Higher dimensional gravity invariant under the AdS group

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A higher dimensional gravity invariant both under local Lorentz rotations and under local Anti de Sitter boosts is constructed. It is shown that such a construction is possible both when odd dimensions and when even dimensions are considered. It is also proved that such actions have the same coefficients as those obtained in ref. [3].

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The most general lagrangian for gravity in d dimensions built up on the same principles as General Relativity (general covariance, second order equations for the metric, and no explicit torsion) is a polynomial of degree \([d/2]\) in the curvature known as the Lanczos-Lovelock lagrangian. Lanczos-Lovelock (LL) theories have the same metric, and no explicit torsion) is a polynomial of degree \([d/2]\) in the curvature known as the Lanczos-Lovelock lagrangian. Lanczos-Lovelock (LL) theories have the same non-linear realizations to gravity, permits constructing an action for Lanczos-Lovelock gravity theory genuinely invariant under the AdS group. It is shown that such a construction is possible both when odd dimensions and when even dimensions are considered. Applications of the theory of non-linear realizations to gravity have been carried out in different ways in previous research, such as for example in ref. [7] where the vierbein field was considered as a Goldstone field related to a nonlinear realization of the group \(GL(4, R)\), or of the affine and conformal groups. In the present work, the Goldstone fields represent a point in an internal AdS space.

The actions invariant under the AdS group are constructed using the vierbein and spin connection 1-forms obtained in ref. [2]. It is also proved that such actions have the same coefficients as those obtained in refs. [2, 8].

The non-linear realizations in de Sitter space can be studied by the general method developed in ref. [2].

\[ S = \int \sum_{p=0}^{[d/2]} \alpha_p L^{(p)} \]

where \(\alpha_p\) are arbitrary constants and \(L^{(p)}\) is given by

\[ L^{(p)} = \varepsilon_{a_1 a_2 \ldots a_p} R^{a_1 a_2} \ldots R^{a_{2p-1} a_{2p}} e^{a_{2p+1}} \ldots e^{a_d}. \]

One important drawback of the Lanczos-Lovelock action is the appearance of a number of dimensionful constants which are not determined from first principles. In ref. [3] it is shown that requiring that the equations of motion uniquely determine the dynamics for as many constants which are not determined from first principles. In ref. [3] it is shown that requiring that the equations of motion uniquely determine the dynamics for as many components of the independent fields as possible fixes the \(\alpha_p\) coefficients (for even as well as for odd dimensions) in terms of the gravitational and cosmological constants.

For \(d = 2n\) the coefficients are

\[ \alpha_p = \alpha_0 (2\gamma)^p \binom{n}{p}, \]

and the action takes a Born-Infeld-like form. With these coefficients, the LL action is invariant only under local Lorentz rotations.

For \(d = 2n - 1\), the coefficients become

\[ \alpha_p = \frac{\alpha_0 (2n-1)(2\gamma)^p}{(2n-2p-1)} \binom{n-1}{p}, \]

where

\[ \alpha_0 = \frac{k}{d!d-1}, \quad \gamma = -\text{sgn}(\Lambda) \frac{l^2}{2}, \]

and, for any dimension \(d, l\) is a length parameter related to the cosmological constant by \(\Lambda = \pm (d-1)(d-2)/2l^2\). With these coefficients [3], the vierbein and the spin connection may be accommodated into a connection for the AdS group, allowing for the lagrangian to become the Chern-Simons form in \(d = 2n + 1\) dimensions, whose exterior derivative is the Euler topological invariant in \(d = 2n \) dimensions,

\[ dL^{(2n-1)} = \frac{k l}{2n} \varepsilon_{A_1 \ldots A_{2n}} R^{A_1 A_2} \ldots R^{A_{2n-1} A_{2n}} \]

with

\[ R^{A B} = \left( R^{a b} + \frac{1}{2} e^{a A} e^{b B} \frac{l^2}{2l^2} T^A \right). \]
Following these references, we consider a Lie group \( G \) and its stability subgroup \( H \).

