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Higher dimensional recombination of intersecting D-branes

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Abstract: We study recombinations of D-brane systems intersecting at more than one angle using super Yang-Mills theory. We find the condensation of an off-diagonal tachyon mode relates to the recombination, as was clarified for branes at one angle in [5]. For branes at two angles, after the tachyon mode between two D2-branes condensed, D2-brane charge is distributed in the bulk near the intersection point. We also find that, when two intersection angles are equal, the off-diagonal lowest mode is massless, and a new stable non-abelian configuration, which is supersymmetric up to a quadratic order in the fluctuations, is obtained by the deformation by this mode.

Keywords: D-branes, Tachyon Condensation.

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Contents

1. Introduction

Low energy dynamics of D-brane is well known by the analysis using super Yang-Mills theories [1]. One of the D-brane systems which can be studied in Yang-Mills theories is intersecting D-brane systems. Fluctuation spectra of intersecting D-branes on a torus are studied by using Yang-Mills action in [2]. In intersecting brane systems, recombination is important in a context of string phenomenology. Standard Model Higgs mechanism is realized as a brane recombination [3]. On the other hand, it plays an important role to realize an inflation in brane world scenario [4]. In Yang-Mills theories, the recombination is studied in [5, 6] for branes at one angle.[1] Fundamental strings stretched between intersecting D-branes are also studied in Yang-Mills theories in [10]. We extend the study of recombination of intersecting D-branes to the case with more than one intersection angle. This enables us to study more complicated intersecting brane systems, which may appear in the context of Standard Model realized by intersecting D-branes. To clarify the mechanism of recombination in higher dimensions is one of the aims of this paper.

When we consider two intersection angles, we find a supersymmetric configuration where two intersection angles are equal. It is known that supersymmetric intersecting branes are embedded into spacetime on a calibrated surface to minimize their worldvol-

[1]It is also discussed in tachyon field theory [7] and Matrix theory [8]. Intersecting D-branes with a separation are discussed in [9].

ume.[2] Calibration equations are realized as BPS conditions in abelian Dirac-Born-Infeld actions [12]. An embedding realized by the calibration geometry is called in [13] as an abelian embedding. The dynamics of multiple D-branes is described by non-abelian Born-Infeld(NBI) actions. In the non-abelian cases, there are other embeddings which cannot be realized by the calibration geometry, and further, which are constructed from the components including off-diagonal elements of the fields in the NBI actions. We call such embeddings as non-abelian embeddings. There are various kinds of the non-abelian embeddings and they include many interesting characters but it has not yet been considered except for a few cases [13]. It is difficult to know the full NBI action because in the non-abelian cases, slowly varying field approximation is meaningless [14] and we must consider derivative terms together with field strengths.[3] But there are some non-abelian embeddings which Yang-Mills analysis is valid for. We believe that the results of such Yang-Mills analysis can be lifted smoothly to full NBI analysis.

In this paper, we study recombination of D2-branes intersecting at two angles. When two angles are not equal, the lowest mode of a Neveu-Schwarz sector is tachyonic, and considering a condensation of this mode, we obtain deformed intersecting branes. For branes at one angle, the recombination occurs locally near the intersection point and we can describe this phenomenon by a condensation of a tachyon mode which is localized at an intersection point [5]. The final state in this decay process is a set of two parallel D2-branes. For branes at two angles, the final state after the recombination is expected as a brane configuration which preserves 1/4 of the supersymmetries [17]. The tachyon mode we consider here is localized at the intersection point even in this case, therefore, the condensation of the tachyon mode describes the first step of the recombination, plays a role of a trigger. After the tachyon mode has condensed, the branes can not be realized as a simple geometrical surface, because two transverse scalar fields can not be diagonalized by any gauge transformation simultaneously. We find D2-brane charge distributed in the bulk near the intersection point after the tachyon condensation. When two intersection angles are equal, we obtain a massless mode which appears in an off-diagonal spectrum in Yang-Mills theory. By considering a deformation by this massless mode, we obtain a configuration which is supersymmetric up to a quadratic order in the fluctuations. In this configuration, D2-brane charge is again distributed in the bulk near the intersection point. The massless deformation considered here does not seem to be expressed by any abelian calibration geometry[4] and this is an interesting example of the non-abelian embeddings.

In section 2, we discuss a recombination of D2-branes intersecting at two angles. When two intersection angles are equal, the brane configuration is supersymmetric and there is an off-diagonal massless mode which is discussed in section 3. The connections with higher orders in field strengths are discussed in section 4. Section 5 is devoted to a conclusion and discussions.

[2]For a recent review, see [11].

[3]We know up to and including F^6 terms of the NBI action now [15, 16].

[4]Calibration of supersymmetric intersecting D-branes in terms of world-volume field theory is discussed in [18].

2. Recombination of D-branes intersecting at two angles

Before performing Yang-Mills analysis, we see a mass spectrum of a fundamental string stretched between two D-branes intersecting at multiple angles $\theta_i (i = 1, \ldots, a)$, which is obtained in the worldsheet analysis as [19, 20][5]

$$m_j^2 = \frac{1}{2\pi\alpha'} \sum_{i=1}^{a} (2n_i - 1)\theta_i \pm 2\theta_j \, , \tag{2.1}$$

where $n_i \in \mathbb{N}$.

