Spontaneous Symmetry Breaking for Massive Spin-2 Interacting with Gravity

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Abstract

An action for a massless graviton interacting with a massive tensor field is proposed. The model is based on coupling the metric tensor to an $SP(4)$ gauge theory spontaneously broken to $SL(2,C)$. The symmetry breaking is achieved by using a Higgs multiplet containing a scalar field and a vector field related by a constraint. We show that in the non-unitary gauge and for the Fierz-Pauli form of the mass term, the six degrees of freedom of the massive tensor are identified with two tensor helicities, two vector helicities of the Goldstone vector, and two scalars present in the Goldstone multiplet. The propagators of this system are well behaved, in contrast to the system consisting of two tensors.

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1 Introduction

It is well known that the action for massive spin-2 field does not have a smooth limit to the massless case [1], [2] and there are many doubts in the literature about the consistency of this theory [3], [4], [5]. The Fierz-Pauli mass term is required to give mass to the five dynamical degrees of freedom expected for a massive spin-2 field [6] . In a covariant formulation, the massive spin-2 field is represented, like the graviton, by a symmetric tensor. However, gauge freedom associated with diffeomorphism invariance guarantees that the graviton is massless. This means that there are six degrees of freedom associated with the symmetric tensor for the massive spin-2 field, as the four gauge degrees of freedom resulting from diffeomorphisms, having been used for the graviton, are not available. The Fierz-Pauli choice for the mass term decouples the independent sixth scalar degree of freedom from the five massive spin-2 degrees of freedom, only at the linearized level [3], [7]. Quantum corrections do not preserve the Fierz-Pauli form of the mass term, and implies that the sixth degree of freedom will have ghost interactions [8]. The renewed interest in massive spin-2 fields (also referred to as massive gravitons) comes from different directions, mainly from brane models where a different metric is taken on each of the branes [9], [10], [11]. As there is only one diffeomorphism to be preserved by the full system, only one combination of the metric tensors can be associated with the massless graviton. The other combination(s) will not have diffeomorphism invariance and will (each) correspond to six degrees of freedom. The mixing of the metrics is governed by the Fierz-Pauli choice, which guarantees that the sixth scalar mode decouples at the linearized level. When quantum corrections are taken into account, one finds that the scalar mode acquires a kinetic term with ghosts. The propagator of the massive field does not have a smooth $m \to 0$ limit, and suffers from the Dam-Veltman discontinuity [1], [2]. Recently, it was shown by Arkani-Hamed, Georgi and Schwartz [12], that the pathological behavior of the massive graviton propagator could be regulated by spontaneously breaking the symmetry of the two diffeomorphisms. This was done by introducing a Goldstone vector field which maps the coordinates of one brane to the other. This is similar to the use of Stueckelberg field for the description of massive gravity [13].

The idea in [12] is closely related to an approach proposed many years ago by Chamseddine, Salam and Strathdee [14] for strong gravity [15]. The system considered in [14] is based on the gauge symmetry $SP(4) \times SP(4)$ broken spontaneously to $SL(2, C)$ through a Higgs field transforming under

both gauge groups[1]. It was shown that in the unitary gauge a Fierz-Pauli term does arise for one combination of the gauge fields associated with the massive graviton, while the other combination gave a massless graviton. The interactions of the Higgs field could be evaluated by working in the non-unitary gauge where the Goldstone mode is kept. The work in [14] depends on the idea that the gravitational field could be formulated as a gauge theory of the $SP(4)$ gauge group where the spin-connection is taken as the gauge field of the $SL(2,C)$ subgroup of $SP(4)$. The vierbein is the gauge field of the four remaining generators of $SP(4)$ [16], [17]. The massive graviton was tuned to have a very heavy mass and very strong coupling. This arrangement is not necessary, and it is now preferable to have a very light graviton with weak coupling. Although this formulation is very elegant, the Higgs structure in the non-unitary gauge is complicated making the analysis of the Goldstone multiplet non transparent. The complication is due to the fact that no space-time metric is used but instead gauge fields are introduced, which then produce a metric as a product of the vierbein gauge fields. To simplify this model and keep track of the Higgs interactions, we shall consider instead the coupling of the metric tensor to a gauge theory based on $SP(4)$ spontaneously broken to $SL(2,C)$ through the Higgs mechanism. In this way the structure of the massive spin-2 field coupled to a massless graviton will be transparent.

