Fermi National Accelerator Laboratory

FERMILAB-Pub-88/75-T
LBL-25515
UCB-PTH-88/16
July 12, 1988

Properties of Supersymmetry Breaking via Gravitino Condensation

Michelangelo MANGANO[1]

Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, IL 60510.

Massimo PORRATI

Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, IL 60510.
and
Dept. of Physics, University of California and LBL,[2]
Berkeley, CA 94720

Abstract

We study the implications of a recently proposed mechanism to break local supersymmetry through gravitino condensation. We find that gravitino condensates destabilize the Minkowski background and generate a large cosmological constant.

[1]Address after September 1. 1988: INFN, Sesione di PISA, S. Piero a Grado (PISA), Italy.
[2]Permanent address.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

Supersymmetric theories of fundamental interactions play a major rôle in any unification scheme, in particular in any theory containing gravity. This is because, in order to solve the hierarchy problem, it is convenient to have an effective Lagrangian describing interactions below the Planck scale M_P endowed with N=1 supersymmetry. Since only spontaneously broken models are consistent with observation, it is important to find general mechanisms to generate supersymmetry breaking in supergravity.

Classical super-Higgs effects can exist in supergravity[1], and yield models with realistic physics and zero cosmological constant (at tree level)[2]. These models, however, strongly depend on the choice of matter content and superpotential and, in theories where the matter and interaction content are fixed, like strings or extended $N > 4$ supergravities, they cannot be reproduced. The only known ways of achieving tree level spontaneous supersymmetry breaking in string theory, through *e.g.* "Scherk-Schwarz" compactifications[3], seem unable to give a breaking at the scale of 1 TeV.

String theories gave importance to the understanding of non-perturbative mechanisms for supersymmetry breaking, obtained for example via gaugino condensates[4]. The trouble with gaugino condensates is that it has not been possible to find, using the effective field-theoretical action of strings, a concrete non-perturbative configuration yielding such a non-zero condensate. It is worth noticing that in globally supersymmetric QCD it is possible to find non-zero gaugino condensates $\langle\bar\lambda\lambda\rangle$, but their relevance for the problem of supersymmetry breaking strongly depends on other properties of the theory. For example, in the presence of massive matter multiplets minimally coupled supersymmetry remains unbroken even if $\langle\bar\lambda\lambda\rangle \neq 0$.

Another possibility recently considered in the context of supergravity theories is to induce the supersymmetry breaking via gravitino condensates. Better, the gauge invariant *v.e.v.* $\langle D_{[\mu}\bar\psi_{\nu]}D^{[\mu}\psi^{\nu]}\rangle$ has been shown to be an order parameter for supersymmetry breaking[5]. For completeness we will illustrate here the argument given in ref. [5].

The chiral $U(1)$ current associated with a chiral matter supermultiplet $Z = (z, \chi)$ satisfies the following equation[5,6,7], in absence of superpotential:

$$D_\mu J^\mu = -\frac{1}{384\pi^2} R_{abcd}\,{}^*R^{abcd} + \ldots \tag{1}$$

where ${}^*R_{abcd} = 1/2\epsilon_{cdef}R_{abef}$.

The left-hand side of equation (1) is the $\theta^2\bar\theta^2$ component of the vector supermultiplet $\bar Z Z = (z^* z, \ldots)$. Equation (1) is therefore related by supersymmetry to other equations, in particular to[8]:

$$\{\bar Q, \bar\chi z\} = \frac{\kappa^2}{384\pi^2}\bar\psi_{ab}\psi^{ab} + \text{terms vanishing on-shell}, \tag{2}$$

where $\psi_{ab} \equiv D_a\psi_b - D_b\psi_a \equiv D_{[a}\psi_{b]}$ and $\kappa = 1/M_P$.

