Strings as Flux Tube and Deconfinement on Branes in Gauge Theories

Hikaru Kawai and Tsunehide Kuroki

Department of Physics, Kyoto University, Kyoto 606-8502, Japan

Abstract

We propose gauge theories in which the unstable branes and the fundamental string are realized as classical solutions. While the former are represented by domain wall like configurations of a scalar field coupled to the gauge field, the latter is by a confined flux tube in the bulk. It is shown that the confined flux tube is really a source of the bulk $B$-field. Our model also provides a natural scenario of the confinement on the brane in the context of the open string tachyon condensation. It is also argued that the fundamental string can be realized as a classical solution in a certain IIB matrix model as in our model.

* e-mail address : hkawai@gauge.scphys.kyoto-u.ac.jp
† e-mail address : kuroki@gauge.scphys.kyoto-u.ac.jp
1 Introduction

It has been believed that string theory can be nonperturbatively defined in terms of a gauge theory. In fact, candidates [1, 2, 3] for the nonperturbative definition of string theory are formulated as gauge theories in lower dimensions and some evidences have been found that these models really contain fundamental strings [4, 5, 6]. Unfortunately, however, we still do not know much about how to extract the fundamental string degrees of freedom explicitly from these gauge theories. From this point of view, it must be important to construct classical solutions corresponding to fundamental strings in gauge theories.

On the other hand, physics of the (open string) tachyon condensation draws much attention recently. There, it is conjectured that after the tachyon condensation, unstable D-branes disappear and end up with a bunch of closed strings. If a gauge theory is supposed to be a nonperturbative definition of string theory, it is desirable to describe such phenomena in terms of a gauge theory.

In light of these situations, it is useful to construct a gauge theory which realizes a fundamental string and an unstable brane as classical configurations. This is exactly the aim of this paper. The idea is the following: let us consider a gauge theory coupled to a scalar field which has an unstable domain wall solution — “brane”. Suppose in the bulk, the gauge theory is in the confinement phase, while on the brane, it is in the Coulomb (deconfinement) phase. Then we have the standard Abelian gauge theory on the brane and the confined flux tube in the bulk plays a role of a fundamental string. The flux tube attached to the brane will provide a deconfined flux on the brane. As the unstable brane decays, this flux tends to be confined as in the bulk, and finally when the brane disappears, a single confined flux tube will be recovered [7]. From this point of view, “confinement on the brane” [8] is automatically realized: it directly follows from the confinement in the bulk.

The organization of this paper is as follows: in the next section, we consider one of such models in which the Abelian gauge field couples to a scalar field through the dielectric ‘constant’ [9]. It turns out that if we choose the potential and the dielectric constant appropriately, our model indeed has the desired properties described above. Moreover, it is shown that the confined flux tube correctly becomes a source for the $B$-field in the bulk. This implies that the flux tube really represents a fundamental string. A relation to the tachyon condensation in string theory is also addressed. In section 3, we propose other gauge theories which have the same properties. In fact, for each confinement mechanism, it is possible to
construct a gauge theory with the desired properties. This suggests that confinement in the bulk and deconfinement on the brane is a quite universal phenomenon. Section 4 is devoted to discussions in which we put an emphasis on a possible relation to a fundamental string solution in a kind of IIB matrix model.

2 Confinement via the Dielectric Effect

2.1 The Model

In this section, we consider the case in which the bulk electric flux is confined by the dielectric effect proposed in [9]. Let us start with the Abelian gauge theory coupled to a scalar field \( \phi \) in \((d + 1)\) dimensions considered in [4]

\[
\mathcal{L} = -\frac{1}{4} \varepsilon(\phi) F_{\mu\nu}^2 + \frac{1}{2} \partial_\mu \phi \partial_\mu \phi - V(\phi) - j^\mu A_\mu + i \bar{\psi} \gamma^\mu \partial_\mu \psi,
\]

where

\[
j^\mu = g \bar{\psi} \gamma^\mu \psi,
\]

is the fermion current. Gauss’ law tells us that

\[
\partial_\mu D^{\mu\nu} = j^\nu,
\]

where

\[
D_{\mu\nu} = \varepsilon F_{\mu\nu}.
\]

The canonical momentum for \( A_\mu \) is given by

\[
\pi_A^\mu = -D^{0\mu} \equiv D^\mu.
\]

In the following, we choose \( A_0 = 0 \) gauge. The Hamiltonian density is given by

\[
\mathcal{H} = \frac{D^2}{2\varepsilon} + \mathcal{H}_\phi,
\]

where we have assumed that there is no magnetic source and dropped the trivial fermion part.