2.1 Fluctuation modes

We start with a worldvolume effective action for two D2-branes, which is obtained by the dimensional reduction of a (9+1) dimensional SU(2) Yang-Mills action,

$$S = -T \, \mathrm{Tr} \int d^2x dt \left[(D_a Y^i)^2 + \frac{1}{2} F_{ab}^2 - \frac{1}{2} [Y^i, Y^j]^2 \right] . \tag{2.2}$$

The indices $a, b = 0, 1, 2$ denote directions along the worldvolume, and $i, j = 3, \ldots, 9$ denote directions of transverse collective coordinates. F_{ab} and $D_a Y^i$ are defined by

$$\begin{aligned} F_{ab} &= \partial_a A_b - \partial_b A_a - i[A_a, A_b] \, , \\ D_a Y^i &= \partial_a Y^i - i[A_a, Y^i] \, , \end{aligned} \tag{2.3}$$

where A_a's are worldvolume gauge fields and Y^i's are transverse scalar fields. Let us consider an intersecting D2-brane system. It is sufficient for us to consider the situation that the two D2-branes are embedded in 4 dimensions and do not extend in other dimensions. We take coordinates of the embedding space as $1, 2, 8$ and 9 and others as $3, \ldots, 7$. We describe intersection angles as θ_1 in x_1-Y^9 plane and θ_2 in x_2-Y^8 plane here. We consider a classical solution representing the intersecting D-branes

$$Y^9 = q_1 x_1 \sigma^3 \, , \qquad Y^8 = q_2 x_2 \sigma^3 \, , \qquad A_a = 0 \, , \tag{2.4}$$

where $q_\alpha (\alpha = 1, 2)$ is related to the intersection angle θ_α as $\theta_\alpha \equiv 2 \tan^{-1}(2\pi\alpha' q_\alpha)$ and we assume $q_\alpha > 0$ here.

Let us consider off-diagonal parts of fluctuations around the solution

$$\begin{aligned} Y^9 &= q_1 x_1 \sigma^3 + f_1(x_a)\sigma^1 - \bar{f}_1(x_a)\sigma^2 \, , \qquad & Y^8 &= q_2 x_2 \sigma^3 + f_2(x_a)\sigma^1 - \bar{f}_2(x_a)\sigma^2 \, , \\ A_1 &= g_1(x_a)\sigma^1 - \bar{g}_1(x_a)\sigma^2 \, , & A_2 &= g_2(x_a)\sigma^1 - \bar{g}_2(x_a)\sigma^2 \, . \end{aligned} \tag{2.5}$$

We adopt a gauge condition $A_0 = 0$. Diagonal parts of the fluctuations are decoupled from the off-diagonal fluctuations at the quadratic order, therefore we neglect the former ones.

[5] See also [21].

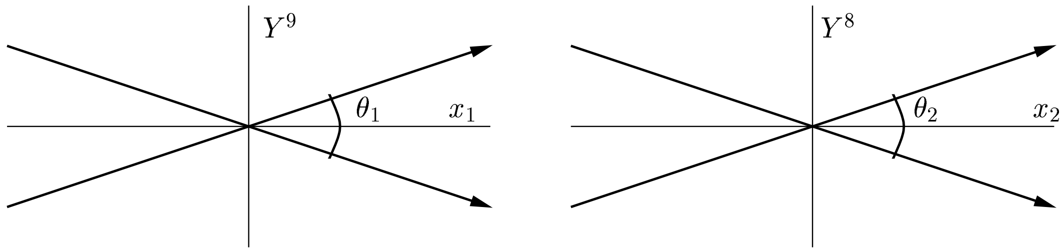

Figure 1: Intersecting D2-branes.

The lagrangian quadratic in the fluctuations is calculated as

$$L = \sum_{i,j=1,2} \Big(-(\partial_0 f_i)^2 - 4q_i f_i \bar{g}_i + (\partial_i f_j + 2\bar{g}_i q_j x_j)^2 - (\partial_0 \bar{g}_i)^2 + (\partial_i \bar{g}_j - \partial_j \bar{g}_i)^2 + \\ + 4(q_i x_i f_j - q_j x_j f_i)^2 + (f_\alpha \to \bar{f}, \bar{g}_\alpha \to g) \Big) . \tag{2.6}$$

The combinations $(f_\alpha, \bar{g}_\alpha)$ and $(\bar{f}_\alpha, g_\alpha)$ are decoupled each other at this order, therefore we neglect the $(\bar{f}_\alpha, g_\alpha)$ pair. From now on, we denote $\bar{g}$ as g for simplicity. The equations of motion for the fluctuations are written as

$$4 \begin{pmatrix} O_{11} & O_{12} \\ O_{21} & O_{22} \end{pmatrix} \begin{pmatrix} \vec{V}_1 \\ \vec{V}_2 \end{pmatrix} = 0 , \tag{2.7}$$

where

$$\begin{aligned} O_{11} &= \begin{pmatrix} \partial_0^2 - \partial_1^2 - \partial_2^2 + 4q_2^2 x_2^2 & -4q_1 - 2q_1 x_1 \partial_1 \\ -2q_1 + 2q_1 x_1 \partial_1 & 4q_1^2 x_1^2 + 4q_2^2 x_2^2 + \partial_0^2 - \partial_2^2 \end{pmatrix} , \\ O_{12} = O_{21} &= \begin{pmatrix} -4q_1 q_2 x_1 x_2 & -2q_1 x_1 \partial_2 \\ 2q_2 x_2 \partial_1 & \partial_1 \partial_2 \end{pmatrix} , \\ O_{22} &= \begin{pmatrix} \partial_0^2 - \partial_1^2 - \partial_2^2 + 4q_1^2 x_1^2 & -4q_2 - 2q_2 x_2 \partial_2 \\ -2q_2 + 2q_2 x_2 \partial_2 & 4q_1^2 x_1^2 + 4q_2^2 x_2^2 + \partial_0^2 - \partial_1^2 \end{pmatrix} , \end{aligned} \tag{2.8}$$

and

$$\vec{V}_1 = \begin{pmatrix} f_1(x_a) \\ g_1(x_a) \end{pmatrix} , \qquad \vec{V}_2 = \begin{pmatrix} f_2(x_a) \\ g_2(x_a) \end{pmatrix} . \tag{2.9}$$

