking of the B-field gauge symmetry in the effective action. This effect takes place due to non-perturbative coupling of the B-field to membrane topological charge. As a result, the B-field is renormalized in the RR backgrounds, making it impossible to obtain consistent non-commutative models with constant B-field. We argue that the gauge invariance should be restored by introducing appropriate external D-brane configuration.

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1. Introduction

Space-time noncommutativity in string theory appears as a result of open string dynamics in a constant B-field background [1], [2]. The case of a special interest is the noncommutativity in the presence of D-branes. Prior to the remarkable work of Polchinski [3] on D-branes it had long been a profound puzzle which objects are the actual sources of Ramond-Ramond charges. In superstring theory charges are usually carried by emission vertices of corresponding physical states, i.e. by BRST invariant and non-trivial vertex operators. For instance, the open superstring vertex operator of a photon:

\[ V_{ph}(k) = A_m(k)(\partial X^m + i(k\psi)\bar{\psi}^m)e^{ikX(z)} \]

is multiplied by the vector gauge potential \(A_m(k)\) of the photon and therefore can be regarded as the source of the U(1) electric charge. In the Ramond-Ramond sector, however, the situation is different. It is well-known that expressions for Ramond-Ramond vertex operators at canonical picture are given by:

\[ V_{RR}(k) = \frac{1}{p!}\gamma_{m_1...m_p}F_{m_1...m_p}(k)c\bar{c}e^{-\frac{1}{2}\phi - \frac{i}{2}\Sigma\bar{\Sigma}}e^{ikX(z, \bar{z})} \]

where \(\gamma_{m_1...m_p}\) is antisymmetrized product of 10d gamma-matrices, \(\Sigma, \bar{\Sigma}\) are spin operators for matter fields, \(\phi\) is bosonized superconformal ghost and the rank \(p\) of the RR field strength \(F\) may be odd or even, depending on the type of superstring theory we consider - that is, type IIA or type IIB.

The standard bozonization formulae [4] for reparametrization and superconformal ghosts \(b, c, \beta\) and \(\gamma\) are given by:

\[ c(w) = e^{\sigma}(w), b(w) = e^{-\sigma}(w), \gamma(w) = e^{\phi-\chi}(w), \beta(w) = e^{\chi-\phi}\partial\chi(w) \]

\[ <\sigma(z)\sigma(w)> = <\chi(z)\chi(w)> = \log(z-w) \]

\[ <\phi(z)\phi(w)> = -\log(z-w) \] (3)

The crucial point here is that the BRST invariance of the operators (2) requires that they couple to RR field strength rather than RR gauge potential. BRST invariance condition is then equivalent to the Maxwell’s equations for the RR field strength \(F(k)\). For this reason, the RR vertex operators (2) cannot be considered as the sources of the RR charges. This puzzle has been resolved in the crucial work by Polchinski where it has been shown that the RR charges are carried by the non-perturbative solitonic objects, the
D-branes. Namely, the RR gauge potentials appear in the WZ terms of DBI actions for Dp-branes, coupling to their worldvolume $p + 1$-forms. D-branes may also be realized as open strings with mixed Dirichlet-Neumann boundary conditions, at least up to massless modes. Remarkably, however, the RR vertex operators carrying the RR charges still can be constructed at non-canonical pictures. Namely, consider the physical RR vertex operator at picture $(-3/2, -3/2)$:

\[ V_{RR}^{(-3/2, -3/2)}(k) = \frac{1}{p!} \gamma_{\alpha \beta}^{m_1 \ldots m_p} R_{m_1 \ldots m_p}(k) c \bar{c} e^{-\frac{i}{2} \phi - \frac{i}{2} \bar{\phi}} \sum_{\alpha} \sum_{\beta} e^{ikX}(z, \bar{z}) \]  

(4)

