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Field Theoretic Branes and Tachyons of the QCD String

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Abstract

Dvali and Shifman have proposed a field-theoretic mechanism for localizing gauge fields to "branes" in higher dimensional spaces using confinement in a bulk gauge theory. The resulting objects have a number of qualitative features in common with string theory D-branes; they support a gauge field and flux strings can end on them. In this letter, we explore this analogy further, by considering what happens when N of these "branes" approach each other. Unlike in the case of D-branes, we find a *reduction* of the gauge symmetry as the "branes" overlap. This can be attributed to a tachyonic instability of the flux string stretching between the branes.

Recently, Dvali and Shifman have considered the possibility of trapping gauge fields on p-branes with $p < 3$ using confining dynamics in a bulk 3+1-dimensional gauge theory [1]. These field theoretic branes are very interesting as their higher dimensional generalizations can be used to construct extensions of the standard model with extra dimensions. In such models gravity and possibly some other fields propagate in higher dimensional space-time whereas the standard model matter and gauge forces are confined to (3+1) dimensional branes [2, 3]. Apart from these potential phenomenological applications field theoretic p-branes also provide a very interesting background for studying the interplay of dynamics in various dimensions. We are particularly intrigued by some obvious similarities between these branes and D-branes [4] in string theory: apart from supporting a gauge field in their world-volume, field theoretic branes also allow color flux-strings to end on them [5]. The aim of this letter is to further explore the analogy between these branes in field theory and in string theory. More concretely, after reviewing the construction of Dvali and Shifman and giving some simple generalizations, we consider what happens when N of these walls are brought on top of each other. From the analogy with D-branes one expects that modes of the QCD flux-string become light and contribute to the effective p dimensional world-volume field theory. In the case of D-branes the lightest modes are spin-1 fields (and their superpartners), and the $U(1)^N$ gauge symmetry of N widely separated D-branes gets enhanced to $U(N)$. In the case of the QCD string the masses of low-lying vibrational modes decrease as

$$m \propto L \, \Lambda^2 \tag{1}$$

when the length L of a long flux tube is reduced. However which modes of the QCD string become light for small L turns out to be different. We will argue from consistency of the $(2+1)$ dimensional low energy field theory on the branes that the lightest such mode is not a spin-1 field. Instead we find a scalar whose mass squared is positive for long stretched strings but becomes negative for very small $L \sim \Lambda^{-1}$ where eq.1 breaks down. Thus for very small separation the scalar condenses and *reduces* the gauge symmetry of the effective field theory on the branes. This understanding of the reduction of gauge symmetry as branes are brought in contact with each other from both a macroscopic effective (2+1)-d theory as well as from a microscopic (3+1)-d picture with QCD strings is our central result. As a bonus we also find that field theoretic branes can be connected by multi-pronged flux tubes corresponding to "baryonic" QCD strings.

To get started we first review the argument of Dvali and Shifman [1] and give a simple generalization before moving on to consider what happens when N of these walls are brought on top of each other.

In the construction of our walls we will frequently assume that a gauge group is broken in some region of space but not in others. We will also have use for matter fields which are very massive in the bulk but light on the walls. In the discussion we will assume that these effects have been arranged by coupling the theory to a "black box" containing appropriate very massive neutral and charged scalars with space dependent vacuum expectation values*.

For simplicity, we first attempt to localize a $U(1)$ gauge field to a region $\mathcal{W}$ in 3+1 dimensional space between $0 < z < l$, which on distances much larger than l would look like a 2+1 dimensional wall supporting a $U(1)$ gauge field.

The most obvious idea is to arrange for the $U(1)$ to be broken outside $\mathcal{W}$ giving the photon a mass $M >> l^{-1}$, but unbroken inside $\mathcal{W}$. Then, since the photon is massive outside $\mathcal{W}$ but massless inside, one may think that there is a massless electric photon in the $(2+1)$-d theory at long distances. This is not the case. To understand this, note that the region outside $\mathcal{W}$ is superconducting while the region inside is normal vacuum. Now place an electric test charge inside $\mathcal{W}$ and examine the field strength at another point in $\mathcal{W}$ a distance $r >> l$ away; if there is a massless photon in the long-distance theory, we should have a $(2+1)$-d Coulomb field in this regime.

