Generalized axion-photon couplings and dual descriptions on brane-worlds

M. O. Tahim and C. A. S. Almeida
Departamento de Física, Universidade Federal do Ceará,
Caixa Postal 6030, 60455-760, Fortaleza, Ceará, Brazil

In this work we obtain topological and dual theories on brane-worlds in several dimensions. Our brane is a solitonic-like hypersurface embedded in a space-time with a specific dimensionality and it appears due to the breaking of a Peccei-Quinn-like symmetry. In the first part of this work, the obtained topological theories are related to a generalization of the axion-photon anomalous interaction in $D = 4$ (in the Abelian case) and to the Wess-Zumino term (in the non-Abelian case). In the second part, we construct dual models on the brane through a mechanism of explicit Lorentz symmetry breaking. The gauge symmetries of such models are discussed within the Stuckelberg formalism.

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I. INTRODUCTION

In current brane-world research the most important subject has to do with (fermionic or bosonic) matter field localization procedures. Its importance resides in the fact that is just this procedure that will give physical information about the universe inside the brane. Such information are interaction energy scales, the range of the interactions, the chiralities of the fields, symmetries, etc. In this direction, hard work have been made. In the case of gravity localization procedures, the standard way commonly accepted is that given by the Randall-Sundrum scenario. This scenario is characterized by treating Einstein gravity in classical level. However, there are a lot of approaches and interpretations to Einstein gravity, one of them is the so called $B \wedge F$-Gravity. Such formalism, strictly related to Loop Gravity, is known by providing a consistent machinery of quantization of gravity. An interesting peculiarity of this type of theory is that it is completely background independent.

Due to these developments, we regard as an important subject look for topological theories in the brane context, with the major objective to add quantum information to the Randall-Sundrum scenario. Pursuing this idea, as the first step, we study topological and equivalent theories on brane-worlds in several dimensions and discuss some mechanisms of dimensional reduction, namely, the naive one and a mechanism involving branes. In the first part of this work we construct topological theories in brane-worlds. The brane-world is regarded as a kink-like soliton and it appears due to a spontaneous symmetry breaking of a Peccei-Quinn-like symmetry. Topological terms are obtained by generalizing to several dimensions the axion-photon anomalous interaction (in the Abelian case) and the Wess-Zumino term (in non-Abelian theories). Indeed, this anomalous interaction is valid in a grand unification scenario with the electroweak angle $\sin^2 \theta_W = 3/8$. Such effect is mainly due to the presence of instantons in the Standard Model. However the generalizations made in this work are simply inspired in this 4D anomalous interaction and they aim in to preserve the field theories basic symmetries in order to open the possibilities of application of these new topological terms in brane-world scenarios. Despite the fact that such theories do not have propagating local degrees of freedom, they are important in dual descriptions, regardless of other applications.

In the second part, we construct dual tensor gauge theories on the brane. Investigations of dual formulations for tensor fields are important for understanding of alternative formulations of known theories like gravity as well as understanding of their role in superstrings. Also, such models are important because, in the non-Abelian case, they describe a low energy region of a QCD system and, since that the topological properties do not depend on small distances (local characteristics) but depend on the global ones, this has a direct application in describing the IR regimen of QCD, where the perturbation theory fails down. Besides this, we point out the fact that these models are equivalent to free and propagating bosonic theories, an interesting characteristic that may be used in the context of localization scenarios. As we will show, this construction involves Lorentz symmetry breaking. We also discuss, at the end, the gauge symmetries of these models within the Stuckelberg formalism. In this case we note the breaking of a gauge symmetry of the type $U(1) \times U(1)$ down to $U(1)$.
II. TOPOLOGICAL THEORIES ON DOMAIN WALL-BRANES (ABELIAN CASE)

In what follows we use capital letters to denote space-time indices in a $D > 4$ space-time, greek indices stand for a $D = 4$ space-time and lower cases (a, b, c, d...) stand for a $D < 4$ space-time. We implement the theory through the following action in $D = 5 + 1$:

$$S = \int d^5x \left( -\frac{1}{2(3!)} H_{MNP} H^{MNP} + g \varepsilon^{MNPQRS} \phi(z) H_{MNP} H_{QRS} + \frac{1}{2} \partial_M \phi \partial^M \phi + V(\phi) \right).$$ (1)