Expanding the fluctuations by the mass eigenfunctions as

$$\vec{V}_1(x_1, x_2, t) = \sum_{n \geq 0} \begin{pmatrix} \tilde{f}_{1n}(x_1, x_2) \\ \tilde{g}_{1n}(x_1, x_2) \end{pmatrix} C_{1n}(t) , \qquad \vec{V}_2(x_1, x_2, t) = \sum_{n \geq 0} \begin{pmatrix} \tilde{f}_{2n}(x_1, x_2) \\ \tilde{g}_{2n}(x_1, x_2) \end{pmatrix} C_{2n}(t) , \tag{2.10}$$

where $C_{\rm in}$ $(i = 1, 2)$'s satisfy the following equations

$$(\partial_0^2 + m_{\rm in}^2) C_{\rm in}(t) = 0 , \tag{2.11}$$

we obtain the differential equations written as

$$4\begin{pmatrix} O'_{11} & O_{12} \\ O_{21} & O'_{22} \end{pmatrix}\begin{pmatrix} \vec{V}'_{1n} \\ \vec{V}'_{2n} \end{pmatrix} = \begin{pmatrix} m^2_{1n}\vec{V}'_{1n} \\ m^2_{2n}\vec{V}'_{2n} \end{pmatrix}, \tag{2.12}$$

where

$$\begin{aligned} O'_{11} &= \begin{pmatrix} -\partial_1^2 - \partial_2^2 + 4q_2^2x_2^2 & -4q_1 - 2q_1x_1\partial_1 \\ -2q_1 + 2q_1x_1\partial_1 & 4q_1^2x_1^2 + 4q_2^2x_2^2 - \partial_2^2 \end{pmatrix}, \\ O'_{22} &= \begin{pmatrix} -\partial_1^2 - \partial_2^2 + 4q_1^2x_1^2 & -4q_2 - 2q_2x_2\partial_2 \\ -2q_2 + 2q_2x_2\partial_2 & 4q_1^2x_1^2 + 4q_2^2x_2^2 - \partial_1^2 \end{pmatrix}, \end{aligned} \tag{2.13}$$

and

$$\vec{V}'_{1n} = \begin{pmatrix} \tilde{f}_{1n} \\ \tilde{g}_{1n} \end{pmatrix}, \qquad \vec{V}'_{2n} = \begin{pmatrix} \tilde{f}_{2n} \\ \tilde{g}_{2n} \end{pmatrix}. \tag{2.14}$$

By solving these differential equations, the eigenfunctions localized at the intersection point are obtained as

$$\begin{aligned} \vec{V}'_{10} &= \begin{pmatrix} \tilde{f}_{10}(x_1,x_2,t) \\ \tilde{g}_{10}(x_1,x_2,t) \end{pmatrix} = e^{-q_1x_1^2 - q_2x_2^2}C_{10}(t)\begin{pmatrix} 1 \\ 1 \end{pmatrix}, \\ \vec{V}'_{20} &= \begin{pmatrix} \tilde{f}_{20}(x_1,x_2,t) \\ \tilde{g}_{20}(x_1,x_2,t) \end{pmatrix} = e^{-q_1x_1^2 - q_2x_2^2}C_{20}(t)\begin{pmatrix} 1 \\ 1 \end{pmatrix}. \end{aligned} \tag{2.15}$$

The mass eigenvalues are obtained as

$$\begin{aligned} m^2_{10} &= 2(q_2 - q_1) \sim \frac{1}{2\pi\alpha'}(\theta_2 - \theta_1)\,, \\ m^2_{20} &= 2(q_1 - q_2) \sim \frac{1}{2\pi\alpha'}(\theta_1 - \theta_2)\,. \end{aligned} \tag{2.16}$$

There are massive and tachyonic modes, which coincide with the mass seen in the lowest mode of (2.1) in the small q region. Thus, we have obtained the correct lowest mode of the string mass spectrum in [19].

A brane system with a tachyon mode in open string theory is unstable and it rolls down to the stable vacuum by the condensation of the tachyon mode[6] [22]. When we consider the condensation of the tachyon mode (2.15), we obtain the configuration written as

$$\begin{aligned} Y^9 &= \begin{pmatrix} q_1x_1 & \tilde{f}_{10} \\ \tilde{f}_{10} & -q_1x_1 \end{pmatrix}, \qquad Y^8 = \begin{pmatrix} q_2x_2 & 0 \\ 0 & -q_2x_2 \end{pmatrix}, \\ A_1 &= \begin{pmatrix} 0 & i\tilde{g}_{10} \\ -i\tilde{g}_{10} & 0 \end{pmatrix}, \qquad A_2 = 0\,, \end{aligned} \tag{2.17}$$

where we consider a case $q_1 > q_2$. We cannot simultaneously diagonalize Y^8 and Y^9 by any gauge transformation, therefore, we do not find the simple recombination effect in the geometrical picture as was studied in [5]. We will discuss more about this point in the next subsection.

