Creation of Fundamental String in M(atrix) Theory

Nobuyoshi Ohta, Takashi Shimizu and Jian-Ge Zhou

Department of Physics, Osaka University,
Toyonaka, Osaka 560, Japan

Abstract

The potential between two D4-branes at angles with partially unbroken supersymmetry is computed, and is used to discuss the creation of a fundamental string when two such D4-branes cross each other in M(atrix) theory. The effective Lagrangian is shown to be anomalous at 1-loop approximation, but it can be modified by bare Chern-Simons terms to preserve the invariance under the large gauge transformation. The resulting effective potential agrees with that obtained from the string calculations. The result shows that a fundamental string is created in order to cancel the repulsive force between two D4-branes at proper angles.

*ohta@phys.wani.osaka-u.ac.jp
†simtak@phys.wani.osaka-u.ac.jp
‡jgzhou@phys.wani.osaka-u.ac.jp, JSPS postdoctral fellow
Recently the creation of a fundamental string when a D0-brane crosses a D8-brane in type IIA theory, and the creation of a longitudinal membrane in M-theory when two longitudinal M5-branes cross each other, have been studied from various points of view [1-7].

In the string theory context, by requiring RR charge conservation, it was first found in [1] that a third brane is created when two certain branes cross. In [2], the anomaly equation was exploited to show that when the two branes in question cross each other an energy level crosses zero and a single particle or hole is created, and this was interpreted as the creation of an open string or brane. It was also suggested that the induced charge on the D8-brane worldvolume indicates the creation of a string when the D0-brane crosses the D8-brane [3]. The one-loop open string calculation in ref. [4] revealed that the potential of the D0-D8 system vanishes due to the cancellation of the forces coming from the dilaton-graviton exchange and a fundamental string created when two such branes cross.

On the other hand, more and more evidence has accumulated to show that M-theory could be described in terms of a matrix model – M(atrix) theory [8]. Many consistency checks have been done, including the calculation of potentials between various D-branes [9-14]. It is then natural to study the brane creation in the context of M(atrix) theory, but only a few papers discuss this phenomenon. The authors of ref. [6] argued that the effective potential between two M5-branes is dominated by the contribution from the chiral fermionic zero mode in the off-diagonal degrees of freedom. On the other hand, ref. [13] calculated the effective potential between D0- and D8-branes, with the approximation of small $b$ and $c_i$ ($b$ is the distance between D0- and D8-branes along $X_9$, and $c_i$’s stand for backgrounds). Since the effective potential is computed only for small $b$ and $c_i$, it is hard to determine from the potential in [13] whether a fundamental string is created or not when D0-brane crosses D8-brane.

The creation of a fundamental string in D0-D8 system can be related to the creation of branes of other dimensions in other systems by sequences of dualities. For example, after T-dualities, when one D4-brane along (1357) directions crosses another D4-brane in (2468) directions, the creation of a fundamental string occurs. But when two parallel D4-branes cross each other, nothing happens. This raises an obvious question: is a fundamental string created when two D4-branes at angles cross each other?

In the present paper, we compute the effective potential without the approximation made in ref. [13] and use the result to discuss the brane creation in M(atrix) theory. In
particular we consider the creation of a fundamental string (a longitudinal M2-brane in M-theory) when two D4-brane (wrapped M5-brane in M-theory) at angles \cite{15} cross each other. The classical configuration we choose is the bound state \{(4+2+2+0)-(4+2+2+0)\}, which is T-dual to two D4-branes at angles \cite{16}.

