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No–scale $N = 4$ supergravity coupled to Yang–Mills: the scalar potential and super– Higgs effect

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Abstract

We derive the scalar potential of the effective theory of type IIB orientifold with 3-form fluxes turned on in presence of non abelian brane coordinates.
$N = 4$ supergravity predicts a positive semidefinite potential with vanishing cosmological constant in the vacuum of commuting coordinates, with a classical moduli space given by three radial moduli and three RR scalars which complete three copies of the coset $(U(1, 1+n)/U(1) \otimes U(1+n))$, together with $6n$ $D3$–branes coordinates, n being the rank of the gauge group G. Implications for the super Higgs mechanism are also discussed.

1 Introduction

Recently it was shown [1] that the bulk sector of the type IIB orientifolds with 3–form fluxes turned on [2], [3] corresponds, for a generic choice of fluxes which give a vanishing vacuum energy, to a gauged $N = 4$ supergravity with a linear action of $SL(2) \times GL(6)$ on the 12 vector potentials (belonging to the representation $(\mathbf{2},\mathbf{6})$) of $N = 4$ supergravity coupled to six $N = 4$ vector multiplets [4].
An important ingredient is that, in order to obtain the correct theory, only a $GL(6)$ subgroup of the global symmetry group $SO(6, 6)$ acts linearly on the gauge potentials, while other $SO(6, 6)$ transformations mix electric with magnetic field strengths [5].
Fluxes turned on correspond, in the supergravity language, to a gauging of translational isometries of $SO(6, 6)$ in the decomposition

$$\mathfrak{so}(6,6) = \mathfrak{gl}(6,\mathbb{R}) + \mathbf{15}^{'+} + \mathbf{15}^{-} \tag{1.1}$$

The only fields which are charged under this gauging are the 15 RR scalars whose dual have a charged coupling to the 12 vectors

$$\nabla_\mu B^{\Lambda\Sigma} = dB^{\Lambda\Sigma} + f^{\Lambda\Sigma\Delta\alpha} A_{\Delta\alpha} \tag{1.2}$$

Here Λ, Σ, Δ are $GL(6)$ indices, the axion $B^{\Lambda\Sigma}$ is an antisymmetric tensor and the "charges" $f^{\Lambda\Sigma\Delta\alpha}$ are the 3–form (RR and NS) fluxes [1], [2], [3].
The corresponding $N = 4$ supergravity sector has a scalar potential given by

$$V = \frac{1}{12}|F^{-\,IJK}|^2 \tag{1.3}$$

where

$$F^{-\,IJK} = \frac{1}{2}\left(F^{IJK} - i\,{}^*F^{IJK}\right) \tag{1.4}$$

and

$$F^{IJK} = L^\alpha f^{IJK}_\alpha \tag{1.5}$$

L^α parametrizes the $SU(1, 1)/U(1)$ coset of the NS and RR dilaton sector as follows: let the generic element of $SU(1, 1)/U(1)$ be given by s,

$$s = \begin{pmatrix} \phi_1 & \overline{\phi}_2 \\ \phi_2 & \overline{\phi}_1 \end{pmatrix} \qquad (\phi_1\overline{\phi}_1 - \phi_2\overline{\phi}_2 = 1) \tag{1.6}$$

introducing the 2-vector

$$L^\alpha \equiv \begin{pmatrix} L^1 \\ L^2 \end{pmatrix} = \frac{1}{\sqrt{2}}\begin{pmatrix} \phi_1 + \phi_2 \\ -i(\phi_1 - \phi_2) \end{pmatrix} \tag{1.7}$$

$$L_\alpha \equiv \epsilon_{\alpha\beta} L^\beta \tag{1.8}$$

the identity $\phi_1\overline{\phi}_1 - \phi_2\overline{\phi}_2 = 1$ becomes

$$L^\alpha \overline{L}^\beta - \overline{L}^\alpha L^\beta = i\epsilon^{\alpha\beta} \tag{1.9}$$

