THE AUXILIARY FIELD / ULTRAVIOLET FINITENESS CONNECTION

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ABSTRACT :

The status of the auxiliary field program for supersymmetric field theories is reviewed, with emphasis on the implications for improved ultraviolet behaviour, and conversely on the implications of ultraviolet behaviour for auxiliary fields.

1. INTRODUCTION

For many years one of the central problems of supersymmetric field theories has been that of auxiliary fields. Generally, supersymmetry transformations that leave invariant a given action will close to form the usual supersymmetry algebra only when the field equations are used. This is called "on-shell supersymmetry". If field equations are not needed to close the algebra then the supersymmetry is "off-shell". For this to happen we must have equal numbers of boson and fermion field components as well as equal numbers of boson and fermion propagating modes (i.e. states). Since fermions always have more components than they propagate modes, a balance in the number of boson and fermion states will usually mean an imbalance in the number of boson and fermion field components. Hence the need for boson "auxiliary" i.e. non-propagating, fields to balance the number of boson and fermion field components off-shell while not disturbing the balance of states. Of course equality of the numbers of bosons and fermions is only a necessary condition for off-shell supersymmetry, not a sufficient one. In general, a solution to an auxiliary field problem will require fermion as well as boson auxiliary fields.

There are several reasons why we should want auxiliary fields. The most obvious one is that if the algebra of transformations includes field equations then these transformations are tied to the particular action that yields these equations. The introduction of auxiliary fields frees the transformation rule of reference to a particular action. This allows, inter alia, the addition of invariant actions to produce new invariant actions. Without auxiliary fields this involves a laborious procedure in which terms are added order by order to the action and transformation rules to obtain a new invariant action with new transformation rules.

A second, and more important, motivation for auxiliary fields is that they allow us to consider superfield perturbation theory in which component field Feynman graph calculations can be done together, and with more ease, as supergraph calculations.

This formalism allows us to deduce immediately, just from the form of the Feynman rules, many "non-renormalization theorems" which have remarkable consequences for the ultraviolet behaviour of the N-extended supersymmetric theories.

Many of the authors that have contributed to our understanding of auxiliary fields have advanced as one of their motivations a better understanding of ultraviolet behaviour while, at the same time, most of the real progress in this field has been technical. Given some recent advances the time now seems ripe for a survey of what we can expect auxiliary fields to say about the quantum theory, and vice versa.

2. AUXILIARY FIELDS

Consider the simplest example of an on-shell supersymmetric theory, the Wess-Zumino model[1],

$$\mathcal{L} = -\frac{1}{2}(\partial_\mu A)^2 - \frac{1}{2}(\partial_\mu B)^2 - \frac{1}{2}\bar{\lambda}\not{\partial}\lambda \tag{2.1}$$

with scalar A, pseudoscalar B and spinor λ_α. One easily establishes that the algebra of supersymmetry transformation rules that leaves $I = \int \mathcal{L}\, d^4x$ invariant, closes to give the usual supersymmetry algebra only if the λ field equation, $\not{\partial}\lambda = 0$, is used. To extend (2.1) to an off-shell supersymmetric model we require auxiliary fields, and a count of components shows that two auxiliary scalar bosons would be the simplest solution. To find these auxiliary fields we can investigate the various irreducible representations of supersymmetry on fields for $N = 1$, which can be found by, e.g. superfield methods. One easily finds a representation with maximum spin 1/2 with the fields (A, B, λ, F, G), corresponding to an N = 1 chiral superfield). F and G are obviously the wanted dimension two auxiliary fields and they will occur in the action as $F^2 + G^2$. The new, F and G dependent, transformation rules will close off-shell, i.e. without the use of field equations.

What makes this example so simple is the fact that the highest propagated spin, i.e. 1/2, is also the highest spin of the off-shell multiplet of fields. In general, for massless theories the maximum (propagating) spin, s, is bounded from below by

$$s \text{ (propagating)} \geqslant \frac{N}{4} \tag{2.2}$$

where N is the number of supersymmetries. For massive on-shell multiplets, or off-shell multiplets of fields, the bound is

$$s \text{ (off-shell)} \geqslant \frac{N}{2} \tag{2.3}$$

which is more stringent. If s (propagating) $\geqslant$ N/2 we can always find an off-shell multiplet of fields with maximum spin s, and this will give a set of auxiliary fields. For those models for which $N/4 \leqslant s\text{(propagating)} < N/2$ the auxiliary field problem is much harder, and probably without a solution for most cases. At any rate, it is clear that in such cases the off-shell multiplet of fields must contain spins higher than the maximum propagated spin. This can happen in two ways.

