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OFF-SHELL $d = 10$, $N = 1$ POINCARÉ SUPERGRAVITY
AND THE EMBEDDIBILITY OF HIGHER DERIVATIVE FIELD THEORIES
IN SUPERSPACE

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A B S T R A C T

We present the constraints on the supertorsion which lead to off-shell $d = 10$, $N = 1$ Poincaré supergravity and identify the superfields which are known to enter the superspace action. We point out that Poincaré off-shell constraints have to be used in a superspace formulation of the low energy effective actions of $N = 1$ superstrings. However, current formulations of $N = 1$ superstrings propagating in arbitrary supergravity backgrounds are incompatible with these off-shell constraints.

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The discovery by Green and Schwarz¹⁾ that the Chapline-Manton theory^{2),3)} can be rid of its anomalies (at one loop) sparked off a tremendous interest in this model and the related $N = 1$ superstrings^{4),5)}. Strictly speaking, however, anomaly freedom was not achieved by adding Wess-Zumino counter-terms to the Chapline-Manton Lagrangian, but rather to a slightly modified Lagrangian. An understanding of the origin of this new Lagrangian and the mechanism by which the anomalies are cancelled in Ref. 1) is obtained by recalling that the Chapline-Manton theory is the low-energy limit of an anomaly-free superstring. This was, in fact, how the mechanism was first discovered^{6),1)}. The necessary modifications of the Chapline-Manton Lagrangian are then seen to be part of the first-order term (pure $N = 1$ supergravity being the zeroth order term) in a power series expansion of the effective action of the $N = 1$ open superstring⁷⁾ or the heterotic string⁵⁾. [For definitions and a quite comprehensive account of various approaches available for calculating field theory (low-energy) effective actions, see Ref. 8).]

Our present knowledge of the form of these effective actions, usually expanded in terms of α' , the inverse string tension, or numbers of derivatives, is unfortunately extremely rudimentary [see, however, Refs. 9) and 8) for some results valid to all orders in α']. There are three approaches for obtaining the effective action usually considered. One may either: (i) expand, in a (covariant) second-quantized approach, the string functional in some basis and eliminate the massive modes in the field theory by means of their field equations^{10),8)}, or (ii) calculate it directly from a two-dimensional σ -model à la Fradkin and Tseytlin¹¹⁾ or deduce it from the β functions as done by Friedan¹²⁾ and Callan et al.¹³⁾, or (iii) determine the necessary terms in the effective field theory of the massless modes by requiring that they reproduce the scattering amplitudes derived from the string^{14),15)}. [Note that the relation between the two approaches mentioned under (ii) is at the moment very poorly understood.] Most $O(\alpha')$ results have so far been obtained by methods (ii) and (iii), while the results to all orders in α' ^{9),8)} were obtained by method (ii). (As far as we know, the first method has not yet been used in practice).

In a recent paper¹⁶⁾, the possibility was considered that in the case of supersymmetric strings a fourth method might exist based on solving superspace Bianchi identities. This idea was substantiated to a certain degree in Ref. 17).

It was also mentioned by Dahr et al. in Ref. 18). In this letter, we will first point out a problem with this method as it was presented in Ref. 17). [Our arguments are, of course, also valid for the slightly different, but equivalent, version of Ref. 18).] We then proceed to show how to resolve it. This involves giving the completely off-shell formulation of $d = 10$, $N = 1$ Poincaré supergravity. [This should be compared to the off-shell results obtained in Ref. 19); see below.] After some comments on the other version of this theory, which involves a six-form potential instead of a two-form potential, we discuss the previous results in relation to Siegel's fermionic symmetry^{20),21)} and to $N = 1$ superstrings propagating in arbitrary backgrounds²²⁾ [see also Refs. 23)-25)]^{*)}.