The Lie group \( G \) has \( n \) generators. Let us call \( \{ V_i \}_{i=1}^{n-d} \) the generators of \( H \). We shall assume that the remaining generators \( \{ A_l \}_{l=1}^d \) are chosen so that they form a representation of \( H \). In other words, the commutator \([V_i,A_l]\) should be a linear combination of \( A_l \) alone. A group element \( g \in G \) can be uniquely represented in the form

\[
g = e^{\xi^i A_i} h \quad (8)
\]

where \( h \) is an element of \( H \). The \( \xi^i \) parametrize the coset space \( G/H \). We do not specify here the parametrization of \( h \). One can define the effect of a group element \( g_0 \) on the coset space by

\[
g_0 g = g_0 (e^{\xi^i A_i} h) = e^{\xi'^i A_i} h'
\]

or

\[
g_0 e^{\xi^i A_i} = e^{\xi'^i A_i} h_1
\]

where

\[
h_1 = h' h^{-1}
\]

\[
\xi' = \xi'(g_0, \xi)
\]

\[
h_1 = h_1(g_0, \xi).
\]

If \( g_0 - 1 \) is infinitesimal, \( 10 \) implies

\[
e^{-\xi^i A_i} (g_0 - 1) e^{\xi^i A_i} - e^{-\xi' A_i} \delta e^{\xi^i A_i} = h_1 - 1.
\]

The right-hand side of \( 14 \) is a generator of \( H \).

Let us first consider the case in which \( g_0 = h_0 \in H \).

Then \( 10 \) gives

\[
e^{\xi'^i A_i} = h_0 e^{\xi^i A_i} h_0^{-1}.
\]

Since the \( A_i \) form a representation of \( H \), this implies

\[
h_1 = h_0; \quad h' = h_0 h.
\]

The transformation from \( \xi \) to \( \xi' \) given by \( 15 \) is linear. On the other hand, consider now

\[
g_0 = e^{\xi^i A_i}.
\]

In this case, eq. \( 10 \) becomes

\[
e^{\xi^i A_i} e^{\xi^i A_i} = e^{\xi'^i A_i} h.
\]

This is a non-linear inhomogeneous transformation for \( \xi \). The infinitesimal form of \( 11 \) is

\[
e^{-\xi^i A_i} \xi_0^i A_i e^{\xi^i A_i} - e^{-\xi^i A_i} \delta e^{\xi^i A_i} = h_1 - 1.
\]

The left-hand side of this equation can be evaluated, using the algebra of the group. Since the results must be a generator of \( H \), one must set equal to zero the coefficient of \( A_I \). In this way one finds an equation from which \( \delta \xi^i \) can be calculated.

The construction of a Lagrangian invariant under local group transformations requires the introduction of a set of gauge fields \( a = a^\mu_{\rho} A_l dx^\mu \), \( \rho = \rho^\mu_{\rho} V_l dx^\mu \), \( p = p^\mu_{\rho} A_l dx^\mu \), \( v = v^\mu_{\rho} V_l dx^\mu \), associated with the generators \( V_l \) and \( A_l \), respectively. Hence \( \rho + a \) is the usual linear connection for the gauge group \( G \), and therefore its transformation law under \( g \in G \) is

\[
g : (\rho + a) \rightarrow (\rho' + a') = \left[ g(\rho + a)g^{-1} - \frac{1}{f} (dg)g^{-1} \right] (20)
\]

where \( f \) is a constant which, as it turns out, gives the strength of the universal coupling of the gauge fields to all other fields.

We now consider the Lie algebra valued-differential 1-forms \( p \) and \( v \) defined by \( 18 \)

\[
e^{-\xi^i A_i} [d + f(\rho + a)] e^{\xi^i A_i} = p + v.
\]

The transformation laws for the forms \( p(\xi, d\xi) \) and \( v(\xi, d\xi) \) are easily obtained. In fact, using \( 17, 18 \) one finds \( 21 \)

\[
p' = h_1 p(h_1)^{-1}
\]

\[
v' = h_1 v(h_1)^{-1} - (dh_1) h_1^{-1}.
\]

Eq. \( 22 \) shows that the differential forms \( p(\xi, d\xi) \) and \( v(\xi, d\xi) \) are transformed linearly by a group element of the form \( 17 \); the former as a tensor and the latter as a connection. The transformation law is the same as that by an element of \( H \), except that now this group element \( h_1(\xi, \xi) \) is a function of the variable \( \xi \). Therefore, any \( H \)-invariant expression, in any dimensions, written with \( \rho + a \) and \( v \), will be \( G \)-invariant if these fields are changed by \( p \) and \( v \), respectively.