Equation (2) shows that a non-zero *v.e.v.* for $\bar{\psi}_{ab}\psi^{ab}$ breaks supersymmetry, and that the Goldstone fermion has a component along $\bar{\chi}z$. In ref.[5] it was observed that, by applying straightforwardly the techniques of ref.[9], and evaluating the contribution of gravitational instantons to the gravitino condensate, one would obtain a non-zero value for $\bar{\psi}_{ab}\psi^{ab}$ equal to:

$$\langle\bar{\psi}_{ab}\psi^{ab}\rangle = \mathrm{const}\cdot\mu^5. \tag{3}$$

In eq. (3) μ is a cut-off, whose explicit appearance is due to the non-renormalizability of gravitational interactions, and which is probably to be fixed around M_P: $\mu \sim M_P$.

The advantage of breaking supersymmetry through gravitino condensation is its model independence: it is sufficient to have a chiral superfield corresponding to some flat direction in order to generate the breaking. Moreover, the fact that it is possible to find explicitly a computational scheme yielding a non-zero $\langle\bar{\psi}_{ab}\psi^{ab}\rangle$ is extremely attractive. The generality of this phenomenon may furthermore extend to string theories, because the relevant gravitational instantons (Eguchi-Hanson metrics) are hyper-Kähler manifolds, so that they might be promoted to exact solutions for the string equations of motion by allowing a non-zero three index tensor $H_{\mu\nu\rho}$. Within the computational framework of ref.[5,9], on the other hand, it is rather difficult to evaluate another key quantity in the description of the supersymmetry breaking, namely the value of the cosmological constant Λ. One has to check indeed whether Λ remains zero even after the breaking, or if it develops instead an unacceptably large value. We think this to be a question worth of investigation and we will try to give it an answer in the rest of this paper.

We take here the effective action approach, which has been proven already to be a powerful tool for investigating globally supersymmetric gauge theories like Super Yang-Mills[10] or Super QCD[11]. In order to justify this approach in supergravity one must assume that a consistent theory of gravitation exists. The important point is that, whatever this theory might be, at low energies it will reduce to Einstein gravity coupled to matter, or, when the supersymmetry breaking scale is below M_P, to a $N = 1$ supergravity coupled to matter. At still lower energies, where the super-Higgs effect occurs, one should write an effective Lagrangian which, still having the form of a $N = 1$ theory, contains only the relevant dynamical degrees of freedom. In our case these are the gravitational supermultiplet and the goldstino superfield. The gravitational supermultiplet contains the graviton (e_μ^a) and the gravitino (ψ_μ), in addition to auxiliary fields, A_μ and $B_{\mu\nu}$, which we choose according to the new-minimal formulation of supergravity[6,12,13]. The goldstino supermultiplet contains the spin 1/2 fermion providing the helicity states $\pm 1/2$ of the massive gravitino. Since the field whose supersymmetry transformation is non-vanishing is composite (see eq. (2)), we should expect a composite superfield to play the rôle of the Goldstino.

Before entering into the details of the nature of this composite field we must constrain as much as possible the form of the effective Lagrangian. To this purpose eq. (2), and in general any anomalous transformation law, proves useful. We will indeed write the effective Lagrangian as a sum of two parts. The first one respects all *classical* symmetries of the theory, and in particular it is invariant under the chiral $U(1)$ transformation of eq. (2). We will call it the kinetic term L_K, and write it as a D-type density (for an introduction to the techniques of super-conformal tensor calculus see *e.g.* ref.[12]):

$$L_K = [V]_D = e\{D + (V_\mu + \frac{i}{2}\bar{\psi}_\mu\gamma_5\lambda)\epsilon_{\mu\nu\rho\sigma}\partial_\nu B_{\rho\sigma} \\ -\frac{1}{2}\bar{\psi}\cdot\gamma\gamma_5\lambda\} + \text{surface terms.} \quad (4)$$