Let us start by briefly discussing superspace as a method of deriving superstring-like field theories. By string-like field theories we will here mean theories whose Lagrangians contain a set, possibly infinite, of higher derivative terms. The method to be described below does seem capable of generating such theories, but it should be emphasized that so far the only connection to string theories is through the Chapline-Manton theory. [However, for the purpose of this paper, we need only assume that the effective actions of the anomaly-free heterotic strings and $N = 1$ open superstring can be formulated in $d = 10$ chiral superspace²⁹⁾; it is not necessary to assume that superspace provides a way of deriving them uniquely.]

The basic idea is to impose constraints on the superfield strengths appearing in the theory, substitute these into the superspace Bianchi identities, and finally solve the equations so obtained. This is by now quite an old technique³⁰⁾ and it has been applied to most of the ordinary supergravity theories. However, the one we are interested in here, namely the Chapline-Manton theory^{2),3)}, was not constructed in this way until quite recently¹⁶⁾⁻¹⁸⁾. This construction is based on the superspace Bianchi identity $dH = \text{tr} F^2$ ¹⁶⁾, where H is the three-form field strength and F is the Yang-Mills field strength. To take the step to string-like theories, one should add a $\text{tr} R^2$ -term together with some terms

*) Neglecting higher derivative terms for the moment, we want to make the following remark. The propagation of $N = 1$ superstrings in a Chapline-Manton background (i.e., with non-zero Yang-Mills fields) is now understood. The key point is that the two-dimensional σ -model action (with Wess-Zumino term) invariant under Yang-Mills transformations (with a non-invariant two-form potential) contains terms involving the square of the Yang-Mills potential. This can be deduced either by checking the invariance explicitly²⁵⁾, or by spontaneous compactification of a higher dimensional theory on $d = 10$ space-time times a group manifold as described in Refs. 26) and 27) [see also Ref. 28)].

involving dH and H^2 ¹⁷⁾ on the right-hand side of this equation, and find the constraints that makes it possible to solve the new equations. As explained in Refs. 17) and 18), these equations must now be solved iteratively since H appears on both sides of the dH equation.

Whether this technique will prove useful is still to be demonstrated. However, there is an immediate problem with this approach as described in Refs. 17) and 18): the torsion constraints used in these two references are incompatible with $(R_{\mu\nu\rho\sigma})^n$ -terms in the effective action (at least for $n > 2$). The reason for this is as follows. As shown in Ref. 19), the supertorsion constraints²⁹⁾

$$T_{ab}{}^\gamma = -i(\gamma^\gamma)_{ab} \quad , \quad (1)$$

$$T_{ab}{}^c = i \delta_{(a}^c T_{b)} \quad , \quad (2)$$

$$T_{a\rho}{}^\gamma = i \delta_\rho^\gamma \hat{T}_a \quad , \quad (3)$$

$$T_{\alpha\beta}{}^\gamma = 0 \quad , \quad (4)$$

together with the torsion Bianchi identity

$$DT^A = -E^B{}_\Lambda R_B{}^A \quad (5)$$

are off-shell save for the (component) equation^{*)}

$$R + 18 \square \phi - 72 (D_\alpha \phi)^2 + 432 Y_{\alpha\beta\gamma} Y^{\alpha\beta\gamma} = 0 \quad . \quad (6)$$

*) The same is true for the constraints used in Ref. 18), although the equation one derives is then slightly different from (6) but has the same implications. Note also that (6) is essentially the condition for vanishing β^ϕ function, which is usually derived by requiring conformal invariance of the corresponding σ -model¹³⁾. The off-shell theory defined by (1) to (4) is also related to conformal supergravity in $d = 10$.

