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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

NON-METRICAL SUPERGRAVITY FOR A SPIN-TWO NONET

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INTERNATIONAL ATOMIC ENERGY AGENCY

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NON-METRICAL SUPERGRAVITY FOR A SPIN-TWO NONET *

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ABSTRACT

The SL(2M,C) gauge invariance of Einstein-Cartan-Weyl Lagrangians of strong gravity is generalized to the supersymmetry group U(2M,2M|1). The spontaneously broken theory, down to SL(2M,C) × U(1) × ^{5}U(1), describes the propagation of M^2 massive spin-2 and M massive spin-$\frac{3}{2}$ quanta. For M = 3, the theory may describe a coloured nonet of strongly interacting f gravitons.

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I. INTRODUCTION

The Einstein-Cartan-Weyl theory of gravity [1] has two types of invariances: i) general co-ordinate transformation invariance and ii) SL(2,C) or local Lorentz gauge invariance. In order to describe a nonet of (possibly coloured) spin-2 quanta in the theory of strong f gravity [2], it was necessary to generalize the SL(2,C) gauge invariance to an SL(6,C) invariance [3]. This generalization incorporates the internal SU(3) symmetry in a relativistic spin-unitary-spin mixing gauge symmetry.

More recently, Einstein's gravitational theory has been generalized into the locally supersymmetric theory of supergravity [4], through the introduction of a spin-$\frac{3}{2}$ fermionic partner to the gravitational field. One attempt to understand the geometric structure of this theory [5] proceeds through the OSP(4,1) gauge symmetry [6),7]. This latter symmetry incorporates the fermionic sector by grading the O(3,2) de Sitter or the $\widetilde{SP}(4)$ extension of SL(2,C). Internal symmetries of the O(N) type can be incorporated [7),8] through the further extension to OSP(4,N). However, in order to incorporate internal symmetries of the U(N) type, one must generalize the full superconformal or the U(2,2|1) group [9),10] to U(2,2|N).

Now each of the OSP(4,N) and the U(2,2|N) theories contains a single spin-2 field. If we wish to describe supersymmetrically a nonet of spin-2 f gravitons, we must consider supersymmetric generalizations of theories of the SL(6,C) type. In this note we study such a generalization.

We find that the supersymmetry group which contains SL(2M,C) is obtained by first extending the latter to U(2M,2M) and then grading it into U(2M,2M|1). The theory describes propagation of M^2 spin-2 and M spin-$\frac{3}{2}$ particles. One may further extend the theory to U(2M,2M|N) group, thus incorporating two distinct unitary groups of dimensions M and N - perhaps colour and flavour (U(M) × U(N)). We shall, however, limit our discussion to the case N = 1.

II. THE U(2M,2M|1) SYMMETRY

The parametrized generators $\mathcal{M}(\Omega^R,\epsilon,\bar\epsilon)$ of the U(2M,2M|1) superalgebra are represented by the (4M + 1) × (4M + 1) matrices,

$$\mathcal{M}(\Omega^R,\epsilon,\bar\epsilon) = \left|\begin{array}{c|c} \mathcal{B}(\Omega^R) & \epsilon \\ \hline \bar\epsilon & 2M\Omega \end{array}\right| , \qquad (1)$$

where

$$\mathcal{B}(\Omega^R) = \frac{1}{2}\Big\{\Omega + \gamma_5{}^5\Omega + \gamma_a\Omega^a + i\gamma_a\gamma_5{}^5\Omega^a + \frac{1}{2}\sigma_{ab}\Omega^{ab} + \lambda_i\Big(\Omega^i + \gamma_5{}^5\Omega^i + \gamma_a\Omega^{ai} + i\gamma_a\gamma_5{}^5\Omega^{ai} + \frac{1}{2}\sigma_{ab}\Omega^{abi}\Big)\Big\} ;$$

$$a = 0,1,2,3 \quad ; \quad i = 1,2,\ldots,(M^2-1) \quad . \qquad (2)$$

Here λ_i are the matrices of SU(M). The $\Omega^R = (\Omega, {}^5\Omega, \Omega^a, \ldots, \Omega^{abi})$ are the $16M^2$ bosonic parameters of the U(2M,2M) subgroup. The ϵ and $\bar\epsilon$ are the 4M fermionic (anticommuting) complex parameters and their Dirac adjoints. They belong to the fundamental representation of the U(2M,2M) subgroup.

