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A D D E N D U M

AN EXCEPTIONAL $N=2$ SUPERGRAVITY WITH FLAT POTENTIAL
AND PARTIAL SUPER-HIGGS

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See also ref. [8].

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A B S T R A C T

We present a new N=2 supergravity model, which is exceptional in the sense that it is not described by the usual tensor calculus. It has flat (i.e., identically vanishing) scalar potential and partial breaking to N=1 supersymmetry ($\Lambda=0$). Our construction is based on a suitable "contraction" of standard AdS supergravity models with matter. We also give the procedure to obtain flat models with hierarchical supersymmetry breakings.

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In ref.[1] it was shown that the general N=2 supergravity models, as are known from the superconformal tensorial calculus [2], cannot have only one broken supersymmetry if space-time is flat. However, there remains open the question about the existence of other "exceptional" classes of N=2 models (not being covered by ref.[2]) which can have partial supersymmetry breaking and vanishing cosmological constant $\Lambda=0$.

In the present letter, we answer the question in the affirmative, and show how a model with the required properties can indeed be obtained starting from the tensorial calculus by a suitable "contraction" procedure.

The resulting model (here we present the complete details only for the very minimal one - a large class of them can be straightforwardly constructed) has a number of remarkable features.

First of all, it is flat in the sense that the scalar potential vanishes identically; in particular, the physical scalars live in the same $SU(1,1)/U(1)$ Kaehler manifold of the standard N=1 flat potential [3], and the field dependence of the gravitino masses is also reminiscent of that of ref.[3]. Thus, we have new flat N=2 supergravities besides those considered in ref.[4]. Recall that the construction of ref.[4], being based on vector multiplets only, always leads to a degenerate gravitino mass-matrix (1). In our case this need not to be so; indeed, our construction can be very easily generalized to the case of N=2 flat supergravity models with any given ratio of the gravitino masses (see below for details). In this way we can have simultaneously a flat potential and hierarchical SUSY breakings (2), a situation that may be of use in the applications.

However, the more interesting features of the model are in the structure of the spin-one sector. The simplest model has two spin-1 particles, of opposite parity, which are massive and are superpartners (with respect to the unbroken supersymmetry) of the massive gravitino.

The only field which is not neutral under the gauge symmetries corresponding to these vectors is a "quaternionic" scalar b^α_i , which is real in the sense that [2]

$$(b^\alpha_i)^* = (b^\alpha_i) = \rho^{\alpha\beta} \epsilon_{ij} b^\beta_j \quad (1)$$

$$\alpha, \beta, i, j = 1, 2 ; \quad \rho^{\alpha\beta} = - \rho^{\beta\alpha}$$

or,

$$b^i_\alpha = \begin{pmatrix} b & a \\ -a & b \end{pmatrix} \quad (2)$$

a, b complex.

The covariant derivative of b^i_α is

$$\hat{\partial}_\mu b^\alpha_i = \partial_\mu b^\alpha_i + ig(A_\mu \sigma_1^\alpha_i - B_\mu \sigma_2^\alpha_i) \quad (3)$$

that is

$$\begin{aligned} \hat{\partial}_\mu b &= \partial_\mu b \\ \hat{\partial}_\mu a &= \partial_\mu a - g(B_\mu - iA_\mu) \end{aligned} \quad (4)$$

and, correspondingly, the gauge symmetry is

$$\begin{aligned} \delta a &= \omega & \omega \text{ complex} \\ \delta B_\mu &= (1/g)\partial_\mu(\text{Re } \omega) \\ \delta A_\mu &= -(1/g)\partial_\mu(\text{Im } \omega) \end{aligned} \quad (5)$$

As it is evident from eqs.(4,5), this is exactly the standard structure of the truly massive vector (not a gauge particle), rewritten in a gauge-invariant way by adding to the Lagrangian the gauge compensator field a , (see ref.[5] for a discussion). As usual, one can define the massive vector fields gauge-invariantly as

$$(B_\mu - iA_\mu) = -(1/g)\hat{\partial}_\mu a, \quad (6)$$

note that B_μ (A_μ) coincides with B_μ (A_μ) in the "unitary" gauge $a=0$. After this redefinition, there is no more trace of the scalar "field" a in the Lagrangian.

Thus, our spin-1 particles turn out to be genuine massive vectors, not gauge connections, when the physical fields are correctly identified. Now, the fact that these particles are superpartners of the massive gravitino naturally suggests the intriguing possibility that we can analogously define a massive spin-3/2 field Ψ_μ , which is covariant with respect to local supersymmetry, by identifying the suitable compensator(s) from the fields present in our "symmetric" form of the Lagrangian. Hence, Ψ_μ may also be a genuine massive spin-3/2 field, not a supersymmetry connection. Unfortunately, it is very hard to test this conjecture because of the very non-linear gravitino supersymmetry transformations.

