A Vanishingly Small Vector Mass from Anisotropy of Higher Dimensional Spacetime

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We consider five-dimensional massive vector-gravity theory which is based on the foliation-preserving diffeomorphism and anisotropic conformal invariance. It does not have an intrinsic scale, and the only relevant parameter is the anisotropic factor $z$ which characterizes the degree of anisotropy between the four-dimensional spacetime and the extra dimension. We assume that the physical scale $M^*$ emerges as a consequence of spontaneous conformal symmetry breaking of the vacuum solution. We demonstrate that a very small mass for the vector particle compared to $M^*$ can be achieved with a relatively mild adjustment of the parameter $z$. At the same time, the motion along the extra dimension is observed to be highly suppressed and the five-dimensional theory can be effectively reduced to four-dimensional spacetime.

Keywords: Photon mass, Extra dimension, Vector-gravity theory, Conformal invariance, Anisotropy of spacetime

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I. INTRODUCTION

The research on the non-vanishing photon mass has a long history [1–3]. That Maxwell’s theory with the massless photon can be extended to a gauge-invariant massive Stueckelberg theory by introducing a scalar field that compensates for the gauge transformation of the vector field is well known. It preserves the unitarity and the renormalizability of the massless theory and describes the electrodynamics of the massive vector field [4]. In a particular gauge where the scalar field is set to zero, the theory reduces to the familiar Proca theory [5]. The deviation of the photon mass from zero, if any, must be extremely small because Maxwell’s theory has been tested to high experimental accuracy. The experimental constraints on the photon mass have considerably increased over the past several decades, putting stringent upper bounds on its mass. Recently, the cosmological implication of a massive photon was investigated in connection with dark energy, and the current acceleration was argued to be provided by a non-vanishing photon mass $m$ governed by the relation $\Lambda \sim m^2$ [6].

Up to now, a method of generating an on extremely small mass of vector particles, in general, has not existed and the mass has to be put in by hand if one starts in four-dimensional spacetime. This requires extreme fine tuning, and from a theoretical point of view, avoiding this would be aesthetically more appealing. One way of bypassing this problem is to consider a higher-dimensional theory in which gravity is coupled with five-dimensional Maxwell-Stueckelberg theory. However, there is a caveat here; the usual higher dimensional theories in which the spacetime is isotropic turn out not to provide a method the suppresses the photon mass to a high degree. The fine-tuning problem still remains. A possible remedy can be shown to necessitate a higher dimensional theory in which the four-dimensional spacetime and the extra dimensions are not treated on an equal footing [7]. On top of that, the theory requires two compatible symmetries of foliation-preserving diffeomorphism (FPD) and anisotropic conformal transformation (ACT). The FPD is implemented in the ADM decomposition of a higher dimensional metric by requiring the foliation-preserving diffeomorphism invariance be adapted to the extra dimensions, thus keeping the general covariance only for the four-dimensional spacetime. Conformal invariance can be incorporated with an extra (Weyl) scalar field and a real parameter $z$ which describes the degree of anisotropy of conformal transformations between the four-dimensional spacetime, and extra dimensional metrics. A cosmological test of $z$ was given, yielding an allowed range of the parameter $z$ [8].

In this paper, we construct five-dimensional Stueckelberg-like vector theory in an anisotropic spacetime background and couple it with gravity theory which has the symmetry of the aforementioned FPD and ACT. The essential point of the construction in the grav-
ity sector is the parameter \( z \) that describes the degree of anisotropy of the conformal transformation between the four-dimensional spacetime and the extra dimension. This construction can be extended to higher dimensional vector theory straightforwardly, and the coupled action describes five-dimensional vector-gravity theory which respects both FPD and ACT. The theory is so constructed as not to contain any intrinsic scale. Thus, the higher dimensional spacetime could be effectively reduced to a four-dimensional spacetime for a certain range of the parameter \( z \). This turns out to possible for a suitable adjustment of the parameter \( z \). The degree of fine-tuning is alleviated to a great extent. Another interesting consequence of the analysis is that the propagation of the massive vector field is alleviated to a great extent. Another interesting consequence of the analysis is that the propagation of the massive vector mass that is vastly suppressed compared to the try breaking scale.