Since the region outside is a conductor, however, the electric field lines emanating from the test charge must end on and be perpendicular to the boundary of $\mathcal{W}$, whereas in order to obtain a $2+1$ dimensional force law these field lines would have to be repelled from the boundary. We can solve for the electric field using the method of images, with an infinite number of image charges of alternating signs. Clearly all the multipole moments vanish for such a configuration, and we are left with an exponentially small field for $r >> l$. Therefore, we conclude that there are no fields lighter than the ultraviolet cutoff l^{-1} of the (2+1)-d theory, coupling to electric charge. It is very easy to see (as we show in detail in the appendix) that instead, there is

*An example of such a black box can be found in [1].

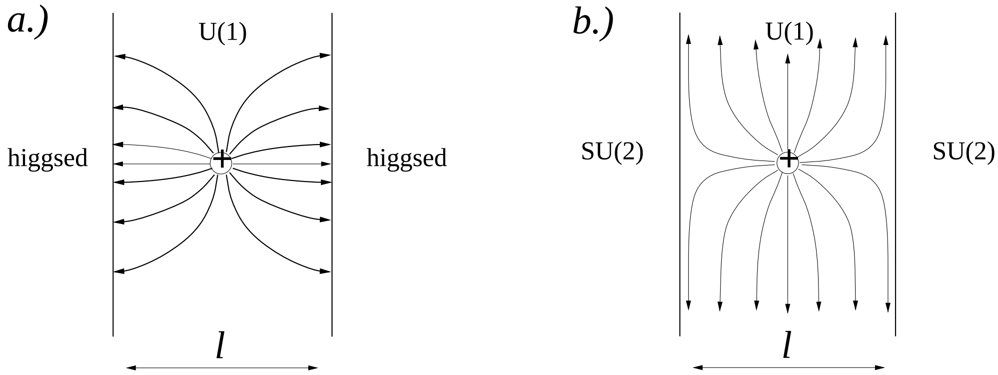

Figure 1: *The figure labeled a.) depicts a domain wall which has no massless electric photon trapped. This can be seen from this figure by noting that the electric field lines are screened by the superconducting Higgs vacuum in the bulk. In figure b.) the bulk is in a confined phase and repels electric flux. As a result, a massless photon coupling to electric charge is trapped.*

a tower of *massive* gauge fields with masses quantized in units of l^{-1}.

This failure suggests the correct way to proceed, however. Suppose that we instead place a magnetic charge g inside $\mathcal{W}$. Now, because of the Meissner effect in the superconducting region, all the magnetic flux lines are repelled from the boundaries and we recover the $(2+1)$-d magnetic Coulomb law. A trivial application of Gauss' law yields the relationship between the effective $(2+1)$-d magnetic charge g_3 and g:

$$\frac{1}{g_3^2} = \frac{1}{g^2} \times l \tag{2}$$

Of course, we actually want to localize electric photons on the wall, this can be accomplished by the 't Hooft-Mandelstam dual of this superconducting picture. Suppose that we begin with a $(3+1)-d$ $SU(2)$ gauge theory, which is broken to a $U(1)$ inside $\mathcal{W}$ by a very massive scalar in the adjoint representation of $SU(2)$. The bulk theory is confining at the scale Λ which we take to be $>> l^{-1}$, whereas the the $U(1)$ inside $\mathcal{W}$ is free. If we now place an electric test charge inside $\mathcal{W}$, confinement expels the electric field lines from the bulk due to the dual Meissner effect, and we recover the $(2+1)$-d Coulomb law for the electric field. This successfully localizes a $U(1)$ gauge field to a $(2+1)$-d wall in a $(3+1)$-d bulk.

There are obvious generalizations of this idea. Suppose we have an $SU(N_c)$ gauge theory with $N_F >> N_C$ flavors, which are given a very large mass outside $\mathcal{W}$ but are massless inside. Then the outside theory is asymptotically free and confines. The theory inside the region $\mathcal{W}$ is infrared free at distances short compared to the wall thickness l where the coupling evolves according to the $(3+1)$-d renormalization group equation. But at length scales long compared to l the theory on the wall is $(2+1)$-dimensional and the coupling evolves according to the $(2+1)$-d renormalization group equation. At the UV cutoff l^{-1} of the low energy theory the (2+1)-d gauge coupling is matched to the higher dimensional coupling as