Properties of this operator are significantly different from those of (2) (this question has also been considered in [5], [6]. It is easy to check that in this case BRST invariance condition imposes no on-shell constraints on the polarization $p$-form $R_{m_1 \ldots m_p}(k)$ and for this reason it cannot be interpreted as a RR field strength. We shall refer to the $p$-form space-time field $R_{m_1 \ldots m_p}(k)$ as Ramond-Ramond prepotential, for reasons which will later become clear. To point out the physical meaning of $R_{m_1 \ldots m_p}(k)$ let us act on it with left and right picture-changing operators $\Gamma$ and $\bar{\Gamma}$ and compare the result with (2). The normal ordered expression for $\Gamma$ is given by

\[ \Gamma(z) =: e^{\phi}(G_{matter} + G_{ghost}) : (z) = \frac{1}{2} e^{\phi} \psi_m \partial X^m(z) + \frac{1}{4} \chi e^{2\phi - \chi} (\partial \sigma + \partial \chi)(z) + e^\chi \partial \chi(z) \]  

and analogously for $\bar{\Gamma}$.

Using the relevant OPE’s, including

\[ \psi^m(z) \Sigma_{\alpha}(w) \sim (z - w)^{-\frac{1}{2}} \gamma^m_{\alpha \gamma} \Sigma_{\gamma}(w) + ... \]
\[ e^{\alpha \phi}(z)e^{\beta \phi}(w) \sim (z - w)^{-\alpha \beta} e^{(\alpha + \beta) \phi}(w) + ... \]
\[ \partial X^m(z)e^{ikX}(w) \sim (z - w)^{-1} i k^m e^{ikX}(w) + ... \]  

(6)

we have:

\[ : \bar{\Gamma} V_{RR}^{(-3/2, -3/2)}(k) := \frac{1}{2} \gamma_{\alpha \beta}^{m_1 \ldots m_{p-1} n} R_{[m_1 \ldots m_{p-1} n](k)} \]
\[ + \gamma_{\alpha \beta}^{m_1 \ldots m_{p-1}} R_{m_1 \ldots m_p}(k) k_{m_p} \times c \bar{c} e^{-\frac{i}{2} \phi - \frac{i}{2} \bar{\phi}} \sum_{\alpha} \sum_{\beta} e^{ikX}(z, \bar{z}) = V_{RR}^{(-3/2, -1/2)}(k) \]  

(7)
where the square brackets imply antisymmetrization over the space-time indices. Next, acting on this expression with $\Gamma(z)$ we get:

$$
: \Gamma \bar{\Gamma} V^{(-3/2,-3/2)}(k) := \frac{1}{4} e^{-\phi - \frac{1}{2} \phi \sum} e^{ikX(z, \bar{z})} \times 2 \Gamma_{\alpha \beta}^{n m_1 ... m_{p-1}} k_m R_{m_1 ... m_{p-1} m_p}(k) \\
\equiv V^{(-1/2,-1/2)}(k)
$$

Comparing this with (1) we deduce the following relations between the RR field strength $F$, RR gauge potential $A$ and RR prepotential $R$:

$$
F_{n m_1 ... m_{p-1}}(k) = k_m R_{n m_1 ... m_{p-1} m_p}(k) \\
A_{m_1 ... m_{p-1}}(k) = k_m R_{m_1 ... m_{p-1} m_p}(k)
$$

In other words, the Ramond-Ramond $(p - 1)$-form gauge potential $A$ is given by the divergence of the $p$-form prepotential $R$, while the RR field strength $F$ is given by the Laplacian of $R$. The $p$-form prepotential $R$ may also be interpreted as a parameter for the Penrose class of solutions to Maxwell’s equations Therefore the RR potential enters the expression for picture $-3/2, -1/2$ (or equivalently $-1/2, -3/2$) RR vertex operator which can now be expressed as

$$
V^{(-3/2,-1/2)}(k) := \frac{1}{2} e^{-\phi - \frac{1}{2} \phi \sum} e^{ikX(z, \bar{z})} (\Gamma_{\alpha \beta}^{m_1 ... m_{p-1}} A_{m_1 ... m_{p-1}}(k) + \Gamma_{m_1 ... m_p}^{m_1 ... m_{p-1}} R_{m_1 ... m_{p-1} m_p}(k))
$$