In eq. (4) V is a real superfield of chiral weight $n = 0$[12] with components:

$$V = (C, \chi, H, K, V_\mu, \lambda, D). \quad (5)$$

The second part is an F-type density[12] and it reproduces, at the classical level, the anomaly equations, while being invariant under the non-anomalous symmetries of the theory. Already from this fact we see that $\bar{Z}Z$ cannot be the light superfield providing the goldstino degrees of freedom. $\bar{Z}Z$ has in fact chiral weight zero. Since, in order to reproduce the anomaly equations, we need a field shifting under a chiral transformation, the natural choice is to assume

$$S = \log\frac{Z^3}{\mu^3} \quad (6)$$

to be the effective composite field. The power of three is a matter of convenience and μ is an energy scale, yet to be determined.

To determine completely the effective Lagrangian let us examine the symmetries of the fundamental theory. The first one, already mentioned, is a chiral $U(1)$ acting in the following way on the fields of the theory:

$$Z \rightarrow e^{i\Lambda/3}Z, \qquad \Lambda = \text{constant}, \quad (7)$$

all the other fields being invariant. The anomalous variation of the action under this $U(1)$ is[7]:

$$\int d^4x\, L \rightarrow -\frac{i}{384\pi^2}\int d^4x\, d^2\theta\, \mathcal{E}\, T_{ab}\,{}^*T^{ab}\,\Lambda \equiv -\frac{i}{384\pi^2}[T_{ab}\,{}^*T^{ab}\Lambda]_F. \quad (8)$$

In eq. (8) $\mathcal{E}$ is the chiral superspace density, while $T_{ab}\,{}^*T^{ab}$ is the supersymmetric extension of the Gauss-Bonnet and Hirzebruch densities in the new-minimal formulation[14]. In other words, $T_{ab}{}^*T^{ab}$ is a chiral superfield whose F and G components are respectively $R_{abcd}{}^*R^{abcd}$ and $R_{abcd}\epsilon^{cdef}\,{}^*R^{ab}_{ef}$.

There is still another $U(1)$ invariance to take into account, namely the chiral R-symmetry of the super-conformal group[12], acting on the various fields in the following way ($Z = (z, \chi, h)$):

$$\delta z = -\frac{2}{3} i z, \quad \delta\chi = \frac{i}{3}\chi, \quad \delta h = \frac{4}{3} i h, \quad \delta\psi_\mu = -i\psi_\mu, \tag{9}$$

all other fields being invariant.

Since:

$$S = \log\frac{Z^3}{\mu^3} = (\log\frac{z^3}{\mu^3},\ \frac{3}{z}\chi,\ \frac{3}{z}h + \frac{3}{z^2}\bar{\chi}_L\chi_L) \equiv (s, \chi_S, h_S), \tag{10}$$

we get:

$$\delta S = (-2i, i\chi_S, 2ih_S). \tag{11}$$

Under the R-symmetry described above, the action transforms anomalously as[6,15]:

$$\begin{aligned} \delta\int d^4x\, L &= \frac{1}{384\pi^2}(7N_{3/2} - N_{1/2})\int d^4x\sqrt{g}\, R_{abcd}{}^*R^{abcd} \\ &= \frac{1}{384\pi^2}(7N_{3/2} - N_{1/2})\int d^2\theta d^4x\, \mathcal{E}\, T_{ab}{}^*T^{ab}, \end{aligned} \tag{12}$$

where in our model $N_{3/2} = N_{1/2} = 1$.

In eq. (12) the first contribution comes from the anomaly of the spin-3/2 fermion, the second from matter fermions. From eqs. (8,12) we can determine the F-term in the effective Lagrangian to be:

$$L_F = [\Sigma]_F = e\{\tilde{h} + \frac{i}{2}\bar{\tilde{\chi}}_L\sigma\cdot\psi_L + \frac{i}{2}\tilde{z}\bar{\psi}_{La}\sigma_{ab}\psi_{Lb}\}, \tag{13}$$

where $\Sigma = (\tilde{z}, \tilde{\chi}_L, \tilde{h})$ is:

$$\Sigma \equiv \frac{1}{384\pi^2}[S\, T_{ab}{}^*T^{ab} - 4T_{ab}{}^*T^{ab}\log\frac{T_{ab}{}^*T^{ab}}{\mu^6}]. \tag{14}$$