Let us assume that \( D \) has a flux tube configuration in the \( x^1 \) direction. From Gauss’ law (2.3), we obtain

\[
f = \int dS \, D_1.
\]
where \( S \) is a \((d - 1)\)-dimensional hyperplane orthogonal to the \( x^1\)-direction, and \( f \) is the flux produced by the charge. Then we should examine whether there is a minimum of the tension (energy of the flux tube per unit length) under the constraint that a given total flux passes through the tube. Therefore, we minimize

\[
T = \int dS \frac{D_1^2}{2\varepsilon} - \lambda \left( \int dS \, D_1 - f \right),
\]

where \( \lambda \) is the Lagrange multiplier implementing the constraint (2.7). It is easy to find that the solution satisfies

\[
D_1 = \lambda \varepsilon.
\]

From this solution, we find that the electric field \( F_{01} \) is simply given by \( \lambda \) and hence is constant. Nevertheless \( D_1 \) can be nontrivial due to the nontrivial dependence of the dielectric effect \( \varepsilon(\phi(x)) \). Note that there is no flux in the region where \( \varepsilon \) is zero as seen from (2.9).

Solving the constraint by (2.9), the energy per unit length can be rewritten as

\[
W = \frac{1}{2} \int \varepsilon \, dS + \int dS \left( \frac{1}{2} (\vec{\nabla} \phi)^2 + \dot{\phi}^2 + V(\phi) \right),
\]

where we have included the contribution from the scalar field. As shown in [9], it is now easy to see that there exits a static configuration of \( \phi \) which minimize \( W \) under a suitable choice for \( \varepsilon \) and \( V \). In [9], they are given by

\[
\varepsilon = \left( \frac{\phi - \phi_0}{\phi_0} \right)^4,
\]

\[
V(\phi) \sim \mu (\phi - \phi_0)^2.
\]

Thus, a nontrivial dependence of \( \varepsilon \) on \( \phi \) allows a flux tube solution for \( D \) which is energetically more favored than the spherically symmetric configuration. Therefore, if we put an electric charge in the bulk, it produces a flux tube which yields a linear potential between charges. This shows that this theory in the bulk is in the confinement phase.

In our model, a new ingredient comes into the choice of the scalar potential \( V(\phi) \). As an example, the potential is chosen to be

\[
V(\phi) = -\phi^2 \log \phi^2,
\]

rather than (2.12) so that it can admit the unstable domain wall solution. In fact, this potential is studied in [10] as a toy model of the tachyon condensation and is known to be
the exact tachyon potential \[11\] in the context of the boundary string field theory (BSFT) \[12\]. As shown in \[10\], the unstable domain wall in this case is simply given as the Gaussian:

\[
\phi(x) = \exp(-x^2/4),
\]  
(2.14)

where we have denoted the one-dimensional transverse coordinate of the domain wall (brane) as \(x\). For illustration, let us take \(\varepsilon\) as

\[
\varepsilon(\phi) = \phi^4,
\]  
(2.15)

which is essentially the same as (2.11) up to an irrelevant constant. In this case, \(\varepsilon(\phi)\) is zero in the bulk except in the flux tube produced by charges put there as described above. On the other hand, since \(\varepsilon\) is non-zero all over the unstable \((d-1)\)-brane, \(D_1\) can be non-zero there. Therefore, it is natural to expect that the standard Abelian gauge theory is recovered on the brane. In the next subsection, this claim is confirmed by examining the existence of the massless mode of the gauge field on the brane.

