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THE VANISHING VOLUME OF $N = 1$ SUPERSPACE

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ABSTRACT

We show that the supergeometry of the new minimal formulation of $N = 1$ supergravity is a limiting case of the non-minimal supergeometry for which the volume of superspace vanishes. The action for supergravity in this case is then no longer the superdeterminant of the vielbein but is analogous to a special form of the Fayet-Iliopoulos term allowed only for the new minimal supergeometry. We discuss the restrictions on matter couplings imposed by the new and non-minimal supergeometries.

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1. INTRODUCTION

It is now known that there are three inequivalent sets of auxiliary fields for $N = 1$ supergravity; the minimal set¹⁾ (A_m, S, P) , the non-minimal set²⁾ $(A_m, S, P, V_m, A'_m, \lambda, \chi)$, and the new minimal set³⁾ (A_m, t_{mn}) (where $t_{mn} = -t_{nm}$ and A_m are both gauge fields)*). Each set of auxiliary fields corresponds to a particular choice of torsion constraints on the curved superspace geometry^{5),6)}. All have in common a set of conventional and representation preserving constraints, which admit a super-Weyl invariance with an unconstrained complex scalar superfield parameter⁷⁾. They are distinguished by their "type (iii)" or conformal constraint⁷⁾ which restricts the super-Weyl parameter in a particular way. In the minimal case the parameter is chiral, in the non-minimal case it is complex linear, and in the new minimal case it is the sum of a real linear and a real general scalar superfield.

The supergeometry of the minimal and non-minimal cases is well understood^{5),6)}. In this letter we show that the new minimal case is a limit of the non-minimal one in which the volume of superspace goes to zero. This fact necessitates a different form of the supergravity action which is analogous to the special form of the Fayet-Iliopoulos (F-I) term allowed only for the new minimal supergeometry. We also discuss the restrictions on the couplings of chiral fields in the new and non-minimal cases. Unlike the minimal case only one power of a chiral field, e.g. ϕ^m , is allowable.

2. LORENTZ SUPERSPACE

We shall first discuss the general supergeometry in the context of the usual superspace for which the tangent space group is just $SL(2, C)$, the Lorentz group. The conventional and representation preserving constraints as the torsions are collectively,

$$\begin{aligned} T_{\alpha\beta}{}^c &= 0 & T_{\alpha\dot{\beta}}{}^c &= -i \delta_{\alpha\dot{\beta}}^c \\ T_{\alpha\beta}{}^{\dot{\gamma}} &= T_{\alpha\dot{\beta}}{}^{\gamma} = 0, & T_{\alpha\dot{\beta}}{}^{\dot{\gamma}} &= -\delta_{\dot{\beta}}^{\dot{\gamma}} T_{\alpha}{}^{\dot{\gamma}} \\ T_{\alpha b}{}^c &= \delta_b^c T_{\alpha}{}^c, & T_{ab}{}^c &= 0 \end{aligned} \quad (2.1)$$

These are invariant under the super-Weyl transformation

$$\delta E_{\alpha}{}^M = -S E_{\alpha}{}^M \quad (2.2)$$

*⁴⁾ The linearized new minimal set had been found earlier⁴⁾.

where S is a complex scalar superfield. (E_a^M and the connections are expressible in terms of E_α^M through the conventional constraints.) It follows from this invariance that the solution to the constraints (2.1) involves a complex scalar superfield A , whose super-Weyl transformation, $\delta A = S$, is responsible for the transformation (2.2). Under this transformation the superdeterminant, $\text{Ber } E_M^A \equiv E$, transforms as

$$\delta E = 2(S + \bar{S})E \quad (2.3)$$

E is the only geometrical invariant of the right dimension to be the supergravity Lagrangian, but from (2.3) and $\delta A = S$, we see that without some further constraint the A field equation would imply $E = 0$. Therefore a further, type (iii), constraint must be imposed to restrict the parameter S , and hence the field A . For the minimal case this constraint is