Corresponding to the parameters $(\Omega^R,\epsilon,\bar\epsilon)$ we introduce the superalgebra generators $(J^R,S,\bar S)$. Then from

$$\left[\mathcal{M}(\Omega^R,\epsilon_1,\bar\epsilon_1),\ \mathcal{M}(\Omega_2^R,\epsilon_2,\bar\epsilon_2)\right] = i\,\mathcal{M}(\Omega_3^R,\epsilon_3,\bar\epsilon_3) \qquad (3)$$

we obtain the superalgebra

$$\{S_\alpha^m\ ,\ S_\beta^n\} = \{\bar S_m^\alpha\ ,\ \bar S_n^\beta\} = 0 \quad ; \quad \begin{matrix}\beta = 1,2,3,4, \\ n = 1,2,\ldots,M,\end{matrix} \qquad (4)$$

$$\{S_\alpha^m\ ,\ \bar S_n^\beta\} = \frac{1}{2M}\,\delta_n^m\Big[J - \gamma_5{}^5J + \gamma_a J^a - i\gamma_a\gamma_5{}^5J^a + \frac{1}{2}\sigma_{ab}J^{ab}\Big]_\alpha^\beta + \frac{1}{4}(\lambda_i)_n^m\Big[J^i - \gamma_5{}^5J^i + \gamma_a J^{ai} - i\gamma_a\gamma_5{}^5J^{ai} + \frac{1}{2}\sigma_{ab}J^{abi}\Big]_\alpha^\beta , \qquad (5)$$

$$[J_R, S_\alpha^m] = \frac{1}{2}(\gamma_R S)_\alpha^m \quad ; \quad [J_R, \bar S_m^\alpha] = (\tfrac{1}{2}\bar S\gamma_R)_m^\alpha \ , \qquad (6)$$

where

$$\gamma_R = \Big\{1_{4M},\ \gamma_5,\ \gamma_a,\ i\gamma_a\gamma_5,\ \sigma_{ab}\ ;\ \lambda_i,\ \lambda_i\gamma_5,\ \lambda_i\gamma_a,\ \lambda_i i\gamma_a\gamma_5,\ \lambda_i\sigma_{ab}\Big\} . \qquad (7)$$

The bosonic sector of the superalgebra, i.e. the U(2M,2M) subalgebra [11)] is given by the commutation relations of the matrices $J_R = \frac{1}{2}\gamma_R$. An element of the group is represented by

$$G(\Omega,\epsilon,\bar\epsilon) = e^{-i\mathcal{M}(\Omega,\epsilon,\bar\epsilon)} \quad . \qquad (8)$$

An object F belonging to the fundamental representation is a (4M+1) dimensional column

$$F = \begin{pmatrix}\xi_\alpha^m \\ \varphi\end{pmatrix} , \qquad (9)$$

where ξ_α^m is a 4M-dimensional U(2M,2M) complex spinor and φ is real. F has a complex adjoint (row)

$$\bar F = (\bar\xi_\alpha^m\ , \varphi) \quad . \qquad (10)$$

F and $\bar F$ transform as

$$F \to GF \quad ; \quad \bar F \to \bar F\,G^{-1} \qquad (11)$$

so that $\bar F F$ is invariant.