### 2.2 Massless Mode on the Brane

Let us examine the existence of the massless mode of the gauge field on the brane. The equation of motion for the gauge field far from a charge is

\[
\partial_\mu (\varepsilon F^{\mu\nu}) = 0.
\]  
(2.16)

It is rather trivial that the \(x\)-independent mode of \(A_\mu\) actually satisfies this equation as the massless mode. However, for completeness, let us solve this equation. Taking the \(A_0 = 0\) gauge and assuming the plane wave solution for \(A_x\) and \(A_i\) with respect to \((t, x^i)\), where \(x^i\) represents the longitudinal direction of the brane, we find that \(\varepsilon A_x\) satisfies the standard wave equation and that the equation of motion for \(A_j\) is reduced to

\[
\partial^2_x A_i + \frac{\partial_x \varepsilon}{\varepsilon} \partial_x A_i + k_x^2 A_i = 0,
\]  
(2.17)

where \(k_x\) is the momentum in the transverse direction of \(A_x\). Substituting (2.14) and (2.15) into this equation, we see that this equation is nothing but the Hermite differential equation and that the eigenvalue is given by

\[
k_x^2 = 2n, \quad n = 0, 1, 2, \ldots.
\]  
(2.18)

Thus the gauge field \(A_i\) on the brane really has the massless mode and there is no tachyon even though the brane itself is unstable. The latter fact can be also seen directly from the equation of motion (2.16).
2.3 Coupling with the $B$-field

In this subsection, we consider the effect of the bulk $B$-field. If the flux tube is really a fundamental string, it must be a source of the $B$-field in the bulk. For the purpose of including the $B$-field into our Lagrangian, let us begin with the dual gauge theory in four dimension in which a magnetic flux is coupled to the bulk $B$-field:

$$L' = -\frac{1}{4} \epsilon^{-1} \tilde{F}_{\mu\nu}^2 - \frac{i}{2} \epsilon^{\mu\nu\lambda\rho} \tilde{F}_{\mu\nu} B_{\lambda\rho}, \quad (2.19)$$

where $\tilde{F}$ is the dual field strength and, correspondingly, the dielectric constant is inverted. Performing the duality transformation of this Lagrangian, we arrive at

$$L = -\frac{1}{4} \epsilon (F_{\mu\nu} - B_{\mu\nu})^2, \quad (2.20)$$

which seems quite natural. Therefore, we expect that in general the coupling with $B$-field is introduced by replacing $F_{\mu\nu}$ with $F_{\mu\nu} - B_{\mu\nu}$. Regarding the $B$-field as a background for the gauge field and repeating the same argument as above, we obtain the Hamiltonian density as follows:

$$H = \frac{D^2}{2\epsilon} + D^\mu B_{0\mu}. \quad (2.21)$$

This shows that our flux tube correctly couples to the $B$-field. It is also easy to verify that $F_{0\mu}$ is again given by $\lambda$ and hence constant for the flux tube configuration. Notice that although $F_{0\mu}$ is constant for the flux tube configuration, $B$ is not necessarily constant. These facts suggest that the confined electric flux tube is a classical configuration in the gauge theory corresponding to the fundamental string. Therefore, if we omit the scalar field part, the complete Lagrangian must be

$$L = -\frac{1}{4} \epsilon (F_{\mu\nu} - B_{\mu\nu})^2 + cH_{\mu\nu\rho}^2, \quad (2.22)$$

where $H$ is the field strength of the $B$-field, and we have not taken account of the effect by gravity.

It is worth pointing out that the statements in this subsection are not restricted to the action of $F_{\mu\nu}^2$ type. In fact, if we start from the ‘dual’ Born-Infeld action

$$L' = -\sqrt{\det(1 + \tilde{F})} - \frac{i}{2} \epsilon^{\mu\nu\lambda\rho} \tilde{F}_{\mu\nu} B_{\lambda\rho}, \quad (2.23)$$

and performing the duality transformation developed in [13], we obtain

$$L = -\sqrt{\det(1 + F - B)}, \quad (2.24)$$
namely, the inclusion of $B$ amounts to making a replacement $F \rightarrow F - B$ as well in the Born-Infeld action. Furthermore, it can be shown that for a general Lagrangian $L((F_{\mu \nu} - B_{\mu \nu})^2)$, the electric field $F_{0\mu}$ becomes constant as above for the flux tube configuration.