$$T_\alpha = 0 \quad (2.4)$$

which implies $D_\alpha(2S + \bar{S}) = 0$, thus leaving a chiral super-Weyl invariance⁸⁾. For the non-minimal case there is a one parameter family of type (iii) constraints⁶⁾:

$$\bar{R} - \zeta D^\alpha T_\alpha - 2\zeta^2 T^\alpha T_\alpha = 0 \quad (2.5)$$

where $\bar{R} = \frac{1}{6} R_{\alpha\beta}^{\alpha\beta}$ is chiral, $D_\alpha \bar{R} = 0$. In this case the parameter S is restricted by

$$\mathcal{D}^2(\zeta) [(1 - 2\zeta)S - \zeta\bar{S}] = 0 \quad (2.6)$$

where,

$$\mathcal{D}^2(a) = D^2 - 2\bar{R} + 2a D^\alpha T_\alpha + 4a^2 T^\alpha T_\alpha + 4a T^\alpha D_\alpha \quad (2.7)$$

For general ζ Eq. (2.6) implies that the super-Weyl parameter is complex linear⁹⁾, but for the special case $\zeta = 1$, $S - \bar{S}$ is unconstrained while $S + \bar{S}$ is constrained to be real linear, i.e., $\mathcal{D}^2(a=1)(S + \bar{S}) = 0$. In fact the equation $\mathcal{D}^2(a)L = 0$ is consistent for real L only if $a = 1$. It is precisely this special case that corresponds to the new minimal auxiliary fields. To see this suppose that the action is $U(1)$ invariant; then only $A + \bar{A} = G$, a real

linear superfield, will appear. G contains a scalar and spinor, which compensate for Weyl and S supersymmetry transformations, and the field strength for the antisymmetric tensor gauge field, t_{mn} . The integral of E is indeed $U(1)$ invariant, as shown by Eq. (2.3), but in this special case it cannot be the action, as it is in both minimal¹⁰⁾ and non-minimal^{6),11)} cases, because $\int E$ vanishes. This can be seen as follows:

A chiral superfield, $\bar{\Sigma}$, may be defined to satisfy

$$\mathcal{D}_\alpha \bar{\Sigma} \equiv (D_\alpha + 2a T_\alpha) \bar{\Sigma} = 0 \quad (2.8)$$

from which it follows that

$$\bar{\Sigma} = \mathcal{D}^2(a) X \quad (2.9)$$

This can be written as

$$\bar{\Sigma} = (D^\alpha + m T^\alpha)(D_\alpha + (4a-m)T_\alpha) X - M X \quad (2.10)$$

where,

$$M = 2\bar{R} - (m-2a)D^\alpha T_\alpha - (m-2a)^2 T^\alpha T_\alpha \quad (2.11)$$

Since

$$\int d^8z E (D^\alpha + 2T^\alpha) V_\alpha = 0$$

by putting $m = 2$ in (2.10), integrating with d^8z , and then setting $a = 0$, we arrive at

$$\int E \bar{\Sigma}(a=0) d^8z = -2(5-1) \int E \{ D^\alpha T_\alpha + 2(5+1) T^\alpha T_\alpha \} X \quad (2.12)$$

But for $a = 0$, $\bar{\epsilon} = 1$ is a particular chiral field. Making this choice we see that fE vanishes with $\zeta = 1$. This means that the new minimal supergravity action cannot be expressed in purely geometrical terms. To discuss the action and verify that it leads to the correct field equations it is simpler to work in a different superspace that is appropriate to this special case; viz. one for which the tangent space group contains a chiral $U(1)$ factor. This superspace has been used previously to discuss conformal supergravity¹²⁾ and also the new minimal Poincaré theory¹³⁾.