[6] Stability of branes at angles is studied by considering the potential between two branes in [23].

Next, we search for the excited mode. We assume the form of the eigenfunctions of the fluctuations as the product of exponential function written as $e^{-q_1x_1^2-q_2x_2^2}$ and linear function of x_a. There are the following three eigenfunctions. The first eigenfunctions are written as

$$\begin{pmatrix} \tilde{f}_{11}(x_1,x_2,t) \\ \tilde{g}_{11}(x_1,x_2,t) \end{pmatrix} = x_2 e^{-q_1x_1^2-q_2x_2^2} C_{11}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix},$$
$$\begin{pmatrix} \tilde{f}_{21}(x_1,x_2,t) \\ \tilde{g}_{21}(x_1,x_2,t) \end{pmatrix} = x_1 e^{-q_1x_1^2-q_2x_2^2} C_{21}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix}. \tag{2.18}$$

The corresponding mass eigenvalues are obtained as

$$m_{11}^2 = 2(3q_2 - q_1) \sim \frac{1}{2\pi\alpha'}(3\theta_2 - \theta_1),$$
$$m_{21}^2 = 2(3q_1 - q_2) \sim \frac{1}{2\pi\alpha'}(3\theta_1 - \theta_1). \tag{2.19}$$

The mass eigenvalues obtained here are found in the string mass spectrum in (2.1). The second eigenfunctions are written as

$$\begin{pmatrix} \tilde{f}_{12}(x_1,x_2,t) \\ \tilde{g}_{12}(x_1,x_2,t) \end{pmatrix} = x_1 e^{-q_1x_1^2-q_2x_2^2} C_{12}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix},$$
$$\begin{pmatrix} \tilde{f}_{22}(x_1,x_2,t) \\ \tilde{g}_{22}(x_1,x_2,t) \end{pmatrix} = x_2 e^{-q_1x_1^2-q_2x_2^2} C_{22}(t) \begin{pmatrix} -1 \\ -1 \end{pmatrix}, \tag{2.20}$$

where the corresponding mass eigenvalues are common in both eigenfunctions, that is, $C_{12}(t) = C_{22}(t)$. The mass eigenvalue is written as

$$m_{i2}^2 = 2(q_1 + q_2) \sim \frac{1}{2\pi\alpha'}(\theta_1 + \theta_2), \tag{2.21}$$

which is found in (2.1).

The third eigenfunctions written by the product of exponential function and linear function of x_i are written as

$$\begin{pmatrix} \tilde{f}_{13}(x_1,x_2,t) \\ \tilde{g}_{13}(x_1,x_2,t) \end{pmatrix} = q_1 x_1 e^{-q_1x_1^2-q_2x_2^2} C_{13}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix},$$
$$\begin{pmatrix} \tilde{f}_{23}(x_1,x_2,t) \\ \tilde{g}_{23}(x_1,x_2,t) \end{pmatrix} = q_2 x_2 e^{-q_1x_1^2-q_2x_2^2} C_{23}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix}, \tag{2.22}$$

where the squared of the mass eigenvalue is obtained as

$$m_1^2 = 0. \tag{2.23}$$

There is no massless mode in the string mass spectrum between intersecting branes at two different angles, except for some fixed angles. This mode is considered as Nambu-

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Goldstone mode of broken U(2) symmetry.[7] To see this explicitly, let us consider a gauge transformation for the intersecting brane solution written as

$$Y^9 = q_1 x_1 \sigma^3\,, \qquad Y^8 = q_2 x_2 \sigma^3\,, \qquad A_1 = A_2 = 0\,. \tag{2.24}$$

When we consider the gauge transformation which is written as

$$\begin{aligned} Y^8 \to \tilde{Y}^8 &= UY^8U^{-1}\,, & Y^9 \to \tilde{Y}^9 &= UY^9U^{-1}\,, \\ A_1 \to \tilde{A}_1 &= UA_1U^{-1} + i(\partial_1 U)U^{-1}\,, & A_2 \to \tilde{A}_2 &= UA_2U^{-1} + i(\partial_2 U)U^{-1}\,, \end{aligned} \tag{2.25}$$

where

$$\begin{aligned} \mathrm{U}(x_1, x_2) &= e^{i\Lambda(x_1,x_2)}\,, \\ \Lambda(x_1, x_2) &= -\frac{C_3}{2} e^{-q_1 x_1^2 - q_2 x_2^2} \sigma_2\,, \end{aligned} \tag{2.26}$$

we obtain the configuration as follows:

$$\begin{aligned} \tilde{Y}^9 &= q_1 x_1 (\sigma_3 + C_3 e^{-q_1 x_1^2 - q_2 x_2^2} \sigma_1) + \mathcal{O}(C_3^2)\,, \\ \tilde{Y}^8 &= q_2 x_2 (\sigma_3 + C_3 e^{-q_1 x_1^2 - q_2 x_2^2} \sigma_1) + \mathcal{O}(C_3^2)\,, \\ \tilde{A}_1 &= -q_1 x_1 C_3 e^{-q_1 x_1^2 - q_2 x_2^2} \sigma_2 + \mathcal{O}(C_3^2)\,, \\ \tilde{A}_2 &= -q_2 x_2 C_3 e^{-q_1 x_1^2 - q_2 x_2^2} \sigma_2 + \mathcal{O}(C_3^2)\,. \end{aligned} \tag{2.27}$$

We find that the off-diagonal part in (2.27) is equivalent with (2.22). An off-diagonal massless mode which is not a Nambu-Goldstone mode is considered in the next section.