To describe the creation of a fundamental string, the one-loop effective action for the present classical background is calculated, from which the potential between two D4-branes at angles can be read off. It is found that when the backgrounds \(c_i\) are arbitrary, the effective action for general \(b\) is quite complicated, and the precise dependence of the potential on general \(b\) cannot be determined. However, for arbitrary \(c_i\), the classical configuration \{(4+2+2+0)-(4+2+2+0)\} breaks supersymmetry \cite{17}, casting suspicion on the validity of the one-loop approximation. When \(c_1 = c_2, c_3 = c_4\) and \(c_i \neq 0\), on the other hand, the resulting configuration preserves 1/8 unbroken supersymmetry, and when \(c_1 = c_2 = c_3 = c_4 \neq 0\), the unbroken supersymmetry is enhanced to 3/16 \cite{16, 18}. In these cases, we may rely on our one-loop calculation. We find that when \(c_1 = c_2, c_3 = c_4\) and \(c_i \neq 0\), the effective Lagrangian is surprisingly simplified, and is given by \(\frac{1}{2}\text{sign}(b)(b+a_0)\) (where we have turned on the background \(a_0\) for gauge field \(A_0\)), completely independent of the backgrounds \(c_i\). The resulting effective Lagrangian is anomalous under the global gauge transformation \cite{19, 20}, but is invariant under \(b \rightarrow -b, a_0 \rightarrow -a_0\) transformation which corresponds to the charge conjugation invariance \cite{19}. One can restore the invariance under the large gauge transformation by adding bare Chern-Simons (CS) terms of the form \(-\frac{1}{2}(b + a_0)\) to the effective Lagrangian, but at the expense of breaking charge conjugation invariance.\footnote{This could be understood as a result of regularization \cite{19}, or as bare terms allowed from symmetries in the theory \cite{21}.}

We find the resulting potential in M(atrix) theory is independent of the backgrounds \(c_i\) related to the angles between two D4-branes (after T-dualities). For \(c_1 = c_2 = c_3 = c_4 = 0\), the potential vanishes, which indicates that when two parallel D4-branes cross each other, no fundamental string is created. For \(c_1 = c_2, c_3 = c_4\), both nonzero, the potential produces a jump (exactly the same amount of string tension) in the force acting on the brane when the above two D4-branes at angles cross each other. In order to maintain the BPS property, that is, for the total potential to vanish on both sides of D4-brane, a fundamental string must be created. It is interesting to note that this occurs independently of the nonzero angles between two D4-branes. Our conclusion is consistent
with the string calculations [3, 4].

Let us start with the M(atrix) theory Lagrangian which is the 10-dimensional $U(N)$ super Yang-Mills Lagrangian reduced to $1 + 0$ dimensions [8]

$$L = \frac{1}{2g_s} \text{Tr} \left( D_t X_i D_t X_i + \frac{1}{2} [X_i, X_j]^2 + i \theta^4 D_t \theta + \theta^4 \gamma_i [X_i, \theta] \right),$$

where we have set $T_s^{-1} = 2\pi \alpha' = 1$, and $X_i (i = 1, \cdots, 9)$ and $\theta$ are bosonic and fermionic hermitian $N \times N$ matrices.

We take the following background configuration:

$$\bar{A}_0 = \begin{pmatrix} a_0 & 0 \\ 0 & 0 \end{pmatrix}, \quad \bar{X}_1 = \begin{pmatrix} Q_1 & 0 \\ 0 & 0 \end{pmatrix}, \quad \bar{X}_2 = \begin{pmatrix} P_1 & 0 \\ 0 & 0 \end{pmatrix},$$
$$\bar{X}_3 = \begin{pmatrix} Q_2 & 0 \\ 0 & 0 \end{pmatrix}, \quad \bar{X}_4 = \begin{pmatrix} P_2 & 0 \\ 0 & 0 \end{pmatrix}, \quad \bar{X}_5 = \begin{pmatrix} 0 & 0 \\ 0 & Q_3 \end{pmatrix}, \quad \bar{X}_6 = \begin{pmatrix} 0 & 0 \\ 0 & -P_3 \end{pmatrix},$$
$$\bar{X}_7 = \begin{pmatrix} 0 & 0 \\ 0 & Q_4 \end{pmatrix}, \quad \bar{X}_8 = \begin{pmatrix} 0 & 0 \\ 0 & -P_4 \end{pmatrix}, \quad \bar{X}_9 = \begin{pmatrix} b \mathbf{1} & 0 \\ 0 & 0 \end{pmatrix},$$