This paper is organized as follows: In Sec. II, we construct the five-dimensional Stueckelberg-like vector-gravity theory with anisotropic conformal invariance. In Sec. III, we search for the Minkowski vacuum solution and consider the effective action of the massive vector field sector. A gauge fixing term is introduced and propagators are presented. Section IV contains conclusions and discussions.

II. MODEL

We start with a formulation of 5D anisotropic conformal gravity. The first part of this section is mostly redrawn from Refs. 7 and 8 to make the paper self-contained. Let us first consider the Arnowitt-Deser-Misner (ADM) decomposition of the five-dimensional metric:

\[
d s^2 = g_{\mu\nu}(dx^\mu + N^\mu dy)(dx^\nu + N^\nu dy) + N^2 dy^2. \tag{1}
\]

Then, the five-dimensional Einstein-Hilbert action without the cosmological constant can be expressed as

\[
S_{EH}^{(5)} = \int d^4x dy N\sqrt{-g} M_5^3 \left[ R - \{K_{\mu\nu} K^{\mu\nu} - K^2 \} \right], \tag{2}
\]

where \( M_5 \) is the five-dimensional gravitational constant, \( R \) is the spacetime curvature and \( K_{\mu\nu} \) is the extrinsic curvature tensor. \( K_{\mu\nu} = (\partial_\mu g_{\nu\rho} - \nabla_\mu N_\nu - \nabla_\nu N_\mu)/(2N) \).

The above action, Eq. (2), can be extended anisotropically by breaking the five-dimensional general covariance down to its foliation-preserving diffeomorphism symmetry given by

\[
x^\mu \rightarrow x'^\mu \equiv x'^\mu(x, y), \quad y \rightarrow y' \equiv y'(y), \tag{3}
\]

\[
g_{\mu\nu}(x', y') = \left( \frac{\partial x'^\rho}{\partial x^\mu} \right) \left( \frac{\partial x^{\nu'}}{\partial x^\rho} \right) g_{\rho\sigma}(x, y), \tag{4}
\]

\[
N'^{\mu}(x', y') = \frac{\partial y^\rho}{\partial y'} N^{\rho}(x, y) - \frac{\partial x'^\mu}{\partial y'} \tag{5}
\]

and its non-uniform conformal transformations given by

\[
g_{\mu\nu} \rightarrow e^{2\lambda(z, y)} g_{\mu\nu}, \quad N \rightarrow e^{z\omega(z, y)} N, \tag{6}
\]

\[
N^\mu \rightarrow N^\mu, \quad \phi \rightarrow e^{-\frac{3z}{2} \omega} \phi, \tag{7}
\]

where a Weyl scalar field \( \phi \) to compensate for the conformal transformation of the metric is introduced. In Eq. (7), a factor \( z \) is introduced in the transformation of \( N \equiv g_{00} \), which characterizes the anisotropy of the spacetime and the extra dimension. The anisotropic Weyl action invariant under Eqs. (3)–(7) for an arbitrary \( z \) can be written as

\[
S_{CG}^{(5)} = \int d^4x dy N \sqrt{-g} \left[ R - \frac{12}{z + 2} \frac{\nabla_\mu \nabla_\nu \phi}{\phi} + \frac{12z}{(z + 2)^2} \frac{\nabla_\mu \phi \nabla_\nu \phi}{\phi^2} \right] - \beta_1 \phi \frac{2iz + 6z}{z + 2} \left\{ B_{\mu\nu} B^{\mu\nu} - \lambda B^2 \right\} + \beta_2 \phi^2 C_\mu C^\mu. \tag{8}
\]

\[
C_\mu = \frac{\partial_\mu N}{N} + \frac{2z - 2}{z + 2} \frac{\partial_\mu \phi}{\phi}. \tag{9}
\]

Several comments are in order. Note that, in accordance with conformal symmetry, the above action in Eq. (8) is built so as not to contain any intrinsic scale. The scale \( M_5 \) of Eq. (2) is supposed to emerge as con-
sequence of spontaneous conformal symmetry breaking. We consider only the case \( z \neq -2 \) because \( \phi \) is not affected under the conformal transformation in Eq. (7) in this case. For \( z = -2 \), an anisotropic scale invariant gravity theory can actually be constructed without the need for the field \( \phi \). In this work, we are interested in Minkowski vacuum solution, and we did not include the potential term in the action in Eq. (8). The isotropic case with \( \beta_1 = \lambda = z = 1 \) and \( \beta_2 = 0 \) leads to five-dimensional Weyl gravity with a zero potential for the field \( \phi \) \([7,8]\). In the anisotropic case, the action in Eq. (8) is, in general, plagued with the perturbative ghost instability coming from the breaking of the full general covariance of 5D. However, this problem can be shown to be cured by constraining the constants \( \beta_1 \) and \( \beta_2 \), especially with \( 0 < \beta_2 < 3/2 \) \([7]\).