To relate this to the standard complex dilaton,

$$S = ie^\varphi + C \tag{1.10}$$

where φ is the string dilaton and C is its RR partner, let us set

$$\frac{\phi_2}{\phi_1} = z = \frac{-i+S}{i+S} \tag{1.11}$$

then we have

$$e^K = \frac{1}{2i(\overline{S}-S)} = \frac{e^{-\varphi}}{4} \tag{1.12}$$

so that

$$\phi_2 = \frac{1}{2}[i(e^{\varphi}-1)+C]e^{-\frac{\varphi}{2}} \tag{1.13}$$

$$\phi_1 = \frac{1}{2}[i(e^{\varphi}+1)+C]e^{-\frac{\varphi}{2}} \tag{1.14}$$

Using the above relations, one easily shows that the potential given in reference [1] agrees with the potential recently discussed in reference [6] since

$$|\phi_2|^2 = \frac{1}{2}(cosh\,\varphi - 1 + \frac{1}{2}e^{-\varphi}C^2) \tag{1.15}$$

The gravitino mass matrix is given by [1], [5]

$$S_{AB} = -\frac{i}{48}\overline{F}^{-IJK}(\Gamma_{IJK})_{AB} \tag{1.16}$$

This is a no–scale model [7], [8] in that the contribution of $|S_{AB}|^2$ to the gravitational potential exactly cancels (even away from the extremum) with the four goldstinos which come from the $\overline{\mathbf{4}}$ in the $\mathbf{6}\times\mathbf{4} = \mathbf{20}+\overline{\mathbf{4}}$ $SU(4)$ decomposition of the gauginos of the six matter multiplets.

2 Coupling to Yang–Mills matter

The above discussion gives the result derived in reference [1] which agrees with the potential obtained by the bulk part of the action of type IIB on $T^6/\mathbb{Z}_2$ orientifold [2], [3], [6].

A less obvious result is the coupling of this system to additional Yang–Mills matter, that is in presence of additional n vector multiplets which are Lie algebra valued on some compact group G. This result, in the superstring framework, comes from the coupling of the Born–Infeld non abelian action to the gravity sector of the ambient space. Some of these couplings have been computed in the literature [9], [10], but their completion, to give an effective theory with $N = 4$ local supersymmetry, is not a straightforward exercise.

This completion was obtained [11] in the particular case of a $N = 1$ sector coupled to Yang–Mills, as one would obtain if all degrees of freedom (in particular three massive gravitinos) in the partial breaking $N = 4 \longrightarrow N = 1$ were integrated out. The result, with only one residual flux breaking $N = 1 \longrightarrow N = 0$, is a no–scale model with a particularly simple structure.

The amazing fact of this result is that the non abelian effective theory does not contain, for example, terms of order a^3, where a are the $D3$–brane coordinates, but only a pure a^4

term, as in the pure Born–Infeld with non gravitational back–reaction corrections. The cubic term is predicted by an explicit calculation [9] but it is known to vanish if the equation $F^{IJK-}=0$ is used. This equation, among other things, stabilizes the (complex) dilaton in terms of flux entries [2], [3].

To have a better insight of these results we now present the full $N=4$ potential where both fluxes and non abelian brane coordinates are present. This potential is completely predicted by $N=4$ supergravity with gauge group $\mathcal{G}$ of dimension $12+n$ ($n=dim\,G$) which is the direct product of 12 translations with the compact Lie group G: $\mathcal{G}=T_{12}\otimes G$. Note that this requires a non standard symplectic embedding [12] of the full $SL(2,\mathbb{R})\times SO(6,6+n)$ symmetry in $Sp(24+2n)$ group because $SL(2,\mathbb{R})$ acts linearly on the first 12 vectors, but acts as an electric–magnetic duality on the remaining 2n field strengths and their dual. This embedding was discussed in reference [5] and the subgroup which acts on the gauge potentials is $GL(6)\times SO(n)$. In this framework the scalar fields $B^{\Lambda\Sigma}$, a_i^Λ ($i=1\dots n$) are treated as tensors of $GL(6)$ and they complete the $GL(6)/SO(6)$ manifold to the full $SO(6,6+n)/SO(6)\times SO(6+n)$. This formulation, in absence of gauge coupling and fluxes, is related by a duality transformation to the $N=4$ action constructed a long time ago in reference [13], [14]. However, when the fluxes and the non abelian couplings of G are turned on, this action is no longer equivalent to any of the previously proposed actions and it allows to obtain a no–scale extended $N=4$ supergravity with non abelian gauge interactions.