We have specified the fields \( p \) and \( v \) as well as their transformation properties, and now we make use of them to define the covariant derivative with respect to the group \( G \):

\[
D = d + v.
\]

The corresponding components of the two-form curvature are

\[
T = Dp
\]

\[
R = dv + vv.
\]

When \( G \) is the AdS lie algebra

\[
[P_a, P_b] = -iam^2 J_{ab}
\]
\[ [J_{ab}, P_c] = i (\eta_{ac} P_b - \eta_{bc} P_a) \quad (28) \]
\[ [J_{ab}, J_{cd}] = i (\eta_{ac} J_{bd} - \eta_{bc} J_{ad} + \eta_{bd} J_{ac} - \eta_{ad} J_{bc}) \quad (29) \]

having as generators \( P_a, J_{ab} \) and the subalgebra \( H \) is the Lorentz algebra \( SO(3, 1) \) with generators \( J_{ab} \), then the equation (21) becomes

\[
\frac{1}{2} i W^{ab} J_{ab} - i V^a P_a = e^{i\xi^a P_a} \left[ d + \frac{1}{2} i \omega^{ab} J_{ab} - i e^a P_a \right] e^{-i\xi^b P_b}. \quad (30)
\]

Using the AdS algebra, we arrive at explicit expressions for \( V^a \) and \( W^{ab} \) which are the basis for the Stelle-West formalism:

\[
V^a = e^a + (\cosh z - 1) \left( \delta^a_b - \frac{\xi_b^a}{\xi^2} \right) e^b + \frac{\sinh z}{z} D\xi^a - \left( \frac{\sinh z}{z} - 1 \right) \left( \frac{\xi^a d\xi^b - \xi^b d\xi^a}{\xi^2} \right). \quad (31)
\]

\[
W^{ab} = \omega^{ab} - \frac{1}{l^2} \frac{\sinh z}{z} \left( \xi^a e^b - \xi^b e^a \right) - \frac{1}{l^2} \left[ \xi^a D\xi^b - \xi^b D\xi^a \right] \left( \frac{\cosh z - 1}{z^2} \right). \quad (32)
\]

with \( z = \frac{1}{4} (\cosh \xi_a)^{1/2} = \frac{1}{4} \xi_a \). Only in the so-called "physical" gauge \( \{ \} \), where \( \xi^a = 0 \), we have \( V^a = e^a \) and \( W^{ab} = \omega^{ab} \). In this gauge the resultant theory is invariant only under the Lorentz group. There is however, an exceptional case which occurs when the odd-dimensional case is considered and when the coefficients are appropriate choices. We shall comment on this below.

On the other hand, under an infinitesimal AdS boost,

\[
\delta\xi^a = \rho^a + \left( \frac{z \cosh z}{\sinh z} - 1 \right) \left( \rho^a - \rho^b \frac{\xi^a \xi^b}{\xi^2} \right) \quad (33)
\]

\[
\delta V^a = \kappa^a_b V^b \quad (34)
\]
\[
\delta W^{ab} = - D\kappa^{ab} \quad (35)
\]

where

\[
\kappa^{ab} = - \frac{1}{l^2} \frac{\cosh z - 1}{z \sinh z} \left( \xi^a \rho^b - \xi^b \rho^a \right). \quad (36)
\]

Then, under the AdS group, \( V^a \) and \( W^{ab} \) behave as a Lorentz vector and connection respectively, but with a non-linear parameter.

Now we proceed to apply the above formalism to build a gauge theory of gravity in any number of dimensions.

The vielbein and the spin connection can be rewritten in the form

\[
V^a = O^a_b e^b + Dz \xi^a, \quad (37)
\]

\[
W^{ab} = \left[ \delta^a_c \delta^b_d - \frac{2}{l^2} \left( \cosh z - 1 \right) \delta^{ab} \xi^f \xi_c \right] \omega^{fg}, \quad (38)
\]

where

\[
O^a_b \equiv \cosh z \delta^a_b - (\cosh z - 1) \frac{\xi^a \xi^b}{\xi^2} - \frac{1}{l^2} \delta^{ab} \xi^f \xi_f, \quad (39)
\]

\[
Dz = \frac{\sinh z z}{z} D\xi^a - \left( \frac{\sinh z}{z} - 1 \right) \frac{dz}{z}. \quad (40)
\]

It is interesting to note that the inverse of the operator \( O^a_b \) is given by

\[
(O^a_b)^{-1} = \frac{1}{\cosh z} \delta^a_b - \left( \frac{1}{\cosh z} - 1 \right) \frac{\xi^a \xi^b}{\xi^2}. \quad (41)
\]