An object $\Phi(Z^R,\eta,\bar\eta)$, belonging to the adjoint representation has the same structure as $M(\Omega^R,\epsilon,\bar\epsilon)$ and transforms as

$$\Phi \to G\,\Phi\,G^{-1} \quad . \qquad (12)$$

The infinitesimal transformations are given by

$$\delta\Phi = -i\left[\mathcal{M}(\Omega,\epsilon,\bar\epsilon)\ ,\ \Phi(Z,\eta,\bar\eta)\right] \quad . \qquad (13)$$

III. INVARIANT LAGRANGIAN

In order to realize the U(2M,2M|1) symmetry locally, we introduce gauge fields through the 1-form

$$\mathcal{A} = \mathcal{M}(V_\mu^R, \psi_\mu, \bar\psi_\mu)\ dx^\mu \quad , \qquad (14)$$

where $\mathcal{M}(V,\psi,\bar\psi)$ has, of course, the same structure as (1). Under the local U(2M,2M|1) transformations, $\mathcal{A}$ transforms inhomogeneously in the adjoint representation as

$$\mathcal{A} \to G\mathcal{A}G^{-1} - i d\, G.G^{-1} \quad . \tag{15}$$

The infinitesimal transformations are

$$\delta\mathcal{A}(V^R,\psi,\bar\psi) = -i\left[\mathcal{M}(\Omega^R,\epsilon,\bar\epsilon),\ \mathcal{A}(V^R,\psi,\bar\psi)\right] - d\mathcal{M}(\Omega^R,\epsilon,\bar\epsilon) \quad . \tag{16}$$

The curvature 2-form is defined as usual by

$$R(\nabla_\wedge V^R,\nabla_\wedge\psi,\nabla_\wedge\bar\psi) = (d - i\mathcal{A})_\wedge \mathcal{A}$$

$$\equiv \nabla_\wedge\mathcal{A}(V^R,\psi,\bar\psi) \tag{17}$$

and transforms homogeneously as,

$$R \to G\,R\,G^{-1} \quad . \tag{18}$$

The components $(\nabla_\wedge V^R,\nabla_\wedge\psi,\nabla_\wedge\bar\psi)$ are given in Table I.

As in the OSP(4,1) theory, one needs the auxiliary (Higgs) multiplet, to which the γ_5 matrix belongs, in order to construct an invariant action. For the U(2M,2M|1) symmetry, this multiplet $\mathcal{N}(\varphi^R,\chi,\bar\chi)$ should belong to the adjoint representation. Of course, the intrinsic parities of the component fields are abnormal. Define the covariant differential

$$\nabla\mathcal{N} = d\mathcal{N} - i\ [\mathcal{A},\mathcal{N}] \quad . \tag{19}$$

The most general U(2M,2M|1) gauge invariant Einstein-Cartan-Weyl action is

$$\begin{aligned} A = \int \mathrm{Tr}\ \Big\{ & R_\wedge R\, f_1(\mathcal{N}) \\ & + \alpha\ R_\wedge \nabla\mathcal{N}_\wedge \nabla\mathcal{N}\, f_2(\mathcal{N}) \\ & + \beta\ \nabla\mathcal{N}_\wedge \nabla\mathcal{N}_\wedge \nabla\mathcal{N}_\wedge \nabla\mathcal{N}\, f_3(\mathcal{N}) \Big\} \\ & + \text{h.c.} \end{aligned} \tag{20}$$

Here α and β are arbitrary constants; $f_1(\mathcal{N})$, $f_2(\mathcal{N})$ and $f_3(\mathcal{N})$ are arbitrary functions of odd powers in $\mathcal{N}$. The trace operation is taken over the U(2M,2M|1) matrices and the integration is over the four-volume determined

by the exterior forms. It is important to realize that the Einstein-Weyl-Cartan theory, represented by (20), is non-metrical and the Lagrangian is severely restricted in form. This restriction makes it impossible for any but the spin-2 and spin-$\frac{3}{2}$ components of $\mathcal{A}$ and $\mathcal{N}$ to propagate.

The action (20) can, in principle, be minimized around a constant value for $f_1(\mathcal{N})$ proportional to

$$\left[\begin{array}{c|c} \gamma_5 & 0 \\ \hline 0 & 0 \end{array}\right] \quad . \tag{21}$$