3. $U(1)$ SUPERSPACE

We now take the action of the tangent space group on a vector $V^A = (V^a, V^\alpha, V^{\dot{\alpha}})$ to be

$$\delta V^A = V^B \hat{L}_B{}^A \quad (3.1)$$

where,

$$\begin{aligned} \hat{L}_\rho{}^\alpha &= L'_\rho{}^\alpha + \delta_\rho{}^\alpha L \\ \hat{L}_b{}^a &= (\delta_b{}^a)_\alpha{}^\beta L'_\beta{}^\alpha + c.c. \\ \hat{L}_{\dot{\rho}}{}^{\dot{\alpha}} &= -(\overline{\hat{L}_\rho{}^\alpha}) ; \quad \bar{L} = -L \end{aligned} \quad (3.2)$$

Here, L is the $U(1)$ parameter and $L'_\alpha{}^\beta$ the $SL(2, C)$ parameter. [Care should be taken in identifying the $U(1)$ charge of a given field since the ϵ tensor used to raise and lower spinor indices is $U(1)$ inert.] Corresponding to the transformations (3.2) the connection has the form

$$\hat{\Omega}_M{}^\alpha{}_\beta = \Omega'_M{}^\alpha{}_\beta + \delta_\alpha{}^\beta \Omega_M \quad (3.3)$$

The curvatures and torsions computed with respect to $\hat{\Omega}$ and Ω' are related by^{*)}

$$\begin{aligned} \hat{R}_{AB}{}^\delta{}_\gamma &= R'_{AB}{}^\delta{}_\gamma + \delta_\gamma{}^\delta F_{AB} \\ \hat{T}_{AB}{}^C &= T'_{AB}{}^C + \Omega_A I_B{}^C - (-1)^{AB} \Omega_B I_A{}^C \end{aligned} \quad (3.4)$$

*) The vielbein $E_M{}^A$ and the Lorentz connection $\Omega'_M{}^\alpha{}_\beta$ are related to those of Section 2 by field redefinitions.

where F_{AB} is the $U(1)$ field strength and $I_A^B = (0, \delta_\alpha^B, \delta_{\dot{\alpha}}^{\dot{B}})$. Imposing conventional and representation preserving constraints the only non-vanishing torsions are

$$\begin{aligned}\hat{T}_{\alpha\dot{\beta}}^{\dot{\gamma}} &= -i\delta_{\alpha\dot{\beta}}^{\dot{\gamma}} \\ \hat{T}_{\alpha\dot{\alpha},\beta,\dot{\beta}}^{\dot{\gamma}} &= i\epsilon_{\alpha\beta}\epsilon_{\dot{\alpha}\dot{\beta}}\hat{R}^{\dot{\gamma}} \\ \hat{T}_{\alpha\dot{\alpha},\beta,\dot{\beta},\gamma} &= -2i\epsilon_{\alpha\gamma}G_{\beta\dot{\beta}} \\ \hat{T}_{\alpha\dot{\alpha},\beta\dot{\beta},\dot{\gamma}} &= \epsilon_{\alpha\beta}[\bar{\psi}_{\dot{\beta}\dot{\gamma}} + \epsilon_{\dot{\beta}\dot{\gamma}}\bar{\psi}_{\dot{\beta}} + \epsilon_{\dot{\beta}\dot{\gamma}}\bar{\psi}_{\dot{\gamma}}] + \epsilon_{\dot{\beta}\dot{\gamma}}\psi_{\alpha\beta}\end{aligned}\quad (3.5)$$

where,

$$\begin{aligned}\bar{D}_{\dot{\alpha}}\hat{R} &= 0 \\ \hat{D}_{(\alpha}G_{\beta)\dot{\beta}} &= -\frac{1}{4}\psi_{\alpha\beta\dot{\beta}} \\ \hat{D}^{\alpha}G_{\alpha\dot{\beta}} &= -\frac{3}{2}\bar{\psi}_{\dot{\beta}} + \bar{D}_{\dot{\beta}}\hat{R} \\ \hat{D}_{\alpha}\bar{\psi}_{\dot{\alpha}\beta\dot{\gamma}} &= 0\end{aligned}\quad (3.6)$$