The plan of this paper is as follows. In section two the action representing the interaction of a metric tensor coupled to $SP(4)$ gauge field and a Higgs multiplet is constructed. In section three the action is expanded in terms of $SL(2,C)$ representations of the $SP(4)$ fields. In section four the action is analyzed in the unitary gauge. In section five the analysis of the action is done in a non-unitary gauge and the degrees of freedom identified. Section six is the conclusion.

2   The action for graviton coupled to a massive spin-2 field

Consider a metric tensor $g_{\mu\nu}$ on a manifold M and the Einstein-Hilbert action associated with it

$$I_g = \frac{1}{4\kappa^2} \int_M d^4x \sqrt{g} R\left(g\right),$$

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[1]In [14] a supergravity system based on gauging the graded algebra $OSP(1;4) \times OSP(1;4)$ was considered. Here we shall only consider the bosonic version.

where $M_{Pl} = \frac{1}{\kappa}$. Consider also the gauge group $SP(4)$ and the gauge field W_μ associated with it. This can be expanded in terms of the $SL(2,C)$ subgroup [17], [14]

$$W_{\mu\alpha}{}^\beta = \left(L_\mu^a \left(\frac{i}{2\kappa_0} \gamma_a \right) + \frac{1}{4} B_\mu^{ab} \gamma_{ab} \right)_\alpha^\beta,$$

which satisfies the symmetry condition

$$(W_\mu C)_{\alpha\beta} = (W_\mu C)_{\beta\alpha},$$

where C is the charge conjugation matrix. It also satisfies the reality condition

$$\gamma_0 W_\mu^\dagger \gamma_0 = -W_\mu,$$

implying that L_μ^a and B_μ^{ab} are real. The $SP(4)$ gauge transformation of the gauge field is given by

$$W_\mu \to \Omega W_\mu \Omega^{-1} + \Omega \partial_\mu \Omega^{-1},$$

where the gauge parameter Ω can be expanded in terms of the $SL(2,C)$ components

$$\Omega = \exp\left(\frac{i}{2} \omega^a \gamma_a + \frac{1}{4} \omega^{ab} \gamma_{ab} \right).$$

The component form of the infinitesimal gauge transformations read

$$\delta L_\mu^a = -\kappa_0 \left(\partial_\mu \omega^a + B_\mu^{ab} \omega_b \right) + \omega^{ab} L_{\mu b},$$

$$\delta B_\mu^{ab} = -\left(\partial_\mu \omega^{ab} + B_\mu^{ac} \omega_c{}^b - B_\mu^{bc} \omega_c{}^a \right) - \frac{1}{\kappa_0} \left(\omega^a L_\mu^b + \omega^b L_\mu^a \right).$$

The gauge covariant field strengths are defined by

$$W_{\mu\nu} = \partial_\mu W_\nu - \partial_\nu W_\mu + [W_\mu, W_\nu],$$

which transforms as $W_{\mu\nu} \to \Omega W_{\mu\nu} \Omega^{-1}$. This can be resolved into

$$W_{\mu\nu} = L_{\mu\nu}^a \left(\frac{i}{2\kappa_0} \gamma_a \right) + \frac{1}{4} B_{\mu\nu}^{ab} \gamma_{ab},$$

where

$$L_{\mu\nu}^a = \partial_\mu L_\nu^a - \partial_\nu L_\mu^a + B_\mu^{ab} L_{\nu b} - B_\nu^{ab} L_{\mu b},$$

$$B_{\mu\nu}^{ab} = \partial_\mu B_\nu^{ab} - \partial_\mu B_\nu^{ab} + B_\mu^{ac} B_{\nu c}{}^b - B_\nu^{ac} B_{\mu c}{}^b - \frac{1}{\kappa_0^2} \left(L_\mu^a L_\nu^b - L_\nu^a L_\mu^b \right).$$