The potential can be computed from the fermion shifts modified by the charge couplings. These shifts can be computed from the $N=4$ gauged theory using superspace Bianchi identities

$$\begin{aligned}
\delta\psi_{A\mu}^{(shift)} &= S_{AB}\gamma_\mu\varepsilon^B=-\frac{i}{48}(\overline{F}^{IJK-}+\overline{C}^{IJK-})(\Gamma_{IJK})_{AB}\gamma_\mu\epsilon^B && (2.17)\\
\delta\chi^{A\,(shift)} &= N^{AB}\epsilon_B=-\frac{1}{48}(\overline{F}^{IJK+}+\overline{C}^{IJK+})(\Gamma_{IJK})^{AB}\epsilon_B && (2.18)\\
\delta\lambda_A^{I\,(shift)} &= Z_A^{IB}\epsilon_B=\frac{1}{8}(F^{IJK}+C^{IJK})(\Gamma_{JK})_A^{\ B}\epsilon_B && (2.19)\\
\delta\lambda_{iA}^{(shift)} &= W_{iA}^{\ \ B}\epsilon_B=\frac{1}{8}L_2\,q^{Jj}q^{Kk}\,c_{ijk}(\Gamma_{JK})_A^{\ B}\epsilon_B && (2.20)
\end{aligned}$$

Here c_{ijk} are the structure constants of G, C^{IJK} are the boosted structure constants defined as

$$C^{IJK}=L_2q^{Ii}q^{Jj}q^{Kk}\,c_{ijk} \tag{2.21}$$

and $q_{Ii}=E_{I\Lambda}a_i^\Lambda$ where $E_{I\Lambda}$ are the coset representatives of $GL(6)/SO(6)$.

The knowledge of the fermion shifts allows us to use the Ward identity of supersymmetry in the Lagrangian [?] to compute the scalar potential which turns out to be

$$V=\frac{1}{12}|F^{IJK-}+C^{IJK-}|^2+\frac{1}{32}|L_2c_{ijk}q^{jJ}q^{kK}|^2 \tag{2.22}$$

Some comments are in order. If we set $q^{iI}=0$ (or commuting) we retrieve the previous potential given in reference [1]. On the other hand, if we set $F^{IJK-}=0$ we retrieve the standard potential (as for instance it comes from the heterotic string compactified on T_6 [18], [19], [20]).

Interestingly, the first term contains an interference term

$$F^{IJK-}C^*_{IJK-}+cc \tag{2.23}$$

which is what was obtained from the Born–Infeld action [9]. However, as it can be seen, this term completes to a perfect square because of the extra (pure gravitational) higher order a^6 term $\frac{1}{12}|C^{IJK-}|^2$. This actually explains why, integrating out the dilaton, such terms may disappear and this is consistent with the $N=1$ reduction studied before [11]. Also note that the a^6 terms would not always disappear if more supersymmetry remains unbroken, for instance, integrating out the dilaton in the $N=4 \longrightarrow N=3$ reduction [1], the non vanishing flux $f_{ijk} = \epsilon_{ijk} f$ just predicts that the pure holomorphic part of C^{IJK-}, i.e. the singlet in the $SU(4) \longrightarrow SU(3)$ decomposition $\mathbf{10} \longrightarrow \mathbf{6} + \overline{\mathbf{3}} + \mathbf{1}$ should not be present. This is indeed true as in the $N=3$ supergravity the a^6 terms are of the form $|zz\overline{z}|^2$ [15], where z^A are holomorphic triplet coordinates in the splitting $\mathbf{6} \longrightarrow \mathbf{3} + \overline{\mathbf{3}}$
Note that a puzzle seems to emerge on the fact that the Yang–Mills contribution to the gravitino mass is proportional to

$$S_{AB} \propto c_{ijk} z^{Ci} z^{Dj} \overline{z}^k_{(B}\epsilon_{A)CD} \tag{2.24}$$

where A, B, C are $SU(3)$ indices; this term is not holomorphic in the z^{Ai} coordinates, while it is holomorphic in the $N=1$ case [11].
The resolution is the fact that the supergravity transformations of the three z^A (for a fixed value of the i index) and the four left–handed λ_A, λ gauginos are

$$\delta z^A = \epsilon^{ABC}\overline{\lambda}_B \epsilon_C + \overline{\lambda}\epsilon^A \tag{2.25}$$

$$\delta\lambda_A = -i\epsilon_{ABC}\gamma^\mu \partial_\mu z^B \epsilon^C \tag{2.26}$$

$$\delta\lambda = i\gamma^\mu \partial_\mu \overline{z}_A \epsilon^A \tag{2.27}$$