$\psi_{\alpha\beta\dot{\gamma}}$ is totally symmetric and $G_{\alpha\dot{\beta}}$ is Hermitian. The ψ 's are the $SL(2,C)$ irreducible parts of the Rarita-Schwinger field strength. These constraints determine algebraically not only the Lorentz connection but also the $U(1)$ connection; in fact $\Omega_{\alpha} = -T_{\alpha}$ where T_{α} is defined in (2.1). The curvatures may be found using the Bianchi identities and were given in Ref. 12). In particular,

$$\begin{aligned}F_{\alpha\beta} &= 0 \\ F_{\alpha\dot{\beta}} &= -3G_{\alpha\dot{\beta}}\end{aligned}\quad (3.7)$$

These constraints are also invariant under a super-Weyl transformation of the form of Eq. (2.2), but the parameter S is now real. The type (iii) constraint is

$$\hat{R} = 0 \quad (3.8)$$

which restricts S to be real linear, $\hat{D}^2S = 0$. The Bianchi identities now tell us that $D^{\alpha\dot{\alpha}}G_{\alpha\dot{\alpha}} = 0$ and this equation implies in turn that $G_{\alpha\dot{\alpha}}|_{\theta=0}$ is the supercovariantized field strength for an antisymmetric tensor gauge field (i.e., $G_{\alpha\dot{\alpha}}$ includes terms quadratic in the spin $3/2$ field). Indeed one can show that there is a superspace two-form potential, t , with field strength

$G = dt$, whose components G_{ABC} satisfy $G_{\alpha\beta\gamma} = G_{\alpha\beta\dot{\gamma}} = G_{\alpha\dot{\beta}\gamma} = 0$, $G_{\alpha\dot{\beta}\dot{\gamma}} = i(\sigma_a)_{\beta\dot{\gamma}} \dot{\gamma}$, $G_{abc} \propto \epsilon_{abcd} G^d$. These constraints on G_{ABC} are consistent with those of (3.5) and (3.8) and with the Bianchi identity $dG = 0$.

Note now that the equation $F_{\alpha\beta} = 0$ implies that

$$\Omega_\alpha = -T_\alpha = -D'_\alpha T \quad (3.9)$$

so that T is the $U(1)$ prepotential, [although it is further constrained by the type (iii) constraint $\hat{R} = 0$]. Associated with T there is a pre-gauge transformation

$$T \rightarrow T + \bar{C}, \quad \hat{D}_\alpha \bar{C} = 0 \quad (3.10)$$

Using T we can write an action for supergravity. It is

$$I = \int d^4z E' (T + \bar{T}) \quad (3.11)$$

where $E' = \text{Ber}(E'_M{}^A)$ ($= E$ of Section 2). This action is invariant under the transformation (3.10) because the integral of a chiral field vanishes. The action (3.11) is analogous to the F-I term which we will discuss in the next section. To show that it is indeed the correct one we solve for the variations of the vielbein and the field T from $\delta T_{AB}^{C(10)}$. Defining

$$H'_A{}^B = E'_A{}^M \delta E'_M{}^B \quad (3.12)$$

one has

$$\delta I = \int E' \left\{ H'_A{}^A (-1)^A (T + \bar{T}) + \delta(T + \bar{T}) \right\} \quad (3.13)$$

After some algebra we find

$$\begin{aligned} H'_A{}^A &= -\frac{1}{2} [\hat{D}_\alpha, \bar{\hat{D}}_{\dot{\alpha}}] V^{\alpha\dot{\alpha}} + 2L \\ \delta(T + \bar{T}) &= -2V^{\alpha\dot{\alpha}} G_{\alpha\dot{\alpha}} - 3L + \frac{1}{2} V^{\alpha\dot{\alpha}} [\hat{D}_\alpha, \bar{\hat{D}}_{\dot{\alpha}}] (T + \bar{T}) \end{aligned} \quad (3.14)$$

where $V^{\alpha\dot{\alpha}}$ is real and L is real linear. Substituting (3.14) into (3.13) we find that the $V^{\alpha\dot{\alpha}}$ variation yields the field equation

$$G_{\alpha\dot{\alpha}} = 0 \quad (3.15)$$

while the L variation yields only a derivative of this as a consequence of the compensated conformal invariance of the action. Equation (3.15) is the equation of motion of new minimal supergravity. The linearized action and field equations in superspace were found previously in Ref. 15).

