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LETTER TO THE EDITOR

The spectrum of the S^5 compactification of the chiral $N=2$, $D=10$ supergravity and the unitary supermultiplets of U(2, 2/4)†

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Abstract. We calculate the spectrum of the S^5 compactification of the chiral $N=2$, $D=10$ supergravity theory. The modes on S^5 fall into unitary irreducible representations of the $D=5$, $N=8$ anti-de Sitter supergroup U(2, 2/4). These unitary supermultiplets involve fields of spin $\leqslant 2$ with quantised 'mass' eigenvalues. The massless multiplet contains fifteen vector fields, six self-dual and six anti-self-dual anti-symmetric tensor fields. The fields of the massless multiplet are expected to be those of a gauged $N=8$ theory in $D=5$ with a local gauge group SU(4).

Recently there has been a surge of interest in Kaluza-Klein type theories in the context of supergravity. The main object of study so far has been the eleven-dimensional simple supergravity of Cremmer *et al* [1] and various compactifications of this theory down to four dimensions have been studied. A compactification is called spontaneous if, for example, the eleven-dimensional theory admits a vacuum solution leading to a four-dimensional spacetime M_4 and a seven-dimensional compact internal manifold M_7 parametrised by the extra seven coordinates. Then from a four-dimensional point of view the spectrum of the theory consists of massless modes and an infinite tower of massive modes at the Planck scale corresponding to the Fourier expansion on M_7.

In this letter we shall study the spectrum of the chiral $N=2$, $D=10$ supergravity theory [2] compactified on S^5 down to five spacetime dimensions. This theory *cannot* be obtained from the eleven-dimensional supergravity by dimensional reduction in contrast to the non-chiral $N=2$, $D=10$ theory which can be. The non-chiral and chiral $N=2$ supergravity theories can be regarded as the massless sector of type IIa and type IIb superstring theories in $D=10$‖. Both of these theories when dimensionally reduced to $D=4$ give the $N=8$ theory.

The fields of the chiral $N=2$, $D=10$ supergravity are: a complex scalar φ, a complex Weyl spinor λ, a complex anti-symmetric second rank tensor $A_{\mu\nu}$, a complex Weyl gravitino ψ_μ, a real graviton ('zehnbein') e^r_μ and a real fourth-rank anti-symmetric tensor $A_{\mu\nu\rho\lambda}$. The field strength $F_{\mu\nu\rho\lambda\sigma}$ of the field $A_{\mu\nu\rho\lambda}$ is self-dual in the free theory. The complex scalar field φ parametrises the non-compact coset space SU(1, 1)/U(1) [4]. The covariant field equations of chiral $N=2$, $D=10$ supergravity were obtained

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‖ For a complete list of references on ten-dimensional string theories see the review article [3].

in reference [4]. So far no compactifying solution of this theory to four dimensions has been found that admits a non-Abelian isometry group G for the six-dimensional internal manifold. However, it is known that this theory can be compactified on S^5 leading to the five-dimensional anti-de Sitter space AdS_5 [5]. The isometry group of this vacuum solution is $SO(4,2)\times SO(6) = SU(2,2)\times SU(4)$ where $SO(4,2)$ is the five-dimensional anti-de Sitter group. This solution has maximal number of supersymmetries in AdS_5 and its full super-invariance group is $U(2,2/4)$ whose even subgroup is $SU(2,2)\times U(4)$. Note that the extra $U(1)_Y$ comes from the ten-dimensional theory itself.

We have calculated the spectrum of the S^5 compactification using the fact that the modes for different spin fields must fall into unitary representations of $U(2,2/4)$ and that the graviton modes fall into symmetric tensor representations of SO(6), corresponding to the scalar modes on S^5. The modes for other spin fields are determined by the fact that together with the graviton modes they must form unitary irreducible representations (UIR) of $U(2,2/4)$. Since the highest spin field in the massive supermultiplets is $s=2$ they correspond to short supermultiplets (in the sense of [6]) with quantised mass eigenvalues.