In other words, the RR vertex operator taken at the mixed $(-3/2, -1/2)$-picture can be considered as a source of the RR-charge (shifted by the exterior derivative of the RR-prepotential, necessary to insure the BRST-invariance). In this sense the $(-3/2, -1/2)$-picture vertex operators are similar to D-branes and should have a non-perturbative nature, while the structure of their scattering amplitudes may be expected to reflect a non-perturbative physics. In this letter we shall observe and discuss one particularly interesting example of non-perturbative effect related to RR scattering amplitudes at non-canonical pictures - the breaking of the B-field gauge invariance in the low energy effective action due to the presence of the non-canonical RR states. Namely, we will show that terms of the form $(B \wedge F^{(p-2)}, F^{(p)})$ appear in the effective Lagrangian (particularly for $p = 4$) due to the interaction of the B-field with non-canonical RR backgrounds in superstring theory. Note that, contrary to the canonical case (related to scattering amplitudes of
the B-field with canonical RR vertices) these terms are not full derivatives (as this is the case for the well-known CS term $B \wedge F^{(4)} \wedge F^{(4)}$ coming from M-theory) and therefore the B-field gauge invariance is broken. Let us stress that this effect is non-perturbative (as the $V_{RR}^{(-3/2,-1/2)}$ vertices correspond to non-perturbative brane dynamics) and of course may lead to significant consequences for issues like non-commutativity occurring in certain B-field backgrounds. Some of these consequences will be discussed in this letter.

2. B-field in RR backgrounds and non-perturbative 2-form state

Before starting the calculation of scattering amplitudes revealing the non-perturbative B-field gauge invariance breaking, we shall comment on some peculiarities of the OPE’s of spin operators at non-canonical pictures to clarify the physics behind the gauge invariance breaking. Consider first the OPE of two canonical Ramond spin operators. We shall be interested in simple poles of these OPE’s (as for 3-point correlators of primary fields only these OPE terms are important). Using the OPE expressions (6) we have:

$$e^{-\frac{1}{2}\phi} \Sigma \alpha(z) e^{-\frac{1}{2}\phi} \Sigma \beta(w) \sim \frac{1}{z-w} e^{-\phi} \psi \gamma^m \Gamma_{\alpha\beta} + \ldots$$

At the same time,

$$e^{-\frac{1}{2}\phi} \Sigma \alpha(z) e^{-\frac{1}{2}\phi} \Sigma \beta(w) \sim \frac{1}{(z-w)^2} e^{-2\phi} \delta_{\alpha\beta} + \frac{1}{2} e^{-2\phi} \psi \gamma^m \psi \gamma^{mn}$$