So, if we pick up $\epsilon^A = (\epsilon^1, 0, 0)$ as in $N=1$ reduction, we see that the left–handed $N=1$ multiplets are

$$(z^3,\, \lambda_2) \quad (z^2,\, \lambda_3) \quad (\overline{z}_1,\, \lambda) \tag{2.28}$$

so that we can identify $\overline{z}$ with the third holomorphic coordinate. The remaining λ_1 is the $N=1$ gaugino.
The condition for the potential to have minimum with vanishing cosmological constant is

$$\delta\lambda_{iA}^{(shift)} = 0 \tag{2.29}$$

$$\delta\chi^{A(shift)} = 0 \tag{2.30}$$

$$\delta\lambda_{IA}^{(20)(shift)} = 0 \tag{2.31}$$

The first equation implies $c_{ijk} q^{iI} q^{jJ} = 0$, that is the q^{iI} are in the Cartan subalgebra of G. This also implies $C^{IJK} = 0$ and then the other equations imply $F^{IJK-} = 0$.
The latter equation implies, for arbitrary values of the gravitino masses, that the complex scalar in L_2 (coordinates of $SU(1,1)/U(1)$) and eighteen of the twenty–one coordinates of $GL(6)/SO(6)$ are stabilized. Twelve axions in $B^{\Lambda\Sigma}$ are eaten by twelve vectors with a Higgs mechanism [1], leaving $6+6n$ ($n = Rank(G)$) flat directions.
By integrating out the massive modes, the space is (at least classically) the product of three $CP^{n+1} = U(1, 1+n)/U(1) \times U(1+n)$ σ–model, as was shown in reference [11].

3 N=2 examples

A particular interesting case to study is also the reduction to $N=2$, where both vector multiplets and hypermultiplets are present and the general for of the non linear σ–model is predicted to be [11]

$$\frac{U(1,1+n)}{U(1)\times U(1+n)}\times\frac{U(2,2+n)}{U(2)\times U(2+n)} \tag{3.32}$$

This theory has two flux parameters corresponding to a gauging of a two translational isometry of the $\frac{U(2,2)}{U(2)\times U(2)}$ coset. They are gauged by the two vectors of the $n=0$ gravitational sector [4]. Special geometry predicts [16] that the $N=2\longrightarrow N=1$ partial breaking is possible only if the $\frac{U(1,1+n)}{U(1)\times U(1+n)}$ special manifold has a symplectic section which cannot be derived from a prepotential function $F(X)$. We now show that this is indeed the case. The argument is very similar to the analogous example, studied in reference [17] for the special cosets $\frac{SU(1,1)}{U(1)}\times\frac{SO(2,n)}{SO(2)\times SO(n)}$.
For the CP_{n+1} special cosets, the holomorphic prepotential in the natural basis where the $U(n)$ symmetry is manifest is

$$F(X)=i(X^0X^1+X^aX^a)\quad a=2,\dots n+1 \tag{3.33}$$

with symplectic (holomorphic) sections $(X^\Lambda,\,F_\Lambda)$, $\Lambda=0,1,\dots n+1$ with

$$F_0=iX^1;\quad F_1=iX^0;\quad F_a=2iX^a \tag{3.34}$$

The corresponding Käler potential is

$$K=-log\,i(\overline{X}^\Lambda F_\Lambda-X^\Lambda\overline{F}_\Lambda)=-log[-2(S+\overline{S}+2x^a\overline{x}^a)] \tag{3.35}$$

where $S=\frac{X^1}{X^0}$ and $x^a=\frac{X^a}{X^0}$. Let us now perform the following symplectic change of holomorphic section, using a symplectic matrix

$$\mathcal{S}=\begin{pmatrix}A & B\\ C & D\end{pmatrix} \tag{3.36}$$

where

$$A=D=\begin{pmatrix}\mathbf{p}_{2\times2} & \mathbf{0}_{2\times n}\\ \mathbf{0}_{n\times2} & \mathbf{1}_{n\times n}\end{pmatrix};\quad -B=C=\begin{pmatrix}\mathbf{q}_{2\times2} & \mathbf{0}_{2\times n}\\ \mathbf{0}_{n\times2} & \mathbf{0}_{n\times n}\end{pmatrix} \tag{3.37}$$

$$\mathbf{p}_{2\times2}=\begin{pmatrix}1 & 0\\ 0 & 0\end{pmatrix};\qquad \mathbf{q}_{2\times2}=\begin{pmatrix}0 & 0\\ 0 & 1\end{pmatrix} \tag{3.38}$$