(i) We may have unconventional field representations with spin "discontinuities" in the massless limit. The simplest case is the gauge antisymmetric tensor $A_{\mu\nu}$ which is spin 1 off-shell but which propagates a massless spin-zero mode (if the kinetic term is the usual $\partial_\rho A_{\mu\nu}\partial^{[\rho}A^{\mu\nu]}$ one).

(ii) We may have high spin (i.e. higher than the physical propagated spin), non-gauge auxiliary fields.

For those cases for which s(propagating) $\geqslant$ N/2 there is a systematic way to construct the relevant off-shell representations by means of "supercurrents"[2]. The best known example is the N = 1 spin 2 conformal supercurrent multiplet, containing the (traceless) energy momentum tensor, $T_{\mu\nu}$, the supersymmetry current S_μ and the axial current j^5_μ . For the model of (2.1), for example, these currents are bilinears in the fields A, B, λ and they are conserved as a consequence of the A, B, λ field equations. The currents themselves ($T_{\mu\nu}$, S_μ , j^5_μ) are not subject to field equations, however, and therefore form an off-shell multiplet. The corresponding "contragredient" multiplet of fields $\{h_{\mu\nu}$, ψ_μ , $A_\mu\}$ is also an off-shell representation with gauge transformations determined by the conservation conditions on the currents via the Noether coupling I(int) = $\int h\cdot J$. That is, δh is whatever leaves I(int) invariant. The multiplet $\{h_{\mu\nu}$, ψ_μ , $A_\mu\}$ is that of N = 1 conformal supergravity.

Given a maximum desired spin, s, e.g. 2 in the above case, the supercurrent construction yields an off-shell multiplet with this spin provided there is a "matter" multiplet of maximum spin s/2 that can be used to form it. Since this maximum spin is bounded by (2.2) we again arrive at the bound s $\geqslant$ N/2 for the off-shell multiplet of currents, or fields. Thus we can say that the really difficult auxiliary field problems start when there is no lower spin "matter" to which the given off-shell model could couple. Precisely, systematic methods, such as supercurrents , yield auxiliary fields for :

(i) matter theories (spin $\leqslant$ 1/2) for N = 1 only

(ii) gauge theories (spin $\leqslant$ 1) for N $\leqslant$ 2

(iii) conformal supergravity (spin $\leqslant$ 2) for N $\leqslant$ 4

The restriction to conformal supergravity in (iii) is because the supercurrent method yields an <u>irreducible</u> multiplet which is what is needed for conformal supergravity. The Poincaré supergravity theories are based on reducible (partially locally reducible) multiplets because such theories couple to a tracefull energy momentum tensor; $T_\mu{}^\mu \neq 0$. To find the Poincaré supergravity theories we must add additional multiplets. For N = 1 and N = 2 these additional multiplets have spin $<$ 2 and the passage from conformal to Poincaré supergravity is relatively straightforward. For N $\geqslant$ 3 there are no lower spin multiplets. Any multiplet added to the spin 2 conformal multiplet would itself have at least spin 2. This is a problem because these higher spin fields will generally be gauge fields and so could not appear easily as auxiliary fields and even if they could they would probably not allow consistent interactions. Because of this we can add to the list of off-shell theories constructible by systematic methods, only

(iv) Poincaré supergravity (spin $\leqslant 2$) for $N \leqslant 2$.

There have been recently two additional successful constructions of off-shell supersymmetric theories that fall outside the scope of the above mentioned methods. They both are auxiliary field solutions with high spin (non gauge) auxiliaries. The first is a set of auxiliary fields for linearized 10-dimensional supergravity[3] which implies, by dimensional reduction, a solution of the auxiliary field problem for linearized N = 4 Poincaré supergravity coupled to 6 N = 4 super-Maxwell multiplets. The maximum spin of the off-shell multiplet is 4. The construction makes essential use of the 10-dimensional supercurrent obtained as a bilinear of on-shell 10-dimensional super-Yang-Mills fields[4], but the final set of fields is more than those of the supercurrent alone. The second[5] is a set of auxiliary fields for the N=2 matter theory with spins 1/2 (sometimes called the hypermultiplet). The problem with a previous off-shell version of this model[6] is that it required one of the spin zero states to be represented by a gauge antisymmetric tensor, or equivalently a conserved vector. In interaction with Y-M fields this becomes a covariantly conserved vector, which is a constraint that cannot be solved in a useful way. The new solution has only conventional field assignments and can be consistently coupled to N = 2 super Y-M theory. The maximum spin of the off-shell multiplet is 1.