Here $A, B, \dots$ are flat superspace indices which split into bosonic ones $\alpha, \beta, \dots$ and fermionic ones $a, b, \dots$ [our conventions are those of Refs. 16), 17) and 19)]. In (6) R is the scalar curvature and ϕ is the scalar field in the supergravity multiplet. $Y_{\alpha\beta\gamma}$ is an arbitrary third-rank antisymmetric tensor field. In the Chapline-Manton theory, there is, of course, a combination of the ϕ field equation and the trace of the Einstein equation that does not contain the square of the Yang-Mills field strength; as can be seen from Ref. 17), this combination is just Eq. (6). However, if one adds $(R_{\mu\nu\rho\sigma})^n$ -terms ($n > 2$) to the Lagrangian, it seems impossible that any linear combination of these two equations could have the structure of (6). Note that additional $(H_{\mu\nu\rho})^n$ -terms could easily be accounted for by the $(Y_{\alpha\beta\gamma})^2$ -term in (6). [This is true for terms containing both H and R as well. For instance, we might have $Y_{\alpha\beta\gamma} \sim H_{\alpha\beta\gamma} + (H^3)_{\alpha\beta\gamma} + (RH)_{\alpha\beta\gamma} + \dots$].

The way out of this dilemma is to find a relaxed set of torsion constraints which leads to a completely off-shell formulation of $N = 1$ Poincaré supergravity [i.e., which does not imply (6)]. The known off-shell (super)field content of the linearized theory³¹⁾ should then emerge. In fact, taking advantage of having this knowledge, it is easy to guess the relaxed constraints. The completely off-shell torsion constraints are [see, however, the footnote following Eq. (17)]:

$$T_{ab}{}^\gamma = -i(\gamma^\gamma)_{ab} + i(\gamma^{\alpha_1 \dots \alpha_5})_{ab} T_{\alpha_1 \dots \alpha_5}{}^\gamma, \quad (7)$$

$$T_{ab}{}^c = 0, \quad (8)$$

$$T_{\alpha\beta}{}^\gamma = 0, \quad (9)$$

leaving for the moment $T_{a\beta}{}^\gamma$ arbitrary. In (7), $T_{\alpha_1 \dots \alpha_5}{}^\delta$ should be constrained to the $SO(1,9)$ representation 1050 contained in the product of 126 and 10 (represented by the self-dual $[\alpha_1 \dots \alpha_5]$ and δ respectively). Furthermore, here we have chosen the conventional constraints²⁹⁾ such that T_a and $\hat{T}_a$ in (2) and (3) now both appear in $T_{a\beta}{}^\gamma$ so that (8) might be imposed. [Note that the torsion components $T_{ab}{}^c$ and $T_{\alpha\beta}{}^\gamma$ do not enter the variation of the σ -model Lagrangian describing the propagation of superstrings in an arbitrary background; see Eq. (24) below.]

The reason for the constraint (7) is the following. We know from Ref. 19) that the solution of (5) using the old constraints (1) to (4) is lacking one of the scalar superfields appearing in the off-shell analysis of the linear theory carried out by Howe et al.³¹⁾. This scalar superfield (denoted by W from now on) has as its last component the auxiliary scalar field which invalidates (6). Furthermore, at the θ level corresponding to T_{ab}^Y (i.e., the 12th θ level), this superfield contains the two representations 1050 and 770 [the complete content of a scalar superfield is given in Ref. 32)]. Of these, only the 1050 fits into T_{ab}^Y , and indeed the Bianchi identities show (see below) that it has to be present in T_{ab}^Y .

We will now indicate how the linear version of (5) is solved completely off-shell by the constraints (7) to (9). Consider first the linearized (abc, δ) component of (5). Using (7) and (8), it reads

$$D_{(a} T_{bc)}^{\underline{1050}}{}^\delta + i(\gamma^\alpha)_{(ab} T_{c)\alpha}{}^\delta = 0. \quad (10)$$