Thus, we obtain the correct string mass of the lowest and first excited part of the spectrum in the approximation of small angles $\mathcal{O}(\theta_i)$. We generalize these analyses to D3-branes intersecting at three angles, which are written in appendix A.

2.2 Tachyon condensation and recombination

We focus on the tachyon mode (2.15) here. By the condensation of the tachyon mode, we obtain the configuration written as

$$\begin{aligned} Y^9 &= q_1 x_1 \sigma_3 + \tilde{f}_{10}\sigma_1\,, & Y^8 &= q_2 x_2 \sigma_3\,, \\ A_1 &= -\tilde{g}_{10}\sigma_2\,, & A_2 &= 0\,. \end{aligned} \tag{2.28}$$

We can not simultaneously diagonalize Y^8 and Y^9 by any gauge transformation, but we can diagonalize Y^8 and Y^9 in some region simultaneously. Let us look for the region where we can diagonalize Y^8 and Y^9 simultaneously. In the region,

$$|x_1| \gg \frac{C_1}{q_1}\,, \tag{2.29}$$

we can diagonalize Y^8 and Y^9 as

$$Y^9 \sim q_1 x_1 \sigma_3\,, \qquad Y^8 \sim q_2 x_2 \sigma_3\,. \tag{2.30}$$

[7] For branes at one angle, it is discussed in [24].

Far away from the intersection point along x_1 direction, the brane configuration remains unchanged, because the tachyon mode, which is due to the brane deformation, is localized near the intersection point. Thus, this result is easily explained. In the region $x_2 \sim 0$, we can also diagonalize Y^8 and Y^9 simultaneously as

$$Y^9 \sim \sqrt{(q_1 x_1)^2 + \tilde{f}_{10}^2}\sigma_3\,, \qquad Y^8 \sim 0\,. \tag{2.31}$$

We find that the intersection point is resolved, and 'recombination' occurs on the x_1-Y^9 plane near the intersection point. The other region can not be diagonalized and the D2-branes diffuse in the $x_1 x_2 Y^8 Y^9$ region near the intersection point. We will see the distribution of the D2-brane charge in the next subsection. Finally, let us remark a comment. If $q_1 \gg q_2$, there are many tachyonic modes with masses $2(2nq_2 + q_2 - q_1)$ $(n = 0, 1, 2, \ldots)$. In the limit $q_2 \to 0$, these infinite number of modes will be summed up by the form written as $f \sim Ce^{-q_1 x_1^2}\sigma_3$, which is equivalent to the fluctuation mode found for branes at one angle in [5].

2.3 D2-brane charge

D2-brane charge density in terms of the D2-brane worldvolume effective action is obtained by [25, 26] as

$$J_{0kl} = \frac{1}{12}\,\mathrm{Tr}(-iF^{ij}[Y^k, Y^l] - D_i Y^k D_j Y^l + D_i Y^l D_j Y^k)\,, \tag{2.32}$$

where i and j are parameterizations of worldvolume directions. D2-brane charge is already integrated along the direction Y^k and Y^l.

We consider the configuration written as

$$\begin{aligned} Y^9 &= q_1 x_1 \sigma_3 + \tilde{f}_{10}\sigma_1\,, \qquad & Y^8 &= q_2 x_2 \sigma_3 + \tilde{f}_{20}\sigma_1\,, \\ A_1 &= -\tilde{g}_{10}\sigma_2\,, A_2 & &= -\tilde{g}_{20}\sigma_2\,. \end{aligned} \tag{2.33}$$

If $q_1 > q_2$, the combination $(\tilde{f}_{10}, \tilde{g}_{10})$ corresponds to the tachyon mode and $(\tilde{f}_{20}, \tilde{g}_{20})$ corresponds to the massive mode. We calculate the charge density of this configuration. By the results,

$$\begin{aligned} -i\,\mathrm{Tr}\,F_{12}[Y^8, Y^9] &= 8(C_2 q_1 x_1 - C_1 q_2 x_2)^2 e^{-2q_1 x_1^2 - 2q_2 x_2^2}\,, \\ -\,\mathrm{Tr}\,D_1 Y^8 D_2 Y^9 &= 8(C_2 q_1 x_1 - C_1 q_2 x_2)^2 e^{-2q_1 x_1^2 - 2q_2 x_2^2}\,, \\ \mathrm{Tr}\,D_1 Y^9 D_2 Y^8 &= 2q_1 q_2 - 4(C_2^2 q_1 + C_1^2 q_2) e^{-2q_1 x_1^2 - 2q_2 x_2^2}\,, \end{aligned} \tag{2.34}$$

we obtain the D2-brane charge density as

$$J_{089} = \frac{1}{6}q_1 q_2 + \left(\frac{2}{3}(C_2 q_1 x_1 - C_1 q_2 x_2)^2 - \frac{1}{3}C_2^2 q_1 - \frac{1}{3}C_1^2 q_2\right) e^{-2q_1 x_1^2 - 2q_2 x_2^2}\,. \tag{2.35}$$