We couple the action in Eq. (8) with the five-dimensional conformal vector field theory given by

\[
S_{MA}^{(5)} = \int d^4 x d y \sqrt{-g} N \phi ^{z/2} \left[ -\frac{1}{4} g ^{\mu \nu } g ^{\sigma \rho } F _{\mu \rho } F _{\nu \sigma } - \frac{1}{2} Z_1 N ^{-2} g ^{\mu \nu } \hat{F}_{5 \mu } \hat{F}_{5 \nu } - \frac{1}{2} Z_2 g ^{\mu \nu } ( A _{\mu } + \nabla _{\mu } \sigma ) ( A _{\nu } + \nabla _{\nu } \sigma ) - \frac{1}{2} Z_3 N ^{-2} ( \hat{A}_5 + \hat{\nabla }_5 \sigma ) ( \hat{A}_5 + \hat{\nabla }_5 \sigma ) \right],
\]

(11)

where \( F _{\mu \nu } = \nabla _{\mu } A _{\nu } - \nabla _{\nu } A _{\mu } = - F _{\nu \mu } \). \( \hat{F}_{5 \mu } = \partial y _{\mu } A _{5} - \partial x _{\mu } \hat{A}_5 \). \( \hat{F}_{5 \mu } = F _{5 \mu } + N ^{\nu } F _{5 \nu } \). \( \hat{\sigma} = \partial y _{\mu } \hat{\sigma} - \partial x _{\mu } \hat{\sigma} \), and \( \hat{A}_5 = A _{5} - N _{\mu } A _{\mu } \). \( \hat{\nabla }_5 \sigma = \partial y _{5} \sigma \). For anisotropic conformal symmetry, the \( Z \)’s are given by

\[
Z_1 = \phi ^{4 + c/2}, \quad Z_2 = \phi ^{2 + c}, \quad Z_3 = \phi ^{4 + 2 + c/2}. \quad (12)
\]

Note that a simple relation exists among \( Z_1, Z_2, \) and \( Z_3 \); \( Z_1 Z_2 = Z_3 \). The action in Eq. (11) respects the following three symmetries: The first is foliation-preserving diffeomorphism; along with Eqs. (3)–(6),

\[
A_\mu (x', y') = \left( \frac{\partial x' }{\partial x } \right) A_\mu (x, y), \quad A_5 (x', y') = \left( \frac{\partial y }{\partial y } \right) A_5 (x, y) + \left( \frac{\partial x' }{\partial y } \right) A_\mu , \quad \sigma' (x', y') = \sigma (x, y). \quad (13)
\]

Under Eq. (13), one can check that

\[
\hat{A}_5 (x', y') = \left( \frac{\partial y }{\partial y } \right) \hat{A}_5 (x, y), \quad \hat{F}_{5 \mu } (x', y') = \left( \frac{\partial x' }{\partial x } \right) \left( \frac{\partial y }{\partial y } \right) \hat{F}_{5 \mu } (x, y). \quad (14)
\]

The second is anisotropic conformal symmetry; along with Eq. (7), with \( A_\mu, A_5, \) and \( \sigma \) being unaffected. The third is gauge symmetry;

\[
A_\mu \rightarrow A_\mu + \nabla _{\mu } \Lambda, \quad A_5 \rightarrow A_5 + \nabla _5 \Lambda, \quad \sigma \rightarrow \sigma - \Lambda. \quad (15)
\]

In the limit \( \phi = \phi _0 = 1 \), all \( Z \)’s are equal to 1. In this case, the action in Eq. (11) reduces to five-dimensional Stueckelberg theory given by \([9]\)

\[
S_{\text{Stueckelberg}}^{(5)} = \int d^4 x \sqrt{\hat{g}^{(5)}} \left[ -\frac{1}{4} G^{MP} G^{NQ} F_{MN} F_{PQ} - \frac{1}{2} G^{MN} ( A_M + \nabla _M \sigma ) ( A_N + \nabla _N \sigma ) \right],
\]

(16)

with unit mass for the vector field. \( G^{MN} \) is the inverse of \( G_{MN} \) given in Eq. (1). The dependence on the parameter \( z \) disappears, and the five-dimensional diffeomorphism invariance is respected.