$$g_3^2(\mu = l^{-1}) = \frac{g_4^2(\mu = l^{-1})}{l} \tag{3}$$

By the same argument as for the $U(1)$ case above, this localizes an $SU(N_c)$ gauge theory on the $(2+1)$ dimensional wall. Notice that unlike the $U(1)$ case, this (2+1)-d theory also confines; however the confinement scale is $\sim g_3^2$ which can be much smaller than the cutoff l^{-1} if g_4^2 is small. This is easy to arrange since the 3+1-d theory inside $\mathcal{W}$ can have a small gauge coupling at its UV cutoff and gets (logarithmically) weaker as it is scaled into the IR towards $\mu = l^{-1}$. Therefore, there is a range of energies $g_4^2/l < E < 1/l$ where we can have an unconfined (2+1)-d $SU(N_c)$ gauge theory. In this manner it is possible to engineer a large variety of field theoretic branes with different gauge theories living on them.

What happens if we move an electric charge from the wall into the confining bulk [5]? The confinement tries to expel the electric field lines, but since a net flux of electric field must be present at large distances by Gauss' law, a string of electric flux forms between the charge in the bulk and the wall as in figure 2. Thus, these walls have a second qualitative feature in common with D-branes: strings can end on them †.

†Note that our electric flux strings ending on a wall of unconfined gauge field are the electric-magnetic dual to cosmic strings (with their associated magnetic flux) ending on a domain wall of unbroken gauge field as described for example in [6]. The microscopic physics allowing strings to end on domain walls here is different from the physics allowing strings to end on domain walls in $N = 1$ supersymmetric QCD [7, 8]. For a recent discussion of domain walls in softly broken $N = 2$ SUSY gauge theories, see [9].

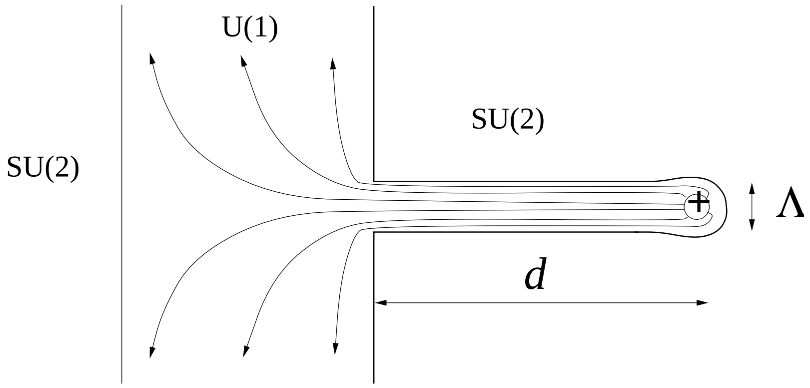

Figure 2: *A test charge which is moved a distance d off the wall drags a flux tube of thickness Λ behind it because electric charge is confined in the bulk.*

We now explore this analogy with D-branes further by considering what happens when we bring two or more of these walls close together. For concreteness, let us take a case with $SU(2)$ Higgsed to a $U(1)$ in two regions $\mathcal{W}_1 = -l_1 < z < 0$ and $\mathcal{W}_2 = d < z < d + l_2$. Let us first consider the case where the walls are very well separated $d >> l_1, l_2$. Then at distances longer than $l_{1,2}$ we have two $(2+1)$-d walls with two separate $U(1)'s$ localized on them. To see that there are really two $U(1)'s$, simply note that the electric field lines emanating from a charge on $\mathcal{W}_1$ can never end on a charge in $\mathcal{W}_2$ because of the confining region separating them. Let us further simplify our description by working in the effective theory at distances $>> d$, where the separation between the walls cannot be discerned. This is then a $(2+1)-$d theory with a $U(1) \times U(1)$ gauge group. For the case of N well-separated walls, this very long distance theory has a $U(1)^N$ gauge symmetry.

Next consider the opposite extreme when two walls are sitting very close to each other $d << 1/\Lambda < l_{1,2}$. At long distances this case is indistinguishable from having just one wall with thickness $(l_1 + l_2)$, and we only localize a single $U(1)$ gauge field in the very long distance theory. Therefore, as the walls are brought close together, the long-distance theory sees a *reduction* of the gauge symmetry from $U(1) \times U(1)$ to $U(1)$. This is opposite to the D-brane case, where the gauge symmetry gets enhanced from $U(1) \times U(1) \rightarrow U(2)$. Nevertheless, as we will see below, the physics of the two situations is very similar.