Let us note that, in the gauge $a=0$, the scalars live in the Kaehler manifold

$$SU(1,1)/U(1) \times SU(1,1)/U(1)$$

typical of $N=1$ supergravity, whereas there is no remainder of the quaternionic nature of the hypermultiplet σ -model, as required in the standard framework [5,6]. This is another hint that our models are really new.

Another property of the present model is the parity assignments of the vectors. The superconformal calculus gives theories which are invariant under a parity operation for which all spin-1 fields are true vectors. Our model has an (unbroken) parity invariance but, now, one of the (massive) vectors must transform as a pseudovector, in agreement with group theory considerations [7]. We feel that unusual parity assignments are crucial for partial supersymmetry breaking with $\Lambda=0$.

Now we turn to the construction of the model.

Recall that in the direct proof of the absence of partial super-Higgs in conventional $N=2$ supergravities [8], crucial points were inequalities of the form

$$\frac{(1/2)|A|^2}{1+(1/2)|A|^2} \neq 1 \quad (7)$$

where $|A|^2 = A^\alpha_i A^\dagger_{\alpha i}$ and A^α_i are the physical hypermultiplet fields [9]. As $|A|^2 \rightarrow \infty$, however, eq.(7) becomes an equality and so the arguments of refs.[1,8] no longer apply. Then, if for the Lagrangians of ref.[2] we are able to find a singular field transformation with the property that the original fields diverge (3) at the regular values of the new ones, while the Lagrangian and supersymmetry transformations remain finite in terms of the new fields, we obtain a new theory which may have a chance for partial super-Higgs.

To implement the above strategy, it is convenient to have partial super-Higgs from the beginning; this is possible if $\Lambda < 0$. An explicit example of partial super-Higgs in AdS space was presented in ref.[10]. That model can be generalized in a number of directions.

The particular model of ref.[10] was based on two vector multiplets (one of which is the compensator multiplet [11]) with minimal coupling [11,12]

$$N_{IJ} = \eta_{IJ} = \text{diag}(1, -1) \quad (8)$$

and one physical hypermultiplet A^α_i ($i, \alpha = 1, 2$) (which satisfy a reality condition similar to eq.(1)). The compensator vector multiplet is used to gauge an $SO(2)$ internal symmetry of the compensator hypermultiplet [2], whereas the matter vector multiplet gauges an $SO(2)$ subgroup of the physical hypermultiplet A^α_i .

In ref.[10] it was shown that the model has non-trivial solutions with cosmological constant

$$\Lambda = - \frac{6g_o'^2 t^2}{1 - t^2} \quad (9)$$

where $t^2 < 1$ is the ratio of the two gauge coupling constants $g_o'^2$ and g_o^2 .

The classical potential does not fix the direction and phases of the scalar v.e.v's and so the gravitino mass-matrix turns out to be sliding [10]. In particular, there is a choice of the indeterminate phases for which one supersymmetry is unbroken, that is [13]

$$m_1 = \sqrt{(-\Lambda/3)}$$

This case corresponds to the maximal parity breaking: both X_1 and the corresponding auxiliary field Y_{ij} are purely imaginary. In this situation the mass of the heavy gravitino is [10]

$$m_2 = \sqrt{2} g_0 t \frac{(1 + t^2)}{(1 - t^2)^{3/2}} \quad (10)$$

From eqs.(9,10), it is obvious that in the limit $t^2 \rightarrow 1$ with

$$g = g_0(1 - t^2)^{-3/2} \quad \text{fixed} \quad (11)$$

the cosmological constant and m_1 vanish, whereas m_2 has a finite limit. The expectation values of $|A|^2$ and $|X_1|^2$ diverge as $(1 - t^2)^{-1}$ in this limit. This is just the situation we need for our construction. It is easier to work in terms of the fields [2]

$$Z = X_0^{-1} X_1 = (1 + |X_1|^2)^{-1/2} X_1 \quad (12)$$

$$B^\alpha_i = (A_0)^{-1} A^\alpha_i = A^\alpha_i / (1 + |A|^2/2)^{1/2}$$

which have a finite limit as $t \rightarrow 1$. However, their kinetic terms are singular in the limit.

We introduce new scalar fields b^α_i (see eq.(1)) and ϕ (complex), by the following transformations

$$B^\alpha_i = \delta^\alpha_i - (1-t^2)b^\alpha_i \quad (13)$$

$$Z = i[1 - (1 - t^2) \phi] \quad (14)$$

which become singular as $t \rightarrow 1$. In eqs.(13,14) the constant terms are chosen to correspond to partial super-Higgs in AdS space for $t \sim 1$; more general choices are allowed, see eq.(32).