Since $\delta(T_{ab}T^{ab})_z = -2i(T_{ab}T^{ab})_z$ the effective action described in eqs. (13,14) reproduces correctly, at the classical level, all of the anomalies of the fundamental theory. In particular, it gives also the trace anomaly with the correct coefficient for the new-minimal formulation[6], where the contribution of the gravitational multiplet to T^μ_μ is $\frac{-7}{384\pi^2}$ [6,15] (in the old-minimal it would be $\frac{41}{384\pi^2}$). A comment here is in order. The effective lagrangian of eq. (13) reproduces correctly, at low energies, all 1PI vertices of the "fundamental" theory. It is this requirement, that fixes the form of the anomalous part. Explicitly, the anomalous variations of all the functional derivatives of the effective action, with respect to the low energy fields, must coincide with the analogous quantities of the exact theory. The absence of zero modes for the spin-1/2 fermions on a gravitational instanton does not affect the derivation, in that it follows from the presence of boundary terms in the Atiyah-Patodi-Singer formula for the gravitational anomaly[15]. Functional derivatives of boundary terms, on the other hand, can not modify (local) 1-PI vertices.

The explicit form of $T_{ab}{}^*T^{ab}$ was given in ref.[14]. In terms of the superconformal field-strengths $\hat{F}^+_{ab}$, $\hat{R}^+_{ab}$ and ψ_{ab} they found:

$$\begin{aligned} T_{ab}{}^*T^{ab} &= (\bar{\psi}_{ab})^*_L(\psi_{cd})_L + \theta(i\sigma_{cd}\hat{R}^+_{cdab} - 2i\hat{F}^+_{ab})^*\psi_{ab} \\ &+\theta^2\left\{2(\bar{\psi}_{ab})_L i\hat{\not{D}}^{*}(\psi_{ab})_R - \hat{F}^+_{ab}{}^*\hat{F}^+_{ab} + \frac{1}{2}\hat{R}^+_{abcd}{}^*\hat{R}^+_{abcd} + \frac{i}{2}\hat{R}^+_{abcd}{}^*\hat{R}^{*+}_{abcd}\right\}. \end{aligned} \quad (15)$$

The definitions of ψ_{ab}, $\hat{R}^+_{ab}$ and $\hat{F}^+_{ab}$ are as follows[14]:

$$\hat{R}^+_{ab} = R^+_{ab} - \frac{i}{4}\bar{\psi}\gamma_5\{\not{H},\sigma_{ab}\}\psi - \frac{1}{4}\bar{\psi}\{\sigma_{ab},\gamma_c\}\psi_{cd}e^d \quad (16)$$

$$\hat{F}^+ = F^+ + \frac{i}{2}\bar{\psi}\not{H}\psi - i\bar{\psi}\gamma_5(r + \frac{1}{2}\gamma\,\gamma\cdot r) \quad (17)$$

$$\psi_{\mu\nu} = D^+_{[\mu}\psi_{\nu]} \quad ,{}^*\psi_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}\psi_{\rho\sigma}. \quad (18)$$

Notice that on shell $\psi_{\mu\nu} + \frac{1}{2}\gamma_5^*\psi_{\mu\nu} = 0$, and that we used the short-hand notation $\psi = \psi_\mu dx^\mu$ and $r_a = \frac{1}{2}\gamma_5\gamma_b{}^*\psi_{ba}$. The derivative D^+_μ is covariantized with respect to the A_μ gauge field, and:

$$F^+_{\mu\nu} = D^+_{[\mu}A_{\nu]}. \quad (19)$$

The fundamental property of eq. (15) is that each component is a *local* total derivative. More precisely, the F component of eq. (15) can be rewritten, according to eq. (14), as[14]:

$$[T_{ab}{}^*T^{ab}]_F = -R^+_{ab}R^+_{ab} - iR^+_{ab}{}^*R^+_{ab} + 2F^{+2}$$

$$-\mathrm{d}\left\{8\overline{(D^+\psi)}\frac{1-\gamma_5}{2}(r + \frac{1}{2}\gamma\,\gamma\cdot r + \not{H}\psi_R)\right\}. \quad (20)$$