Let us first try to understand what is going on purely in the long-distance theory. As the parameter d in the theory is varied, we go from having a $U(1) \times U(1)$ symmetry for $d >> l_{1,2}$ to just a $U(1)$ symmetry for $d = 0$. The most plausible interpretation is that the $U(1) \times U(1)$ symmetry is Higgsed somewhere in the transition where $d \sim \Lambda$. Since neither of the walls is special, we expect that $U(1) \times U(1)$ must be broken to the diagonal $U(1)$. This satisfies an interesting consistency check. From the microscopic viewpoint, when the walls merge to give a new wall of thickness $(l_1 + l_2)$, the (2+1)-d coupling of the single $U(1)$ should be

$$\frac{1}{g_3^2} = \frac{(l_1 + l_2)}{g_4^2} \tag{4}$$

On the other hand, the gauge coupling determined by Higgsing $U(1) \times U(1)$ to the diagonal subgroup is

$$\frac{1}{g_{3diag}^2} = \frac{1}{g_{3,1}^2} + \frac{1}{g_{3,2}^2} = \frac{l_1}{g_4^2} + \frac{l_2}{g_4^2} \tag{5}$$

as required.

Therefore, purely from considerations of the very low-energy theory, we conclude that some new state becomes light when $d \sim \Lambda$, and acquires a condensate to spontaneously break $U(1)_1 \times U(1)_2 \to U(1)_{diag}$. The condensate must of course be a Lorentz scalar, and must be charged under both $U(1)'s$ to break to the diagonal subgroup. The simplest possibility is that as d is reduced and becomes smaller than $\sim \Lambda$ a scalar field $\phi^{+,-}$ of charge $(+,-)$ under $U(1)_1 \times U(1)_2$ becomes light and then tachyonic, triggering the non-zero condensate $\langle \phi^{+,-} \rangle$.

We stress that the existence of such a condensate was deduced by the requirement of a consistent low-energy effective theory. But we can easily identify a natural candidate for $\phi^{+,-}$ in the microscopic theory. For $d >> l$, there is a stable configuration corresponding to the QCD string stretching between the walls as shown in figure 3.

One can imagine forming this string as follows. Place very heavy test quarks $q, \bar{q}$ inside the confining medium between the walls; a QCD string of confined color electric flux will stretch between them. Now, move $q(\bar{q})$ until it is just inside region $\mathcal{W}_{1(2)}$. This will cost a great deal of energy $\sim \Lambda^2 d$, but the resulting string is stable: it can not break since there are no dynamical

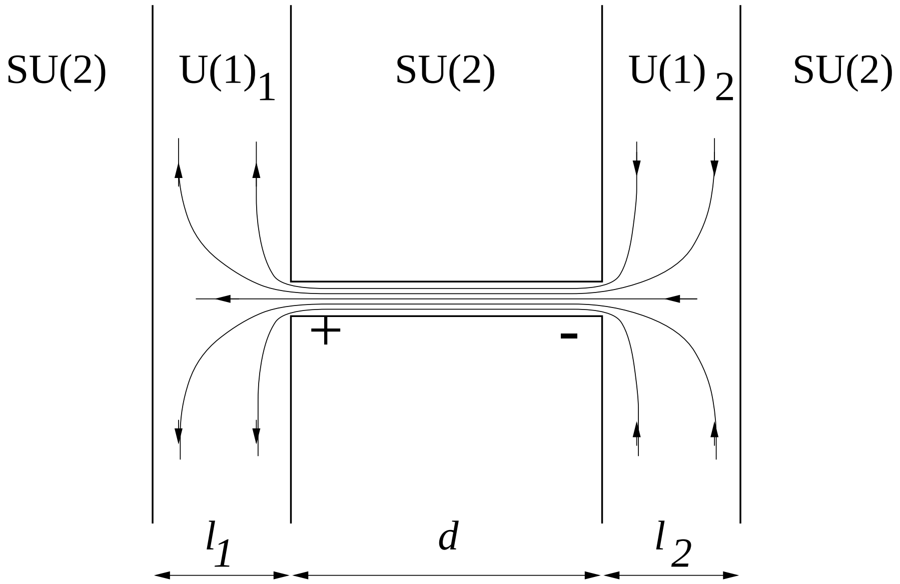

Figure 3: *A color electric flux tube connecting two walls. In terms of the low energy effective theory on the walls this flux tube is described by a field charged under both $U(1)$'s.*