We will now return to the Lorentz superspace of Section 2 in order to compare matter couplings with new minimal fields to those with minimal or non-minimal.

4. MATTER COUPLING

Flat space matter models are described either by chiral, ϕ , or real linear, G , superfields for (propagating) superspin 0, or by a real scalar gauge superfield, V , for (propagating) superspin 1/2. Those actions that are expressed as integrals over the whole superspace, such as the kinetic terms for ϕ or G , are easily extended to curved superspace by covariantization of derivatives and multiplication by the density E . In the new and non-minimal cases there is an ambiguity in the definition of $\int d^8z E \phi \bar{\phi}$ corresponding to an arbitrariness in the definition of a covariantly chiral field, i.e., $(D_\alpha + 2bT_\alpha)\bar{\phi} = 0$ for arbitrary b . Different choices of the constant b correspond to different choice for the coefficients of improvement terms in the matter stress tensor. These coefficients are also charged by the addition to the action of quantities that vanish in flat space such as $\int d^8z E(\phi^2 + \bar{\phi}^2)$. There is no such arbitrariness in the meaning of $\int E G^2$, however, because the curved space definition of G is uniquely $D^2(a=1)G = 0$.

For a vector superfield V the F-I term $\int d^8z V$ can be directly extended to $\int d^8z EV$ but the invariance under the gauge transformation $V \rightarrow V + i(\Lambda - \bar{\Lambda})$, $\bar{D}_\alpha \Lambda = 0$, is not necessarily preserved. $\int d^8z EV$ will be gauge invariant if the integral of a chiral field (defined with $b = 0$) vanishes. From Eq. (2.12) we see that this is the case for the new minimal constraint ($\zeta=1$), as has been noted previously in a tensor calculus for the new minimal fields¹³⁾, and only in this case. In the minimal case, although $\int EV d^8z$ is not gauge invariant, the action $\int E e^V d^8z$ is invariant under a combined gauge and super-Weyl transformation¹⁵⁾. This is possible because the super-Weyl parameter is chiral in this case. In the non-minimal case, however, the super-Weyl parameter is complex linear and so this escape is not possible.

To consider curved space extensions of mass and interaction terms for chiral fields which are expressed as integrals over the chiral subspace we need a chiral measure. Consider a chiral Lagrangian Σ , satisfying (2.8) such that it can be written as $\bar{\Sigma} = \mathbb{D}^2(a)X$. An invariant is

$$I_{\Sigma} = \int d^8z \bar{\Sigma} X + c.c. \quad (4.1)$$

but we must be able to express this locally in terms of Σ . To this end note that since $\bar{\Sigma}$ can be written in terms of X as in Eq. (2.10) we have, dividing by M ,

$$X = -\frac{\bar{\Sigma}}{M} - \frac{(\mathbb{D}^{\alpha} + (4a-m)T^{\alpha})X (\mathbb{D}_{\alpha} + (m-2)T_{\alpha})M}{M^2} \quad (4.2)$$

+ surface term

Hence

$$I_{\Sigma} = - \int d^8z \frac{\bar{\Sigma}}{M} \Sigma + c.c. \quad (4.3)$$

with M given by (2.11), provided that $(\mathbb{D}_{\alpha} + (m-2)T_{\alpha})M$ vanishes. This condition is satisfied with the new or non-minimal type (iii) constraint [or without using any type (iii) constraint] for arbitrary m only if $a = 1$. Thus, we require that Σ in the chiral action (4.3) satisfy

$$(\mathbb{D}_{\alpha} + 2T_{\alpha})\bar{\Sigma} = 0 \quad (4.4)$$

Imposing the constraint (2.5), M can be written as

$$M = (2\zeta + 2-m) [\mathbb{D}^{\alpha} T_{\alpha} + (2\zeta - 2+m) T^{\alpha} T_{\alpha}] \quad (4.5)$$

so that for any choice of ζ there is always a choice of m such that M is non-vanishing. Imposing instead the constraint $T_{\alpha} = 0$ yields $M = 2\bar{R}$ so that we recover the usual chiral measure of the minimal case^{16),6)}.