Discarding for a moment the index δ , Eq. (10) is identical to the (abc) component of the super Maxwell Bianchi identity^{33),*}

$$D_{[A} F_{BC]} + t_{AB}{}^E F_{[E|C]} = 0, \quad (11)$$

where $t_{AB}{}^E$ is the supertorsion of flat superspace [whose only non-zero component $t_{ab}{}^Y = -i(\gamma^Y)_{ab}$]. From Ref. 33), we know that the solution to the (abc) component of (11) is, with $F_{ab} = i(\gamma^{\alpha_1 \dots \alpha_5})_{ab} A_{\alpha_1 \dots \alpha_5}$,

$$D_a A_{\alpha_1 \dots \alpha_5} = 10(\gamma_{\alpha_1 \alpha_2})_a{}^b B_{b\alpha_3\alpha_4\alpha_5} + 5(\gamma_{\alpha_1 \dots \alpha_4})_a{}^b B_{b\alpha_5}, \quad (12)$$

$$F_{a\alpha} = 30 \cdot 56 B_{a\alpha} + (\gamma_\alpha)_a{}^b G_b, \quad (13)$$

*) This might hint at the existence of an off-shell formulation of the coupled Einstein-Yang-Mills system, at least for gauge groups with dimension $8n$ where n is a positive integer. I am grateful to Paul Townsend for reminding me that the existing off-shell no-go theorems in $d = 4$ break down for $6 \bmod(8)$ $N = 4$ vector multiplets coupled to $N = 4$ supergravity.

where the B spinor-tensors are γ -traceless. We now transfer this solution to Eq. (10) and demand that $T_{\alpha_1 \dots \alpha_5}^\delta$ contains only the component field in the superfield W that transforms under the SO(1,9) representation 1050. We find that

$$D_\alpha T_{\alpha_1 \dots \alpha_5}^\delta = \frac{1}{336} (\gamma_{[\alpha_1 \dots \alpha_4})_\alpha{}^b T_{b\alpha_5]}^\delta, \quad (14)$$

where $T_{a\beta}^\gamma$ is given by

$$T_{a\beta\gamma} = i(\eta_{\beta\gamma} T_a + (\gamma_{\beta\gamma})_a{}^b \tilde{T}_b) + W_{a\beta\gamma}.$$

Here $W_{a\beta\gamma}$ is γ -traceless and antisymmetric in $\beta\gamma$ in order to avoid introducing superfields not present in the off-shell version of the linearized theory³¹⁾. Clearly, (14) implies that

$$T_{a\beta\gamma} = -i(\gamma_{\beta\gamma})_a{}^b T_b + W_{a\beta\gamma}, \quad (15)$$

where $W_{a\beta\gamma}$ is identified with the independent component of the scalar superfield W at the 13th θ level.

The continuation of the analysis of the Bianchi identities shows the role played by the components at the last three θ levels of W, namely $W_{[\alpha\beta\gamma]}$, W_a and W. We discover that the expansion of $T_{ab}^{c\ 29)}$ now reads

$$T_{ab}^c = -\frac{1}{2}(\gamma^{\rho\gamma})_b{}^c Y_\rho + (\gamma^{\rho\gamma})_b{}^c Y_{\alpha\beta\gamma} + (\gamma_{\alpha}{}^{\rho\gamma\delta})_b{}^c \hat{Y}_{\rho\gamma\delta} \quad (16)$$

where Y_β is proportional to $(\gamma_\beta)^{ab} D_a T_b$ and, more importantly, $Y_{\alpha\beta\gamma}$ and $\hat{Y}_{\alpha\beta\gamma}$ are independent three forms. In fact,

$$2(Y_{\alpha\beta\gamma} + 6\hat{Y}_{\alpha\beta\gamma}) = 9W_{\alpha\beta\gamma}, \quad (17)$$

where

$$D_a W_{b\alpha\beta} = i(\gamma_{\alpha\beta}{}^{\gamma\delta\epsilon})_{ab} W_{\gamma\delta\epsilon}.$$

Finally, W and W_a are seen to be proportional to, respectively, $R + 18\Box\phi$ and its fermionic analogue $(\gamma^{\alpha\beta\gamma})_{ab} T_{\beta\gamma}^b$ where $T_{\beta\gamma}^a$ is the Rarita-Schwinger field strength. The theory is thereby completely off-shell (i.e., it describes off-shell Poincaré supergravity), and we have succeeded in removing the above-mentioned obstacle to constructing manifestly supersymmetric effective actions [involving $(R_{\mu\nu\rho\sigma})^n$ -terms] for $N = 1$ superstrings^{*}).