We have constructed the UIRs of $U(2,2/4)$ using the oscillator methods developed in [7–8]†. This method has recently been applied to the calculation of the spectra of S^4 and S^7 compactifications of the eleven-dimensional supergravity [10–11] and agrees with the calculations previously done using coset space techniques. The advantage of the oscillator method is that it yields automatically the irreducible unitary supermultiplets with quantised mass eigenvalues into which the spectrum must fit. The basic idea of the oscillator method is to identify a compact sub-supergroup K of a non-compact supergroup G and construct the UIRs of G in terms of finite dimensional unitary representations of K. For the non-compact supergroup $U(2,2/4)$ the relevant sub-supergroup is $U(2/2)\times U(2/2)\times U(1)$ with respect to which the Lie superalgebra $U(2,2/4)$ has a Jordan structure [8], i.e.,

$$L = L^- \oplus L^0 \oplus L^+ \tag{1}$$

where L^0 represents the generators of $U(2/2)\times U(2/2)\times U(1)$ and $L^-\oplus L^+$ represents the non-compact generators of $U(2,2/4)$ such that

$$\begin{aligned} &[L^0, L^+] = L^+, \qquad [L^0, L^-] = L^-, \\ &[L^+, L^-] = L^0, \qquad [L^+, L^+] = 0 = [L^-, L^-]. \end{aligned} \tag{2}$$

The Lie superalgebra $U(2,2/4)$ is realised in terms of bilinears of boson and fermion annihilation and creation operators ξ_A, $\xi^A = (\xi_A)^\dagger$ and η_M, $\eta^M = (\eta_M)^\dagger$ transforming covariantly and contravariantly under the two $U(2/2)$ subgroups of $U(2,2/4)$, respectively:

$$\begin{aligned} &\xi_A = \begin{pmatrix} a_i \\ \alpha_\mu \end{pmatrix} \qquad \xi^A = \begin{pmatrix} a^i \\ \alpha^\mu \end{pmatrix} \qquad \begin{matrix} i, j = 1,2 \\ \mu, \nu = 1,2 \end{matrix} \\ &[a_i, a^j] = \delta_i^j \qquad \{\alpha_\mu, \alpha^\nu\} = \delta_\mu^\nu \\ &\eta_M = \begin{pmatrix} b_r \\ \beta_x \end{pmatrix} \qquad \eta^M = \begin{pmatrix} b^r \\ \beta^x \end{pmatrix} \qquad \begin{matrix} r, s = 1,2 \\ x, y = 1,2 \end{matrix} \\ &[b_r, b^s] = \delta_r^s \qquad \{\beta_x, \beta^y\} = \delta_x^y. \end{aligned} \tag{3}$$

† For a review of the oscillator method and a discussion of its applicability to the calculation of the spectra of supergravity theories see [9].

Then the generators of U(2, 2/4) are given in terms of the oscillators as

$$L^- = \vec{\xi}_A \cdot \vec{\eta}_M, \qquad L^0 = \vec{\xi}^A \cdot \vec{\xi}_B \oplus \vec{\eta}^M \cdot \vec{\eta}_N, \qquad L^+ = \vec{\xi}^A \cdot \vec{\eta}^M \tag{4}$$

where the arrow over $\vec{\xi}$ and $\vec{\eta}$ means that we are taking an arbitrary number p of oscillators and the dot represents the summation over the generation index running from 1 to p. Note that the three grading, (4), (Jordan structure) is given by the number operator of all the oscillators that generate an Abelian U(1) belonging to L^0. Then starting from a 'ground state' $|\Omega\rangle$ in the super Fock space that transforms irreducibly under K and is annihilated by all operators in L^- one can generate a UIR of G by repeated application of the operators in L^+:

$$R = \{|\Omega\rangle, L^+|\Omega\rangle, (L^+)^2|\Omega\rangle, \ldots\}. \tag{5}$$

Note that the irreducibility of the unitary representation of U(2, 2/4) follows from the irreducibility of $|\Omega\rangle$ under $U(2/2) \times U(2/2) \otimes U(1)$. The anti-de Sitter group SO(4, 2) has a Jordan structure with respect to its maximal compact subgroup $SU(2)_{j_1} \times SU(2)_{j_2} \times U(1)_E$. Positive energy UIRs of SO(4, 2) can be obtained by the oscillator method using the bosonic oscillators as above. They are also uniquely determined by a lowest state annihilated by L_{ir} and which transforms irreducibly under $SU(2)_{j_1} \times SU(2)_{j_2} \times U(1)_E$. The $SU(2)_{j_1} \times SU(2)_{j_2} \simeq SO(4)$ is the rotation group in $D=5$ and $U(1)_E$ denotes the Abelian group generated by the energy operator. Similarly the internal symmetry group $SO(6) = SU(4)$ has a Jordan structure with respect to its $SU(2)_{k_1} \times SU(2)_{k_2} \times U_I(1)$ subgroup. Its representations can all be constructed by the oscillator method using fermion bilinears as above. The UIRs of SU(4) are also uniquely determined by the transformation properties of their lowest states annihilated by $L_{\mu x}$.