(2)

where $[Q_1, P_1] = ic_1, [Q_2, P_2] = ic_2, [Q_3, P_3] = ic_3, [Q_4, P_4] = ic_4$. Note that the background $a_0$ has been turned on for the gauge field $A_0$. In the M(atrix) theory language, the D4-brane is described by a configuration corresponding to a $U(N)$ instanton and this background configuration may be interpreted as $(4 + 2 + 2 + 0) - (4 + 2 + 2 + 0)$ [17]. After T-dualities ($T_{1357}$), it can also be interpreted as one D4-brane lying in the $(1357)$ plane, the other being rotated off the $(1357)$ plane by the rotations in the $(12), (34), (56)$ and $(78)$ planes [14]. The angles $\theta_{12}, \theta_{34}, \theta_{56}$ and $\theta_{78}$ mix the directions $(12), (34), (56)$ and $(78)$, and $\tan \theta_{12}, \tan \theta_{34}, \tan \theta_{56}$ and $\tan \theta_{78}$ are proportional to $c_1, c_2, c_3$ and $c_4$, respectively.

In order to compute the effective Lagrangian, we expand Lagrangian (1) to quadratic order in the fluctuations around the above background ($A_0 = \bar{A}_0 + Y_0, X_i = \bar{X}_i + Y_i$), and integrate out the off-diagonal matrix elements which correspond to the degrees of freedom of the virtual strings stretched between two D4-branes at angles. The off-diagonal blocks can be chosen as [10]

$$Y_0 = \begin{pmatrix} 0 & \phi_0 \\ \phi_0^\dagger & 0 \end{pmatrix}, Y_i = \begin{pmatrix} 0 & \phi_i \\ \phi_i^\dagger & 0 \end{pmatrix}, \theta = \begin{pmatrix} 0 & \psi \\ \psi^\dagger & 0 \end{pmatrix},$$

(3)

Following refs. [10, 12, 21], we can integrate out the off-diagonal modes to get the following
determinants (with \( \tau = it \)):

**Bosons:** \[ \det^{-2}[-(\partial_\tau + a_0)^2 + H] \]
\[ \det^{-1}[-(\partial_\tau + a_0)^2 + (H - 2c_1)] \det^{-1}[-(\partial_\tau + a_0)^2 + (H + 2c_1)] \]
\[ \det^{-1}[-(\partial_\tau + a_0)^2 + (H - 2c_2)] \det^{-1}[-(\partial_\tau + a_0)^2 + (H + 2c_2)] \]
\[ \det^{-1}[-(\partial_\tau + a_0)^2 + (H - 2c_3)] \det^{-1}[-(\partial_\tau + a_0)^2 + (H + 2c_3)] \]
\[ \det^{-1}[-(\partial_\tau + a_0)^2 + (H - 2c_4)] \det^{-1}[-(\partial_\tau + a_0)^2 + (H + 2c_4)], \quad (4) \]

**Ghost:** \[ \det^2[-(\partial_\tau + a_0)^2 + H], \quad (5) \]

**Fermions:** \[ \det[(\partial_\tau + a_0) + m_f], \quad (6) \]

where
\[ H = \sum_{i=1}^{4}(Q_i^2 + P_i^2) + b^2, \]
\[ m_f = Q_1\gamma_1 + P_1\gamma_2 + Q_2\gamma_3 + P_2\gamma_4 - Q_3\gamma_5 + P_3\gamma_6 - Q_4\gamma_7 + P_4\gamma_8 + b\gamma_9. \]

The \( Q_i^2 + P_i^2 \) terms in \( H \) are a collection of simple harmonic oscillator Hamiltonians with eigenvalues \( c_i(2n_i + 1), n_i = 0, 1, \cdots \). Since there is no tenth \( 16 \times 16 \) matrix which anticommutes with all the \( \gamma_i \)'s, the fermionic determinant cannot be converted into Klein-Gordon form by the usual method. We can still compute \( \frac{\partial \Gamma_F(b,a_0)}{\partial a_0} \) to find from (3)
\[ i \frac{\partial \Gamma_F(b,a_0)}{\partial a_0} = \text{Tr} \left( \frac{1}{\partial_\tau + a_0 + m_f} \right) \quad (8) \]