In this action, $H_{MNP} = \partial_M B_{NP} + \partial_N B_{PM} + \partial_P B_{MN}$ $(M, ..., S = 0, ..., 5)$ is the field strength of an antisymmetric tensor gauge field $B_{MN}$. The field $B_{MN}$ has an important function in string theory: it couples correctly with the string world-sheet in a very similar way to the coupling of a gauge vector field $A_M$ to the universe line of a point particle. The field $\phi$ is a real scalar field, and $V(\phi)$ is a potential that provides a phase transition:

$$V(\phi) = \lambda (1 - \cos \phi).$$ (2)

In fact, the solution for the specified potential, is

$$\phi(z) = \pi + 2 \arcsin(\tanh \sqrt{x} z),$$ (3)

for $\phi \in (0, 2\pi)$ when $z \to \pm \infty$. This solution corresponds to the topological sector which contains just one soliton [1, 2]. The second term in the action (1) is a term that generalizes the coupling that appears from the anomaly of the Peccei-Quinn quasisymmetry [2] in $D = 3 + 1$, namely, $\phi \to \phi + 2\pi$. This symmetry is broken by the potential term written with the real scalar field $\phi$. In a phase transition, only the later field acquires a nonzero vacuum expectation value (VEV). For the case of the action (1) the phase transition favors the creation of domains of different phases: in fact, the vacuum is made of many disconnected points, i. e., the potential is minimized when $\phi_{\text{vacuum}} = 2\pi n$, where $n$ is an integer number. The domains created are separated by dynamical soliton hypersurfaces topologically stable. In this sense, these objects are brane-like objects [3, 4]. For such, the space-time dimension is $D = 5 + 1$ and the hypersurface is a $D = 4 + 1$ world. Now we can work with the second term of Eq. (1) (considering that $\phi$ only depends on the $z$-coordinate) in order to obtain new terms by integration by parts. The first term obtained is a total derivative and may be disregarded. Another term is identically null because of the antisymmetry of the Levi-Civita symbol. As the field $\phi$ depends only on the $z$-coordinate, the topological term may be rewritten as follows:

$$S_{\text{top.}} = 3\alpha \int d^5x \varepsilon^{3NPQRS} \partial_M \phi(z) B_{NP} H_{QRS}. $$ (4)

Considering that the $B_{MN}$ field weakly depends on the $z$-coordinate, Eq. (4) is rewritten as:

$$S_{\text{top.}} = \int d^5x \varepsilon^{3NPQRS} k B_{NP} H_{QRS}. $$ (5)

This last equation shows that over the hypersurface an effective topological term appears with a coupling constant $k$ that have canonical dimension of mass. This coupling constant is quantized in various ways [5, 10]. The theory over the hypersurface is completely five-dimensional. This term is very similar to the Chern-Simons term [10], which is written in $D = 2 + 1$ with a gauge vector field $A_\mu$:

$$S_{\text{cs}} = g \int d^3x \varepsilon^{abc} A_a F_{bc}. $$ (6)

Nevertheless, the Eq. (5) is written only with tensorial antisymmetric fields $B_{MN}$. Such term have been used to explain some peculiarities of the Cosmic Microwave Background Radiation (CMBR) [11] within the Randall-Sundrum scenario [12]. It is interesting now to observe the properties of the action (1) in lower dimensional space-times using dimensional reduction. Thus, supposing that the fields of the action (1) are independent of the coordinate $x_M \equiv x_5$ which is not the argument of the field $\phi(z)$ and defining

$$B_{P6} = V_P, $$
$$B_{6P} = -V_P, $$
$$V_{MN} = \partial_M V_N - \partial_N V_M, $$

$$S_{\text{top.}} = \int d^5x \varepsilon^{3NPQRS} k B_{NP} H_{QRS}. $$ (5)
the action \( S \) becomes:

\[
S = \int d^5x \left( -\frac{1}{4} V_{MN} V^{MN} - \frac{1}{2} \epsilon^{MNPQR} \phi(z) V_{MN} H_{PQR} + \frac{1}{2} \partial_M \phi \partial^M \phi + V(\phi) \right).
\] (8)

This action in \( D = 4 + 1 \) has now a vectorial gauge field \( V_M \) reminiscent of the reduction, and contains yet the real scalar field \( \phi \), that again may give rise to the formation of a lower dimension domain wall-brane. In this case, the space-time dimension is \( D = 4 + 1 \) and the hypersurface is a \( D = 3 + 1 \) universe. If we observe the theory over the solitonic hypersurface we will obtain that, rewriting the topological term of the action (8) as made in Eqs.(4), (5) and (6):

\[
S_{\text{top.}} = \int d^4x \left( k \epsilon_{\mu\nu\alpha\rho} \phi(z) \partial_\mu \partial_\nu \partial_\alpha \partial_\rho \right).
\] (9)

We can note that the theory on the domain wall is strictly four-dimensional. If the field \( V_\mu \) is identified with the potential four-vector \( A_\mu \) then we obtain the action for the \( B \wedge F \) model \[12\] on the domain wall-brane. This action, under certain conditions, can give rise to a mechanism of topological mass generation for the field \( A_\mu \) or for the field \( B_{\mu\nu} \).