quark states to pop out of the vacuum. Once $q, \bar{q}$ are inside their respective walls, they are no longer in a confining medium and can be moved off to large distances [‡]. The resulting configuration is just a flux tube with field lines coming in from infinity on $\mathcal{W}_1$, through the tube and back out to infinity on $\mathcal{W}_2$ (see figure 3). Note that an observer on $\mathcal{W}_1$ sees this as a state of charge $+1$ under $U(1)_1$, while his friend on $\mathcal{W}_2$ measures this same state to have charge -1 under $U(1)_2$.

Thus at least for $d >> \Lambda$, we have identified a stable state $\phi^{+,-}$, the lowest scalar vibrational mode of the QCD string stretching between the walls, with the quantum numbers we are after. Furthermore, it is clear that for large d the mass of this mode decreases as $\Lambda^2 d$ as the walls are brought closer. It is now tempting to speculate that as the walls come very close together, this state gets lighter and lighter until it becomes tachyonic somewhere around $d \sim \Lambda$ and condenses. In other words, we imagine that the mass for ϕ as a

[‡]Of course since the $(2+1)$-dimensional theory is itself confining at much longer distances, the energy to move $q, \bar{q}$ to infinity diverges logarithmically in the IR.

function of d has the form

$$m_\phi^2(d) = -c_0\Lambda^2 + c_1(\Lambda^2 d)^2 \tag{6}$$

where c_0, c_1 are $O(1)$ constants. The contribution proportional to c_1 has a classical origin and dominates when the string is long, while the first term reflects a (presumably quantum-mechanical) tachyonic instability of the unstretched QCD string. Of course, since the flux tube has a thickness $O(\Lambda)$, in the interesting region it is as long as it is thick so a "string" picture is necessarily heuristic. In terms of the microscopic $(3+1)$-d description the condensate of the tachyonic scalar can be understood as a spontaneous deconfinement transition of the QCD vacuum between the two walls due to a condensate of flux tubes.

It is amusing that while the end result of bringing these "branes" together is very different from the case of bringing D-branes together, the physics has a similar interpretation: the strings stretching between the branes become light and donate their lowest excitations to the effective theory. In the case of D-branes, the lowest-lying excitations of open strings contain gauge fields which enhance the gauge symmetry. On the other hand, our flux-strings have no massless gauge fields so there is no enhancement of gauge symmetry. Instead, the lightest excitation that is donated is tachyonic and further breaks the gauge group! It is also interesting that the tachyonic instability of the QCD string here does not imply that the theory is sick and should be discarded; it simply means that the correct vacuum, where the strings have condensed, must be chosen.

Before we move on note that we can obtain some information about the dynamics of flux tubes by simply translating $(2+1)$ dimensional results into our microscopic description. For example, the fact that the $(2+1)$-d $U(1) \times U(1)$ theory confines tells us that flux tubes in our picture are confined. A stable finite energy configuration is a spinning bound state of two flux tubes of opposite flux. In the case of large wall separation when the flux tubes are long and heavy, this is a non-relativistic bound state but as we tune the distance between the two walls such that the scalar mode $\phi^{+,-}$ becomes light the bound state becomes non-relativistic. It is amusing that these "bound states" of flux are spinning closed flux-strings which overlap both walls.

We will not discuss at length an obvious generalization to N walls with flux tubes stretching between any pair of neighboring walls. These strings are charged under neighboring $U(1)$ and donate the necessary scalar fields to break $U(1)^N \to U(1)_{diag}$ as all N walls are merged.

Instead, we cannot resist the temptation to describe an interesting generalization with strings corresponding to baryons of the confined gauge theory in the bulk. First note that the above construction of domain walls with trapped gauge fields generalizes to branes of dimension $(1+1)$. To construct such a "string" or "1-brane" consider a patch $\mathcal{W}$ in the $y-z$ plane where the bulk non-abelian gauge group is broken to $U(1)$ as depicted in figure 4.

