Gravity in Dynamically Generated Dimensions

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Abstract

A theory of gravity in $d+1$ dimensions is dynamically generated from a theory in $d$ dimensions. As an application we show how $N$ dynamically coupled gravity theories can reduce the effective Planck mass.

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The idea that a gauge theory in \( (1,d) \) dimensions appears as the low energy limit of a \( (1,d-1) \) theory with many fields has been recently put forward [1,2]. One starts with \( (1,d-1) \) gauge fields \( A_\mu(x,i), \mu = 0, 1, \cdots, d-1 \) with the range of the discrete index \( i \) being either infinite of finite and periodic \( (A_\mu(x,1) = A_\mu(x,N)) \). Interactions are chosen as to insure that a field \( A_{i,i+1}(x) \) is generated dynamically and whose interactions, in the low energy limit, mimic \( A_d \) in a \( (1,d) \) dimensional space with \( i \) turning into the discrete extra dimension. A feature of this approach is that the \( (1,d-1) \) theory has the desired properties of renormalizability and asymptotic freedom.

In this work, we extend this approach to gravity, namely we generate a \( (1,d) \) dimensional gravity from a theory in \( (1,d-1) \) dimensions. Gravity will be described in a moving frame formalism as an \( SO(1,d-1) \) gauge theory on a \( (1,d-1) \) manifold. Unlike the situations discussed in [1,2], where one started with a gauge theory in some dimension and generated the same gauge theory in a space with one higher dimension, in order to generate gravity in the higher dimension we have to start with an extended gravity in the lower dimensional space. Namely, we start with an \( SO(1,d) \) gauge theory on a \( (1,d-1) \) manifold and dynamically generate the \( SO(1,d) \) theory on a \( (1,d) \) manifold [3]. Of course, the lower dimensional theory includes gravity in the \( SO(1,d-1) \) subgroup of \( SO(1,d) \).

In this case we cannot appeal to renormalizability or to asymptotic freedom to justify this approach as the lower dimensional gravity or extended gravity is unlikely to be renormalizable or asymptotically free. What our construction insures is that the dynamically generated higher dimensional theory is no more singular that the lower dimensional one and that coordinate invariance and local \( SO(1,d) \) invariance are maintained at each step. Details, as well as a discussion of \( SO(1,d) \) invariant interactions in \( (1,d-1) \) dimensional spaces, are presented in Sec. II. In Sec. III we apply this approach to a scenario where in four dimensions the existence of \( N \) dynamically coupled gravity theories decreases the effective gravitational coupling by a factor of \( N \).
II. DYNAMICAL GENERATION OF GRAVITY

A. Gravity in \( d + 1 \) Dimensions

We shall first discuss gravity theory on a \((1, d)\) dimensional manifold with coordinates \(x_\mu; \mu = 0, 1, \cdots, d\), which we wish to obtain through the dynamical generation of a dimension in some theory on a \((1, d - 1)\) dimensional manifold. Our goal is the \((1, d)\) Einstein-Hilbert Lagrangian, which we express in the moving frame or d-ad formalism.

\[
L_{EH} = M^{d-1}_{(d+1)} \epsilon^{\mu_1 \mu_2 \cdots \mu_{d+1}}_{a_1 a_2 \cdots a_{d+1}} \left[ B_{\mu_1 \mu_2}^{a_1 a_2} (x) + M^2_{(d+1)} \lambda_{(d+1)} e^a_{\mu_1} (x) e^a_{\mu_2} (x) \right] \\
\times e^a_{\mu_3} (x) e^a_{\mu_4} (x) \cdots e^a_{\mu_{d+1}} (x);
\]

(1)

in the above the \(e^a_{\mu} (x)\)'s, with flat space, Minkowski, indexes \(a = 0, 1, \cdots, d\), are the \((d+1)\)-ads and the \(\omega_{\mu}^{ab} (x)\)'s are the spin-connections. \(M_d\) is the \(d\) dimensional Planck mass, \(M^2_{d} \lambda_d\) is a cosmological constant and the curvature tensor \(R_{\mu_1 \mu_2}^{a_1 a_2} (x)\) is related to the spin-connections by

\[
R_{\mu \nu}^{ab} (x) = \partial_{\mu} \omega_{\nu}^{ab} (x) - \omega_{\mu, \nu}^{a} (x) \omega_{\nu}^{b} (x) - (\mu \leftrightarrow \nu).
\]

(2)

In order to see what theory in a \((1, d - 1)\) space we should start with, we foliate the \((1, d)\) dimensional manifold into \((1, d - 1)\) dimensional ones. Specifically we single out the last coordinate \(x_d\). Coordinates in the \((1, d)\) dimensional space are written as \((x_\mu, x_d)\), where now \(\mu = 0, 1, \cdots, d - 1\). We leave the first \(d\) d-ads as they were but separate out the “shift” vector [4]

\[
e^a_{d} (x) = N^a (x);
\]

(3)

the Minkowski index, \(a\), still ranges over \((d+1)\) values. In terms of this shift vector, eq.(1) becomes [5]

\[
L_{EH} = L_A + L_B,
\]