Multiplying the denominator by \(- (\partial_\tau + a_0) + m_f \) and using Schwinger’s proper-time representation, we have
\[ i \frac{\partial \Gamma_F}{\partial a_0} = \text{Tr} \left( \left\{ - (\partial_\tau + a_0) + Q_1\gamma_1 + P_1\gamma_2 + Q_2\gamma_3 + P_2\gamma_4 - Q_3\gamma_5 + P_3\gamma_6 - Q_4\gamma_7 + P_4\gamma_8 + b\gamma_9 \right\} \right. \]
\[ \times \left. \int_0^\infty ds \exp \left\{ - s \left[ - (\partial_\tau + a_0)^2 + H + ic_1\gamma_{12} + ic_2\gamma_{34} + ic_3\gamma_{56} + ic_4\gamma_{78} \right] \right\} \right). \quad (9) \]

Using the Fourier representation for the \( \tau \) variable, one can take its trace. After the shift of the integration variable, one finds that the first term \((\partial_\tau + a_0)\) in the first curly bracket vanishes, and the next 8 terms are zero because of the traces of odd numbers of \( \gamma_i \). We thus find that eq. (3) is reduced to
\[ i \frac{\partial \Gamma_F}{\partial a_0} = \int_0^\infty ds \int_{-\infty}^{\infty} \frac{dk}{2\pi} \int d\tau b \text{Tr} \left( \gamma_9 \exp \left\{ - s \left[ k^2 + H + ic_1\gamma_{12} + ic_2\gamma_{34} + ic_3\gamma_{56} + ic_4\gamma_{78} \right] \right\} \right) \quad (10) \]
By choosing proper representation for $\gamma_i$, we can arrange $\gamma_{12}$, $\gamma_{34}$, $\gamma_{56}$, $\gamma_{78}$ and $\gamma_9$ to take the form \[18\]

\[
\begin{align*}
\gamma_{12} &= i \text{ diag.}(1_{8\times8}, -1_{8\times8}), \\
\gamma_{34} &= i \text{ diag.}(1_{4\times4}, -1_{4\times4}, 1_{4\times4}, -1_{4\times4}), \\
\gamma_{56} &= i \text{ diag.}(1_{2\times2}, -1_{2\times2}, \cdots, 1_{2\times2}, -1_{2\times2}), \\
\gamma_{78} &= i \text{ diag.}(1, -1, \cdots, 1, -1), \\
\gamma_9 &= \gamma_{12345678} = \text{ diag.}(1, -1, -1, 1, -1, 1, -1, 1, -1, 1, -1, 1), \tag{11}
\end{align*}
\]

where $1_{n\times n}$ is the $n \times n$ identity matrix. Substituting (11) into (10) and performing the integrations yield

\[
\frac{\partial \Gamma_F}{\partial a_0} = \int dt \frac{|b|}{2b}, \tag{12}
\]

for $c_i \neq 0$. When any one (two or three) of $c_i$ is (are) zero, we find

\[
\frac{\partial \Gamma_F}{\partial a_0} = 0. \tag{13}
\]

It is easy to repeat a similar calculation for the contributions from the bosons and ghosts, and we find that they vanish. Consequently eqs. (12) and (13) are the whole results for the derivative of the total effective action:

\[
\frac{\partial \Gamma}{\partial a_0} = \int dt \frac{|b|}{2b}, \quad \text{for } c_i \neq 0, \tag{14}
\]

and

\[
\frac{\partial \Gamma}{\partial a_0} = 0, \quad \text{when at least one of } c_i \text{ is zero}, \tag{15}
\]

with

\[
\Gamma = \Gamma_B + \Gamma_G + \Gamma_F. \tag{16}
\]

Here we point out that in deriving the results (14) and (13), we have not imposed any restriction on $c_i$.