Under the transformations \( e^a \rightarrow e^a + \delta e^a \), \( \omega^{ab} \rightarrow \omega^{ab} + \delta\omega^{ab} \), the vielbein and the connection change as

\[
\delta_e V^a = O^a_b \delta e^b, \quad (42)
\]

\[
\delta_\omega V^a = \frac{\sinh z}{z} \xi^b \delta\omega^b, \quad (43)
\]

\[
\delta_e W^{ab} = \frac{1}{l^2} \frac{\sinh z}{z} \delta^{ab} \xi^f \delta e^g, \quad (44)
\]

\[
\delta_\omega W^{ab} = \delta \omega^{ab} - (\cosh z - 1) \delta^{ab} \xi^f \xi^g \xi_f \xi_g. \quad (45)
\]

The action (11) can now be written in the form

\[
S = \int \sum_{p=0}^{k} \alpha_p \xi_{a_1 a_2 \ldots \cdot a_d} R^{a_1 a_2 \ldots \cdot a_d} \ldots \ldots R^{a_{2p-1} a_{2p} \ldots \cdot a_{2p+1} \ldots \cdot a_d}, \quad (46)
\]
where now
\[ R^{ab} = dW^{ab} + W^c_{\, \, \, c} W^{cb}. \] (47)

The space-time torsion \( T^a \) is given by
\[ T^a = DV^a \] (48)
where \( D \) is the covariant derivative in the connection \( W^{ab} \).

This action (46) is invariant under general coordinate transformations and under \( AdS \) transformations (34), (35). The interesting result is that the action (46) is invariant under \( AdS \) transformations, not only for the odd-dimensional case \( d = 2n - 1 \), but also for the even-dimensional case \( d = 2n \).

Now we consider the variations of the action with respect to \( \xi^a \), \( e^a \), \( \omega^{ab} \). The variations of the action (46) with respect to \( e^a \), \( \omega^{ab} \) lead to the following equations:

\[
\sum_{p=0}^{[d-1]/2} \frac{2 \, \text{senh} \, z}{l^2 z} p(d-2p)\alpha_p \varepsilon_{a_1 \ldots a_d} \xi^{a_1} T^{a_2} \]
\[ \times R^{a_3 a_4} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} \]
\[ + \sum_{p=0}^{[d-1]/2} (d-2p)\alpha_p \varepsilon_{a_1 \ldots a_{d-1}} f R^{a_1 a_2} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} \xi_{a_{d-1}} = 0. \] (49)

These equations reproduce the equations of motion of Lanczos-Lovelock gravity theory. In fact, taking the product of Eq. (49) with \((O^2_n)^{-1}\), and of Eq. (50) with \( \xi^{a_{d-1}} \), we obtain

\[
\sum_{p=0}^{[d-1]/2} \frac{2 \, \text{tanh} \, z}{l^2 z} p(d-2p)\alpha_p \varepsilon_{a_1 \ldots a_{d-1}} \xi^{a_1} T^{a_2} \]
\[ \times R^{a_3 a_4} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} \]
\[ + \sum_{p=0}^{[d-1]/2} (d-2p)\alpha_p \varepsilon_{a_1 \ldots a_{d-1}} f R^{a_1 a_2} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} = 0. \] (51)

Taking the addition of Eq. (51) to Eq. (52), we have

\[
\sum_{p=0}^{[d-1]/2} (2 \, \text{coth} \, z - 3) \frac{p(d-2p)\alpha_p \varepsilon_{a_1 \ldots a_d} \xi^{a_1} T^{a_2}}{l^2 \text{senh} \, z} \]
\[ \times R^{a_3 a_4} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} \]
\[ - \sum_{p=0}^{[d-1]/2} (d-2p)\alpha_p \varepsilon_{a_1 \ldots a_d} f R^{a_1 a_2} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} = 0. \] (52)

and therefore we can write

\[
\sum_{p=0}^{[d-1]/2} (d-2p)\alpha_p \varepsilon_{a_1 \ldots a_d} f R^{a_1 a_2} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}} V^{a_{2p+1}} \ldots V^{a_{d-1}} = 0, \] (53)

From (51), (52) and (53) we obtain:

\[
\sum_{p=0}^{[d-1]/2} \frac{p(d-2p)\alpha_p \varepsilon_{a_{d-1} \ldots a_d} R^{a_3 a_4} \ldots \ldots \ldots R^{a_{2p-1} a_{2p}}}{l^2 \text{senh} \, z} \]
\[ \times V^{a_{2p+1}} \ldots V^{a_{d-1}} V^f \xi_{a_{d-1}} = 0. \] (54)
These equations have the same form as those obtained for Lanczos-Lovelock theory with the usual fields $e^a$ and $\omega^{ab}$ replaced by $V^a$ and $W^{ab}$. The field equation corresponding to the variation of the action with respect to $\xi^a$, is not an independent equation. In fact, taking the covariant derivative operator $D$ of equation (10), we obtain the same equation that one obtains by varying the action with respect to $\xi^a$. This is to be expected, since the Goldstone field $\xi^a$ has no dynamical degrees of freedom.