The coupled action \( S_{MA}^{(5)} + S_{\text{Stueckelberg}}^{(5)} \) does not contain any dimensional parameter and exhibits the following scaling symmetry:

\[
x \rightarrow b^{-1} x, \quad y \rightarrow b^{-z} y, \quad \phi \rightarrow b^{2+ z} \phi, \quad N^\mu \rightarrow b^{-1} N^\mu, \quad A_\mu \rightarrow b A_\mu, \quad A_5 \rightarrow b^2 A_5, \quad \sigma \rightarrow \sigma.
\]

(17)

and the engineering dimensions are given as

\[
[x] = -1, \quad [y] = -z, \quad [\phi] = \frac{2 + z}{2}, \quad [N^\mu] = z - 1, \quad [A_\mu] = 1, \quad [A_5] = z, \quad [\sigma] = 0. \quad (18)
\]

We note that some of the quantities must carry \( z \)-dependent scaling dimensions in order to respect the symmetries. The anisotropic conformal invariance also determines the conformal weights of each field given by

\[
[g_{\mu \nu}]_c = 2, \quad [N]_c = z, \quad [\phi]_c = \frac{2 + z}{2}, \quad [N^\mu]_c = 0, \quad [A_\mu]_c = [A_5]_c = [\sigma]_c = 0. \quad (19)
\]

We assume that a scale appears as a consequence of spontaneous conformal symmetry breaking of the vacuum solution, and the canonical dimensions of the fields and spacetime are recovered.
III. VACUUM SOLUTION

In this section, we search for the solutions of the equation of motion derived from \( S^{(5)}_{\text{CG}} + S^{(5)}_{\text{MA}} \). First, we introduce a scale \( M_* \), which sets the scale for the conformal symmetry breaking of the vacuum solution. We also assume that \( y, \phi, A_\mu, A_5 \), and \( \sigma \) recover their canonical dimensions with the vacuum solution; therefore, we re-scale them via

\[
y \to M_*^{-z+1} y, \quad \phi \to M_*^{z+1} \phi, \quad A_\mu \to M_*^{-\frac{1}{2}} A_\mu, \\
A_5 \to M_*^{-\frac{3}{2}} A_5, \quad \sigma \to M_*^{-\frac{3}{2}} \sigma,
\]

which yields the following effective action:

\[
S = \int d^4x dy \left[ -\frac{1}{4} F^2 - \frac{1}{2} Z_1(0) F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} Z_2(0) M_*^2 \left( A_\mu + \frac{1}{M_*} \nabla_\mu \sigma \right) \left( A^\mu + \frac{1}{M_*} \nabla^\mu \sigma \right) \\
- \frac{1}{2} Z_3(0) M_*^2 \left( \eta + \frac{1}{M_*} \nabla_5 \sigma \right) \left( \eta + \frac{1}{M_*} \nabla^5 \sigma \right) \right],
\]

where the constant \( \phi_0^{2z/(z+2)} \) has been reabsorbed into \( A_\mu, \eta, \) and \( \sigma \). Moreover, \( Z_i(0) \equiv Z_i(\phi = \phi_0) \) \( (i = 1, 2, 3) \) of Eq. (12). Using the gauge invariance of the action in Eq. (22), we add a gauge-fixing term of the form

\[
S_{\text{GF}} = \int d^4x dy \left[ -\frac{1}{2} \zeta_1 \partial_\mu A^\mu + \zeta_2 Z_1(0) \partial_5 \eta + \zeta_3 Z_2(0) M_* \sigma \right]^2.
\]