Moreover, we rescale all other fields (4) as follows:

$$A_\mu = \frac{\sqrt{2}}{(1-t^2)^{1/2}} A_\mu \Big|_{\text{NEW}} \quad (15)$$

$$B_\mu = \frac{\sqrt{2}}{(1-t^2)^{1/2}} B_\mu \Big|_{\text{NEW}} \quad (16)$$

$$\xi^\alpha = (1-(1/2)|B|^2)\xi^\alpha = (1-t^2)^{1/2}\xi^\alpha \Big|_{\text{NEW}} \quad (17)$$

$$i \Lambda_i = i(1 - |z|^2)\Omega_i = (1-t^2)^{1/2}\Lambda_i \Big|_{\text{NEW}} \quad (18)$$

where B_μ is the graviphoton and A_μ is the matter photon.

(In the following, we'll drop the subscript $|_{\text{NEW}}$).

With these redefinitions, as $t \rightarrow 1$, we get new covariant and supercovariant derivatives which are very different from the original ones; they are given in table I. In table II we give the new form of some objects which enter the supersymmetry transformations. We emphasize that these quantities can be used only in the field transformations and not in the Lagrangian, since in $\mathcal{L}$ they are multiplied by a coefficient of order $(1-t^2)^{-1}$ and so corrections of order $(1-t^2)$ to tables I, II contribute to $\mathcal{L}$ as $t^2 \rightarrow 1$.

The new supersymmetry transformations are ($t^2 \rightarrow 1$)

$$\delta e^a_\mu = \bar{\epsilon}^i \gamma^a \psi_{\mu i} + \bar{\epsilon}_i \gamma^a \psi_\mu^i \quad (19)$$

$$\delta \gamma = (\phi + \bar{\phi})^{1/2} \bar{\epsilon}^i \Lambda_i \quad (20)$$

$$\delta b^\alpha_i = -2(b + \bar{b})^{1/2} [\bar{\xi}^\alpha \epsilon_i + \rho^{\alpha\beta} \epsilon_{ij} \xi_\beta \epsilon^j] + i g (\alpha_1 \sigma_1^\alpha_i - \alpha_2 \sigma_2^\alpha_i) \quad (21)$$

$$\begin{aligned} \delta(B_\mu - i A_\mu) &= \\ &= \sqrt{2} \left[\frac{2 \bar{\epsilon}_i \psi_{\mu i} \epsilon^{ij}}{(\phi + \bar{\phi})^{1/2}} - \frac{\bar{\epsilon}_i \gamma_\mu \lambda_i \epsilon^{ij}}{(\phi + \bar{\phi})} \right] + \partial_\mu (\alpha_2 - i \alpha_1) \end{aligned} \quad (22)$$

$$\begin{aligned} \delta \psi^i_\mu &= 2 \nabla_\mu \epsilon^i + i \mathcal{A}_\mu \epsilon^i + \mathcal{V}_\mu^i \epsilon^j - (1/8) \sigma^{\sigma\rho} T_{\sigma\rho} \epsilon^{ij} \gamma_\mu \epsilon_j \\ &\quad - \gamma_\mu \eta^i - (i/2) \Lambda_A \psi_\mu^i + \Lambda^i_j \psi_\mu^j \end{aligned} \quad (23)$$

$$\delta \Lambda_i = \frac{2 \not{D} \phi \epsilon^i}{(\phi + \bar{\phi})^{1/2}} - \frac{2\sqrt{2} g}{(b + \bar{b})} (\delta_{ij} + \sigma_{3ij}) \epsilon^j$$

$$\begin{aligned}
& -\sigma^{\mu\nu}\{(\sqrt{2}/2)(\phi + \bar{\phi})(\hat{G}_{\mu\nu}^- - i\hat{F}_{\mu\nu}^-) \\
& - 2[(\phi + \bar{\phi})^{1/2}/(b + \bar{b})] \rho^{\alpha\beta}\bar{\xi}_{\alpha}\sigma_{\mu\nu}\xi_{\beta}\}_{\varepsilon_{ij}}\varepsilon^j - \\
& -(i/2)\Lambda_A\Lambda_i + \Lambda_i^j\Lambda_j + (\delta\phi)\Lambda_i/(\phi + \bar{\phi})
\end{aligned} \tag{24}$$

$$\begin{aligned}
\delta\xi^{\alpha} = & \frac{-\not{D}b^{\alpha}_i\varepsilon^i}{|b + \bar{b}|^{1/2}} + \frac{\sqrt{2}g}{(\phi + \bar{\phi})^{1/2}(b + \bar{b})^{1/2}}(\delta^{\alpha i} + \sigma_3^{\alpha i})\varepsilon_i - \\
& -(i/2)\Lambda_A\xi^{\alpha} + (b + \bar{b})^{-1}(\delta b^{\alpha}_j)\delta^j_{\beta}\xi^{\beta}
\end{aligned} \tag{25}$$

where α_i ($i=1,2$) are compensating gauge transformations needed in the presence of a gauge-fixing term.