2.4 Stability of the Flux Tube

In this section, we examine the stability of the flux tube against expanding it. For this purpose, let us examine how the total energy for the static flux tube

$$W = \frac{1}{2} \left( \int \varepsilon \, dS \right)^2 + \int dS \left( \frac{1}{2} \vec{\nabla} \phi \right)^2 + V(\phi),$$

changes under a transformation

$$\phi(x^i) \rightarrow \lambda \phi(x^i),$$

where $x^i$ are the $d-1$ directions in which a section of $D$ has a support: $\int d^{d-1}x D = \int dS D = f$. This transformation for $\lambda < 1$ corresponds to fattening the flux tube. More precisely, let us consider the transformation

$$\phi(x^i) \rightarrow \phi'(x^i) = \lambda^\alpha \phi(\lambda x^i),$$

where we choose $\alpha$ such that

$$\int \varepsilon(\phi') \, dS = \int \varepsilon(\phi) \, dS,$$

namely, the above transformation decreases the range of $\phi$ simultaneously so that the integration of the dielectric term will keep invariant. Suppose $\varepsilon(\phi) \sim \phi^m$, then a simple calculation shows that $\alpha = (d - 1)/m$. If the potential behaves like $V(\phi) \sim \phi^n$, the kinetic term and potential term for the scalar field in (2.25) become under this transformation

$$\int \frac{1}{2} (\vec{\nabla} \phi')^2 dS = \lambda^{3-d+2\alpha} \int \frac{1}{2} (\vec{\nabla} \phi)^2 dS,$$

$$\int V(\phi') \, dS = \lambda^{1-d+\alpha n} \int V(\phi) \, dS,$$

where

$$3 - d + 2\alpha = 3 - d + \frac{2(d - 1)}{m},$$

$$1 - d + \alpha n = (d - 1) \left( \frac{n}{m} - 1 \right).$$

Therefore, as long as $m$ is larger than the smaller value of $n$ and $2(d - 1)/(d - 3)$, expanding the soliton costs more energy of the gauge field. It is now evident that our choice (2.13), (2.15) agrees with this condition. For $\lambda > 1$, higher derivative terms possibly stabilize the flux tube. This establishes the stability of the fundamental string in the gauge theory (2.1).
2.5 Relation to the Tachyon Condensation

It is conjectured that when the open string tachyon around the unstable D-brane condenses, the gauge field on the brane forms a confined flux tube \([8]\) which plays a role of a piece of a closed string \([7]\). It is further pointed out that in this case the tachyon potential plays a role of the dielectric constant \([14, 15]\). In order to confirm this, let us apply the previous arguments to the Lagrangian

\[
L = -\frac{V(\phi)}{g_s} \sqrt{-\det(g + F)} - \frac{1}{g_s} \sqrt{-\det(g + F)} G_{\mu\nu}^S \partial_\mu \phi \partial_\nu \phi, \quad (2.31)
\]

where \(\phi\) is the tachyon field, \(G_{\mu\nu}^S\) is the symmetric part of \((g_{\mu\nu} + F_{\mu\nu})^{-1}\), and \(V(\phi)\) is given as in (2.13). This Lagrangian is derived by using the boundary string field theory \([11]\).

As before, we assume the flux tube solution along the \(x^1\)-direction and concentrate only on \(F_{01} = E_1\). The Gauss’ law constraint (2.7) can be solved by

\[
D_1 = f \frac{\tilde{V}}{\tilde{V}dS}, \quad (2.32)
\]

and the minimized energy per unit length is given as

\[
W = \sqrt{f^2 + (\int \tilde{V}dS)^2} + \frac{\sqrt{f^2 + (\int \tilde{V}dS)^2}}{\tilde{V}dS} \int (\partial_\mu \phi)^2 dS, \quad (2.33)
\]

where \(\tilde{V} = V/g_s\). Around the tachyonic vacuum\([3]\), \(\tilde{V} << 1\) and

\[
W = f + \frac{(f \tilde{V}dS)^2}{2f} + \frac{f}{\tilde{V}dS} \int (\partial_\mu \phi)^2 dS. \quad (2.34)
\]

As seen from (2.32), the flux tube can exist only in the region where \(\tilde{V} \neq 0\). In this sense, \(\tilde{V}\) indeed plays a similar role to the dielectric constant near the tachyonic vacuum. Note that even if \(\tilde{V} << 1\), \(D_1\) can remain finite according to (2.32). Thus we expect that the flux tube configuration satisfies both the minimum energy condition and the Gauss’ law constraint.