The OPE (12) differs from (11) substantially. While the r.h.s. of (11) contains only the usual vector field at picture $-1$, the r.h.s. of the OPE (12) involving the non-canonical spin field $e^{-\frac{1}{2}\phi} \Sigma \alpha$ contains the two-form term given by $\Phi_{mn} = e^{-2\phi} \psi \gamma^{mn}$. The origin of this two-form has been discussed in [7]. It has been shown that this corresponds to the membrane topological charge [8], appearing as a two-form central term in picture-changed space-time SUSY algebra. This intermediate state may also be interpreted as a two-form $\Phi$-parameter associated with a choice of regularization in the worldsheet path integral for a string theory with the B-field [9],[10]. Indeed, it is easy to see that for the space-time conjugate momentum operator at picture -1 $P_m = \oint \frac{dz}{2\pi i} e^{-\phi} \psi \gamma^m (z)$ one has $[P_m, P_n] = \oint \frac{dz}{2\pi i} \Phi_{mn} \sim [\partial_m, \partial_n]$. At nonzero momenta this two-form gives rise to physical vertex operator given by $\sim e^{-2\phi} \psi \gamma^m \psi \gamma^{mn} e^{ikx}$ which does not correspond to any perturbative open string excitation (such as a photon) but which describes the non-perturbative membrane dynamics. There is no version of this vertex at picture zero; this operator is BRST-nontrivial if the momentum $k$ is directed along any of 8 space-time directions transverse...
to its indices \(m\) and \(n\). Now, because of the form of the OPE (12) these two-form vertices appear as an intermediate state (both in left and right sectors) in all amplitudes involving the RR vertices at non-canonical pictures. The crucial point is that it is the interaction of this intermediate two-form membrane-like vertex with the axionic state that makes the amplitude picture-dependent and plays the crucial role in breaking the gauge invariance of the B-field in the low-energy effective action. This effect is therefore non-perturbative. In the following section we shall demonstrate it by direct computation of scattering amplitude.

3. Interaction of the B-field with non-canonical RR backgrounds

In this section we compute the interaction of the B-field with 2-form RR field strength and the 4-form RR-prepotential taking place in the type IIA theory, showing that it gives rise to the anomalous term in the low-energy effective action.

The relevant correlator to compute is given by

\[
A_{FFB} = \langle V_{RR}^{(-3/2,-3/2)}(p)V_{RR}^{(-1/2,-1/2)}(k)V_B^{(0,0)}(q) \rangle
\]

where \(V_B^{(0,0)}(p)\) is the axionic vertex at picture \((0, 0)\), given by

\[
V_B^{(0,0)}(q) = \bar{c}c(\partial X^m + i(q\psi)\psi_m)(\tilde{\partial} X^m + i(q\bar{\psi})\bar{\psi}_n)e^{iqX}(z, \bar{z})B_{mn}(q)
\]

Let us start with computing the correlator of RR-vertices with the purely fermionic part of \(V_B\), i.e. the one biquadratic in \(\psi\) and \(\bar{\psi}\). We have:

\[
A_{FFB}^{(1)}(p, k, q) = -\frac{1}{2!4!} \langle c\bar{c}e^{-\frac{1}{2}\phi - \frac{3}{2}\hat{\phi}\Sigma_\alpha_1\Sigma_\beta_1}e^{ipX(z_1, \bar{z}_1)}c\bar{c}e^{-\frac{1}{2}\phi - \frac{3}{2}\hat{\phi}\Sigma_\alpha_2\Sigma_\beta_2}e^{ikX}(z_2, \bar{z}_2)
\]

\[
\times q^s q^t c\bar{c} : \psi_s\psi_m\bar{\psi}_t\bar{\psi}_n : (z_3, \bar{z}_3) > \gamma_{m_1...m_4}^{n_1...n_2} \gamma_{\alpha_1\beta_1}^{\alpha_2\beta_2} R_{m_1...m_4}(p) F_{n_1n_2}(k) B_{mn}(q)
\]

(15)

(the minus sign here is due to the total \(i^2\) factor in the fermionic part of \(V_B\)). Computing this correlator using the OPE expressions (6) we obtain

\[
A_{FFB}^{(1)}(p, k, q) = -\frac{1}{(2!4!)} \text{Tr}(\gamma_{m_1...m_4}^{n_1...n_2} \gamma_{\alpha_1\beta_1}^{\alpha_2\beta_2} R_{m_1...m_4}(p) F_{n_1n_2}(k) B_{mn}(q))
\]

\[
\times (p + k + q)
\]

(16)