$C_1(t)$ is an exponentially growing function of t. We consider the situation where the tachyon mode condenses here, therefore, we take $C_2 = 0$ because C_2 is the coefficient of the massive mode. D2-brane charge is rewritten as

$$\frac{1}{6}q_1 q_2 + \left(\frac{2}{3}(C_1 q_2 x_2)^2 - \frac{1}{3}C_1^2 q_2\right) e^{-2q_1 x_1^2 - 2q_2 x_2^2}\,. \tag{2.36}$$

The first term is background D2 charge. The second term is induced by the tachyon mode and localized near the intersection point. Total D2-brane charge is obtained by the integration of x_i as

$$\int dx_1 dx_2 e^{-2q_1x_1^2-2q_2x_2^2}\left(\frac{2}{3}(C_2q_1x_1 - C_1q_2x_2)^2 - \frac{1}{3}C_2^2q_1 - \frac{1}{3}C_1^2q_2\right) = 0\,. \tag{2.37}$$

Thus, we confirm that D2-brane charge is still conserved after tachyon mode has condensed.

3. Supersymmetry

In this section, we consider the case of equal intersection angles $q_1 = q_2 \equiv q$ here. Now the mode (2.15) becomes massless and the corresponding eigenfunctions are written as

$$\begin{aligned}
\vec{V}'_{10} &= \begin{pmatrix} \tilde{f}_{10}(x_1,x_2,t) \\ \tilde{g}_{10}(x_1,x_2,t) \end{pmatrix} = e^{-qr^2}C_{10}\begin{pmatrix} 1 \\ 1 \end{pmatrix}, \\
\vec{V}'_{20} &= \begin{pmatrix} \tilde{f}_{20}(x_1,x_2,t) \\ \tilde{g}_{20}(x_1,x_2,t) \end{pmatrix} = e^{-qr^2}C_{20}\begin{pmatrix} 1 \\ 1 \end{pmatrix},
\end{aligned} \tag{3.1}$$

where $r^2 \equiv x_1^2 + x_2^2$. C_{10} and C_{20} are some numerical constant. It is known that two intersecting D-branes at equal two angles preserve 1/4 of the supersymmetries.[8] Therefore, let us check the supersymmetry of the configuration written as

$$\begin{aligned}
Y^9 &= q_1x_1\sigma_3 + \tilde{f}_{10}\sigma_1\,, \qquad & Y^8 &= q_2x_2\sigma_3 + \tilde{f}_{20}\sigma_1\,, \\
A_1 &= \tilde{g}_{10}\sigma_2\,, & A_2 &= \tilde{g}_{20}\sigma_2\,.
\end{aligned} \tag{3.2}$$

The supersymmetric variation of gaugino is written as

$$\delta\psi = F_{\mu\nu}\Gamma^{\mu\nu}\epsilon\,. \tag{3.3}$$

The configuration (3.2) satisfies the following BPS conditions up to the quadratic order in the fluctuations as

$$\begin{aligned}
&F_{12} + i[Y^8, Y^9] = 0\,, \\
&D_1Y^8 + D_2Y^9 = \sigma^3\mathcal{O}(C^2)\,, \\
&D_1Y^9 - D_2Y^8 = \sigma^3\mathcal{O}(C^2)\,,
\end{aligned} \tag{3.4}$$

and therefore, we obtain

$$\delta\psi = \mathcal{O}(C^2)\epsilon\,. \tag{3.5}$$

Thus, this configuration is supersymmetric up to the quadratic order in the fluctuations. The configuration with the supersymmetry up to quadratic order and where all supersymmetries are broken beyond this order is studied in [28].

[8] Black hole entropy in this system is studied in [27].

We can diagonalize Y^8 and Y^9 simultaneously in the region $x_1 \sim x_2 \sim 0$ as

$$Y^9 \sim \tilde{f}_{10}\sigma_3 , \qquad Y^8 \sim \tilde{f}_{20}\sigma_3 . \tag{3.6}$$

Thus, the intersection point is resolved, which was also seen in the previous section. Note that in the previous section, only one mode, which is tachyonic, condenses. On the other hand, we consider the deformation by two massless modes here. Finally, we calculate the distribution of the D2-brane charge. D2-brane charge is written in (2.35) as

$$J_{089} = \frac{1}{6}q^2 + \left(\frac{2}{3}(C_2 x_1 - C_1 x_2)^2 q^2 - \frac{1}{3}C_2^2 q - \frac{1}{3}C_1^2 q\right) e^{-2qr^2} . \tag{3.7}$$

Total charge is obtained as

$$\int dx_1 dx_2 \left(\frac{1}{6}q^2 + \left(\frac{2}{3}(C_2 x_1 - C_1 x_2)^2 q^2 - \frac{1}{3}C_2^2 q - \frac{1}{3}C_1^2 q\right) e^{-2qr^2}\right) = \frac{q^2}{6} \cdot (\text{area of } x_1\text{-}x_2 \text{ plane}) . \tag{3.8}$$

Thus, total D2-brane charge is again conserved.