Starting from Eq.(8), the discussion for lower dimensions (\( D = 3 + 1 \) and \( D = 2 + 1 \)), using the same methods, will lead to the following topological action:

\[
S_{\text{top.}} = \int d^4x \left( k \epsilon_{\mu\nu\alpha\rho} \phi(z) \partial_\mu \partial_\nu \partial_\alpha \partial_\rho \right) + \epsilon_{\mu
u\alpha\rho} \phi(z) F_{\mu
u} W_{\alpha\rho}.
\] (10)

The fields \( \phi \) and \( W_{\alpha\rho} = \partial_\alpha W_\rho - \partial_\rho W_\alpha \) emerge as degrees of freedom reminiscent of the reduction. These fields are defined as in Eq.(7). If we work with the first term of Eq.(10) on the domain wall, we will find a different topological theory \[13\], namely:

\[
S = \int d^3x \left( g \epsilon^{abc} \partial_a \phi B_{bc} \right).
\] (11)

Identifying again in the second term of Eq.(10), the vector field \( W_\mu \) as the gauge field \( A_\mu \), we will obtain the anomalous interaction term between the real scalar field \( \phi \) and the field \( A_\mu \). This term, rearranged on the domain wall, reduces to the Chern-Simons term given by Eq.(6).

### III. TOPOLOGICAL THEORIES ON DOMAIN WALL-BRANES (NON-ABELIAN CASE)

Non-Abelian theories provide alternative mechanisms of treating and quantization of gravity \[10\]. Recent interest have appeared in the study of metric independent gravity theories \[2\]. As applications of the results discussed above, we will show how to obtain non-Abelian topological terms on domain wall-branes. These type of terms are important for studies about topological gravity trapped on brane-worlds.

In this case, the brane emerges from the same mechanism discussed above. In our description, we consider the following action:

\[
S = \int d^4x \left( \frac{1}{2} \partial_\mu \phi \partial^\mu \phi + V(\phi) + k \epsilon^{\mu\nu\alpha\rho} \phi(z) F_{\mu
u}^i F_{\alpha\rho}^i \right).
\] (12)

In this action, \( F_{\mu\nu}^i = \partial_\mu A_\nu^i - \partial_\nu A_\mu^i + g f_{ijk} A_\mu^i A_\nu^j \), where \( i, j, k = 1, ..., n \) and \( \phi \) is a real scalar field. We consider \( i, j, k, l, ... \) as indices of a finite dimensional semi-simple Lie group \( G \). As already discussed, writing the third term of Eq.(12) over the hypersurface, we obtain the following effective action:

\[
S = k \int d^3x \epsilon^{\mu\nu\alpha} \left( \partial_\mu A_\nu^i A_\alpha^i + g f_{ijk} A_\mu^i A_\nu^j A_\alpha^k \right).
\] (13)

This term is rather similar to the non-Abelian Chern-Simons term. As discussed by Deser and Jackiw \[10\], the non-Abelian Chern-Simons term is not invariant under large gauge transformations. However, this behavior is avoided if
we consider quantization of the coupling constant of this theory. The term found in Eq.\(^{(13)}\) may be used to describe gravity in \(D = 3\) in the same way as in the references \(^{2,10}\).

Another model may be obtained starting from the following action in \(D = 5\):

\[
S = \int d^5x \left( \frac{1}{2} \partial_M \phi \partial^M \phi + V(\phi) + k \epsilon^{MNPQR} \phi \partial_M B_{NP} \right),
\]

where \(H^i_{MNP} = \partial_M B^i_{NP} + \partial_N B^i_{PM} + \partial_P B^i_{MN} + g' \epsilon^{ijk} A^j_M B^k_N \) \((M,N,... = 0,...,5)\). In this last case, after simple calculations we obtain an action that contains an effective topological term of the non-Abelian \(B \wedge F\) type:

\[
S = k \int \! d^4x \epsilon^{\mu
u\rho\sigma} B_{\mu\nu} F_{\rho\sigma}.
\]

This action is regarded as basic for treating metric independent gravity theories in \(D = 4\) as topological constrained field theories \(^3\).