which satisfy the symplectic conditions

$$A^TD-C^TB=A^2+C^2=1;\quad A^TC=B^TD=0 \tag{3.39}$$

We have in the new basis

$$\begin{pmatrix}\widetilde{X}\\ \widetilde{F}\end{pmatrix}=\begin{pmatrix}A & B\\ C & D\end{pmatrix}\begin{pmatrix}X\\ F\end{pmatrix} \tag{3.40}$$

where

$$\widetilde{X}^\Lambda=(X^0,\,-iX^0,\,X^a);\quad \widetilde{F}_\Lambda=(iX^1,\,X^1,\,2iX^a) \tag{3.41}$$

Since $\widetilde{X}^\Lambda$ does not contain X^1 this basis is singular and no prepotential $\widetilde{F}(\widetilde{X})$ exists. The matrix $\mathcal{N}_{\Lambda\Sigma}$ can be obtained from the special geometry relations [17]

$$\begin{aligned} & F_\Lambda = \mathcal{N}_{\Lambda\Sigma} X^\Sigma, \qquad \mathcal{D}_{\bar{l}} \overline{F}_\Lambda = \mathcal{N}_{\Lambda\Sigma} \mathcal{D}_{\bar{l}} \overline{X}^\Sigma \\ & \mathcal{D}_i X^\Lambda = \partial_i X^\Lambda + \partial_i K X^\Lambda, \quad \mathcal{D}_i F_\Lambda = \partial_i F_\Lambda + \partial_i K F_\Lambda \end{aligned} \tag{3.42}$$

The vector kinetic matrix $\mathcal{N}_{\Lambda\Sigma}$ turns out to be holomorphic and is given in this basis by

$$\widetilde{\mathcal{N}}_{\Lambda\Sigma} = (\widetilde{\mathcal{N}}_{ab}, \widetilde{\mathcal{N}}_{a0}, \widetilde{\mathcal{N}}_{a1}, \widetilde{\mathcal{N}}_{01}, \widetilde{\mathcal{N}}_{00}, \widetilde{\mathcal{N}}_{11}) \tag{3.43}$$

with

$$\begin{aligned} & \widetilde{\mathcal{N}}_{ab} = -2i\delta_{ab}, \widetilde{\mathcal{N}}_{a0} = 2ix_a, \widetilde{\mathcal{N}}_{a1} = -2x_a \\ & \widetilde{\mathcal{N}}_{01} = x^a x_a, \widetilde{\mathcal{N}}_{00} = iS - ix^a x_a, \widetilde{\mathcal{N}}_{11} = iS + ix^a x_a \end{aligned} \tag{3.44}$$

So that the vector kinetic term takes the form

$$Im(\mathcal{N}_{\Lambda\Sigma} F^{+\Lambda}_{\mu\nu} F^{+\Sigma\mu\nu}) \tag{3.45}$$

The above result show that the $N = 2$ model is compatible with partial breaking of $N = 2$ supersymmetry to $N = 1, 0$ ($N = 2$ models with partial supersymmetry breaking have been considered in the literature [21], [22], [23], [24]). The moduli space of such $N = 2 \longrightarrow N = 1 \longrightarrow N = 0$ theories is given by three copies ofthe coset $\frac{SU(1,1)}{U(1)}$. One is the original $\frac{SU(1,1)}{U(1)}$ in the vector multiplet sector, while $\left(\frac{SU(1,1)}{U(1)}\right)^2 \subset \frac{U(2,2)}{U(2)\times U(2)}$ comes from the quaternionic manifold. Note that this is an extension of the minimal model based on the coset $\frac{Usp(2,2)}{Usp(2)\times Usp(2)}$ of reference [16].
A detailed analysis of the above situation will be given elsewhere.

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References

[1] R. D'Auria, S. Ferrara and S. Vaula, "N = 4 gauged supergravity and a IIB orientifold with fluxes,"New J. of Phys. 4:71,2001; arXiv:hep-th/0206241.

[2] A. R. Frey and J. Polchinski, "N = 3 warped compactifications". hep-th/0201029.

[3] S. Kachru, M. Schulz and S. Trivedi, "Moduli Stabilization from Fluxes in a Simple IIB Orientifold". hep-th/0201028.

[4] L. Andrianopoli, R. D'Auria, S. Ferrara and M. A. Lledo, "Duality and spontaneously broken supergravity in flat backgrounds," Nucl. Phys. B **640** (2002) 63 [arXiv:hep-th/0204145].