Moreover, as mentioned in the previous subsection, this flux tube correctly couples to the \(B\)-field. These facts seem to suggest that the flux tube is exactly the closed string at the tachyonic vacuum where the flux tube is confined via the dielectric effect caused by the tachyon potential. However, the expression of \(W\) implies that the flux tube in this case is unstable under the ‘fattening’ transformation described in the previous subsection. In fact, it is easy to see that if we make the transformation (2.27) satisfying (2.28) with \(\varepsilon\) replaced

\(^3\)The phrase tachyonic vacuum refers to the vacuum after the tachyon condensates.
by $\tilde{V}$, the first and second term in (2.34) are invariant, while the last kinetic term decreases. Therefore, the configuration of $\phi$ spreads and eventually becomes flat. The flux tube is unstable. Since the tachyon potential (2.13) is known to be exact [11], this fact suggests that the kinetic term should be modified if the flux tube really plays a role of the fundamental string at the tachyonic vacuum. Indeed, compared to the potential term, there is no good reason yet why the kinetic term can be still represented in terms of the open string metric as in (2.31) even near the tachyonic vacuum. In fact, the trouble in (2.34) originates from the fact that $\tilde{V}$ plays both roles of the potential and of the dielectric constant. Thus one of the resolutions of this problem may be a modification in the kinetic term in (2.31).

3 Other Models

In this section, we construct other gauge theories in which the fundamental string and the branes are realized as classical configurations. The string is described by a confined flux tube solution in the bulk and is deconfined on the branes. As origins of confinement other than the dielectric effect, we employ the non-Abelian gauge field and the vortex line. It turns out that for each confinement mechanism, it is possible to construct a gauge theory with this property.

3.1 Confinement via the Non-Abelian Gauge Field

Let us construct a four-dimensional non-Abelian gauge theory in which a confined flux tube in the bulk becomes deconfined on a domain wall of the scalar field coupled to the gauge field. For example, the Lagrangian is given by

$$
\mathcal{L} = -\frac{1}{4g^2} F_{\mu \nu}^a F^{a\mu \nu} + |D_\mu \Phi|^2 - \left( \frac{\nu^2}{2} + \eta^2 \right) |\Phi|^2 - \frac{\kappa}{2} (|\Phi|^2)^2 + \frac{1}{2} (\partial_\mu \eta)^2 - \lambda (\eta^2 - \nu^2)^2,
$$

(3.1)

where $F_{\mu \nu}^a$ is the $SU(2)$ gauge field strength, $\Phi$ is a complex scalar field in the adjoint representation of $SU(2)$, and $\eta$ is a real (neutral) scalar field. This model has already been considered in [16]. As shown in [16], $\eta$ has a stable domain wall solution

$$
\eta_0 = v \tanh(\sqrt{2\lambda} \nu x).
$$

(3.2)

In the bulk away from the domain wall, $\eta = \pm \nu$. In this case the potential for $\Phi$ has a positive mass term. Thus the bulk gauge field is the standard non-Abelian one and is in the confinement phase. Therefore, the string can be realized as a confined flux tube. On the
other hand, at the core of the domain wall, \(-v/\sqrt{2} < \eta < v/\sqrt{2}\), \(\Phi\) has the negative mass term and consequently the \(SU(2)\) gauge symmetry is spontaneously broken down to \(U(1)\). Thus the flux is deconfined on the brane.

In this example, we have a stable brane (3.2). It is also possible to construct an unstable brane solution by changing the form of the potential for \(\eta\). For example, if we take an \(\eta^3\) potential, there exists a lump solution. Then it is possible to adjust parameters in such a way that \(-v/\sqrt{2} < \eta < v/\sqrt{2}\) is satisfied only inside the unstable brane.