(4)

with

3
\[ L_A = M_{(d+1)}^a \epsilon^{a_{1}a_{2}...a_{d+1}} \left[ (d-1) R_{\mu_1 \mu_2}^a (x) + (d+1) M_{(d+1)}^2 \epsilon_{\mu_1}^a (x) \epsilon_{\mu_3}^a (x) \right] \]
\[ \times \epsilon_{\mu_3}^a (x) \epsilon_{\mu_4}^a (x) \cdots \epsilon_{\mu_d}^a (x) N_{a_{d+1}} (x), \]
\[ (5) \]
\[ L_B = 2M_{(d+1)}^a \epsilon^{a_{1}a_{2}...a_{d+1}} R_{\mu_1 \mu_2}^a (x) \epsilon_{\mu_1}^a (x) \epsilon_{\mu_3}^a (x) \epsilon_{\mu_4}^a (x) \cdots \epsilon_{\mu_d}^a (x). \]
\[ (6) \]

We note that \( L_A \) does not contain any derivatives in the \( x_d \) direction nor any terms involving \( \omega_{d \mu} (x) \) while \( L_B \) does but, in turn, does not involve the shift vectors \( N^a (x) \). \( L_A \) will determine the \((1,d-1)\) theory we start out with, and our goal will be to generate dynamically \( L_B \).

**B. SO(1,d) Gauge Theory in \( d \) Dimension**

The Lagrangian \( L_A \) in eq. (5) describes an \( SO(1,d) \) gauge theory in on a \((1,d-1)\) manifold; as \( SO(1,d-1) \) is a subgroup of \( SO(1,d) \), this model includes gravity, an \( SO(1,d-1) \) gauge theory on a \((1,d-1)\) manifold. We shall now obtain some of the properties of this extended model; these properties are needed by other fields and their interactions. Following eq. (5) we rewrite the Lagrangian as

\[ L_{SO(1,d)} = M_{(d+1)}^a \epsilon^{a_{1}a_{2}...a_{d+1}} \left[ R_{\mu_1 \mu_2}^a (x) + M_{d}^2 \lambda_{d} \epsilon_{\mu_1}^a (x) \epsilon_{\mu_2}^a (x) \right] \]
\[ \times \epsilon_{\mu_3}^a (x) \epsilon_{\mu_4}^a (x) \cdots \epsilon_{\mu_d}^a (x) N_{a_{d+1}} (x). \]
\[ (7) \]

As \( \epsilon_{\mu}^a (x) \) is not a square matrix we may not define an \( \epsilon_{a}^\mu (x) \) as its inverse. We can, however, introduce a \( d \times d \) metric tensor

\[ g_{\mu \nu} (x) = \epsilon_{\mu a} (x) \epsilon_{\nu}^a (x) \]
\[ (8) \]

as well as its inverse \( g^{\mu \nu} (x) \), thus allowing us to raise and lower the curved space coordinate indexes. Using the \( SO(1,d) \) Clifford algebra

\[ \{ \gamma_a, \gamma_b \} = 2\eta_{ab} \]
\[ (9) \]

and its associated spin matrices \( \Sigma_{ab} = [\gamma_a, \gamma_b] / 2i \) the Dirac Lagrangian for an \( SO(1,d) \) spinor field \( \psi(x) \) is
\[ L_D = \epsilon^{\mu_1 \mu_2 \cdots \mu_d} \epsilon_{a_1 a_2 \cdots a_{d+1}} \bar{\psi}(x) \gamma^a D_{\mu_1} \psi(x) \epsilon^{a_2}_{\mu_2} \cdots \epsilon^{a_d}_{\mu_d} N^{a_{d+1}}(x), \]  \hspace{1cm} (10)

with the covariant derivative

\[ D_{\mu} \psi(x) = \left( \partial_{\mu} - \frac{1}{2} \Sigma_{ab} \omega_{\mu}^{ab} \right) \psi(x). \]  \hspace{1cm} (11)

### C. Dynamical Generation of Gravity in \( d + 1 \) Dimensions

In order to generate an extra dimension we study many mutually non-interacting \( SO(1, d) \) theories described by d-ads \( e^a_{\mu}(x, i) \), spin connections \( \omega_{\mu}^{ab}(x, i) \) and fields \( N^a(x, i) \); at this point the range of the \( i \)'s need not be specified. The Lagrangian for this collection of theories is

\[ L_0 = M^d - 2 \sum_i \epsilon^{\mu_1 \mu_2 \cdots \mu_d} \epsilon_{a_1 a_2 \cdots a_{d+1}} \bigg[ R^a_{\mu_1 \mu_2} (x, i) + M^d_\lambda \epsilon_{\mu_1}^{a_1} (x, i) \epsilon_{\mu_2}^{a_2} (x, i) \bigg] \times \epsilon_{\mu_3}^{a_3} (x, i) \epsilon_{\mu_4}^{a_4} (x, i) \cdots \epsilon_{\mu_d}^{a_d} (x, i) N^{a_{d+1}}(x, i). \]  \hspace{1cm} (12)

It is invariant under the product group \( \cdots \times SO^i(1, d) \times SO^{(i+1)}(1, d) \times \cdots \).