Consider now the case that $\Sigma = f(\phi)$ so that $(D_\alpha + 2T_\alpha)\bar{f}(\bar{\phi}) = 0$. This is consistent with the assignment of a definite chiral weight, b , to ϕ , i.e.,

$$(D_\alpha + 2bT_\alpha)\bar{\phi} = 0 \quad (4.6)$$

only if $f(\phi) = \phi^n$ with $nb = 1$. If we require that the kinetic term for ϕ be $U(1)$ invariant then it follows that ϕ must have a definite chiral weight. Hence for the new minimal case where $U(1)$ invariance is required only one power of ϕ is allowed in a chiral action. For example, in order to have a mass term ϕ^2 , we must choose $b = 1$, which then excludes self-interactions.

It might appear that the non-minimal fields avoid this restriction on chiral actions because $U(1)$ invariance is not required. However, this is not the case as one sees from the following argument. Define a field ϕ_0 satisfying $D_\alpha \bar{\phi}_0 = 0$. In terms of ϕ_0 , an $f(\phi)$ satisfying $(D_\alpha + 2T_\alpha)\bar{f}(\bar{\phi}) = 0$ is given by

$$\bar{f}(\bar{\phi}) = e^{-2T} \bar{f}(\bar{\phi}_0) \quad (4.7)$$

where T is defined through $T_\alpha = D_\alpha T$. The chiral action for $f(\phi)$ is therefore

$$\int d^4x \frac{E}{M} e^{-2T} \bar{f}(\bar{\phi}_0) + \text{c.c.} \quad (4.8)$$

But T is not a covariant object and transforms under pre-gauge transformations by a chiral field $T \rightarrow T + \bar{C}$, $D_\alpha \bar{C} = 0$. Hence the chiral action (4.8) will not be pre-gauge invariant unless ϕ_0 also transforms. However, this would spoil the pre-gauge invariance of the kinetic terms unless f is a monomial. On the other hand if the action is not pre-gauge invariant then an additional chiral field appears in the theory, but only in the chiral action (4.8), thus imposing the constraint $f(\phi_0) = 0$. Therefore the only allowed chiral interactions in both the new minimal and non-minimal cases are monomials. Of course, other restrictions will apply if there is more than one chiral field.

5. COMMENTS

We have seen that the new minimal and non-minimal supergravity fields are both very restrictive in the matter couplings that they allow, in contrast to the minimal fields. This makes the minimal fields more useful for most purposes. In particular, in the quantum theory the $U(1)$ invariance of the new minimal formulation will be broken by chiral anomalies so that it is not likely to be useful for quantum calculations. However, the restrictions on couplings for the non-minimal fields might have interesting consequences for the quantum theory if one can quantize covariantly. For example, suppose we have a self-interaction term ϕ^4 for a chiral field ϕ , then one would expect all higher powers of ϕ to appear as counterterms. But there are no such counterterms that are invariant off-shell, although of course they are invariant on-shell. This is therefore an example of an on-shell invariant that cannot be extended off-shell (within the given quantization scheme, i.e., auxiliary field set).

We have also seen that the new minimal fields are really a special case of the non-minimal fields where the parameter $\zeta = 1$. It is likely that in the quantum theory this parameter will be renormalized, in general by an infinite amount. The special value $\zeta = 1$ would be partially, (i.e., at one-loop) stabilized by the $U(1)$ invariance of the new minimal case, but only partially because of the chiral anomaly. Thus in the quantum theory the new minimal and non-minimal cases are essentially equivalent.

The fact that the superspace volume vanishes for the new minimal supergeometry is an interesting echo of the vanishing volume of $N = 2$ superspace¹⁷⁾, suggests that perhaps the known $N = 2$ supergravity auxiliary fields are also a limiting case of a larger set. The non-geometrical form of the new minimal action also underlines the importance of the alternative approaches to superspace geometry¹⁸⁾.

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