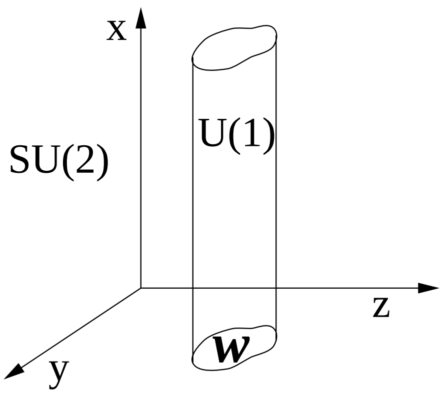

Figure 4: *A photon can be localized on a 1+1 dimensional "1-brane" by embedding the $U(1)$ into a confining $SU(2)$ in the bulk.*

Again, the bulk non-abelian gauge theory is chosen to confine at distances Λ^{-1} taken to be much shorter than the square root of the area A of $\mathcal{W}$. This traps a $(1+1)$ dimensional $U(1)$ gauge theory with gauge coupling $g_2^2 \sim g_4^2/A$ on the string. Given two such 1-branes with areas $A_{1,2}$ and associated $U(1)$ gauge theories one can consider bringing the two 1-branes in contact. Again, the low energy theory sees a reduction of the gauge group from $U(1)\times U(1)$ to $U(1)_{diag}$ with the interpretation of Higgsing of the gauge group via a scalar which becomes light and tachyonic as the two regions are brought within distances of order Λ^{-1}. Evidence for this interpretation is the matching of $U(1)$ couplings which in this case reads

$$\frac{1}{g_{2diag}^2} = \frac{(A_1 + A_2)}{g_4^2} = \frac{A_1}{g_4^2} + \frac{A_2}{g_4^2} = \frac{1}{g_{2,1}^2} + \frac{1}{g_{2,2}^2} \ . \tag{7}$$

Just as in the case of domain walls a QCD string connecting A_1 with A_2 has the correct quantum numbers to supply this scalar.

Consider now a situation with a confining gauge group $SU(N)$ in the bulk and N patches with roughly equal areas A_i for $i = 1, .., N$ on which the running of the $SU(N)$ coupling has been slowed down. (For example, this could be arranged by adding matter fields in the adjoint representation of $SU(N)$ which have very large masses in the bulk but are light on the 1-branes.) Then the four dimensional gauge coupling remains small on the N patches, and below the matching scale $\mu \sim A_i^{-1/2}$ the long distance physics is described by a $(1+1)$ dimensional $SU(N)^N$ gauge theory. As before we can create flux tubes connecting any pair of the various 1-branes by placing a pair of heavy test quarks q and $\bar{q}$ in the confined bulk, pulling them apart to form a flux string and then moving the two quarks in two separate regions A_i and A_j. After removal of the test quarks we are left with a flux tube which transforms in the fundamental representation of $SU(N)_i$ and an antifundamental of $SU(N)_j$.

What happens if we start with N test quarks in a color singlet state corresponding to a baryon of the $SU(N)$ bulk gauge group? We can now move each of the test quarks into a different one of the patches A_i; the flux tubes from each of these quarks meet at a common junction in the bulk where N units of flux combine into a color singlet. After removal of the test quarks we are left with a baryonic N-pronged flux tube connecting the N regions as in figure 5.

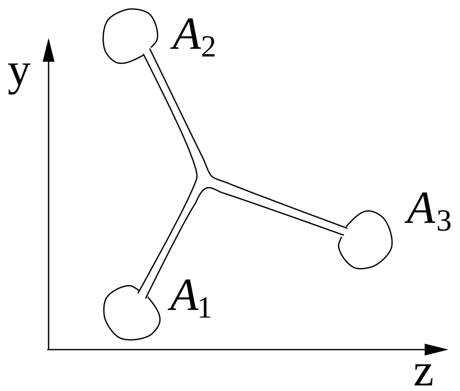

Figure 5: *A 3-pronged baryonic flux tube connecting three 1-branes.*

This baryonic string is stable and has a mass of order $N\Lambda^2$ times the characteristic distance between the various patches. In the low energy effective

$SU(N)^N$ theory its lowest vibrational mode would be described by a massive field which transforms in the fundamental representation of each of the $SU(N)$'s. As we bring all of the N patches close together, the baryonic string as well as all the "mesonic" strings become light. As we bring the 1-branes very close to each other we expect a condensate of strings which deconfines the vacuum in the region between the 1-branes, causing them to merge. In the long distance theory this is described by a condensate of scalars which breaks $SU(N)^N \to SU(N)$.