The short unitary supermultiplets whose highest spin is two, can all be constructed by taking as the ground state $|\Omega\rangle$ the Fock vacuum $|0\rangle$. Even though the vacuum $|0\rangle$ is a singlet under $SU(2/2) \times SU(2/2)$ subgroup, it has definite $U(1)_E \times U(1)_I$ quantum numbers that depend on p. Therefore by varying the generation index p we obtain inequivalent unitary supermultiplets of U(2, 2/4). By acting on the vacuum $|0\rangle$ with supersymmetry generators L^{ix} and $L^{\mu r}$ one generates new lowest states of both SO(4, 2) and SU(4), i.e., states that are annihilated by both L_{ir} and $L_{\mu x}$. Acting on such lowest states with L^{ir} and $L^{\mu x}$ one generates UIRs of SO(4, 2) with definite SU(4) transformation properties. Such lowest states are labelled by their transformation properties under the even subgroup $SU(2)_{j_1} \times SU(2)_{k_1} \times SU(2)_{j_2} \times SU(2)_{k_2} \times U(1)_E \times U(1)_Y \otimes U(1)_I$. For example, for $p = 1$ we list below the possible lowest states, their $SU(2)_{j_1} \times SU(2)_{k_1} \times SU(2)_{j_2} \times SU(2)_{k_2}$ transformation properties (in Young tableau notation) and the SU(4) representations one obtains by the action of $L^{x\mu}$ on them.

Lowest States	$SU(2)_{j_1} \times SU(2)_{k_1} \times SU(2)_{j_2} \times SU(2)_{k_2}$	$SU(2)_{j_1} \times SU(2)_{j_2}$	SU(4)
$\|0\rangle =$	$\|(1, 1, 1, 1)\rangle$	(0, 0)	6
$L^{xi}\|0\rangle =$	$\|(\square, 1, 1, \square)\rangle$	$(\frac{1}{2}, 0)$	$\bar{4}$
$(L^{xi})^2\|0\rangle =$	$\|(\square\square, 1, 1, \begin{smallmatrix}\square\\\square\end{smallmatrix})\rangle$	(1, 0)	1
$L^{r\mu}\|0\rangle =$	$\|(1, \square, \square, 1)\rangle$	$(0, \frac{1}{2})$	4
$(L^{r\mu})^2\|0\rangle =$	$\|(1, \begin{smallmatrix}\square\\\square\end{smallmatrix}, \square\square, 1)\rangle$	(0, 1)	1

The basis of a positive energy UIR of SO(4, 2) can be taken to be the Fourier modes of a field in AdS_5. The spacetime properties of this field is determined by the $SU(2)_{j_1} \times SU(2)_{j_2} \otimes U(1)_E$ transformation properties of the lowest state of the UIR. The anti-de Sitter 'mass' of a field of spin (j_1, j_2) is determined by the eigenvalue of the $U(1)_E$ generator on the lowest state. Therefore for $p=1$ the unitary supermultiplet of $U(2, 2/4)$ consists of six scalar fields $\varphi^{ab} = -\varphi^{ba}$ $(a, b = 1, 4)$, four complex left-handed spinors λ_a^+ and four complex right-handed spinors λ^{-a} and an antisymmetric self-dual tensor field $A_{\mu\nu}$ and an anti-self-dual tensor field $\tilde{A}_{\mu\nu}$. The number of bosonic and fermionic degrees of freedom do not match in this supermultiplet. We shall call this supermultiplet the doubleton of $U(2, 2/4)$ in analogy with the doubleton in $D=7$ [10] and singleton in $D=4$ [11, 12, 13]. The doubleton representation does not have a Poincaré limit and it decouples from the spectrum of S^5 compactification. This can best be understood by the fact that the six scalars of the doubleton multiplet are the conformal scalars (their gradients give conformal Killing vectors) and can be gauged away by appropriate conformal gauge choice. As in the case of $D=4$ singleton [13] and $D=7$ doubleton [10] one may be able to construct a field theory of $D=5$ doubleton in one lower spacetime dimension, i.e., in $D=4$. The AdS group $SU(2, 2)$ will then be the conformal group in $D=4$ and the resulting field theory will be conformally invariant. In fact, the unique candidate for this doubleton theory in $D=4$ is the $N=4$ Yang–Mills theory which is known to be conformally invariant. The study of the doubleton field theory in $D=4$ will be given elsewhere [14].