Straightforward evaluation of the gamma-matrix trace gives:

\[
\text{Tr}(\gamma_{m_1...m_4}^{n_1...n_2}) = -(96g^{m_1n_1}g^{s_1m_2}g^{t_3m_4} + \text{perm}([n_1 \leftrightarrow n_2], [m \leftrightarrow s]))
\]

\[
-(96g^{m_1t_2}g^{s_1n_2}g^{t_3m_4} + \text{perm}([n_1 \leftrightarrow t], [m \leftrightarrow s]))
\]

\[
+(96g^{s_1m_1}g^{n_2m_2}g^{t_3m_4} + \text{perm}([n_1 \leftrightarrow n_2], [m \leftrightarrow s]))
\]

\[
+(96g^{s_1m_1}g^{n_2m_2}g^{t_3m_4} + \text{perm}([n_1 \leftrightarrow t], [m \leftrightarrow s]))
\]

(17)
where $g^{mn}$ is Minkowski tensor and permutations imply antisymmetrizations over the appropriate indices. Contracting the obtained expression for the trace with the space-time fields and using the on-shell conditions:

$$q^m B_{mn}(q) = 0$$
$$k^m F_{mn}(k) = 0$$
$$k_{[m} F_{np]}(k) = 0$$
$$q^2 = k^2 = p^2 = 0$$

we obtain the result for this part of the amplitude:

$$A^{(1)}_{FPB}(p, k, q) = \frac{192}{2!^4 4!} q_s R_{mnnn_1}(p) q_t F_{tn_1}(k) B_{mn}(q) = -A^{RR}_{mnn_1}(p) p_t F_{tn_1}(k) B_{mn}(q) = -F_{tmnn_1}(p) F_{tn_1}(k) B_{mn}(q)$$

(18)

The next contribution is from the correlator involving the $X$-part of the axionic vertex:

$$A^{(2)}_{FPB}(p, k, q) = \frac{1}{2!^4} < c \bar{c} e^{-\frac{1}{2} \phi - \frac{1}{2} \phi} \Sigma_{\alpha_1} \bar{\Sigma}_{\beta_1} e^{ipX(z_1, \bar{z}_1)} c \bar{c} e^{-\frac{1}{2} \phi - \frac{1}{2} \phi} \Sigma_{\alpha_2} \bar{\Sigma}_{\beta_2} e^{i k X(z_2, \bar{z}_2)} >$$

$$\times c \bar{c} \partial X^m \bar{\partial} X^n e^{iqX} > \times \bar{\gamma}_{\alpha_1 \beta_1} \gamma_{\alpha_2 \beta_2} R_{m_1 ... m_4 \gamma n_1 n_2} F_{n_1 n_2}(k) B_{mn}(q)$$

$$= k_n p_m R_{m_1 ... m_4 (p)} F_{n_1 n_2}(k) B_{mn}(q) Tr(\gamma_{m_1 ... m_4} \gamma_{n_1 n_2}) = 0$$

(19)

e. this contribution is zero as the gamma-matrix trace vanishes. Finally, the cross-term contribution is given by:

$$A^{(3)}_{FPB}(p, k, q) = \frac{i}{2!^4} < c \bar{c} e^{-\frac{1}{2} \phi - \frac{1}{2} \phi} \Sigma_{\alpha_1} \bar{\Sigma}_{\beta_1} e^{ipX(z_1, \bar{z}_1)} c \bar{c} e^{-\frac{1}{2} \phi - \frac{1}{2} \phi} \Sigma_{\alpha_2} \bar{\Sigma}_{\beta_2} e^{i k X(z_2, \bar{z}_2)} >$$

$$\times c \bar{c} \partial X^m \bar{\partial} X^n e^{iqX} > \times \bar{\gamma}_{\alpha_1 \beta_1} \gamma_{\alpha_2 \beta_2} R_{m_1 ... m_4 (p)} F_{n_1 n_2}(k) B_{mn}(q)$$

$$= -2 \times \frac{i}{(2!)^2 4!} R_{m_1 ... m_4 (p)} F_{n_1 n_2}(k) B_{mn}(q) q_s q_n Tr(\gamma_{m_1 ... m_4} \gamma_{n_1 n_2} \gamma_{m_1 ... m_4})$$