### 3.2 Confinement via the Vortex Line

The second model is based on the Abelian Higgs model in which magnetic charges are confined by the vortex line proposed in [17]. The Lagrangian is

\[
\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + |D_\mu \Phi|^2 - (\frac{v^2}{2} - \eta^2)|\Phi|^2 - \frac{\kappa}{2} (|\Phi|^2)^2 + \frac{1}{2} (\partial_\mu \eta)^2 - \lambda (\eta^2 - v^2)^2. \tag{3.3}
\]

This model is different from the one in the last subsection in that \(F_{\mu\nu}\) is the Abelian gauge field and the sign in front of the mass term for \(|\Phi|\) is opposite. In this case in the bulk the magnetic charges are confined via the vortex line. This flux tube can be identified as a fundamental string in the bulk. Inside the brane the potential for \(\Phi\) is the stable one, hence we have the standard Abelian gauge field in the Coulomb phase.

In the dual picture, the electric flux is confined in the bulk and deconfined on the brane. This model is proposed in [18]. However, in this case the electric flux is dual to the magnetic one which is introduced by hand in order to cancel a singularity arising from a singular gauge transformation. In this sense, the electric flux tube is not a classical solution, but a kind of background.

We conclude this section by making a remark on the confinement in the bulk and deconfinement on the brane. In the context of the tachyon condensation in open string theory, it is most likely that the bulk confinement is realized by the dielectric effect as described by the previous section. However, a variety of models in this section which realize the same situation suggest that this is a quite universal phenomenon.
4 Discussions

In this section we discuss a relation between our model and the IIB matrix model. Motivated by (2.1), let us consider a variant of the IIB matrix model

\[ S = -\frac{1}{4} \text{Tr} \varepsilon(Y)[A_\mu, A_\nu]^2 - \frac{1}{2} \text{Tr} (\bar{\psi} \Gamma^\mu [A_\mu, \psi]) + V(Y), \]

where \( A_\mu \) and \( Y \) are bosonic \( N \times N \) Hermitian matrices, and \( \psi \) is a fermionic \( N \times N \) matrix. \( \varepsilon(Y) \) and \( V(Y) \) are assumed to be given as (2.13) and (2.13) respectively. We see in (4.1) that \( \varepsilon(Y) \) plays a role of the dielectric function and the dynamics of \( Y \) is governed by the potential \( V(Y) \). Of course the action is for the Yang-Mills field, but we may still expect that we have a confined Abelian flux tube in this model due to \( \varepsilon(Y) \). For example, as shown in [19], if we expand (4.1) around a following classical solution for \( A_\mu \):

\[ [\hat{A}_\mu, \hat{A}_\nu] = iB_{\mu\nu}, \]

then the matrix model becomes the noncommutative \( U(1) \) gauge theory with the dielectric function \( \varepsilon(Y) \). In this theory, \( \varepsilon(Y) \) is expected to confine a flux tube as well as in the commutative model considered in section 2. This confined tube should be also stable by the argument given in section 2.4. Thus we have a classical solution corresponding to a fundamental string in a kind of IIB matrix model. Electric and magnetic flux tube solutions in the noncommutative \( U(1) \) gauge theories have been also obtained in e.g. [14, 20], but the confinement problem has not been fully addressed.

This type of model with \( \varepsilon(Y) \sim Y^{-1} \) and \( V(Y) \sim Y \) has been proposed in [22] as a nonperturbative regularization of the Schild action [21] of type IIB superstring. In their formulation, \( Y \) is introduced to play the same role as \( \sqrt{g} \) in the Schild action. It is also pointed out in [22] that \( Y^{-1} \) can be regarded as the dielectric function proposed in [9]. In this sense, our model is quite similar to the one in [22] although the potentials are different because of the different motivations.

It must be emphasized that the above interpretation of a classical flux tube solution as the fundamental string in the IIB matrix model is conceptually different from the one in [3]. It would be interesting to clarify their direct relationship.

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