In order to cancel the \( A_\mu, \sigma \) and the \( \eta, \sigma \) mixing terms in Eq. (22) which are first-order in derivative, we choose \( \zeta_1 = \zeta_2 = \zeta_3 = 1 \equiv \zeta \). Adding to the action in Eq. (22) yields the following effective action:

\[
S_{\text{eff}} = M_* \int d^4x dy \left[ -\frac{1}{2} A_M K^{MN} A_N + \frac{1}{2} \sigma L_\sigma \right],
\]

where

\[
K^{\mu\nu} = -\left[ \left( \Box + Z_1(0) \right) \partial^2 - Z_2(0) M_*^2 \right] \eta^{\mu\nu} - \left( 1 - \xi^2 \right) \partial^\mu \partial^\nu, \\
K^{\mu5} = K^{5\mu} = Z_1(0) \left( 1 - \xi^2 \right) \partial^\mu \partial^5, \\
K^{55} = -Z_1(0) \left[ \Box + \xi^2 Z_1(0) \partial_5^2 - Z_2(0) M_*^2 \right],
\]

and \( L = Z_2(0) \left[ \Box + Z_1(0) \partial_5^2 - \xi^2 Z_2(0) M_*^2 \right] \). Here, we used the relation \( Z_i(0) Z_2(0) = Z_i(0) \).

Calculation of the inverse matrix \( K^{MN} \) gives the following propagator in momentum space (with \( k_5 \equiv \sqrt{Z_1(0) k_5} \), \( M_* \equiv \sqrt{Z_2(0) M_*} \)):

\[
G_{\mu\nu} = \frac{1}{k^2 + k_5^2 + M_*^2} \left[ \eta^{\mu\nu} - \frac{\xi^2 - 1}{\left( \xi^2 k^2 + \xi^2 k_5^2 + M_*^2 \right)} k_\mu k_\nu \right],
\]

\[
G_{\mu\eta} = \frac{1 - \xi^2}{\left( \xi^2 k^2 + \xi^2 k_5^2 + M_*^2 \right)} k_\mu k_5, \\
G_{\eta\eta} = \frac{\xi^2 k^2 + k_5^2 + M_*^2}{\left( \xi^2 k^2 + \xi^2 k_5^2 + M_*^2 \right)} k_5^2, \\
G_{\sigma\sigma} = \frac{1}{\left( k^2 + k_5^2 + \xi^2 M_*^2 \right)}.
\]

Here, we rescaled \( \eta \to \eta/\sqrt{Z_1(0)} \) and \( \sigma \to \sigma/\sqrt{Z_2(0)} \) to recover the canonical form for the kinetic energy. If we take \( \xi = 0 \), we find that Eqs. (26)–(28) combine into (with \( p_M = (k_\mu, k_5) \))

\[
G_{MN} = \frac{1}{p^2 + M_*^2} \left[ \eta_{MN} + \frac{\eta_{MP} \eta_{PN}}{M_*^2} \right],
\]

which describes the five-dimensional Proca field with mass \( M_*^2 \), \( G_{\sigma\sigma} \) decoupled; and this is the unitary gauge. Note, however, that the conjugate momentum of the extra coordinate \( y \) is \( k_5 \) and that the dispersion relation \( p^2 = k^2 + Z_1(0) k_5^2 \) suggests that the motion along the fifth-direction is suppressed by a factor of \( \sqrt{Z_1(0)} \). An extremely small value of \( Z_1(0) \) could effectively localize the motion to four-dimensional spacetime. For example,
with \( \phi_0 = 10 \) and \( z = -2.1 \), \( Z_4(0) \) of Eq. (12) becomes as small as \( 10^{-124} \). In this case, the mass of the vector field is \( M^2_\sigma = 10^{-40}M^2 \). Recently, the value of \( z \) has been tested in comparison with cosmological data, and a range of values of the parameter \( z \) that can address the current dark energy density compared to the Planck energy density was given [8]. The upshot is that a tiny number can be generated with the anisotropic factor \( z \), which does not require extreme fine-tuning. The propagators become particularly transparent in the Feynman gauge with \( \xi = 1 \). In this case, \( G_{\mu\nu} = 0 \) and the gauge field \( A_\mu \) decouples from the scalar field \( \eta \). All fields propagate with the same mass \( M^2_\sigma \), and \( \eta \) supplies longitudinal polarization to \( A_\mu \) field. The choice \( \xi = \infty \) gives a transverse propagator in the sense that \( p^M G_{MN} = 0 \). The mass term of the \( \sigma \) field disappears.