In eq. (20) we have made manifest the decomposition of $[T_{ab}{}^*T^{ab}]_F$ into a part which is a *local* total derivative and a part which can be *globally* written as the divergence of a well defined quantity. The integral of the latter, *but not of the former*, vanishes. Notice that in eq. (20) we have used the form-notation, so that $F^{+2} \equiv F^+ \wedge F^+$ (*i.e.* $\mathrm{d}A \wedge F^+$).

To evaluate the effective Lagrangian coming from eqs. (5), (14) and (15) we make the choice $V = F(S,\bar{S})$ in eq. (5), where, for the positivity of the potential, we have to assume[12,14] $\frac{\partial^2 F}{\partial s \partial \bar{s}} > 0$. F will for the rest remain unspecified, a detailed computation of it being a non-trivial dynamical problem. For our purposes its precise form will not be needed.

To evaluate the F-term in eq. (15) we use the formula for the multiplication of chiral multiplets:

$$f(\Sigma) = \left(f(z),\ \chi^i f_i,\ h^i f_i - \frac{1}{4}\bar{\chi}^i_L \chi^j_L f_{ij}\right), \tag{21}$$

with $\Sigma^i = z^i + \theta\chi^i + 1/2\theta^2 h^i$ and $f_i = \partial f/\partial z^i$.

From eq. (21) the F-density for $S\,T_{ab}\,{}^*T^{ab}$ reads:

$$\begin{aligned}
[S\,T_{ab}\,{}^*T^{ab}]_F \;=\; & s\left(-R^+_{ab}R^+_{ab} - iR^+_{ab}\,{}^*R^+_{ab} + 2F^{+2} - \mathrm{d}\left\{8\overline{(D^+\psi)}\frac{1-\gamma_5}{2}(r + \frac{1}{2}\gamma\,\gamma\cdot r + \not{H}\psi_R)\right\}\right) \\
& + h_S(\bar{\psi}_{ab})_L\,{}^*(\psi_{ab})_L - \frac{1}{4}\bar{\chi}_S(T_{ab}\,{}^*T^{ab})_\chi + \frac{i}{2}\bar{\chi}_L\gamma^\mu\psi_\mu(\bar{\psi}_{ab})_L\,{}^*(\psi_{ab})_L.
\end{aligned} \tag{22}$$

In a similar way we obtain:

$$\begin{aligned}
[T_{ab}\,{}^*T^{ab}\log\frac{T_{ab}\,{}^*T^{ab}}{\mu^5}]_F = & -\log\frac{(\bar{\psi}_{ab})_L(\psi_{ab})_L}{\mu^5}[R^+_{ab}R^+_{ab} + iR^+_{ab}\,{}^*R^+_{ab} - 2F^{+2} \\
& + \mathrm{d}\left\{8\overline{(D^+\psi)}\frac{1-\gamma_5}{2}(r + \frac{1}{2}\gamma\,\gamma\cdot r + \not{H}\psi_R)\right\}] + \frac{i}{2}(\overline{T_{ab}\,{}^*T^{ab}})_\chi\gamma^\mu\psi_\mu \\
& - \frac{1}{(\psi_{ab})_L\,{}^*(\psi^{ab})_L}(\overline{T_{ab}\,{}^*T^{ab}})_\chi(T_{ab}\,{}^*T^{ab})_\chi + (T_{ab}\,{}^*T^{ab})_h.
\end{aligned} \tag{23}$$