Acknowledgements

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Appendix

In this appendix we formalize the conclusions regarding the trapping massless electric or magnetic photons in the theory where a $U(1)$ gauge field is higgsed away from $\mathcal{W}$ but is unbroken inside $\mathcal{W}$. Let us consider formally the limit where $Ml \to \infty$, so that the photon outside is really infinitely heavy. The $U(1)$ then really only exists inside $\mathcal{W}$; with a Lagrangian given by

$$\mathcal{L} = \int d^3x \int_{-l/2}^{l/2} dz \frac{1}{g_4^2} F^{\mu\nu} F_{\mu\nu}. \tag{8}$$

This just looks like the compactification of a $U(1)$ gauge theory from $4 \to 3$ dimensions on an interval of length l. However, the spectrum of the theory at energies beneath the compactification scale l^{-1} depends crucially on the boundary conditions imposed on $F^{\mu\nu}$ at $z = +l/2, -l/2$. This is because massless states in the low energy theory must be zero modes in the z direction, and are therefore sensitive to the boundary conditions, which may or may not eliminate them. In the present case, the region outside the wall is superconducting, so the appropriate boundary conditions are that the electric field is perpendicular to the wall (true for any conductor) and that the magnetic field is parallel to the wall (which is only true for a superconductor); that is

$$E_x = E_y = 0, \quad B_z = 0 \tag{9}$$

which can be written more covariantly as

$$F^{ab} = 0, \quad a, b = t, x, y. \tag{10}$$

This makes it clear that a massless photon coupling to electric charge is not present in the low energy theory, it is projected out of the usual Kaluza-Klein spectrum by the boundary conditions. The usual KK scalar, corresponding to F^{a3}, remains in the massless spectrum, but does not couple to electric charge. Rather, a massless magnetic photon has been trapped. Indeed, the boundary conditions can also be written as

$$\tilde{F}^{a3} = 0 \tag{11}$$

which leaves the zero mode of $\tilde{F}^{ab}$ in the massless spectrum. This of course works because in $(2+1)$ dimensions, a scalar is dual to a vector field.

Of course in both cases, we also have a tower of massive states. This follows from the standard Kaluza-Klein analysis with the boundary conditions appropriately imposed. But we can also see it in another way. Let us generalize to the case of an n dimensional wall of thickness l in an $n+1$ dimensional space. Let $\vec{x}$ be the n-dimensional coordinates, and y the $n+1$'th coordinate. Placing an electric charge at the origin, let us compute the electric potential at the point $\vec{x}, y = 0$ on the wall. We can enforce the boundary conditions by placing an infinite sequence of image charges of charge $(-1)^q$ at $\vec{x} = 0, y = ql$. The potential is then

$$V(\vec{x}) = \sum_{q=-\infty}^{+\infty} \int \frac{d^n k}{(2\pi)^n} \frac{dk'}{(2\pi)} \frac{e^{i(\vec{k}\vec{x} + k'lq - \pi q)}}{\vec{k}^2 + k'^2} \tag{12}$$

where we have used the expression for the $(n+1)$ dimensional Coulomb potential in terms of its Fourier transform. If we now use the familiar Poisson resummation identity

$$\sum_{q=-\infty}^{\infty} e^{iq\theta} = 2\pi \sum_{s=-\infty}^{\infty} \delta(\theta - 2\pi s) \tag{13}$$

we can perform the integral over k', leaving

$$V(\vec{x}) = \sum_{s=-\infty}^{\infty} \int \frac{d^n k}{(2\pi)^n} \frac{e^{i\vec{k}\vec{x}}}{\vec{k}^2 + (2\pi/l)^2(s+1/2)^2}. \tag{14}$$

Note that the integrand is just the n-dimensional Yukawa potential for a field of mass $2\pi/l \times (s+1/2)$. Therefore, we have shown that the potential can be expressed in terms of a sum over a tower of n dimensional massive states. If we instead place a magnetic charge at the origin, an infinite sequence of magnetic image charges of the *same* sign enforce the boundary condition, and the πq term in eqn.(12) disappears. Therefore, the potential is due to a sum over n dimensional massive fields of mass $2\pi/l \times s$, which for $s=0$ includes the massless magnetic photon we expect in this case.

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