3

Now introduce the Goldstone field G_α^β in the antisymmetric representation of $SP(4)$:

$$(GC)_{\alpha\beta} = -(GC)_{\beta\alpha},$$

also subject to the reality and tracelessness conditions

$$\gamma_0 G^\dagger \gamma_0 = -G, \qquad G_\alpha^\alpha = 0.$$

Therefore we can decompose G in the form

$$G_\alpha^\beta = (\varphi\,(i\gamma_5) - v_a\,(\gamma_a\gamma_5))_\alpha^\beta,$$

where φ and v_a are real fields. In order to isolate the massive graviton degrees of freedom, we can eliminate one scalar degree of freedom by imposing a gauge invariant constraint on the multiplet G [19]

$$Tr\left(G^2\right) = -4a^2,$$

which in component form reads

$$\varphi^2 + v_a v^a = a^2.$$

The gauge transformation of G is $G \to \Omega G \Omega^{-1}$, so that the covariant derivative is given by

$$\nabla_\mu G = \partial_\mu G + [W_\mu, G],$$

transforming as $\nabla_\mu G \to \Omega \nabla_\mu G \Omega^{-1}$. The simplest action for the fields W_μ and G and not involving the space-time metric is

$$I_{W-G} = \int\limits_M d^4 x \epsilon^{\mu\nu\kappa\lambda} Tr \left(\frac{\alpha}{32a^3}\left(G\nabla_\mu G\nabla_\nu G + \nabla_\mu G\nabla_\nu G G\right)W_{\kappa\lambda}\right.$$

$$\left. + \frac{\beta}{96a^5}G\nabla_\mu G\nabla_\nu G\nabla_\kappa G\nabla_\lambda G\right).$$

Notice that both terms with the α coefficients are needed for the action to be Hermitian. One can also add to this action the term

$$\int\limits_M d^4 x \epsilon^{\mu\nu\kappa\lambda} Tr\left(GW_{\mu\nu}W_{\kappa\lambda}\right),$$

but in the unitary gauge this will only change, apart from adding a Gauss-Bonnet topological term, the coefficients of the terms already present.

The next step is to add a mixing terms between the two sectors. To do this define the field [12]

$$H_{\mu\nu} = g_{\mu\nu} + \frac{\kappa_0^2}{4a^2} Tr\left(\nabla_\mu G \nabla_\nu G\right),$$

and add the interaction term

$$I_{g-H} = m^4 \int_M d^4x \sqrt{g} g^{\mu\rho} g^{\nu\sigma} \left(H_{\mu\nu} H_{\rho\sigma} + (b-1) H_{\mu\rho} H_{\nu\sigma}\right).$$

It is important to note that the multiplet G has the following gauge transformations in component form

$$\delta\varphi = -\omega^a v_a,$$
$$\delta v_a = \omega_a \varphi + \omega_{ab} v^b.$$

This implies that by an appropriate choice of ω_a it is possible to use the unitary gauge

$$v_a = 0.$$

so that the constraint simplifies to $\varphi^2 = a^2$, and corresponds to a non-linear realization of the symmetry breaking from $SP(4)$ to $SL(2,C)$ [18]

3   The action in component form

To derive the action in component form we first write

$$\nabla_\mu G = \nabla_\mu \varphi \left(i\gamma_5\right) - \nabla_\mu v^a \gamma_a \gamma_5,$$

where

$$\nabla_\mu \varphi = \partial_\mu \varphi - \frac{1}{\kappa_0} L_\mu^a v_a,$$
$$\nabla_\mu v^a = \partial_\mu v^a + B_\mu^{ab} v_b + \frac{1}{\kappa_0} L_\mu^a \varphi.$$