**IV. DUAL THEORIES ON DOMAIN WALL-BRANES**

In this section we discuss a procedure for construction of dual theories on domain wall-branes. We start with the following action in \(D = 5\):

\[
S = \int d^5x \left[ \frac{1}{2} \partial_M \phi \partial^M \phi - V(\phi) - \frac{k}{2} \epsilon^{MNPQR} \phi \partial_M B_{NP} \partial F_{QR} - \frac{k'}{2} B^{MN} B_{MN} \partial V^P \right].
\]

The field \(\phi\) yet provides the emergence of the domain wall-brane, in this case, a 3-brane. The field \(F_{MN}\) is the field strength for the gauge field \(A_M\), i.e., \(F_{MN} = \partial_M A_N - \partial_N A_M\) (we are treating the Abelian case) and \(B_{MN}\) is an antisymmetric tensor field, the Kalb-Ramond field. The vector \(V^M\) in the fourth term represents an additional parameter of the theory, i.e., it represents another gauge freedom of this theory. This freedom can be fixed by choosing \(V^M\) pointing along a preferential direction. With a choice like this we can break the \(SO(1,4)\) Lorentz symmetry of the model. Similar procedure have been made in the context of topological gravity theories \(^{14}\). Another interesting attempts have been made in scenarios with Lorentz symmetry breaking \(^{13}\).

In the background of a domain wall-brane \(\phi = \phi(x_4)\), and we choose \(V^\mu = (0,0,0,0,\phi)\) in such a way that the last terms of the action \(^{(10)}\) can be rearranged:

\[
S \sim \int \! d^4x \left[ \frac{k}{2} \epsilon^{4MN} \partial_4 \phi(x_4) B_{MN} \partial F_{QR} + \frac{k'}{2} \partial_4 \phi(x_4) B^{MN} B_{MN} \right].
\]

We note that this last action is yet invariant under \(\phi \rightarrow \phi + 2\pi\). Making a thin wall approximation, i.e., \(\partial_4 \phi(x_4) = \delta(x_4)\) and defining \(g^2 = \frac{k'}{k}\) we obtain

\[
S \sim k \int \! dx_4 \delta(x_4) \int \! d^4x \left[ \frac{1}{2} \epsilon^{\nu\alpha\beta\lambda} B_{\nu\alpha} F_{\beta\lambda} + \frac{1}{2} g^2 B^{\alpha\beta} B_{\alpha\beta} \right].
\]

where we have made the identification \(\epsilon^{4\nu\alpha\beta\lambda} = \epsilon^{\nu\alpha\beta\lambda}\). The conclusion is that we obtain the \(B \wedge F\)-Maxwell model in \(D = 4\), namely

\[
S \sim \int \! d^4x \left[ \frac{1}{2} \epsilon^{\nu\alpha\beta\lambda} B_{\nu\alpha} F_{\beta\lambda} + \frac{1}{2} g^2 B^{\alpha\beta} B_{\alpha\beta} \right].
\]

This model, as it is well known, is equivalent to the free gauge invariant (non-massive) Maxwell theory \(^7\). We arrive at this result by explicitly breaking the \(SO(1,4)\) Lorentz symmetry down to \(SO(1,3)\) through the \(V^\mu\) choice and the supposition \(\phi \equiv \phi(x_4)\).

The procedure explained above can be applied to other sort of theories. In particular, we can obtain similar results for lower dimensional theories. For example, we consider the following models in \(D = 4\):

\[
S = \int \! d^4x \left[ \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) - k \epsilon^{\mu\nu\alpha\beta} \phi \partial_\mu W_\nu F_{\alpha\beta} + k' W^\mu W_\mu \partial_\lambda V^\lambda \right]
\]

and

\[
S' = \int \! d^4x \left[ \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) - k \epsilon^{\mu\nu\alpha\beta} \phi \partial_\mu B_{\nu\alpha} \partial_\beta \phi - \frac{k'}{2} B^{\alpha\beta} B_{\alpha\beta} \partial_\lambda V^\lambda \right].
\]
In these cases we have more fields: $W^\mu$ is an Abelian gauge field in the first model and $\varphi$ is a real scalar field in the second model. These two theories have $SO(1,3)$ Lorentz symmetry. Again, choosing $V^\mu = (0,0,0,\phi)$ and $\phi \equiv \phi(x_3)$ we break $SO(1,3)$ down to $SO(1,2)$. In the thin wall approximation we are lead to the following theories:

\[ S \sim \int d^3x [\varepsilon^{abc} W_a F_{bc} - g^2 W^a W_a] \quad (22) \]

and

\[ S' \sim \int d^3x [\varepsilon^{abc} B_{ab} \partial_c \varphi + \frac{1}{2} g^2 B_{ab} B_{ab}]. \quad (23) \]

The first model is the $B \wedge F$-Maxwell model in $D = 3$ which is dual to a free and non-massive Maxwell theory. The last model is called $B \varphi$-Klein-Gordon model \[13\] which is also dual to the free and non-massive Maxwell theory or yet to the free and non-massive Klein-Gordon theory.