Similarly, when $c_i \neq 0$, we find

\[
\frac{\partial \Gamma}{\partial b} = \frac{b}{2\sqrt{\pi}} \int dt \int_0^\infty dss^{-1/2}e^{-b^2s} \sum_{n_1, \cdots, n_4=0}^{\infty} e^{-s} \sum_{i=1}^{4} c_i (2n_i + 1)
\]

\[
\times \left\{ e^{-s(c_1+c_2-c_3-c_4)} + e^{-s(c_1+c_2-c_3+c_4)} + e^{-s(c_1+c_2+c_3-c_4)} + e^{-s(c_1+c_2+c_3+c_4)} \\
+ e^{-s(c_1-c_2-c_3-c_4)} + e^{-s(c_1-c_2-c_3+c_4)} + e^{-s(c_1-c_2+c_3-c_4)} + e^{-s(c_1-c_2+c_3+c_4)} \\
+ e^{-s(-c_1+c_2-c_3-c_4)} + e^{-s(-c_1+c_2-c_3+c_4)} + e^{-s(-c_1+c_2+c_3-c_4)} + e^{-s(-c_1+c_2+c_3+c_4)} \\
+ e^{-s(-c_1-c_2-c_3-c_4)} + e^{-s(-c_1-c_2-c_3+c_4)} + e^{-s(-c_1-c_2+c_3-c_4)} + e^{-s(-c_1-c_2+c_3+c_4)} \\
- 2e^{-2sc_1} - 2e^{2sc_1} - 2e^{-2sc_2} - 2e^{2sc_2} - 2e^{-2sc_3} - 2e^{2sc_3} - 2e^{-2sc_4} - 2e^{2sc_4} \right\}. \tag{17}
\]
Eq. (17) looks so complicated that it seems impossible to get a definite result for $\partial \Gamma / \partial b$ without the approximation for $c_i$ or $b$, as has been done in ref. [13]. However, we are interested in the classical BPS-brane configurations in M(atrix) theory with partially unbroken supersymmetry. For the present classical configuration, we know that it has 1/8 unbroken supersymmetry when $c_1 = c_2$, $c_3 = c_4$ and $c_i \neq 0$ [16]. With this condition, all the contributions from nonzero modes cancel out from eq. (17), and we find $c_i$-independent result

$$\frac{\partial \Gamma}{\partial b} = \int dt \frac{|b|}{2b}.$$  

(18)

Note that when $c_1 = c_2$, $c_3 = c_4$, but $c_1$ or $c_3$ is zero, eq. (17) reduces to

$$\frac{\partial \Gamma}{\partial b} = 0.$$  

(19)

From eqs. (14)-(19), the effective action for the BPS-saturated background with $c_1 = c_2$ and $c_3 = c_4$ can be obtained:

$$\Gamma(b, a_0) = 0, \quad \text{for } c_1 \text{ or } c_3 \text{ is zero},$$  

and

$$\Gamma(b, a_0) = \int dt \frac{1}{2} \text{sign}(b)(b + a_0), \quad \text{for } c_1, c_3 \neq 0,$$  

(21)

where the irrelevant integration constants (independent of $b$ and $a_0$) have been dropped.

The potential obtained from the effective action (21) gives a jump in the force at $b = 0$ by a string tension. This has been interpreted in terms of half strings in refs. [3, 4]. We believe that the following modification gives a better interpretation.

Note that eq. (21) possesses the charge conjugation invariance under the transformation $b \rightarrow -b, a_0 \rightarrow -a_0$, but it is anomalous due to the presence of the term: $\frac{1}{2} \text{sign}(b)a_0$. The factor $\frac{1}{2}$ indicates that the partition function is not invariant under the global gauge transformations, leading to a global anomaly [13, 21]. The origin of this global anomaly is the chiral fermionic zero mode we noted in the preceding footnote. However, the invariance under the global gauge transformations can be restored by adding bare CS terms to the Lagrangian, but they break the charge conjugation invariance. From the viewpoint

---

2If one expands eq. (17) in small $c_i$ up to fourth orders, one finds the result reported in ref. [13].