The action is then given by

$$
I = \frac{1}{4\kappa^2} \int_M d^4x \sqrt{g} R(g)
$$

$$
+ \frac{\beta}{24a^5} \int_M d^4x \epsilon^{\mu\nu\kappa\lambda} \epsilon_{abcd} \left(\varphi \nabla_\mu v^a \nabla_\nu v^b \nabla_\kappa v^c \nabla_\lambda v^d - 4 \nabla_\mu \varphi v^a \nabla_\nu v^b \nabla_\kappa v^c \nabla_\lambda v^d \right)
$$

$$
- \frac{\alpha}{16a^3} \int_M d^4x \epsilon^{\mu\nu\kappa\lambda} \epsilon_{abcd} \left(\left(\varphi \nabla_\mu v^a \nabla_\nu v^b - 2 \nabla_\mu \varphi v^a \nabla_\nu v^b \right) B_{\kappa\lambda}^{cd} + \frac{2}{\kappa_0} v^a \nabla_\mu v^b \nabla_\nu v^c L_{\kappa\lambda}^d \right)
$$

$$
+ m^4 \int_M d^4x \sqrt{g} g^{\mu\rho} g^{\nu\sigma} \left(H_{\mu\nu} H_{\rho\sigma} + (b-1) H_{\mu\rho} H_{\nu\sigma} \right),
$$

where

$$
H_{\mu\nu} = g_{\mu\nu} - \frac{\kappa_0^2}{a^2} \left(\nabla_\mu \varphi \nabla_\nu \varphi + \nabla_\mu v^a \nabla_\nu v_a \right).
$$

4   The action in unitary gauge

It is easy to analyze the action in the unitary gauge $v^a = 0$ as this implies

$$
\varphi = a,
$$
$$
\nabla_\mu \varphi = 0,
$$
$$
\nabla_\mu v^a = \frac{a}{\kappa_0} L_\mu^a,
$$
$$
H_{\mu\nu} = g_{\mu\nu} - l_{\mu\nu},
$$

where $l_{\mu\nu} = L_\mu^a L_{\nu a}$. In this case the action simplifies to

$$
I = \frac{1}{4\kappa^2} \int_M d^4x \sqrt{g} R(g) + \frac{(3\alpha - \beta)}{\kappa_0^4} \int_M d^4x \det L_\mu^a
$$

$$
- \frac{\alpha}{16\kappa_0^2} \int_M d^4x \epsilon^{\mu\nu\kappa\lambda} \epsilon_{abcd} L_\mu^a L_\nu^b R_{\kappa\lambda}^{cd}(B)
$$

$$
+ m^4 \int_M d^4x \sqrt{g} g^{\mu\rho} g^{\nu\sigma} \left(h_{\mu\nu} h_{\rho\sigma} - (1-b) h_{\mu\rho} h_{\nu\sigma} \right),
$$

where

$$
h_{\mu\nu} = \left(g_{\mu\nu} - l_{\mu\nu} \right),
$$
$$
R_{\mu\nu}^{ab}(B) = \partial_\mu B_\nu^{ab} - \partial_\mu B_\nu^{ab} + B_\mu^{ac} B_{\nu c}{}^b - B_\nu^{ac} B_{\mu c}{}^b.
$$

We can write

$$\epsilon^{\mu\nu\kappa\lambda}\epsilon_{abcd}\, L_\mu^a L_\nu^b \, R_{\kappa\lambda}^{cd}\left(B\right) = -4 \det L \, L_a^\mu L_b^\nu \, R_{\mu\nu}^{ab}\left(B\right),$$

where we have defined $L_\mu^a L_a^\nu = \delta_\mu^\nu$. It is also well known that when the field B_μ^{ab} is eliminated by its equations of motion the above term reduces to

$$\frac{\alpha}{4\kappa_0^2} \int_M d^4x \det L \, R\left(l\right)$$

where $R\left(l\right)$ is the curvature of the metric tensor $l_{\mu\nu}$. Without any loss in generality we can set $\alpha = 1$. This system is known to give the coupling of a massive graviton to a massless graviton with the Fierz-Pauli choice $b = 0$. The massive graviton has mass of the order $\kappa_0 m^2$, which can be arranged to be small by an appropriate choice of m. The choice $b = 0$ is not stable under quantum corrections and the propagator does not have a smooth $m \to 0$ limit. When $b \neq 0$ there is a ghost mode in the tensor $l_{\mu\nu}$ which will propagate [3]. To find the order of the quantum corrections it is essential to the study the coupled system in other gauges, and to separate the six degrees of freedom of the massive tensor into two massless tensor polarizations of $l_{\mu\nu}$, the two vector polarizations in v^a and two scalar modes.