This induces the spontaneous breakdown of the U(2M,2M|1) symmetry down to SL(2M,C) $\times$ U(1) $\times$ ^{5}U(1). This can be understood as follows. From the infinitesimal transformations (13) of the Higgs multiplet $(\varphi^R,\chi,\bar\chi)$, we find that as $^5\varphi$, the coefficient of γ_5, develops a vacuum expectation value $\langle{}^5\varphi\rangle$, then $(\varphi^a,\ {}^5\varphi^a,\ \varphi^{ai},\ {}^5\varphi^{ai},\ \chi,\ \bar\chi)$ become Goldstone fields. The inhomogeneous terms in their transformations are

$$\begin{aligned} \delta\varphi^a &= -{}^5\Omega^a \cdot \langle{}^5\varphi\rangle \quad , & \delta\,{}^5\varphi^a &= -\Omega^a\ \langle{}^5\varphi\rangle \ ; \\ \delta\varphi^{ai} &= -{}^5\Omega^{ai}\,\langle{}^5\varphi\rangle \quad , & \delta\,{}^5\varphi^{ai} &= -\Omega^{ai}\langle{}^5\varphi\rangle , \\ \delta\chi &= \frac{i}{2}\gamma_5\epsilon\,\langle{}^5\varphi\rangle \quad , & \delta\bar\chi &= -\frac{i}{2}\bar\epsilon\gamma_5\ \langle{}^5\varphi\rangle \ . \end{aligned} \tag{22}$$

These Goldstone fields correspond to the spontaneously broken generators $(J^a,\ {}^5J^a,\ J^{ai},\ {}^5J^{ai},\ S,\ \bar S)$. The generators $(J^i,\ {}^5J^i,\ J_{ab},\ J^i_{ab})$ of the SL(2M,C) algebra as well as J and ^5J, the generators of U(1) and ^{5}U(1), respectively, are left unbroken. The Goldstone fields can, of course, be gauged away (the unitary gauge). The remaining fields $(\varphi,\ {}^5\varphi,\ \varphi^{ab},\ \varphi^i,\ {}^5\varphi^i,\ \varphi^{abi})$ of the Higgs multiplet do not have any propagation character in this theory *)

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*) It may be thought that one can introduce a group-invariant metric of the form $g_{\mu\nu} = \mathrm{Tr}[\nabla_\mu\mathcal{N}\ \nabla_\nu\mathcal{N}]$ and use it to construct invariant terms containing the usual spin-0 and spin-1 kinetic terms, $\frac{1}{2}(\partial_\mu\varphi)^2$ and $-\frac{1}{4}(F_{\mu\nu})^2$. However, this approach is beset with the following difficulties: a) the appearance of negative-norm states associated with the non-compact parameters of the U(2M,2M|1) group, unless further constraint equations are used; b) the appearance of kinetic terms of second order for the fermions (ghosts) whenever second-order supersymmetric Lagrangians are constructed, and c) the loss of conjugacy between the connection fields $(V^{ab}_\mu, V^{abi}_\mu)$ and the vierbeins (V^a_μ, V^{ai}_μ). A set of new particles described by the connections acquire independent propagation.

The first term in (20), proportional to the vacuum expectation value (21), yields the Lagrangian

$$\mathcal{L} = e^{\mu\nu\lambda\rho} \Big[\frac{1}{8} \epsilon_{abcd} \nabla_\mu V_\nu^{ab} \nabla_\lambda V_\rho^{cd} + \frac{1}{4M} \epsilon_{abcd} \nabla_\mu V_\nu^{abi} \nabla_\lambda V_\rho^{cd}{}_i$$

$$+ \nabla_\mu V_\nu \nabla_\lambda {}^5V_\rho + \frac{2}{N} \nabla_\mu V_\nu^i \nabla_\lambda V_\rho{}^i$$

$$+ \frac{1}{2M} \nabla_\mu \bar{\psi}_\nu \gamma_5 \nabla_\lambda \psi_\rho\Big] + \text{h.c.} \qquad (23)$$

The Lagrangian (23) by itself, without including the terms of (20) involving the remaining Higgs fields, is, of course, invariant under the surviving $SL(2M,C) \times U(1) \times {}^5U(1)$ symmetry. A closer investigation of (23) shows the following:

1) The first two terms of (23) contain the usual SL(2M,C) theory and provide, via the conjugate fields (V_μ^a, V_μ^{ab}) and $(V_\mu^{ai}, V_\mu^{abi})$, the propagation of M^2 massive spin-2 quanta. The masses come from the "cosmological" terms involving V_μ^a and V_μ^{ai}. It should be pointed out that since the theory contains "cosmological" terms, it must be linearized around a curved (de Sitter) background. The $(\text{magnitude})^2$ of the cosmological term is of the same order as that for the spin-2 coupling.