4. Higher order corrections of F

We consider the effect of the higher order corrections of F. The F^4 terms are the first nontrivial contribution to the fluctuation analysis, therefore, we consider the mass spectra and eigenfunctions including F^4 terms. The symmetrized traced lagrangian in [29] is written as

$$L = \mathrm{Str}\sqrt{-\det(\eta_{\mu\nu} + 2\pi\alpha' F_{\mu\nu})} , \tag{4.1}$$

and F^4 terms are obtained by the expansion of this action as

$$L = \mathrm{Str}(2\pi\alpha')^2 \left(\frac{1}{8}F_{\mu\nu}F_{\nu\lambda}F_{\lambda\sigma}F_{\sigma\mu} - \frac{1}{32}F_{\mu\nu}F_{\nu\mu}F_{\lambda\sigma}F_{\sigma\lambda}\right) . \tag{4.2}$$

By taking the T-duality along $8, 9$ directions and considering the classical solutions (2.4) and fluctuations (2.5), we obtain quadratic parts in the fluctuations as

$$\begin{aligned}
L = & -Q_1\left((\partial_0 f_1)^2 + (\partial_0 g_1)^2\right) - Q_2\left((\partial_0 f_2)^2 + (\partial_0 g_2)^2\right) + \\
& + Q_3\left((\partial_1 f_1 + 2g_1 q_1 x_1)^2 - 4q_1 f_1 g_1\right) + Q_4\left((\partial_2 f_2 + 2g_2 q_2 x_2)^2 - 4q_2 f_2 g_2\right) + \\
& + Q_5\left((\partial_1 f_2 + 2g_1 q_2 x_2)^2 + (\partial_2 f_1 + 2g_2 q_1 x_1)^2 + 4(q_1 x_1 f_2 - q_2 x_2 f_1)^2 + (\partial_1 g_2 - \partial_2 g_1)^2\right) ,
\end{aligned}$$

$$\begin{aligned}
Q_1 &\equiv \left(1 - \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right)\left(1 + \frac{1}{6}\frac{q_2^2}{(2\pi\alpha')^2}\right), \qquad Q_2 \equiv \left(1 + \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right)\left(1 - \frac{1}{6}\frac{q_2^2}{(2\pi\alpha')^2}\right), \\
Q_3 &\equiv \left(1 - \frac{1}{2}\frac{q_1^2}{(2\pi\alpha')^2}\right)\left(1 + \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right), \qquad Q_4 \equiv \left(1 + \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right)\left(1 - \frac{1}{2}\frac{q_1^2}{(2\pi\alpha')^2}\right), \\
Q_5 &\equiv \left(1 - \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right)\left(1 - \frac{1}{6}\frac{q_1^2}{(2\pi\alpha')^2}\right).
\end{aligned} \tag{4.3}$$

By solving the equations of motions for the fluctuations, we obtain the eigenfunctions of the lowest mode as

$$\begin{aligned}\vec{V}'_{10} &= \begin{pmatrix}\tilde{f}_{10}(x_1,x_2,t)\\ \tilde{g}_{10}(x_1,x_2,t)\end{pmatrix} = n_{10}e^{-q_1x_1^2-q_2x_2^2}C_{10}(t)\begin{pmatrix}1\\1\end{pmatrix},\\ \vec{V}'_{20} &= \begin{pmatrix}\tilde{f}_{20}(x_1,x_2,t)\\ \tilde{g}_{20}(x_1,x_2,t)\end{pmatrix} = n_{20}e^{-q_1x_1^2-q_2x_2^2}C_{20}(t)\begin{pmatrix}1\\1\end{pmatrix}.\end{aligned} \tag{4.4}$$

The squared of the mass eigenvalue is obtained as

$$m_{i0}^2 = \pm(2(q_1-q_2) - \frac{2}{3(2\pi\alpha')^2}(q_1^3-q_2^3)) = \pm\frac{(\theta_1-\theta_2)}{2\pi\alpha'} + \mathcal{O}(\theta_i^5)\,. \tag{4.5}$$

We obtain the correct mass eigenvalue of tachyon mode at order θ^3. The eigenfunction of tachyon mode is the product of gaussian functions of x_1 and x_2. In the analysis [16], the eigenfunctions of tachyon mode remain gaussian function at order F^6, therefore we expect that the eigenfunctions remain the product of gaussian functions of x_1 and x_2 even in two angle's case in the action including the F^6 terms, and further, more higher order F terms.

5. Conclusion and discussion

We have considered the recombinations of D-branes intersecting at more than one angle using SU(2) super Yang-Mills theory. There are tachyon modes in the off-diagonal fluctuations and the condensation of this mode triggers the recombinations. On the other hand, in supersymmetric intersecting brane systems, there are two kinds of nontrivial massless deformations, the deformations into calibration geometry and the recombinations by the condensation of the off-diagonal fluctuation modes. The former is obtained by solving a minimal surface problem of membranes and described by U(1) DBI actions. On the other hand, to describe the latter phenomenon, we need to know the full knowledge of the non-abelian Born-Infeld action. In particular configurations, we can analyze this phenomenon even in Yang-Mills theory. We have studied the condensation of the tachyon mode and we have found that there are D2-branes distributed in the bulk near the intersection point after tachyon mode had condensed. Tachyon condensation diffuses D2-brane charge at the intersection point at first stage, and after that, the localization might be relaxed and recombined D2-branes, which preserve 1/4 of the supersymmetries, would emerge as a final state. In the case that two intersection angles are equal, the lowest off-diagonal fluctuation mode is massless. The configuration deformed by this mode is supersymmetric up to the quadratic fluctuations, and the intersection point is resolved. This is an interesting example of non-abelian embeddings. In this case, abelian mode which governs the calibration geometry and non-abelian mode which governs the non-abelian embeddings. The region that Yang-Mills analysis is appropriate is also up to the quadratic order in the fluctuations. In string theory, or NBI action, this supersymmetric non-abelian embedding might be valid beyond the quadratic order in the fluctuations. To discuss this point more deeply, it might be interesting to consider the higher order fluctuations. F^4 corrections are also studied