In a maximally symmetric background, where the bilinear $(\bar{\psi}_{ab})_L(\psi_{ab})_L$ acquires a non-zero vacuum expectation value, hereafter called C, it is easy to evaluate the scalar potential for the effective supergravity Lagrangian. The only relevant terms in eqs. (5), (22) and (23) are in fact[14]:

$$\int d^4x L = \int d^4x\ e\{\frac{1}{4}e^{2\sigma}R + \frac{1}{2}F_{s\bar{s}}h_s h_{\bar{s}} + h_s C + \text{h.c.}\}, \tag{24}$$

with $e^{2\sigma} = 1 - sF_s$ and $F_s \equiv \partial F/\partial s$.

The Einstein term in the action has been retained to find the correct rescaling of the potential term, corresponding to the normalization $\int d^4x\, eR/2$.

After solving for h_s and performing a Weyl rescaling of the action in eq. (24) according to $e^a_\mu \to e^\sigma e^a_\mu$, we arrive at the action:

$$\int d^4x L = \int d^4x\ e\{\frac{1}{2}R - (F_{s\bar{s}})^{-1}e^{-2\sigma}|C|^2 + \ldots\}. \tag{25}$$

Eq. (25) is the main result of our investigation; it shows that by breaking supersymmetry through the introduction of a gravitino field-strength condensate one develops a cosmological constant of the order of $|C|^2/M_P^6$. This result is independent of the specific form of $F(S,\bar{S})$. The only requirement, namely the positivity of $F_{s\bar{s}}$, being indispensable in order to guarantee the positivity of the kinetic terms for the scalar multiplet S [12,14]. The fact that our result is independent of F is a most welcomed feature, since the effective Lagrangian approach cannot fix uniquely the exact form of F.

The introduction of more general interactions is not going to change this picture, as long as the breaking is completely due to the non-perturbative gravitino condensate. The only change is in the definition of S, that in a more general case will be the (composite) Goldstino superfield.

The only way to keep the cosmological constant small would be by introducing a superpotential at tree level, giving a negative contribution to the vacuum energy which could be cancelled by the gravitino condensate.

The situation analyzed in this paper refers to the simple case of pure supergravity with a single sliding chiral field. In order to extend our considerations to strings several points have to be clarified. First of all, one should find viable way of computing instanton effects in strings and show that the particular metric chosen to produce the non-perturbative effect (Eguchi-Hanson) is indeed a solution of the string equations of motion. This is probably possible, since the gravitational instantons of Eguchi and Hanson[16] are hyper-Kähler manifolds, and therefore probably conformal field theories on the sphere[17].

Secondly, the rôle of the antisymmetric field $B_{\mu\nu}$ and its field strength $H_{\mu\nu\rho}$ has to be clarified. On a non-trivial background, in fact, it is not possible to set $H_{\mu\nu\rho}=0$. One should, instead, solve the equation of motion for $B_{\mu\nu}$ consistently together with Einstein's equations. The presence of a non-zero $H_{\mu\nu\rho}$ seems to introduce, at least for large instantons, of size $a \gg M_P^{-1}$, an exponential enhancement in the scale of the gravitino condensation, because of the weight factor,

$$e^{+\int d^4x \hat{H}_{\mu\nu\rho}\hat{H}^{\mu\nu\rho}}, \tag{26}$$

in the functional integral yielding $\langle\bar{\psi}_{ab}\psi^{ab}\rangle$. $\hat{H}_{\mu\nu\rho}$ is the value of $H_{\mu\nu\rho}$ given by the equations of motion, which are quadratic in $B_{\mu\nu}$. This phenomenon would signal that, even in the string effective-supergravity lagrangian, a nonvanishing gravitino condensate could destabilize the perturbative vacuum[1] .

Gravitino condensation is at the moment an interesting way of breaking supersymmetry, owing to the model-independent character of the mechanism; still, to make it fit into a standard low-energy scenario, with $m_{3/2} \sim \mathcal{O}(1TeV)$ and $\Lambda << 10^{-4}(eV)^4$, other phenomena must be invoked.