IV. CONCLUSION AND DISCUSSIONS

In this paper, we constructed five-dimensional massive vector-gravity theory with an anisotropic conformal invariance and demonstrated the possibility of obtaining a vector particle with a very slight mass compared with the physical mass scale which emerges as a consequence of spontaneous conformal symmetry breaking. At the same time, the built-in anisotropy can suppress highly the motion along the extra dimension, effectively localizing it to four spacetime.

This approach can be contrasted to the mainstream research on the theory of higher spacetime dimensions which treats the four-dimensional spacetime and the extra dimensions on an equal footing. The isotropic spacetime is more appealing from the aesthetical point of view, and we note that the anisotropic approach can be reconcile if we consider the flow of the anisotropic factor \( z \). For example, the theory starts from \( z \neq 1 \) at high energy, but it relaxes to \( z = 1 \) at low energy with a full spacetime symmetry. These aspects were considered in the critical behavior of gravity [10,11].

The five-dimensional operator \( \Box + Z_4 \partial_5^2 - Z_2 M^2 \) with a \( Z_4(< 1) \) factor suppresses propagation along the fifth dimension. For example, the static Coulomb potential with \( M_* = 0 \) gives

\[
V(r, y) \sim \int d^3 k e^{ik y} \int \frac{e^{ikr}}{k^2 + Z_1 k^2} \sim \frac{\sqrt{Z_1 r}}{r(Z_1 r^2 + y^2)}. \tag{31}
\]

In the isotropic case with \( Z_1 = 1 \), the above expression reduces to the four space-dimensional Coulomb potential. On the other hand, in the limit \( Z_1 \to 0 \), it becomes

\[
V(r, y) \to \frac{1}{r} \delta(y), \tag{32}
\]

completely obviating the possible motion along the fifth dimension. Therefore, the last expression in Eq. (31) naturally interpolates between the four and five dimensions. It gives an interesting possibility that for a very small value of \( Z_1 \), the effective motion can be practically considered to be four-dimensional. This can render the extra dimension almost completely obsolete, making its presence fall into desuetude.

In effective four-dimensional theory, the size \( l_E \) of the extra dimension can be shown to satisfy the following relation from Eq. (8) after the conformal symmetry breaking:

\[
l_E = l_* \phi_0^{-2} \left( \frac{M_{pl}}{M_*} \right)^2, \tag{33}
\]

where \( l_* \) is the size associated with the symmetry breaking scale \( M_* \) and \( M_{pl} \) is the four-dimensional Planck scale. When \( M_* \) is close to the Planck scale, \( l_E \sim l_* \sim l_{pl} \) for \( \phi_0 \sim \mathcal{O}(1) \). For a value of \( \phi_0 \sim 10^{-7} \), \( M_* \sim 10^3 \text{ GeV} \) yields \( l_E \) to be on the order of a millimeter. A possible connection with the localization in the brane physics [12] deserves further investigation.

We conclude with a remark on the possible application to the physics of gravitational waves [13]. Recently, whether the gravitational wave is leaking into the extra dimension, was investigated and the analysis was found to negate the existence of the extra dimension [14,15]. See also Refs. 16, 17 in which the question of the observed gravitational wave leaking into the extra dimension was investigated in the brane-world model. In the anisotropic approach, the gravitational wave equation is governed by the anisotropic five-dimensional d’Alembertian where the isotropic operator is replaced with \( \Box + Z_4 \partial_5^2 \), where \( Z_4 \) can be given as \( Z_4(\phi_0) \propto \phi_0^{(2(z-1))/(z+2)} \) in Eq. (8). This could again suppress the gravitational wave along the extra dimensions and the effect is the same as to keep the wave from leaking into the extra dimension. Note that \( Z_4(\phi_0), Z_2(\phi_0), \) and \( Z_1(\phi_0) \) all can produce a very small number near \( z < -2 \) and \( \phi_0 \sim 10 \), or \( z > -2 \) and \( \phi_0 \sim 0.1 \). Detailed comparisons with observation would be interesting.

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