[5] V. A. Tsokur and Y. M. Zinoviev, "Spontaneous supersymmetry breaking in N = 4 supergravity with matter," Phys. Atom. Nucl. **59** (1996) 2192 [Yad. Fiz. **59N12** (1996) 2277] [arXiv:hep-th/9411104].

[6] A. R. Frey and A. Mazumdar, "3-form induced potentials, dilaton stabilization, and running moduli," arXiv:hep-th/0210254.

[7] E. Cremmer, S. Ferrara, C. Kounnas and D. V. Nanopoulos, "Naturally Vanishing Cosmological Constant In N=1 Supergravity," Phys. Lett. B **133** (1983) 61.

[8] J. R. Ellis, A. B. Lahanas, D. V. Nanopoulos and K. Tamvakis, "No - Scale Supersymmetric Standard Model," Phys. Lett. B **134** (1984) 429. J. R. Ellis, C. Kounnas and D. V. Nanopoulos, "No Scale Supersymmetric Guts," Nucl. Phys. B **247** (1984) 373. For a review see A. B. Lahanas and D. V. Nanopoulos, "The Road To No Scale Supergravity," Phys. Rept. **145** (1987) 1.

[9] R. C. Myers, "Dielectric-branes," JHEP **9912** (1999) 022 [arXiv:hep-th/9910053].

[10] J. Polchinski and M. J. Strassler, "The string dual of a confining four-dimensional gauge theory," arXiv:hep-th/0003136.

[11] S. Ferrara and M. Porrati, "N = 1 no-scale supergravity from IIB orientifolds," Phys. Lett. B **545** (2002) 411 [arXiv:hep-th/0207135].

[12] M. K. Gaillard and B. Zumino, "Duality Rotations For Interacting Fields," Nucl. Phys. B **193** (1981) 221.

[13] E. Bergshoeff, I. G. Koh and E. Sezgin, "Coupling Of Yang-Mills To N=4, D = 4 Supergravity," Phys. Lett. B **155** (1985) 71.

[14] M. de Roo and P. Wagemans, "Gauge Matter Coupling In N=4 Supergravity," Nucl. Phys. B **262** (1985) 644.

[15] L. Castellani, A. Ceresole, S. Ferrara, R. D'Auria, P. Fre and E. Maina, "The Complete N=3 Matter Coupled Supergravity," Nucl. Phys. B **268** (1986) 317.

[16] S. Ferrara, L. Girardello and M. Porrati, "Minimal Higgs Branch for the Breaking of Half of the Supersymmetries in N=2 Supergravity," Phys. Lett. B **366** (1996) 155 [arXiv:hep-th/9510074].

[17] A. Ceresole, R. D'Auria, S. Ferrara and A. Van Proeyen, "Duality transformations in supersymmetric Yang-Mills theories coupled to supergravity," Nucl. Phys. B **444** (1995) 92 [arXiv:hep-th/9502072].

[18] E. Witten, "Dimensional Reduction Of Superstring Models," Phys. Lett. B **155** (1985) 151.

[19] S. Ferrara, C. Kounnas and M. Porrati, "General Dimensional Reduction Of Ten-Dimensional Supergravity And Superstring," Phys. Lett. B **181** (1986) 263.

[20] J. H. Schwarz and A. Sen, "Duality Symmetries Of 4-D Heterotic Strings," Phys. Lett. B **312** (1993) 105 [arXiv:hep-th/9305185].

[21] S. Cecotti, L. Girardello and M. Porrati, "Some Features Of Susy Breaking In N=2 Supergravity," Phys. Lett. B **151** (1985) 367.

[22] Y. M. Zinovev, "Soft Breaking Terms From N=2 Supergravity," Sov. J. Nucl. Phys. **46** (1987) 724 [Yad. Fiz. **46** (1987) 1240].

[23] Y. M. Zinovev, "Spontaneous Symmetry Breaking In N=2 Supergravity," Sov. J. Nucl. Phys. **46** (1987) 540 [Yad. Fiz. **46** (1987) 943] [arXiv:hep-th/9512041].

[24] P. Fre, L. Girardello, I. Pesando and M. Trigiante, "Spontaneous N $= 2 \rightarrow$ N $= 1$ local supersymmetry breaking with surviving compact gauge groups," Nucl. Phys. B **493** (1997) 231 [arXiv:hep-th/9607032].