In connection with Eq. (16), it should be emphasized that since $Y_{\alpha\beta\gamma}$ and $\hat{Y}_{\alpha\beta\gamma}$ are no longer proportional to each other, the supersymmetry transformation rule for the Rarita-Schwinger field $E_\mu^a(\theta=0) = \phi_\mu^a$ is not the same as in the Chapline-Manton theory [except possibly for $O(\alpha')$ corrections as suggested by the results of Refs. 37) and 38)]. In fact, once the $(R_{\mu\nu\rho\sigma})^n$ terms are included the components of the auxiliary superfield W must be solved for from their (algebraic) field equations and any non-zero result for $W_{\alpha\beta\gamma}$ will enter T_{ab}^c [recall Eq. (17)] and therefore $\delta\phi_\mu^a$. However, considering solutions to the field equations and the possibility of supersymmetric compactifications, the new terms in ϕ_μ^a are probably relevant only for situations with non-zero $H_{\mu\nu\rho}$ (at least when $\phi = \text{constant}$).

In giving the constraints on the supertorsion tensor which lead to an off-shell formulation of $d = 10$, $N = 1$ Poincaré supergravity, we have also identified the position in T_{AB}^C of the superfields in terms of which one may construct a superspace action along the lines of Ref. 31). In the component version of this action there is a six-form gauge field rather than the two-form one usually used. The on-shell theory based on the two-form can be dualized and re-expressed in terms of the six-form^{34),35)}. However, it is by looking at the structure of the off-shell theory that one learns which one is more fundamental. Both the work by Howe et al.³¹⁾ on the linearized off-shell theory and the superspace results presented in Ref. 19) show that the formulation involving a six-form is the correct one [see also Ref. 36)]. Here we just want to point out that there does not exist an off-shell set of constraints for the system of superspace Bianchi identities consisting of (5) and $dH = 0$ where H is the three-form field strength of a two-form potential. [This theory was analyzed in

^{*}) In this analysis we have consistently ignored terms involving space-time derivatives of components of the auxiliary superfield W . Since it is probably inconsistent to put all these derivative terms to zero, one should relax the torsion constraints still further. For this purpose it is probably sufficient to remove the 1050 restrictions on $T_{\alpha_1\beta_1\gamma_1}^{\beta_2\gamma_2\delta_2}$. As a consequence there will appear new terms in, for instance, (15) and (16).⁵

Ref. 29) but the above comment was not made since it depends on the results of Ref. 19).] Furthermore, this statement is probably not true if we instead combine (5) with $dH = 0$ where H is a seven-form field strength of a six-form potential. In fact, the off-shell constraints should then be essentially (7), (8) and (9) together with

$$H_{\alpha_1 \dots \alpha_5 ab} = i (\gamma_{\alpha_1 \dots \alpha_5})_{ab} + i (\gamma^{\beta_1 \dots \beta_5})_{ab} \tilde{W}_{\alpha_1 \dots \alpha_5, \beta_1 \dots \beta_5},$$

$$H_{\alpha_1 \dots \alpha_n a_1 \dots a_{7-n}} = 0 \quad n < 5, \quad (18)$$

where the $\tilde{W}$ -tensor is constrained to contain representations appearing at the 12th θ level in W . Interestingly enough, both the 1050 and 770 can be incorporated into $\tilde{W}$; this was not possible in the case of T_{ab}^{γ} .