For $p=2$ one obtains the massless supermultiplet which contains the massless graviton. For $p \geq 3$ one gets massive supermultiplets of $U(2, 2/4)$. We give the complete spectrum of the S^5 compactification of chiral $N=2$, $D=10$ theory in table 1. Note that there are new massive towers of $U(2, 2/4)$ supermultiplets starting at $p=3$ and $p=4$. For $p>4$ there are no new towers starting and one gets only the massive excitations of the towers starting at $p \leq 4$†. In table 1 we also list the transformation properties of the lowest states under the rotation group $SO(4) = SU(2)_{j_1} \times SU(2)_{j_2}$ and the corresponding fields. λ_+ and λ_- denote 'left- and right-handed' complex spinors. The chirality or handedness refers to the anti-de Sitter group $SU(2, 2)$ or to the rotation group $SO(4)$. In five-dimensional Minkowski space one cannot define chiral spinors. Instead one can have symplectic Majorana spinors [16]. The *complex* anti-symmetric tensors $A_{\mu\nu}$ and $\tilde{A}_{\mu\nu}$ are 'self-dual' and 'anti-self-dual' anti-symmetric tensor fields in the sense of [17]‡. They satisfy the duality conditions

$$A_{\mu\nu} = (\mathrm{i}/120m)\varepsilon_{\mu\nu\rho\sigma\lambda}F^{\rho\sigma\lambda}$$

$$\tilde{A}_{\mu\nu} = (-\mathrm{i}/120m)\varepsilon_{\mu\nu\rho\sigma\lambda}\tilde{F}^{\rho\sigma\lambda}$$

where $F^{\rho\sigma\lambda}(\tilde{F}^{\rho\sigma\lambda})$ is the field strength of the field $A^{\sigma\lambda}(\tilde{A}^{\sigma\lambda})$ and m is the 'mass' of the respective field.

The irreducible unitary supermultiplet for $p=2$ is called the massless supermultiplet since it contains the massless graviton. In addition to the singlet graviton it contains 42 scalar fields, $48 s=\frac{1}{2}$ fields, 15 vector fields, 6 $A_{\mu\nu}$, 6 $\tilde{A}_{\mu\nu}$ and 8 gravitini. If one were to count the degrees of freedom of this supermultiplet in the limit the anti-de Sitter group is contracted to the Poincaré group in $D=5$ then one finds that they match precisely those of the ungauged $N=8$ supergravity. The only subtlety in this limit involves the equivalence, via duality, of a vector field to an anti-symmetric tensor field. However, the fields that are dual to each other in Minkowski space are, in general, not equivalent to each other in the corresponding anti-de Sitter space. This

† The spectrum of the bosonic sector of the S^5 compactification has been independently obtained by Romans using differential geometrical techniques and agrees with our results [15].

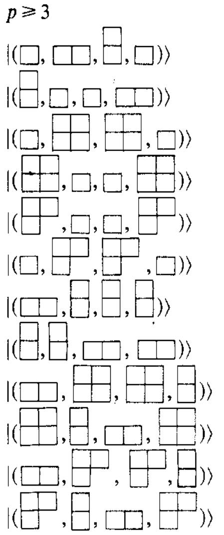
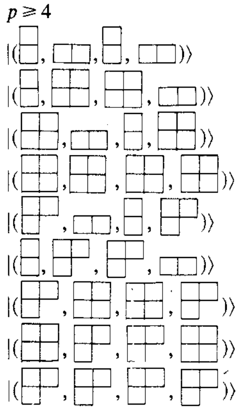
‡ For previous work on self-duality see [18].

phenomenon has already been observed in $D=7$ where anti-symmetric tensor fields of rank 2 and rank 3 are dual to each other in Minkowski space. However, a gauged $N=4$ supergravity in $D=7$ has been constructed involving the anti-symmetric tensor fields of rank 3 [19] whereas the construction using the fields of rank 2 has failed [20]. In fact, the gauged $N=4$ theory in $D=7$ involving tensors of rank 3 has no Poincaré limit [19]. Thus, in analogy with $D=7$, we expect that there be a *gauged* maximally extended simple supergravity in $D=5$ whose fields are precisely those of the massless $(p=2)$ supermultiplet of $U(2,2/4)$. The kinetic energy terms for the scalars of this theory will be a σ-model corresponding to the symmetric space $E_{6(6)}/Sp(8)$ as in the