3 The first 16 terms in eq. (17) are the contributions from fermions. We can easily see that there is only one zero eigenvalue in the spectrum (in the thirteenth term) for $b = 0$ and all $n_i = 0$, which is the chiral zero mode identified in ref. [3].
of the world-line theory of the D-particle, when \( b \ll l_s \), our D4-D4 system can be approximately described by the open string model, and the resulting theory is supersymmetric quantum mechanics with 4 supercharges. Imposed this symmetry, the bare CS terms can be chosen as \[ [19, 20, 4] \] \[
\Gamma_{cs} = - \int dt \frac{1}{2} (b + a_0). \] (22)

Thus the effective Lagrangian is given by
\[
\mathcal{L}_{eff} = \frac{1}{2} \text{sign}(b)(b + a_0) - \frac{1}{2}(b + a_0), \tag{23}
\]
which restores the invariance under the large gauge transformation at the expense of breaking charge conjugation invariance: \( b \to -b, \ a_0 \to -a_0 \). \[ A \] Actually this modification can be interpreted as different choices of the regularization schemes \[ [19] \], for example, using Pauli-Villars regularization.

We also note that the theory is anomaly-free on the backgrounds of D0-brane interacting with D2-, D4- and D6-branes, as can be seen from eq. (15).

Performing T-dualities (\( T_{1357} \)), the \( \{(4 + 2 + 2 + 0) - (4 + 2 + 2 + 0)\} \) background configuration can be interpreted as one of the D4-brane lying on (1357) plane, and the other rotated away from (1357) plane along (12) and (34) directions with an angle \( \theta_1 \) and along (56) and (78) directions with an angle \( \theta_2 \). Under T-dualities mapping the configuration to that of two D4-branes at angles, the distance \( b \) is not changed. Hence after choosing the \( a_0 = 0 \) gauge, the interaction potential between two D4-branes at angles at distance \( b \) is simply given by \[ [13] \]
\[
\Gamma(b, a_0 = 0) = - \int dt V(b). \tag{24}
\]

From eqs. (23) and (24), the effective potential between two D4-branes at angles can be read off as
\[
V(b) = -\frac{1}{2} T_s (|b| - b), \tag{25}
\]
where we have switched on the string tension \( T_s = (2\pi\alpha')^{-1} \). An important point is that the jump in the force at \( b = 0 \) is unchanged by this procedure.

We emphasize that the jump in the force is due to the contribution of the fermionic zero mode. Also note that the potential in eq. (23) is independent of the angles between

\[ A \]

\[ A \] A similar discussion of the effective action is given in somewhat simplified manner for \( N = 8 \) supersymmetric quantum mechanics in ref. \[ [20] \] and is later exploited in refs. \[ [13, 14] \].
two D4-branes, and is the same as that for two orthogonal D4-branes with eight ND-directions.

Eq. (20) shows that when $c_1$ or $c_3$ is zero, the potential vanishes. Suppose $c_3 = 0$. We have the following physical picture: the first D4-brane lies on (1357) directions, and the second is rotated off (13) plane along (12) and (34) directions with two common directions (57) with the first D4-brane. This configuration corresponds to two D2-branes at angle and the rotation is a real element of $SO(2)$ [15]. Since the potential vanishes, when two such D4-branes cross each other, no fundamental string is created.

Let us next consider $c_1 = c_2 = c_3 = c_4 = c$ case where the unbroken supersymmetry is enhanced from $1/8$ to $3/16$ [16, 18], and in the resulting configuration the orientations of the two D4-branes are related by a rotation in $Sp(2)$ subgroup of $SO(8)$ commuting with multiplication by a quaternion [18]. If $c = 0$ i.e., two D4-branes are parallel and lie on (1357) directions, the unbroken supersymmetry is further enhanced from $3/16$ to $1/2$ [22, 23]. Since the potential vanishes, no fundamental string is created when two parallel D4-branes cross each other. If $c \neq 0$, the potential gets contributions only from the fermionic zero mode (and bare CS terms) and is given by eq. (25) which shows that there is no force on one side and a repulsion of a single string tension occurs when two D4-branes at angles cross. This force is canceled by a fundamental string created between them in order to maintain the BPS property. The origin of this mechanism will be further discussed below. Since eq. (25) is independent of $c_1(c_3)$, the above conclusion is valid for finite angles. In particular the case of the angles equal to $\pi/2$ corresponds to two orthogonal D4-branes with 8 ND-directions related to that extensively discussed in the literature [21, 22, 23, 24, 25].