Another approach to the auxiliary field problem is to allow off-shell central charges[7]. Here one allows field equations to remain in the supersymmetry algebra if they can be interpreted as new central charge transformations. For massless theories these transformations must vanish on-shell, but need not off-shell. This approach is less restrictive, allowing for example a set of auxiliary fields for the N = 4 Maxwell (i.e. non-interacting) theory, and general methods have been developed for finding such auxiliary field sets[7,8]. But for superspace perturbation theory such auxiliary field solutions are not useful because they cannot be used to develop an unconstrained superfield formalism. The importance of this point is what concerns us next.

3. ULTRAVIOLET FINITENESS

A set of auxiliary fields, without an off-shell central charge, is the first and critical step towards an unconstrained superfield formulation. One proceeds first to a constrained superfield formulation according to one of several well established techniques. For example, each field, or field strength in the case of a gauge field, can be considered as the $\theta = 0$ component of a superfield. The transformation rules can then be used to determine the higher θ-components of these superfields. Of course only the lowest dimension field strength superfields constructed in this way will be independent. The higher dimension field strengths will occur in the θ-expansion of the lower dimension field strength superfields. These superfields are subject to various constraints, or rather identities if one arrives at them in the above fashion, analogous to the constraint $\partial_\mu {}^*F^{\mu\nu} = 0$ for the Maxwell field strength.

If these constraints are solved, e.g. $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ in the above illustration, one obtains an unconstrained superfield formulation in terms of a "prepotential" (the analogue of A_μ). Thus, in the simplest well-known case of the N = 1 (1, 1/2) multiplet the lowest dimension covariant field is the spinor λ_α which becomes the $\theta = 0$ component of the field strength superfield $W_\alpha(x, \theta)$. These constraints on W_α are solved by $W_\alpha = \bar{D}^2 D_\alpha V$, thereby introducing the unconstrained prepotential V.

These constraints can easily be solved for the free theory but with more difficulty for the interacting theory, particularly for $N>1$. For N = 1 the solutions to both the Yang-Mills and the supergravity constraints are known, but for N = 2 neither is yet known in supersymmetric form. This is probably not a serious problem in principle. One can obtain a solution to the constraints of the interacting theory as a perturbation series about the known solution to the free theory constraints. By solving these constraints either in closed form or perturbatively we arrive finally at an unconstrained superfield formulation of the interacting theory. This is the starting point for the development of superfield perturbation theory.

There is one exception to the need for unconstrained superfields to do superfield perturbation theory, and that is the N = 1 chiral field ϕ , satisfying $\bar{D}_{\dot\alpha}\phi = 0$ (two component spinor notation). The solution to this constraint, for the free field, is $\phi = \bar{D}^2 X$, introducing a complex scalar prepotential X . The chirality constraint is sufficiently simple that with a few modifications to the usual methods of extracting Feynman rules one can deal directly with ϕ without having to use X . Indeed the introduction of the prepotential X illustrates a problem that can be avoided for N = 1 but is endemic to $N>1$ theories[9]. The prepotential X has a gauge transformation $X \to X + \bar{D}_{\dot\alpha}\bar{\Lambda}^{\dot\alpha}$, but the parameter $\bar{\Lambda}^{\dot\alpha}$ is only determined up to a transformation $\bar{\Lambda}^{\dot\alpha} \to \bar{\Lambda}^{\dot\alpha} + \bar{D}_{\dot\beta}\xi^{(\dot\beta\dot\alpha)}$ because this transformation leaves X invariant. Hence, the ghosts associated with the gauge transformation of X, themselves have a gauge transformation with parameter $\bar{\Lambda}^{\dot\beta\dot\alpha} = \bar{\Lambda}^{\dot\alpha\dot\beta}$. But the secondary ghosts associated with this gauge transformation will again have an even larger invariance, with a symmetric tri-spinor parameter, and so on ad infinitum. This is analogous to the series of ghosts associated with gauge antisymmetric tensors. In that case, however, the series involves an increasing or decreasing number of antisymmetrized vector indices and so terminates.