[1]In the Gross-Neveu model, for instance, the gauge invariant quantities behave as e^{1/g^2} for small coupling constant g in the symmetric (unstable) vacuum.

Acknowledgements:
We thank K. Konishi, N. Magnoli, H. Panagopoulos, S. Ferrara and G. Veneziano for discussions on this subject and comments to the manuscript. M.P. would like to thank Fermilab for the hospitality extended to him while this work was done.

References

[1] S. Deser and B. Zumino, Phys. Rev. Lett. 38 (1977) 1433;
E. Cremmer, S. Ferrara, L. Girardello and A. van Proeyen, Nucl. Phys. B212 (1983), 413.

[2] H.P. Nilles, Phys. Rep. 110 (1984), 1.

[3] R. Rohm, Nucl. Phys. B237 (1984), 553;
C. Kounnas and M. Porrati, Preprint UCB-PTH-88/1, LBL-24789, LPTENS-88/3.

[4] H.P. Nilles, Phys. Lett. 115B (1982) 193;
S. Ferrara, L. Girardello and H.P. Nilles, Phys. Lett. 125B (1983), 459;
J. P. Derendinger, L.E. Ibáñez and H.P. Nilles, Phys. Lett. 155B (1985), 65;
M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. 156B (1985), 55;
C. Kounnas and M. Porrati, Phys. Lett. 191B (1987), 91.

[5] K. Konishi, N. Magnoli and H. Panagopoulos, Genoa Preprint 88-0162 (1988).

[6] S. J. Gates, M.T. Grisaru, M. Roček and W. Siegel, "Superspace: One Thousand and One Lessons in Supersymmetry", (Frontiers of Physics, Vol 58), Benjamin, New York (1983).

[7] R. Delbourgo and A. Salam, Phys. Lett. 40B (1972), 381;
T. Eguchi and P.G.O. Freund, Phys. Rev. Lett. 37 (1976), 1251.

[8] K. Konishi, Phys. Lett. 135B (1984), 439.

[9] S.W. Hawking and C.N. Pope, Nucl. Phys. B146 (1978), 385.

[10] G. Veneziano and S. Yankielowicz, Phys. Lett. 113B (1982), 231.

[11] I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B241 (1984), 493;
T.R. Taylor, G. Veneziano and S. Yankielowicz, Nucl. Phys. B218 (1983), 493.

[12] A. van Proeyen, in "Supergravity 1983", (ed. B. Milewski, World Scientific), 1983.

[13] S. Ferrara and P. Van Nieuwenhuizen, Phys. Lett. 74B (1978), 333;
K.S. Stelle and P.C. West, Phys. Lett. 74B (1978), 330;
M. Sohnius and P.C. West, Nucl. Phys. B198 (1981), 493;
V.P. Akulov, D.V. Volkov and V.A. Soroka, Theor. Math. Phys. 31 (1977), 285.

[14] S. Ferrara and S. Sabharwal, Preprint CERN-TH.4995/88, UCLA/88/TEP/11.

[15] M. Duff and S.M. Christensen, Phys. Lett. 76B (1978), 571; Nucl. Phys. B154 (1979), 301;
M. Grisaru in "Proc. NATO Adv. Study Inst. on Gravitation, Cargèse 1978";
P. Van Nieuwenhuizen, M.T. Grisaru, H. Röhmer and N. Nielsen, Nucl. Phys. B140 (1978), 477;
A. Niemi and G. Semenoff, Nucl. Phys. B291 (1986) 131.

[16] T. Eguchi and A.J. Hanson, Phys. Lett. 74B (1978), 249;
T. Eguchi, P. Gilkey and A.J. Hanson, Phys. Rep. 66 (1980), 213.

[17] C.M. Hull, Nucl. Phys. B267 (1986), 266.