We thus find that the potential obtained from the effective Lagrangian in M(atrix) theory misses the term from $R(-1)^F$ sectors, but it can be modified to agree with that obtained from the string calculations in [4].

From the above discussion, we see that the brane creation is closely related to the existence of a chiral fermionic zero mode in the off-diagonal degrees of freedom, which in the present case manifests itself in the form of the global anomaly, and the created string is needed to cancel this global anomaly completely. This can be understood from the point of view of the worldvolume theory. Consider two D4-branes: $D^{(1)4}$ and $D^{(2)4}$ with their worldvolumes $B_1, B_2$, and the intersection between them $B_{12} = B_1 \cap B_2$ is I-brane [22].

\footnote{A similar observation was also made in refs. [4, 13].}
The worldvolume actions of $D^{(1)4}$, $D^{(2)4}$ and I-brane are denoted by $S_1$, $S_2$ and $S_I$. Under a gauge transformation, the variation of $S_1$ and $S_2$ has a boundary piece localized at $B_{12}$, which precisely cancels the anomalous variation of $S_I$ [24, 2]. In our case, the effective action calculated from the off-diagonal blocks in M(atrix) theory corresponds to the action $S_I$ of the worldvolume theory. Under large gauge transformation, the effective action has a global anomaly indeed. Such anomaly from $S_I$ can be canceled by the variation of the bulk actions. Thus the inflow of charge that is required for the absence of global anomaly can be regarded as the created fundamental string.

In M(atrix) theory, to implement such anomaly cancellation, one would add explicit CS terms $\int C \wedge \text{tr}F$ to the M(atrix) theory Lagrangian, and the anomaly can be canceled by the anomaly inflow from the bulk mediated by these CS terms [24, 26]. Indeed, such CS terms combined with the D8-brane background were used in ref. [3] to cancel the potential. In [25], such CS couplings were suggested to be an effective description of a more microscopic mechanism where the supergravity background could be generated by integrating out certain heavy matrix modes. The effective result could be expected as the source for the creation of the fundamental string and extra degrees of freedom corresponding to the created string might be found from $D^{(1)4}$-$D^{(1)4}$ and $D^{(2)4}$-$D^{(2)4}$ strings. On the other hand, the recent work in refs. [27] suggests that there may be more degrees of freedom to M(atrix) theory than just 0-branes. It would be interesting to see if these CS terms have any implications for “corrections” to the original model [8]. Work along this line is under investigation.

In conclusion, the creation of a fundamental string when two D4-branes at angles cross each other has been discussed in the context of M(atrix) theory. When $c_1 = c_2$, $c_3 = c_4$ and $c_i \neq 0$, the background possesses $1/8$ unbroken supersymmetry, the effective action is surprisingly simplified. We have found that the potential obtained from the effective Lagrangian possessing the invariance under the large gauge transformation is independent of $c_i$ related to the angles between two D4-branes (after T-dualities) and exhibits a jump in the force by the amount of a string tension. This result indicates that a fundamental sting is created indeed when two D4-branes at angles (preserving $1/8$ or $3/16$ unbroken supersymmetry) cross each other in order to cancel the force. Such a string creation is related to the existence of the global anomaly, and we have interpreted the created string as the object required to cancel this global anomaly completely. Our results are consistent with string calculations [3, 4].
Acknowledgments: We would like to thank I. Klebanov and T. Nakatsu for helpful comments. This work was supported in part by Grand-in-aid from the Ministry of Education, Science, Sports and Culture No. 96208. J.-G. Zhou thanks the Japan Society for the Promotion of Science for the financial support.

References

[1] A. Hanany and E. Witten, Nucl. Phys. B492 (1997) 152, hep-th/9611230.

[2] C. Bachas, M. Douglas and M. Green, preprint, hep-th/9705074.