Table 1. We give the spectrum of the S^5 compactification of the chiral $N=2$, $D=10$ supergravity. The states for a given p, together with their anti-de Sitter excitations, form an irreducible unitary supermultiplet of $U(2,2/4)$. The $p=1$ supermultiplet corresponds to the doubleton supermultiplet and decouples from the spectrum. The $p=2$ supermultiplet is the 'massless' multiplet containing the singlet graviton. We also give the anti-de Sitter energies E_0 of the lowest states of a given UIR of SO(4, 2). The 'mass' of the corresponding field is determined by E_0 and the SO(4) quantum numbers of the lowest state. The last column gives the $U(1)_Y$ quantum numbers of the fields.

$SU(2)_{j_1}\times SU(2)_{k_1}$ $\times SU(2)_{j_2}\times SU(2)_{k_2}$ Young tableau of the lowest states	SO(4) Labels of the lowest states	Anti-de Sitter energy E_0 of the lowest states	Fields of the UIR of U(2, 2/4)	SU(4) Transformation properties of the fields (Dynkin labelling)	$U(1)_Y$ Quantum numbers of the fields
$p\geqslant 1$					
$\lvert 0\rangle$	$(0,0)$	$2p$	$\varphi^{(1)}$	$(0,p,0)$	0
$\lvert(\ydiagram{1},1,1,\ydiagram{1})\rangle$	$(\frac{1}{2},0)$	$2p+1$	$\lambda_+^{(1)}$	$(0,p-1,1)$	$\frac{1}{2}$
$\lvert(1,\ydiagram{1},\ydiagram{1},1)\rangle$	$(0,\frac{1}{2})$	$2p+1$	$\lambda_-^{(1)}$	$(1,p-1,0)$	$-\frac{1}{2}$
$\lvert(\ydiagram{2},1,1,\ydiagram{1,1})\rangle$	$(1,0)$	$2p+2$	$A_{\mu\nu}^{(1)}$	$(0,p-1,0)$	1
$\lvert(1,\ydiagram{1,1},\ydiagram{2},1)\rangle$	$(0,1)$	$2p+2$	$\bar{A}_{\mu\nu}^{(1)}$	$(0,p-1,0)$	-1
$p\geqslant 2$					
$\lvert(\ydiagram{1,1},1,1,\ydiagram{2})\rangle$	$(0,0)$	$2p+2$	$\varphi^{(2)}$	$(0,p-2,2)$	1
$\lvert(1,\ydiagram{2},\ydiagram{1,1},1)\rangle$	$(0,0)$	$2p+2$	$\bar{\varphi}^{(2)}$	$(2,p-2,0)$	-1
$\lvert(\ydiagram{2,2},1,1,\ydiagram{2,2})\rangle$	$(0,0)$	$2p+4$	$\varphi^{(3)}$	$(0,p-2,0)$	2
$\lvert(1,\ydiagram{2,2},\ydiagram{2,2},1)\rangle$	$(0,0)$	$2p+4$	$\bar{\varphi}^{(3)}$	$(0,p-2,0)$	-2
$\lvert(\ydiagram{2,1},1.1,\ydiagram{2,1})\rangle$	$(\frac{1}{2},0)$	$2p+3$	$\lambda_+^{(2)}$	$(0,p-2,1)$	$\frac{3}{2}$
$\lvert(1,\ydiagram{2,1},\ydiagram{2,1},1)\rangle$	$(0,\frac{1}{2})$	$2p+3$	$\lambda_-^{(2)}$	$(1,p-2,0)$	$-\frac{3}{2}$
$\lvert(\ydiagram{1},\ydiagram{1},\ydiagram{1},\ydiagram{1})\rangle$	$(\frac{1}{2},\frac{1}{2})$	$2p+2$	$A_\mu^{(1)}$	$(1,p-2,1)$	0
$\lvert(\ydiagram{2},\ydiagram{1},\ydiagram{1},\ydiagram{1,1})\rangle$	$(1,\frac{1}{2})$	$2p+3$	$\psi_{+\mu}^{(1)}$	$(1,p-2,0)$	$\frac{1}{2}$
$\lvert(\ydiagram{1},\ydiagram{1,1},\ydiagram{2},\ydiagram{1})\rangle$	$(\frac{1}{2},1)$	$2p+3$	$\psi_{-\mu}^{(1)}$	$(0,p-2,1)$	$-\frac{1}{2}$
$\lvert(\ydiagram{2},\ydiagram{1,1},\ydiagram{2},\ydiagram{1,1})\rangle$	$(1,1)$	$2p+4$	$h_{\mu\nu}$	$(0,p-2,0)$	0