5 The non-unitary gauge

To examine the degrees of freedom in a different gauge, we shall keep the field v^a present, and use the gauge degrees of freedom ω^a to impose a gauge choice of the form

$$g^{\mu\nu}\left(\partial_\mu L_\nu^a + B_\mu^{ab} L_\nu^b\right) = 0,$$

or something equivalent. This will guarantee that the only propagating degrees of freedom present in L_μ^a are the two tensor degrees of helicities $+2$ and -2. This is easy to see because the system for L_μ^a is identical to the massless graviton with the only difference being that the diffeomorphism parameters ζ^μ are used instead of the parameters ω^a to impose the above gauge condition. From the explicit form of the action it is evident that all derivatives on the field v^a coming from the terms with coefficients α and β are antisymmetrized, implying that these terms do not give kinetic energy for v^a. The only terms that contain second order derivatives for v^a come from the mass mixing term with coefficient m^4. To see this we first express the field $H_{\mu\nu}$

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in terms of the component fields

$$H_{\mu\nu} = g_{\mu\nu} - \frac{\varphi^2}{a^2}l_{\mu\nu} - \frac{\kappa_0\varphi}{a^2}\left(\nabla_\mu v_\nu + \nabla_\nu v_\mu\right) - \frac{1}{a^2}\left(v_\mu - \kappa_0\partial_\mu\varphi\right)\left(v_\nu - \kappa_0\partial_\nu\varphi\right)$$
$$- \frac{\kappa_0^2}{a^2}l^{\rho\sigma}\nabla_\mu v_\rho\nabla_\nu v_\sigma + \cdots$$

where $v_\mu = L_\mu^a v_a$ and $\nabla_\mu v_\nu = \partial_\mu v_\nu - \Gamma_{\mu\nu}^\rho\left(l\right)v_\rho$ and we have used the relations

$$-\left(D_\mu L_\nu^a + D_\nu L_\mu^a\right) = 2\Gamma_{\mu\nu}^\rho\left(l\right)v_\rho + \cdots$$
$$D_\mu v_a D_\nu v^a = l^{\rho\sigma}\nabla_\mu v_\rho\nabla_\nu v_\sigma + \cdots$$

after substituting the B_μ^{ab} equation of motion. The fields φ and v_a are related by a constraint, and it is possible to express $\varphi = \sqrt{a^2 - v_a v^a}$ or simply constrain one degree of freedom in v_a. It is then clear that to leading order

$$H_{\mu\nu} = h_{\mu\nu} - \frac{\kappa_0}{a}\left(\nabla_\mu v_\nu + \nabla_\nu v_\mu\right) + \cdots$$

Then the mass mixing term gives, to lowest order, the following contribution to the kinetic energy of v_a

$$\kappa_0^2 m^4 \int\limits_M d^4x\sqrt{g}g^{\mu\rho}g^{\nu\sigma}\left(\left(\partial_\mu v_\nu + \partial_\nu v_\mu\right)\left(\partial_\rho v_\sigma + \partial_\sigma v_\rho\right)\right.$$
$$\left. -(1-b)\left(\partial_\mu v_\rho + \partial_\rho v_\mu\right)\left(\partial_\nu v_\sigma + \partial_\sigma v_\rho\right)\right).$$

After integration by parts and setting $b = 0$, this can be rewritten in the form

$$\kappa_0^2 m^4 \int\limits_M d^4x\sqrt{g}g^{\mu\rho}g^{\nu\sigma}\left(\left(\partial_\mu v_\nu - \partial_\nu v_\mu\right)\left(\partial_\rho v_\sigma - \partial_\sigma v_\rho\right)\right).$$