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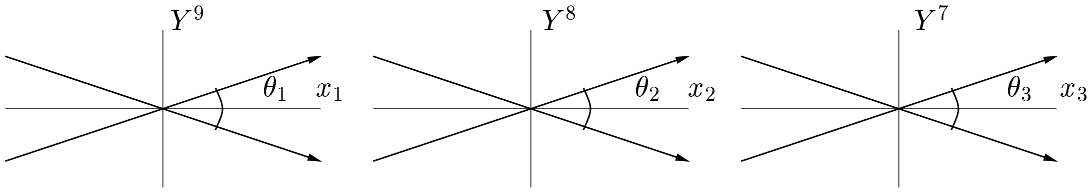

Figure 2: Intersecting D3-branes.

and the higher order corrections of θ are obtained. It is straightforward to discuss the recombination of Dp-branes intersecting at more angles. The extension to three intersection angles is studied in appendix A. In future direction, it might be interesting to study the Higgs mechanism of the Standard model from this direction.

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A. Recombination of D3-branes intersecting at three angles

D3-branes effective action is written as

$$S = -T \ \mathrm{Tr} \int d^2x dt \left[(D_a Y^i)^2 + \frac{1}{2} F_{ab}^2 - [Y^i, Y^j]^2 \right] . \tag{A.1}$$

where the indices $a, b = 0, 1, 2, 3$ and $i, j = 4, \ldots, 9$. Two D3-branes are embedded in 6 dimensions and do not extend in other dimensions. The embedded directions are chosen as $1, 2, 3, 7, 8$ and 9 and others as $4, 5$ and 6. A classical solution we consider here is written as

$$Y^9 = q_1 x_1 \sigma^3 , \qquad Y^8 = q_2 x_2 \sigma^3 , \qquad Y^7 = q_3 x_3 \sigma^3 , \qquad A_a = 0 , \tag{A.2}$$

which describes the two intersecting D3-branes.

Let us turn on the off-diagonal fluctuations as

$$\begin{aligned} Y^9 &= q_1 x_1 \sigma^3 + f_1(x_a)\sigma^1 - \bar{f}_1(x_a)\sigma^2 , & Y^8 &= q_2 x_2 \sigma^3 + f_2(x_a)\sigma^1 - \bar{f}_2(x_a)\sigma^2 , \\ A_1 &= g_1(x_a)\sigma^1 - \bar{g}_1(x_a)\sigma^2 , & A_2 &= g_2(x_a)\sigma^1 - \bar{g}_2(x_a)\sigma^2 . \end{aligned}$$

The lagrangian quadratic in the fluctuations is calculated as

$$\begin{aligned} L = \sum_{i,j=1,2,3} \Big(&- (\partial_0 f_i)^2 - 4 q_i f_i \bar{g}_i + (\partial_i f_j + 2 \bar{g}_i q_j x_j)^2 - (\partial_0 \bar{g}_i)^2 + (\partial_i \bar{g}_j - \partial_j \bar{g}_i)^2 + \\ &+ 4 (q_i x_i f_j - q_j x_j f_i)^2 + (f_\alpha \to \bar{f}, \bar{g}_\alpha \to g) \Big) . \end{aligned} \tag{A.4}$$

The combinations $(f_\alpha, \bar{g}_\alpha)$ and $(\bar{f}_\alpha, g_\alpha)$ are decoupled each other in the quadratic fluctuations, therefore we neglect $(\bar{f}_\alpha, g_\alpha)$. From now on, we denote $\bar{g}$ as g.

By solving the equations of motion for the fluctuations, we obtain the eigenfunction of the lowest mode of the fluctuations as

$$\begin{aligned}
\begin{pmatrix} \tilde{f}_{10}(x_1,x_2,t) \\ \tilde{g}_{10}(x_1,x_2,t) \end{pmatrix} &= \sum_{i=1,2,3} n_1 e^{-q_i x_i^2} C_{10}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix}, \\
\begin{pmatrix} \tilde{f}_{20}(x_1,x_2,t) \\ \tilde{g}_{20}(x_1,x_2,t) \end{pmatrix} &= \sum_{i=1,2,3} n_2 e^{-q_i x_i^2} C_{20}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix}, \\
\begin{pmatrix} \tilde{f}_{30}(x_1,x_2,t) \\ \tilde{g}_{30}(x_1,x_2,t) \end{pmatrix} &= \sum_{i=1,2,3} n_3 e^{-q_i x_i^2} C_{30}(t) \begin{pmatrix} 1 \\ 1 \end{pmatrix}. \qquad (A.5)
\end{aligned}$$

The lowest mode remains the product of gaussian functions of x_i. The squared of the mass eigenvalue is obtained as[9]

$$\begin{aligned}
m_{10}^2 &= \pm 2(-q_1 + q_2 + q_3) \sim \frac{1}{2\pi\alpha'}(-\theta_1 + \theta_2 + \theta_3)\,, \\
m_{20}^2 &= \pm 2(q_1 - q_2 + q_3) \sim \frac{1}{2\pi\alpha'}(\theta_1 - \theta_2 + \theta_3)\,, \\
m_{30}^2 &= \pm 2(q_1 + q_2 - q_3) \sim \frac{1}{2\pi\alpha'}(\theta_1 + \theta_2 - \theta_3)\,. \qquad (A.6)
\end{aligned}$$

This result is consistent with the worldsheet analysis in [19]. The parameter region of the intersection angles which corresponds to the tachyonic configuration is considered in [32].

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