It is interesting to study the gauge symmetries of these models. It is clear, for example, that the action \[16\] is not invariant under the usual gauge transformations (regarding $V_M$ as a gauge field):

\[ \delta \phi = 0, \quad \delta A_M = \partial_M \alpha, \quad \delta V_M = \partial_M \overline{\alpha}, \quad \delta B_{MN} = \partial_{[M} \Lambda_{N]} \quad (24) \]

In order to recover the gauge symmetry we can make use of the Stuckelberg formalism \[16\] simply redefining the fields of the model:

\[ A^M \rightarrow A^M + \frac{1}{g} \partial^M \theta, \quad V^M \rightarrow V^M + \frac{1}{g} \partial^M \overline{\theta}, \quad B_{MN} \rightarrow B_{MN} + \partial_{[M} \Lambda_{N]} \quad (25) \]

Defining the new gauge transformations

\[ \delta \phi = 0, \quad \delta \Lambda = \partial_M \alpha, \quad \delta \theta = -g \alpha, \quad \delta V_M = \partial_M \overline{\alpha}, \quad \delta B_{MN} = \partial_{[M} \Lambda_{N]}, \quad \delta \Lambda_N = -\Lambda_N \quad (26) \]

we recover the gauge symmetry of the model. An interesting characteristic is that the action \[16\] contains two gauge fields, namely $A_M$ and $V_M$ and, therefore, the gauge symmetry for the vector fields is of the type $U(1) \times U(1)$. A symmetry like this have been discussed in the context of models for superconductors \[17\]. The important point here is that when we make the choice for the direction of the vector $V^M$ we also break the $U(1) \times U(1)$ gauge symmetry down to $U(1)$, i.e., the Stuckelberg formalism can be applied in the domain wall-brane to restore this $U(1)$ gauge symmetry. The same arguments are valid for other sort of theories.

V. DISCUSSIONS AND OUTLOOK

In this work we have discussed how to construct topological and dual theories in domain wall-branes in several dimensions. In the first part, we have discussed the appearance of several Chern-Simons-like topological terms, in Abelian and in non-Abelian theories, by a procedure of dimensional reduction (the naive one). We have seen that the brane breaks the Lorentz symmetry of these models. In order to do this, the domain wall-brane has been simulated by a kink-like soliton embedded in a higher dimensional space-time and it has emerged due to a spontaneous symmetry breaking of a specific discrete symmetry, namely, a Peccei-Quinn-like symmetry.

In the second part of this work, we have developed a way to build dual models in domain wall-branes. We have constructed, from a $D = 5$ theory, the $B \wedge F$-Maxwell model in $D = 4$ and, from a $D = 4$ theory, the $B \wedge F$-Maxwell model and the $B \varphi$-Klein-Gordon both in $D = 3$. The procedure adopted consist in the explicit breaking of the Lorentz symmetry due to the choice of a preferential direction for the additional vector parameter $V^\lambda$ of the theories. Interestingly, in the case of the $B \wedge F$-Maxwell model, we can implement through the Stuckelberg formalism a $U(1) \times U(1)$ gauge symmetry that is broken down to $U(1)$ in the domain wall-brane due to our choice. Another characteristic is that we always obtain in the domain wall-brane, theories that are dual to free and gauge invariant models. Indeed, this fact is compatible with the idea of localization of fields in brane worlds, where the zero modes (described by non-massive theories) mimics the standard model fields. In that case, the existence of bosonic modes is guaranteed if we assume the existence of kinetic terms in the world-volume as well as in the bulk of the brane. In branes which contain matter, this can be turned out by perturbative corrections \[18\]. On the other hand, we argue in this work that those kinetic terms can be generated using tensor like dualities. Hence its importance for mechanisms of field localization in branes. A complete analysis of such idea should be interesting. We consider also that the
generalization of these dual models for the non-Abelian case should be important. In this case, applications to gravity theories would be possible.

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