An infinite series of ghosts, apparently unavoidable for $N>1$, presents a problem because the effective action from which Feynman rules are read off cannot be written down in closed form. However, one can arrange for all but a finite number of ghosts to decouple, or, in a background field approach, for all but a finite number to couple only to the background field[10]. In this case the infinite series of ghosts can only contribute at one loop. This makes one loop a special case that must be dealt with separately in the background field method. It can then be shown, with the exception of one loop and on the assumption that <u>an N-extended unconstrained superfield formalism</u>

exists, that all counterterms in N-extended superfield form must be such that

(i) they are full superspace integrals $I = \int d^Dx\, d^{4N}\theta\, \mathcal{L}$ (counterterm)

(ii) $\mathcal{L}$(counterterm) is a local product of superfields

(iii) Only "covariant" superfields occur in $\mathcal{L}$ and in such a way that $\mathcal{L}$ is background gauge invariant.

The first two criteria[11] follow directly from the Feynman rules and the usual properties of counterterms. The third criterion[10] is more subtle and depends on the deployment of the background field method. The precise formulation of this method is what determines the meaning of "covariant". It means pre-gauge invariant, which means that the superspace dimension 1/2 potential $\mathcal{A}_\alpha^i$ for super-Yang-Mills theories, or the vielbein $E_M{}^A$ and super-connections $\Omega_{MA}{}^B$ for supergravity, may appear in gauge covariant quantities, but not the unconstrained prepotentials.

If we choose the Y-M action in D dimensions to be $g^{-2}\int d^Dx\,(F_{\mu\nu})^2$ such that $\dim(A_\mu) = 1$ then $\dim(g^2) = 4 - D$ (in units of mass). The counterterm at ℓ loops has the form

$$(g^2)^{\ell-1} \int d^Dx\, d^{4N}\theta\, \mathcal{L} \tag{3.1}$$

and the same result holds for both super-matter (spins $\leq 1/2$) and conformal supergravities if the coupling constant is again chosen to appear as a g^{-2} factor multiplying a g-independent $\mathcal{L}$. For Poincaré supergravity at ℓ loops we have a counterterm of the form

$$(\kappa^2)^{\ell-1} \int d^Dx\, d^{4N}\theta\, \mathcal{L} \tag{3.2}$$

where κ is the gravitational constant of dimension (1 - D/2). The minimum dimension for $\mathcal{L}$ in (3.1) is 2 for super Y-M theories, corresponding to the choice $\mathcal{L} = \mathcal{A}_\alpha^i \lambda^\alpha_i$; 2 for super-matter, corresponding to $\mathcal{L}$ = (physical scalar)^2 ; zero for conformal supergravity, corresponding to $\mathcal{L} = E = \mathrm{Ber}(E_M{}^A)$. Similarly, the minimum dimension for $\mathcal{L}$ in (3.2) is zero.

Since counterterms have dimension zero, (after account is taken of the integration measure and factors of coupling constants), the minimal dimensions of $\mathcal{L}$ imply the absence of allowed counterterms at ℓ loops in D dimensions if

$$D < 4 + \frac{(2N-2)}{\ell} \quad \text{for super-Y-M and super-matter} \tag{3.3a}$$

$$D < 2 + \frac{(2N-2)}{\ell} \quad \text{for Poincaré supergravity} \tag{3.3b}$$

$$D < 4 + \frac{(2N-4)}{\ell} \quad \text{for conformal supergravity} \tag{3.3c}$$

N is the number of supersymmetries and ℓ the loop order. This is meant to hold beyond one loop. These results would imply[10] that both N = 2 and N = 4 super Y-M theories in D = 4 are finite beyond one loop, which appears to be the case up to three loops[12,13]. The N = 4 theory happens to be also one loop finite, while the N = 2 theory is not.

For Poincaré supergravity (3.3b) would imply[10] finiteness through 6 loops in $D = 4$. For conformal supergravity (3.3c) would imply[14] finiteness of the $N \geqslant 3$ theories in $D = 4$ beyond one loop (and it appears that the $N = 4$ theory is finite also at one loop[15]. In fact, we can do somewhat better for conformal supergravity because the counterterm available for $N = 2$ is $\int d^4x\, d^8\theta\, E$ which vanishes[16], so the $N = 2$ conformal supergravity should also be finite beyond one loop.