This describes the two spin-1 vector polarizations of the massive spin-2 field. There is a spin-0 polarization which can represented by the mass term of v_a in addition to the scalar field φ. This is so because three components of the vector v_a propagate, so the constraint can be thought to be a restriction on the fourth component of v_a. In this way φ is represented as

$$\varphi = a + \overline{\varphi},$$

where $\overline{\varphi}$ are fluctuations. Although φ does not have a direct kinetic term it occurs as a scaling factor for the curvature scalar $R(l)$. Isolating the relevant contributions we have

$$\int\limits_M d^4x \det L_\mu^a \, \varphi^3 \left(\frac{1}{4\kappa_0^2}R\left(l\right) + \frac{\beta - 3\alpha}{\alpha\kappa_0^4}\right),$$

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where we have used the constraint on φ and v_a to simplify $\beta\varphi^3\left(\varphi^2 + v_a v^a\right)$ to $\beta a^2 \varphi^3$. By a Weyl scaling of the field L_μ^a the field φ acquires a kinetic energy term, just as the dilaton in string theory, or when gravity is compactified to lower dimensions. By substituting the constraint one finds that the mass term for v_a is

$$-\frac{3\left(\beta - 3\alpha\right)}{2a^2\kappa_0^4}v_a v^a.$$

To conclude, in the unitary gauge, the tensor $h_{\mu\nu}$ has six degrees of freedom, five for the massive graviton coupled to an additional scalar degree of freedom. In the non-unitary gauge, there are two tensor polarizations of helicities ±2 (corresponding to the tensor $l_{\mu\nu}$), two polarizations of helicities ±1 corresponding to the transverse components of v_a and two scalar polarization of spin-0 corresponding to the field φ and the longitudinal component of v_a. The ill behavior of the massive graviton propagator can be avoided by considering instead of the $g_{\mu\nu}$, $l_{\mu\nu}$ system, the coupled constrained system of $g_{\mu\nu}$, L_μ^a, φ and v_a. The instability of the Fierz-Pauli choice of $b = 0$ would occur at the quantum level, but as explained in [12], the corrections would occur at a cut-off energy, where the ghost mode would start to propagate. At energies much lower than the cut-off scale, the corrections could be ignored, and the system is well behaved.

6   Conclusions

We have shown that it is possible to formulate an action for a massless graviton interacting with a massive spin-2 field as a theory obtained by coupling a metric tensor to a gauge theory of $SP(4)$ spontaneously broken to $SL(2, C)$. The symmetry breaking is done through a Higgs multiplet containing a scalar field and a vector field related by a constraint, employing a non-linear realization for the symmetry breaking. The action consists of three parts. The first part is the Einstein-Hilbert action for the metric tensor $g_{\mu\nu}$ on a space-time manifold M. The second part is metric independent and $SP(4)$ gauge invariant. The third is a mixing term between the metric and the gauge sectors. In the unitary gauge the Goldstone vector field v_a is set to zero and the field φ to a constant. The action reduces to the Fierz-Pauli form of a massless graviton interacting with a massive tensor. In the non-unitary gauge, the degrees of freedom of the massive spin-2 field and the scalar are given by the tensor polarizations of helicities ±2 in L_μ^a, the vector polarizations of helicities ±1 in v_a and two scalars of helicities 0 in φ and v_a. In this form the propagators of the separate modes are well behaved and have a smooth $m \to 0$ limit. Quantum corrections to the Fierz-Pauli choice of

the mass term would be damped by the cut-off energy. At energies much below the cut-off the action is well behaved. What remains to be seen is an explicit computation to show how the massless limit is attained. More importantly is to have a variant of this action where the system is treated more symmetrically with the two tensors corresponding to the two metrics on the separate sheets of two membranes. This can be done either by considering a purely metric theory as in [12], or by considering a gauge theory based on the gauge group $SP(4) \times SP(4)$ broken to $SL(2, C)$. This is indeed possible as shown in [14], but is complicated by the fact that the Higgs field transforms under $SP(4) \times SP(4)$ and has 16 components requiring the introduction of a good number of constraints. The resulting model is very similar to what is described here and is fairly straightforward in the unitary gauge. But the analysis is more complicated in the non-unitary gauge as one has to keep track of all independent components of the Goldstone fields.

7   Acknowledgments

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