Table 1 (continued)

$SU(2)_{j_1} \times SU(2)_{k_1} \times SU(2)_{j_2} \times SU(2)_{k_2}$ Young tableau of the lowest states	SO(4) Labels of the lowest states	Anti-de Sitter energy E_0 of the lowest states	Fields of the UIR of U(2, 2/4)	SU(4) Transformation properties of the fields (Dynkin labelling)	$U(1)_Y$ Quantum numbers of the fields
$p \geqslant 3$					
	$(\frac{1}{2}, 0)$	$2p+3$	$\lambda_+^{(3)}$	$(2, p-3, 1)$	$-\frac{1}{2}$
	$(0, \frac{1}{2})$	$2p+3$	$\lambda_-^{(3)}$	$(1, p-3, 2)$	$\frac{1}{2}$
	$(\frac{1}{2}, 0)$	$2p+5$	$\lambda_+^{(4)}$	$(0, p-3, 1)$	$-\frac{3}{2}$
	$(0, \frac{1}{2})$	$2p+5$	$\lambda_-^{(4)}$	$(1, p-3, 0)$	$\frac{3}{2}$
	$(\frac{1}{2}, \frac{1}{2})$	$2p+4$	$A_\mu^{(2)}$	$(1, p-3, 1)$	1
	$(\frac{1}{2}, \frac{1}{2})$	$2p+4$	$\bar{A}_\mu^{(2)}$	$(1, p-3, 1)$	-1
	$(1, 0)$	$2p+4$	$A_{\mu\nu}^{(2)}$	$(2, p-3, 0)$	0
	$(0, 1)$	$2p+4$	$\tilde{A}_{\mu\nu}^{(2)}$	$(0, p-3, 2)$	0
	$(1, 0)$	$2p+6$	$A_{\mu\nu}^{(3)}$	$(0, p-3, 0)$	-1
	$(0, 1)$	$2p+6$	$\tilde{A}_{\mu\nu}^{(3)}$	$(0, p-3, 0)$	1
	$(1, \frac{1}{2})$	$2p+5$	$\psi_{+\mu}^{(2)}$	$(1, p-3, 0)$	$-\frac{1}{2}$
	$(\frac{1}{2}, 1)$	$2p+5$	$\psi_{-\mu}^{(2)}$	$(0, p-3, 1)$	$\frac{1}{2}$
$p \geqslant 4$					
	$(0, 0)$	$2p+4$	$\varphi^{(4)}$	$(2, p-4, 2)$	0
	$(0, 0)$	$2p+6$	$\varphi^{(5)}$	$(0, p-4, 2)$	-1
	$(0, 0)$	$2p+6$	$\bar{\varphi}^{(5)}$	$(2, p-4, 0)$	1
	$(0, 0)$	$2p+8$	$\varphi^{(6)}$	$(0, p-4, 0)$	0
	$(\frac{1}{2}, 0)$	$2p+5$	$\lambda_+^{(5)}$	$(2, p-4, 1)$	$\frac{1}{2}$
	$(0, \frac{1}{2})$	$2p+5$	$\lambda_-^{(5)}$	$(1, p-4, 2)$	$-\frac{1}{2}$
	$(\frac{1}{2}, 0)$	$2p+7$	$\lambda_+^{(6)}$	$(0, p-4, 1)$	$-\frac{1}{2}$
	$(0, \frac{1}{2})$	$2p+7$	$\lambda_-^{(6)}$	$(1, p-4, 0)$	$\frac{1}{2}$
	$(\frac{1}{2}, \frac{1}{2})$	$2p+6$	$A_\mu^{(3)}$	$(1, p-4, 1)$	0

ungauged $N=8$ theory [21, 16]. However, both the $E_{6(6)}$ and Sp(8) symmetries will be broken by the other terms in the Lagrangian. It will have a local non-Abelian gauge group SU(4) whose gauge fields will be the 15 vector fields of the massless multiplet. We are presently working on the construction of this theory [22].

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