All these restrictions on counterterms assume that an N-extended unconstrained superfield formalism exists. This critical assumption is very probably true for the conformal supergravity theories because their auxiliary fields are all known[17], but appears to be false[18,19], apart from a few exceptions, for the $N>2$ super Y-M and (Poincaré) supergravity theories. In this case, we may make a weaker assumption :

(A) There exists an unconstrained M-extended superfield formalism (and hence M-extended auxiliary fields) for N-extended super Y-M or Poincaré supergravity, with $M \leqslant N$ to be determined.

With this assumption the formulae (3.3) are replaced, in the relevant cases, by

$$D < 4 + \frac{(2M - 2)}{\ell} \qquad \text{for N-extended super Y-M} \tag{3.4a}$$

$$D < 2 + \frac{(2M - 2)}{\ell} \qquad \text{for N-extended Poincaré supergravity} \tag{3.4b}$$

But what is the correct value of M ? For super-matter and conformal supergravity we have, or expect, $M = N$, and the same is true for $N = 2$ Y-M theory. Given the restrictions on auxiliary fields found in refs.(18, 19) and given the recent progress in the auxiliary field search of refs.(3, 5) it would appear that we already have auxiliary fields for almost all those cases for which they can be found, and that M is restricted by

$$M \leqslant N/2 \quad \text{for } N = 4 \text{ Y-M, } N = 4 \text{ \underline{pure} Poincaré supergravity}, \; N = 8 \text{ Poincaré supergravity} \tag{3.5}$$

(By "pure" supergravity I mean with no matter multiplets). I say "almost" because in order to have an $M = 4$ formulation of $N = 8$ Poincaré supergravity one needs an off-shell version of the $N = 4$ "spin 3/2 multiplet" which is not yet known. But let us be optimistic and assume that the bound of (3.5) is saturated. Then we arrive at the conclusion that, for $\ell > 1$ there is no allowed counterterm in D dimensions at ℓ loops if

$$D < 4 + \frac{2}{\ell} \qquad N = 4 \text{ Y-M theory} \tag{3.6a}$$

$$D < 2 + \frac{6}{\ell} \qquad N = 8 \text{ Poincaré supergravity} \tag{3.6b}$$

We observe that (3.6a) is still sufficient to ensure finiteness[5]. It is interesting to compare (3.6) with similar results obtained from the superstring approach[20] by extrapolation from one loop :

$$D < 4 + \frac{4}{\ell} \qquad N = 4 \text{ Y-M theory} \qquad (3.7a)$$

$$D < 2 + \frac{6}{\ell} \qquad N = 8 \text{ Poincaré supergravity} \qquad (3.7b)$$

The results (3.6b) and (3.7b) agree. They also agree with the result of yet another approach[21] and imply that an ultraviolet divergence can be expected to appear at three loops. The relevant three loop counterterm is also known[22], so a calculation for supergravity at three loops remains the critical test of whether we can expect further "miraculous" cancellations in these theories.

There is a discrepancy between (3.6a) and (3.7a). This is not necessarily a contradiction because (3.6a) is valid only beyond one loop while (3.7a) is known to be valid only at one loop. If (3.7a) is valid beyond one loop also then it would appear that divergence cancellations are slightly better than one has a right to expect from superspace perturbation theory. I remark that if M = 3 were possible for the N = 4 Y-M theory, then (3.4a) would yield a formula in agreement with (3.7a), but this possibility is ruled out according to ref. 19.

In all of the above discussion some kind of supersymmetric regularization has been assumed. But is this reasonable ? The scheme used successfully in practice is dimensional regularization by dimensional reduction but its internal consistency is questionable[23]. For general arguments (if not explicit calculations) it is preferable to avoid it. There is an alternative scheme ; higher derivative regularization. For gauge theories this fails at one loop but Slavnov has shown[24] how a two stage regularization method can be devised in which one loop is regulated separately and supersymmetrically. When this method is applicable its consistency is not in question, but the trouble is that it is not applicable unless we have auxiliary fields. The reason is that the higher derivative terms used to regulate the theory propagate additional massive ghost modes (whose mass goes to infinity as the regulator is switched off), and these modes will be in supermultiplets only if auxiliary fields were initially present. This may seem to crush hopes of applying this method to N = 4 Y-M and N = 8 supergravity theories, but fortunately auxiliary fields for M = N/2 is just sufficient to allow the method to work. The point is that in this case the ghost modes of the regulated theory will be in N-supersymmetric multiplets with a central charge equal to the mass, and so the N-extended supersymmetry will not be broken by the regulator. It has been checked that the ghost modes for the N = 4 theory with M = 2 do indeed form N = 4 multiplets with a central charge[5].