2) The remaining bosonic fields of the gauge multiplet, namely $(V_\mu, {}^5V_\mu, {}^5V_\mu^a, V_\mu^i, {}^5V_\mu^i, {}^5V_\mu^{ai})$ do not acquire propagation character at the classical level. These fields together with the remaining non-propagating fields of the Higgs multiplet, however, may pick up propagation through quantum corrections (quantum completion), as has been suggested [1)] for the usual SL(6,C) theory. *)

3) The last term in (23) provides propagation for M massive spin-$\frac{3}{2}$ quanta. Other contributions to the masses of these quanta come from the first two terms.

We conclude this note by giving the Lagrangian term describing the coupling to matter. Taking the fundamental representation F, given by (9), to represent matter, the non-metric geometric Lagrangian term can be written as

*) See footnote, p.6.

$$\mathcal{L}_{matter} = \bar{F}\, \nabla\mathcal{N}_\wedge \nabla\mathcal{N}_\wedge \nabla\mathcal{N}_\wedge \nabla F + \text{h.c.} \qquad (24)$$

This Lagrangian provides for the propagation of the fermionic components ξ of F, but not for the scalar φ. As noted before, this is a general property of all such Lagrangians (i.e. propagation of spins 2, $\frac{3}{2}$ and $\frac{1}{2}$ only)*). In the unitary gauge, we have

$$\nabla\mathcal{N} \rightarrow \langle {}^5\varphi \rangle \left| \begin{array}{c|c} V^a\, i\gamma_a\gamma_5 & 0 \\ \hline 0 & 0 \end{array} \right| . \qquad (25)$$

Hence, the propagation of spin-$\frac{1}{2}$ matter is described by a term of the form

$$\bar{\xi}\, \gamma_5\, V^a\gamma_{a\wedge} V^b\gamma_{b\wedge} V^c\gamma_{c\wedge} \nabla\xi + \text{h.c.} , \qquad (26)$$

which is proportional to the usual term

$$(\det V_{\mu a})\, \bar{\xi}\, V^{\mu a}\gamma_a \nabla_\mu \xi + \text{h.c.} \qquad (27)$$

*) See footnote, p.6.

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Table I

Table of curvature 2-form components

$$\nabla_\wedge V = d_\wedge V - \frac{i}{2M}\,\bar\psi_\wedge\psi$$

$$\nabla_\wedge {}^5V = d_\wedge {}^5V + V^a{}_\wedge {}^5V_a + \frac{2}{M} V^{ai}{}_\wedge {}^5V_{ai} + \frac{i}{2M}\bar\psi\gamma_{5\wedge}\psi$$

$$\nabla_\wedge V^a = d_\wedge V^a + {}^5V_\wedge {}^5V^a + V^{ab}{}_\wedge V_b + \frac{2}{M}\left({}^5V^i{}_\wedge {}^5V^a_i + V^{abi}{}_\wedge V_{bi}\right) - \frac{i}{2M}\bar\psi\gamma^a{}_\wedge\psi$$

$$\nabla_\wedge {}^5V^a = d_\wedge {}^5V^a + {}^5V_\wedge V^a + V^{ab}{}_\wedge {}^5V_b + \frac{2}{M}\left({}^5V^i{}_\wedge V^a_i + V^{abi}{}_\wedge {}^5V_{bi}\right) + \frac{i}{2M}\bar\psi i\gamma^a\gamma_{5\wedge}\psi$$

$$\begin{aligned}\nabla_\wedge V^{ab} = {} & d_\wedge V^{ab} - V^a{}_\wedge V^b + {}^5V^a{}_\wedge {}^5V^b - V^{ac}{}_\wedge V^b_c \\ & + \frac{2}{M}\left(-V^{ai}{}_\wedge V^b_i + {}^5V^{ai}{}_\wedge {}^5V^b_i - V^{aci}{}_\wedge V^b_{ci}\right) - \frac{i}{2M}\bar\psi\,\sigma^{ab}{}_\wedge\psi\end{aligned}$$