We would now like to discuss the off-shell torsion constraints in the context of $N = 1$ supersymmetric strings propagating in arbitrary backgrounds. Thus, consider the action²²⁾

$$S = \int d^2 \xi \left[\frac{1}{2} V \gamma^{ij} E_i^{\alpha} E_j^{\beta} \eta_{\alpha\beta} + \frac{1}{2} \epsilon^{ij} E_i^B E_j^A B_{AB} \right], \quad (19)$$

where $i, j, \dots$ are two dimensional world indices, γ^{ij} is the world sheet metric, $V = (\det \gamma_{ij})^{\frac{1}{2}}$, ϵ^{ij} is the Levi-Civita tensor density. The supergravity background enters via

$$E_i^A \equiv \left(\frac{\partial}{\partial \xi^i} Z^M(\xi) \right) E_M^A(Z) \quad (20)$$

and B_{AB} , where this latter field and E_M^A are the super two-form potential and the super vielbein, respectively. They constitute the basic superfields of the superspace formulation of $N = 1$, $d = 10$ supergravity²⁹⁾.

For consistency^{20), 21)}, it is necessary that (19) be invariant under the transformations²²⁾

$$\delta E^{\alpha} = 0, \quad (21)$$

$$\delta E^a = 2 (\gamma_a)^{ab} E_i^{\alpha} \eta_b^i, \quad (22)$$

$$\delta(V\gamma^{ij}) = X^{a(i} \eta_a^{j)}, \quad (23)$$

where $\delta E^A \equiv (\delta Z^M) E_M^A(Z)$, η_a^i are self-dual parameters and X^{ai} are to be determined. The variation of (19) reads

$$\delta S = \int d^2\zeta \left[\frac{1}{2} \delta(V\gamma^{ij}) E_i^\alpha E_j^\beta \eta_{\alpha\beta} + \right. \\ \left. + V\gamma^{ij} E_i^\alpha E_j^\beta \delta E^a T_{ab}{}^\gamma \eta_{\alpha\gamma} + \frac{1}{2} \varepsilon^{ij} \delta E^d E_i^B E_j^A H_{ABd} \right], \quad (24)$$

where T_{AB}^C and H_{ABC} are the supertorsion and three-form field strength (of B_{AB}), respectively. As has been demonstrated by many authors^{22),25),39)} δS vanishes for the constraints given in Refs. 29) and 22), provided X^{ai} is suitably chosen. However, as explained above, these constraints do not permit the low energy effective field theories of the heterotic strings and the $SO(32)$ open superstring to be embedded in superspace. Therefore, we must ask if $\delta S = 0$ is consistent with the (Poincaré) off-shell constraints given in (7), (8), (9) and (15) together with the necessary generalizations of the constraints on H_{ABC}^* . However, the answer is no!

The interpretation of this unsatisfactory situation is not clear. Nor is it straightforward to see how to reformulate the σ -model to circumvent it. What one is looking for is a formulation of the σ -model which does not have Eq. (6) as a consequence. Of course, one complication is that the presence of the two-form potential in (19) [and H_{ABC} in (24)] will, as explained above, put the theory on-shell. Part of the resolution to this problem might involve the use of the six-form potential (in a non-trivial way) instead of the two-form. Or, it may be that the Lorentz Chern-Simons term present in $H_{\alpha\beta\gamma}$ means that there should be a (Lorentz connection)²-term in the σ -model in analogy with the Yang-Mills case discussed above. Another (less likely) possibility is that the superspace used here is not the most natural one for this problem. [One could consider enlarging it by introducing "central" charges in the 126 representation of $SO(1,9)$; this might also lead to a more elegant description of the six-form potential as the corresponding gauge field.]

*) In fact, it would be interesting to find the most general constraints that satisfy $\delta S = 0$ ⁴⁰⁾.

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NOTE ADDED

While writing this paper we became aware of Ref. 41) which also discusses superspace as an approach to construct field theories with higher derivative terms. That paper also provides the information that the constraints (1) to (4) do not need be relaxed if one considers $O(\alpha')$ corrections only (as was anticipated above).

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