REFERENCES

1. J. Wess and B. Zumino, Nucl. Phys. B49, 52 (1974)
2. S. Ferrara and B. Zumino, Nucl. Phys. B87, 207 (1975) ; M.F. Sohnius, Phys. Lett. B81, 8 (1979) ; E. Bergshoeff, M. de Roo, and B. de Wit, Nucl. Phys. B182, 173 (1981) ; P.S. Howe, K.S. Stelle and P.K. Townsend, Nucl. Phys. B192, 332 (1981)
3. P.S. Howe, H. Nicolai and A. van Proyen, Phys. Lett.B 112, 446 (1982)

4. E. Bergshoeff, M. de Roo, B. de Wit and P; van Nieuwenhuizen, Nucl. Phys. B195, 97 (1982) ; E. Bergshoeff and M. de Roo, Phys. Lett.112, 53 (1982)
5. P.S. Howe, K.S. Stelle and P.K. Townsend, CERN preprint TH 3271 (1982) to appear in Nucl. Phys.B
6. P. Breitenlohner and M.F. Sohnius, Nucl. Phys. B165, 483 (1980) ; M.F. Sohnius K.S. Stelle and P.C. West in "Supergravity and Superspace" eds. S.W. Hawking and M. Rocek, C.U.P. (1981)
7. M.F. Sohnius, K.S. Stelle and P.C. West, Phys. Lett. B92, 123 (1980) ; W. Siegel, Nucl. Phys. B173, 51 (1980)
8. M.F. Sohnius, K.S. Stelle and P.C. West, Nucl. Phys. B173, 127 (1980) ; J.G. Taylor in "Quantum Structure of Spacetime", eds. M.J. Duff and C.J. Isham, C.U.P. to appear
9. W. Siegel and S.J. Gates, Nucl. Phys. B189, 295 (1981)
10. M.T. Grisaru and W. Siegel, Nucl. Phys. B201, 292 (1982)
11. M.T. Grisaru, W. Siegel and M. Rocek, Nucl. Phys. B159, 429 (1979)
12. O.V. Tarasov and A.A. Vladimirov, Phys. Lett. B96, 94 (1980) ; M.T. Grisaru, M. Rocek and W. Siegel, Phys. Rev. Lett. 45, 1063 (1980) ; W. Caswell and D. Zanon Phys. Lett. B100, 152 (1980)
13. L.V. Adveev and O.V. Tarasov, Phys. Lett. B112, 356 (1982)
14. K.S. Stelle and P.K. Townsend, Phys. Lett. B113, 25 (1982)
15. E.S. Fradkin and A.A. Tseytlin, Nucl. Phys. B203, 157 (1982)
16. E. Sokatchev, Phys. Lett. B100, 466 (1981)
17. E. Bergshoeff, M. de Roo and B. de Wit, Nucl. Phys. B182, 173 (1981)
18. W. Siegel and M. Rocek, Phys. Lett. B105, 278 (1981)
19. V.O. Rivelles and J.G. Taylor, "Off-shell no-go theorems for higher dimensional supersymmetries and supergravities", Kings College preprint (1982)
20. M. Green, J. Schwarz and L. Brink, Nucl. Phys. B198, 474 (1982); for a review see J. Schwarz, preprint CALT 68-911 (1982)
21. T. Curtright, Phys. Rev. Lett. 48, 1704 (1982)
22. R. Kallosh, Phys. Lett. B99, 122 (1981) ; P.S. Howe, K.S. Stelle and P.K. Townsend Nucl. Phys. B191, 445 (1981)
23. See W. Siegel, P.K. Townsend and P. van Nieuwenhuizen in "Superspace Supergravity" eds. S.W. Hawking and M. Rocek (C.U.P.)
24. A.A. Slavnov in "Superspace and Supergravity" eds. S. Hawking and M. Rocek (C.U.P.)