In order to couple theories at different \( i \)'s we have to introduce several more fields. For each pair \((i, i+1)\) there is a nonabelian gauge field \( A_{\mu}^{i,i+1}(x) \); the only requirement on the group \( G \) under which these fields transforms and the strength of the gauge coupling is that certain fermion condensates, to be discussed below, are induced. In addition, for each \( i \), we have two Weyl fermion fields. One, \( \psi^i(x) \), couples to \( A_{\mu}^{i,i+1}(x) \) as a fundamental under \( G \) while the other one, \( \chi^i(x) \), couples as an anti-fundamental under \( A_{\mu}^{i-1,i}(x) \). We assume that the \( SO^i(1, d) \times SO^{(i+1)}(1, d) \) symmetry is broken by a condensate

\[ \langle \psi^i(x) \chi^{i+1}(x) \rangle \sim f_G^{d-1} \exp \left[ \frac{i}{2} \Sigma_{ab} O_{i,i+1}^{ab}(x) \right]. \]  \hspace{1cm} (13)

\( O_{i,i+1}^{ab}(x) \) is an \( SO(1, d) \) Lorentz transformation matrix and \( f_G \) parameterizes the strength of the condensate. The low energy effective theory for the fields \( O_{i,i+1}^{ab}(x) \) is governed by the Lagrangian

\[ L_1 = 2 f_G^{d-1} \sum_i \epsilon^{\mu_1 \mu_2 \cdots \mu_d} \epsilon_{a_1 a_2 \cdots a_{d+1}} C_{i,i+1}^{a_1} (x) D_{\mu_1} C_{i,i+1}^{a_2} (x) \epsilon_{\mu_2}^{a_3} (x, i) \epsilon_{\mu_3}^{a_4} (x, i) \cdots \epsilon_{\mu_d}^{a_{d+1}} (x, i); \]  \hspace{1cm} (14)
the covariant derivative is
\[
D_\mu \mathcal{O}_{i,i+1}^{ab}(x) = \left[ \partial_\mu \mathcal{O}_{i,i+1}^{ab}(x) + \omega^a_{\mu n}(x,i) \mathcal{O}_{i,i+1}^{nb}(x) - \mathcal{O}_{i,i+1}^{an}(x) \omega^b_{\mu n}(x,i+1) \right] .
\] (15)

In the continuum limit we may expand \( \mathcal{O}_{i,i+1}^{ab}(x) \)
\[
\mathcal{O}_{i,i+1}^{ab}(x) = \eta^{ab} + \frac{1}{a} \omega^{ab}_{i,i+1}(x) + \cdots,
\] (16)
where \( a \) is the lattice separation. With the following identifications we recover the discrete version of \( L_B \) (eq. (6)):
\[
a = M^{d-2} / f_{G}^{d-1},
\]
\[
f_{G} = (d-1)M_{d+1},
\]
\[
M_{d}^{a} \lambda_{d} = (d+1)M_{d+1} \lambda_{d+1}.
\] (17)

III. EXTRA DISCRETE DIMENSIONS

Recently, extensive research has been carried out on the possibility that extra compact but large dimensions may account for the apparently large value of the Planck mass [6,7]. The present work shows how to formulate a discrete version of such schemes. In the continuum case phenomenology demands that we have more than one extra dimension. For simplicity we shall discuss only one extra “large” discrete dimension [8]. We envisage a four dimensional manifold with many \( SO(1,4) \) theories. The Lagrangian is the the sum of eq. (12) and (14) with \( d = 3, \ i = 0, 1, \ldots, N - 1 \), and periodic conditions on the discrete index \( i \), \( e_{\mu}^{a}(x, N) = e_{\mu}^{a}(x, 0), \ldots \). All other non-gravity fields appear only once and couple only to \( e_{\mu}^{a}(x, 0) \); in the continuum language this would indicate that these extra fields do not propagate into the extra dimension. The dynamical mechanism discussed in Sec. (II) generates a fifth dimension of circumference \( Na \). Using techniques similar to those discussed in ref. [1] we find that the potential for two masses coupled only to the \( i = 0 \) gravity is
\[
V(r) = \frac{m_1 m_2}{NM_{4}^2} \frac{1}{r} \sum_{m=0}^{N-1} \exp\left(-\frac{2r}{a} \sin \frac{m\pi}{N}\right) .
\] (18)
For \( r >> Na \) we recover the \( 1/r \) potential with an effective Planck mass \( M_{P} = NM_{4}^2 \).
REFERENCES


[3] A. Sugamoto, hep-ph/0104241, generated a four dimensional gravity from a three dimensional space by relying on the observation that $SO(4) = SO(3) \times SO(3)$. This technique cannot be extended to higher dimensions.

[4] Usually such a separation is done for the time coordinate where the spatial part of $e^a_t$ is referred to as the shift vector and $e^0_t$ is the lapse, see C. W. Misner, K. S. Thorne and J. A. Wheeler, *Gravitation* (Freeman, San Francisco, 1973). We refer to the whole last d-ad, $e^a_d$ as the shift vector.


[8] The simultaneous dynamical generation of more than one dimension presents fine-tuning problems.