$$\nabla_\wedge V^i = d_\wedge V^i + \frac{1}{2} f^{ijk}\left(V_{j\wedge}V_k - {}^5V_{j\wedge}{}^5V_k + V^a_{j\wedge}V_{ak} - {}^5V^a_{j\wedge}{}^5V^a_k + \frac{1}{2}V^{ab}_{j\,\wedge}V_{abk}\right) - \frac{i}{4}\bar\psi\lambda^i{}_\wedge\psi$$

$$\begin{aligned}\nabla^5{}_\wedge V^i = {} & d_\wedge {}^5V^i + V^{ai}{}_\wedge {}^5V_a - {}^5V^{ai}{}_\wedge V_a + d^{ijk}\,V^a_{j\wedge}{}^5V_{ak} + \\ & + f^{ijk}\left(V_{j\wedge}{}^5V_k + \frac{1}{8}\epsilon_{abcd}\,V^{ab}_{j\,\wedge}V^{cd}_k\right) + \frac{i}{4}\bar\psi\lambda^i\gamma_{5\wedge}\psi\end{aligned}$$

$$\begin{aligned}\nabla_\wedge V^{ai} = {} & d_\wedge V^{ai} + {}^5V^i{}_\wedge {}^5V^a - {}^5V^{ai}{}_\wedge {}^5V + V^{abi}{}_\wedge V_b - V^{ab}{}_\wedge V^i_b \\ & + d^{ijk}\left({}^5V_{j\wedge}{}^5V^a_k + V^{ab}_{j\,\wedge}V_{bk}\right) + f^{ijk}\left(V_{j\wedge}V^a_k - \frac{1}{2}\epsilon^{abcd}\,{}^5V_{bj\wedge}V_{cdk}\right) - \frac{i}{4}\bar\psi\gamma^a\lambda^i{}_\wedge\psi\end{aligned}$$

$$\begin{aligned}\nabla_\wedge {}^5V^{ai} = {} & d_\wedge {}^5V^{ai} + {}^5V^i{}_\wedge V^a - V^{ai}{}_\wedge {}^5V + V^{abi}{}_\wedge {}^5V_b + V^{ab}{}_\wedge {}^5V^i_b \\ & + d^{ijk}\left({}^5V_{j\wedge}V^a_k + V^{ab}_j{}_\wedge {}^5V_{bk}\right) + f^{ijk}\left(V_{j\wedge}{}^5V^a_k - \frac{1}{2}\epsilon^{abcd}\,V_{bj\wedge}V_{cdk}\right) + \frac{i}{4}\bar\psi i\gamma_a\gamma_5\lambda^i{}_\wedge\psi\end{aligned}$$

$$\begin{aligned}\nabla_\wedge V^{abi} = {} & d_\wedge V^{abi} - V^{ai}{}_\wedge V^b + V^{bi}{}_\wedge V^a + {}^5V^{ai}{}_\wedge {}^5V^b - {}^5V^{bi}{}_\wedge {}^5V^a \\ & - V^{aci}{}_\wedge V^b_c + V^{bci}{}_\wedge V^a_c - d^{ijk}\left(V^a_{j\wedge}V^b_k - {}^5V^a_{j\wedge}{}^5V^b_k + V^{ac}_{j\,\wedge}V^b_{ck}\right) \\ & + f^{ijk}\left(V_{j\wedge}V^{ab}_k - \frac{1}{2}\epsilon^{abcd}\left[{}^5V_{j\wedge}V_{cdk} + 2\,V_{cj\wedge}{}^5V_{dk}\right]\right) - \frac{i}{4}\bar\psi\,\sigma^{ab}\lambda^i{}_\wedge\psi\end{aligned}$$

$$\nabla_\wedge\psi = \left(d - \frac{i}{2}V^R\gamma_R\right)_\wedge\psi$$

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