[3] U. Danielsson, G. Ferretti and I.R. Klebanov, Phys. Rev. Lett. 79 (1997) 1984, hep-th/9705084; I.R. Klebanov, preprint, hep-th/9709160; U. Danielsson and G. Ferretti, preprint, hep-th/9709171.

[4] O. Bergman, M. Gaberdiel and G. Lifschytz, preprint, hep-th/9705130.

[5] S.P. de Alwis, preprint, hep-th/9706142.

[6] P.-M. Ho and Y.-S. Wu, preprint, hep-th/9708137.

[7] Y. Imamura, preprint, hep-th/9710026.

[8] T. Banks, W. Fischler, S.H. Shenker and L. Susskind, Phys. Rev. D55 (1997) 5112, hep-th/9610043.

[9] O. Aharony and M. Berkooz, Nucl. Phys. B491 (1997) 184, hep-th/9611215.

[10] G. Lifschytz and S.D. Mathur, preprint, hep-th/9612087; G. Lifschytz, preprint, hep-th/9612223.

[11] G. Lifschytz, preprint, hep-th/9703201.

[12] I. Chepelev and A.A. Tseytlin, preprints, hep-th/9704127; hep-th/9705120; hep-th/9709087; E. Keski-Vakkuri and P. Kraus, preprints, hep-th/9706196; hep-th/9709122.

[13] J.M. Pierre, preprint, hep-th/9705110.

[14] R. Gopakumar and S. Ramgloom, preprint, hep-th/9708022.
[15] M. Berkooz, M.R. Douglas and R.G. Leigh, Nucl. Phys. B480 (1996) 256, hep-th/9606139; V. Balasubramanian and R.G. Leigh, Phys. Rev. D55 (1997) 6415, hep-th/9611163.

[16] N. Ohta and J.-G. Zhou, preprint, hep-th/9709065.

[17] T. Banks, N. Seiberg and S. Shenker, Nucl. Phys. B490 (1997) 91, hep-th/9612157.

[18] J.P. Gauntlett, G.W. Gibbons, G. Papadopoulos and P.K. Townsend, Nucl. Phys. B500 (1997) 133, hep-th/9702202; P.K. Townsend, preprint, hep-th/9708074; N. Ohta and P.K. Townsend, preprint, hep-th/9710129.

[19] S. Elitzur, E. Rabinovici, Y. Frishman and A. Schwimmer, Nucl. Phys. B273 (1986) 93.

[20] T. Banks, N. Seiberg and E. Silverstein, Phys. Lett. B401 (1997) 30, hep-th/9703052.

[21] M.R. Douglas, D. Kabat, P. Pouliot and S.H. Shenker, Nucl. Phys. B485 (1997) 85, hep-th/9608024.

[22] J. Polchinski, Phys. Rev. Lett. 75 (1995) 4724, hep-th/9510017; preprint, hep-th/9611050.

[23] G. Papadopoulos and P.K. Townsend, Phys. Lett. B380 (1996) 273, hep-th/9603087; A.A. Tseytlin, Nucl. Phys. B475 (1996) 149, hep-th/9604033; J.P. Gauntlett, D.A. Kastor and J Traschen, Nucl. Phys. B478 (1996) 544, hep-th/9604179; J.-G. Zhou, H.J.W. Müller-Kirsten, J.-Q. Liang and F. Zimmerschild, Nucl. Phys. B487 (1997) 155, hep-th/9611148; R. Argurio, F. Englert and L. Houart, Phys. Lett. B398 (1997) 61, hep-th/9701042; N. Ohta, Phys. Lett. B403 (1997) 218, hep-th/9702164.

[24] M.B. Green, J.A. Harvey and G. Moore, Class. Quant. Grav. 14 (1997) 49, hep-th/9605033.

[25] P. Hořava, preprint, hep-th/9705055.

[26] D. Kabat and S.-J. Rey, preprint, hep-th/9707093.

[27] A. Sen, preprint, hep-th/9709220; N. Seiberg, preprint, hep-th/971009; M. Dine and A. Rajaraman, preprint, hep-th/9710174; M. Douglas and H. Ooguri, preprint, hep-th/9710178.