(20)

(21)

Evaluating the gamma-matrix trace as before and using the on-shell conditions for the space-time fields along with momentum conservation we get

$$A^{(3)}_{FPB}(p, k, q) = -\frac{8 \times 4!}{(2!)^2 4!} q_s R_{s m n_1 n_2}(p) F_{n_1 n_2}(k) p_n B_{mn}(q)$$

$$= 2 A^{RR}_{mnn_1 n_2}(p) F_{n_1 n_2}(k) p_n B_{mn}(q) = 2 F_{mnn_1 n_2}(p) F_{n_1 n_2}(k) B_{mn}(q)$$

$$= -2 F_{mnn_1 n_2}(p) F_{n_1 n_2}(k) B_{mn}(q)$$

(22)
Physically, the factor of 2 in this contribution is related to sum of the contributions from the left and the right sectors (interaction with left and right intermediate two-form states). Adding all the contributions together, we get

$$A_{FFB}(p, k, q) \equiv A_{FFB}^{(1)}(p, k, q) + A_{FFB}^{(2)}(p, k, q) + A_{FFB}^{(3)}(p, k, q)$$

$$= -3 F_{mn_1n_2}(p) F_{n_1n_2}(k) B_{mn}(q)$$

(23)

This concludes our calculation of the 3-point correlator. It corresponds to the term

$$S_{BFF} \sim \int d^{10} x F_{mn_1n_2} F_{n_1n_2} B_{mn}$$

(24)

in the low-energy effective action, apparently breaking the B-field gauge symmetry. As we have already remarked above, the mechanism of this symmetry breaking originates from the interaction of the B-field vertex with the membrane-like intermediate states (left and right), described by two-form vertex operators at the picture $-2$, i.e. this effect is non-perturbative, even though technically it involves only the perturbative string amplitudes. In the next section we shall discuss the relevance of this result to the space-time non-commutativity problem.

4. Gauge symmetry breaking and non-commutativity in RR backgrounds

The phenomenon of the gauge invariance breaking, caused by the non-perturbative interaction of the B-field with non-canonical RR-states, particularly raises questions of how consistent are the models of space-time noncommutativity based on open strings in a constant B-field background. In these models the presence of the B-field modifies the two-point propagator $< X^m(z) X^n(w) >$ on the worldsheet boundary. As a result the modified propagator acquires the antisymmetric part giving rise to non-commutativity of space-time coordinates after the regularization at coincident points z and w. The non-commutativity parameter, $\theta_{mn}$, is then given by the function of the $B_{mn}$ axionic field, originating from a rank 2 antisymmetric massless mode of a closed string. In these models the two-form B-field is treated as a static background, without any regard to a closed string. In the perturbative string-theoretic framework such a consideration is valid and non-contradictory since perturbatively the constant $B$-field background plays no role in closed string dynamics. This is because perturbatively the B-field is a gauge field, entering the low-energy effective action through its 3-form field strength $H = dB$. The only exception is the CS term $B \wedge F^{(4)} \wedge F^{(4)}$ where $F^{(4)}$ is the RR 4-form field strength. This topological term, originating from M-theory, is gauge-invariant and does not affect
the equations of motion, playing the role analogous to the \( \theta \)-term. For constant \( B_{mn} \)-field the B-field strength is zero and it can be gauged away by suitable transformation. The situation changes, however, if one takes into account the non-perturbative interaction of the B-field with the RR-sector (taking place through the intermediate NS 2-form state \( e^{-2\phi} \psi_m \psi_n e^{ikX} \)). Due to the non-vanishing s-matrix element of the B-field with two RR-states. As a result, the gauge symmetry is broken and the B-field background becomes dynamical and gets renormalized by the Ramond-Ramond space-time fields. Due to this renormalization, the space-time profile of the B-field is no longer arbitrary but is related to the profile of the RR-fields so as to satisfy the condition of the worldsheet conformal invariance. To see the relation consider the worldsheet RG flow equations involving the B-field and the RR-states, taking into account their non-perturbative interaction (13). Using the result of our calculation of the 3-point correlator (23) it is easy to write down equations for the beta-functions:

\[
\begin{align*}
\beta_B &\equiv \frac{dB_{mn}}{d(\log \Lambda)} \sim -k^2 B_{mn} - 3F_{mnpq}^{RR(4)} F_{pq}^{RR(2)} + \ldots \\
\beta_{F^{(2)}} &\equiv \frac{dF_{pq}^{RR(2)}}{d(\log \Lambda)} \sim -k^2 F_{pq}^{RR(2)} - 3F_{mnpq}^{RR(4)} B_{mn} + \ldots \\
\beta_{F^{(4)}} &\equiv \frac{dF_{mnpq}^{RR(4)}}{d(\log \Lambda)} \sim -k^2 F_{mnpq}^{RR(4)} - 3F_{[pq}^{RR(2)} B_{mn]} 
\end{align*}
\]

Upon the Fourier transform, the conformal invariance condition \( \beta_B = 0 \) particularly implies the equation of motion for the B-field:

\[
\nabla^2 B_{mn} = -3F_{mnpq}^{RR(4)} F_{pq}^{RR(2)}
\]

reflecting the non-perturbative breaking of the gauge symmetry. Thus we see that for general RR backgrounds \( F^{RR} \neq 0 \) constant B-field is not a solution of this equation. In other words, standard non-commutativity models based on open superstring theory in a constant B-field are not conformally invariant on the worldsheet (on the non-perturbative level) in the presence of the RR-fields. One way to cure this problem and to restore the conformal symmetry along the gauge symmetry is to introduce appropriate D-brane configuration which would screen the RR-charges. In terms of the effective action this would correspond to shifting the B-field as \( B \to B + dA \) where \( A \) is the D-brane’s U(1) field. As a result, the terms with the B-field will be transformed into those with the the \( B_{mn} + F_{mn} \) - type structure; the terms that were breaking the B-field gauge symmetry before the introduction of
D-branes would evolve into those of the type \((C^{RR(p)} + C^{RR(p-2)} \wedge (B + dA))^2\) which are gauge invariant under the appropriate combination of the U(1) and B field gauge transformations. In other words, the gauge symmetry will be restored by the D-brane’s U(1) field. Therefore the introduction of D-brane backgrounds is indispensable to preserve the B-field gauge symmetry on the non-perturbative level, as well as to consistently formulate non-commutative theories in Ramond-Ramond backgrounds, based on a constant B-field.

5. Conclusions

In this letter we have considered the Ramond-Ramond vertex operators at non-canonical pictures, showing that they may be considered as sources of the RR gauge potential. In this way these operators have a non-perturbative character (similarly to D-branes) and correlation functions involving these operators contain essential information about the non-perturbative physics of strings. We also have discussed one particularly interesting non-perturbative effect related to these vertices - the breaking of the B-field gauge symmetry. The only way to restore this gauge symmetry is to introduce D-branes so that the presence of the D-brane U(1) field compensates for the B-field gauge non-invariance. Therefore the presence of the non-canonical vertex operators in superstring spectrum automatically entails the introduction of D-branes to insure the gauge symmetry. The gauge symmetry restoration may also be understood in another way. Namely, the dynamics of various D-branes can be represented in terms of the NS and NS-NS brane-like vertex operators similar to those appearing as intermediate extra states in the amplitude (23). As we have seen, these intermediate extra states are in fact those leading to the gauge symmetry violation. Therefore introducing D-brane backgrounds (described by the brane-like states) is equivalent to screening these intermediate poles. Such a screening insures that the gauge symmetry associated with the B-field is restored.

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