Following the same procedure of ref. 53, one can see that the equations (57), lead, in the even-dimensional case, to the coefficients given by eq. (3) and, in the odd-dimensional case, to the coefficients given by eq. (4).

Therefore, the use of the SW-formalism does not change the coefficients $\alpha_p$ of the action already obtained in refs. 7, 8. In these refs, can be found a more detailed discussion of the coefficients $\alpha_p$.

We have shown in this Letter that the Stelle-West formalism for non-linear gauge theories allows the construction of an off-shell AdS-invariant higher dimensional gravity. No boundary terms are added to the lagrangian when the AdS gauge transformations are performed. This is accomplished in any dimensions, no matter whether it be even or odd. It is also proved that such actions have the same coefficients as those obtained in refs. 3, 4.

We emphasize that the action (40) is invariant under the AdS group and that, when one picks the physical gauge $\xi^a = 0$, the theory becomes indistinguishable from the usual one, and the AdS symmetry is broken down to the Lorentz group. The only exception to this rule occurs in odd dimensions when the coefficients (3) are chosen. In this case, and for any value of $\xi^a$, it is possible to show that the Euler-Chern-Simons action written with $e^a$ and $\omega^{ab}$ differs from that written with $V^a$ and $W^{ab}$ by a boundary term. As a matter of fact, the defining relation for the non-linear fields $V^a$ and $W^{ab}$ given in (40), represents a gauge transformation for the linear connection $A = \frac{1}{2} \xi^{ab} J_{ab} - i e^a P_a$, which can be written in the form

$$A \rightarrow \tilde{A} = g^{-1} (d + A) g,$$

where $g = e^{-i \xi^a P_a}$ and $\tilde{A} = \frac{1}{2} W^{ab} J_{ab} - i V^a P_a$. This means that the linear and non-linear curvatures $F = dA + A^2$ and $\tilde{F} = d\tilde{A} + \tilde{A}^2$ are related by

$$\tilde{F} = g^{-1} F g.$$

Just as the usual Euler-Chern-Simons lagrangian, the odd-dimensional non-linear lagrangian, with the special choice of coefficients given in eq. (4), satisfies

$$dL_{\text{VW}}^{(2n-1)} = \langle \tilde{F}^n \rangle,$$

where $\langle J_{a_1 a_2} \cdots J_{a_{2n-3} a_{2n-2}} P_{a_{2n-1}} \rangle = \frac{1}{n!} \varepsilon_{a_1 \cdots a_{2n-1}}$.

Then, eq. (57) implies that

$$dL_{\text{CS}}^{(2n-1)} = \langle \tilde{F}^n \rangle = \langle F^n \rangle = dL_{\text{CS}}^{(2n-1)},$$

and hence we see that both lagrangians may locally differ only by a total derivative. The same arguments lead to the conclusion that, in general, any Chern-Simons lagrangian written with non-linear fields differs from the usual one by a total derivative.

When the linear lagrangian is invariant only under the stability subgroup, as it happens in even dimensions, the introduction of the non-linear fields brings in new terms in the action, which cannot be written as boundary contributions. We are left with a new action which is invariant under the full group.

It is perhaps interesting to note that, if one considers $g_{\mu\nu} = V^a V^b \eta_{ab}$, one can write the lagrangian of the action (40) in the form that was written in ref. 3. This means that, if one considers the theory to be constructed in terms of the space-time metric $g_{\mu\nu}$, ignoring the underlying formulation, the theory described in our manuscript is completely equivalent to the theory developed in refs. 2, 4. No trace of the new structure of the vierbein existing in the underlying formulation of the theory can be found at the metric level.

The interest in the study of the Lanczos-Lovelock gravity theory invariant under AdS transformations takes root in the fact that, in recent years, M theory has become the preferred description for the underlying structure of string theory 12, 13. Some of the expected features of M-theory are (i) its dynamics should somehow exhibit a superalgebra in which the anticommutator of two supersymmetry generators coincides with the AdS superalgebra in eleven dimensions 14, (ii) the low-energy regime should be described by an eleven dimensional supergravity of a new type which should stand on a firm geometric foundation in order to have an off-shell local supersymmetry 15, (iii) the perturbation expansion for graviton scattering in M-theory has recently led to a conjecture that the new supergravity Lagrangian should contain higher powers of curvature 16.

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