From eqs.(19,-,25) and tables I,II, we see that, indeed, the susy transformations are finite in terms of the new fields as $t \rightarrow 1$. From general principles, we know that this implies that the Lagrangian is finite in the same limit. However, it is so in a rather subtle way.

The potentially divergent parts of $e^{-1}\mathcal{L}$ come from the following terms of the general Lagrangian (see eq.(4.13) of the first paper in ref.[2]).

$$\begin{aligned}
& (1/8)N_{IJ}\hat{F}_{\mu\nu}^I\hat{F}^{J\mu\nu} + \{(-1/64)N_{IJ}X^IX^J(\tau^+_{\mu\nu ij}\varepsilon^{ij})^2 - \\
& (1/4)N_{IJ}\bar{\Omega}^{il}\gamma_{\mu}\psi_{\nu}^j\varepsilon_{ij}\hat{F}^{\mu\nu J} - (1/4)N_{IJ}X^I\bar{\psi}_{\mu i}\psi_{\nu j}\varepsilon^{ij}\hat{F}^{\mu\nu J} - \\
& (1/16)N_{IJ}e^{-1}\varepsilon^{\mu\nu\rho\sigma}\bar{\psi}_{\mu i}\gamma_{\nu}\Omega^I_j\varepsilon^{ij}\bar{\psi}_{\rho k}\gamma_{\sigma}\Omega_m^J\varepsilon^{km} - \\
& (1/8)N_{IJ}e^{-1}\varepsilon^{\mu\nu\rho\sigma}\bar{\psi}_{\mu i}\psi_{\nu j}\varepsilon^{ij}\bar{\psi}_{\rho k}(\gamma_{\sigma}\Omega_m^I X^J \\
& + (1/2)\psi_{\rho m}X^IX^J)\varepsilon^{km} + \text{h.c.}\}
\end{aligned} \tag{26}$$

(in eq.(26) all fields are $|_{\text{on shell}}$). With some algebra, and the equation of motion for the auxiliary fields, it is easy to see that all terms of order $(1-t^2)^{-1}$ in eq.(26) combine in a total derivative

$$(i/4(1-t^2))\nabla_{\mu}[e^{-1}\varepsilon^{\mu\nu\rho\sigma}B_{\nu}\nabla_{\rho}A_{\sigma}] \tag{27}$$

and so, after integration, the action is finite as $t^2 \rightarrow 1$.

The final Lagrangian reads

$$\begin{aligned}
e^{-1}\mathcal{L} = & -(1/2)R - \frac{\partial_\mu \bar{\phi} \partial^\mu \phi}{(\phi + \bar{\phi})^2} - \frac{\hat{\partial}_\mu b^\alpha{}_i \hat{\partial}^\mu b^i{}_\alpha}{(b + \bar{b})^2} \\
& - (1/4) (\bar{\Lambda}^i \hat{\not{\partial}} \Lambda_i) / (\phi + \bar{\phi}) - 2(1/(b + \bar{b})) \bar{\xi}^\alpha \partial \xi_\alpha \\
& + \{ (e^{-1}/2) \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_{\mu i} \gamma_\nu \hat{\partial}_\rho \psi_\sigma{}^i - (\phi/4) \hat{G}^+_{\mu\nu} \hat{G}^{+\mu\nu} \\
& - (\phi/4) \hat{F}^+_{\mu\nu} \hat{F}^{\mu\nu} + \text{h.c.} \} \tag{28} \\
& + \{ -(2\sqrt{2})g \frac{\bar{\Lambda}^i (\delta_{i\alpha} + \sigma_{3i\alpha}) \xi^\alpha}{(\phi + \bar{\phi})(b + \bar{b})^{3/2}} + (2\sqrt{2})g \frac{\bar{\xi}^\alpha (\delta_{\alpha\beta} + \sigma_{3\alpha\beta}) \xi^\beta}{(\phi + \bar{\phi})^{1/2}(b + \bar{b})^{3/2}} \\
& - \bar{\psi}^i_\mu \gamma^\mu \left[\frac{2\sqrt{2} g}{(\phi + \bar{\phi})^{1/2}(b + \bar{b})^{3/2}} (\delta_{i\alpha} + \sigma_{3i\alpha}) \xi^\alpha - \right. \\
& - \frac{g}{\sqrt{2}(b + \bar{b})(\phi + \bar{\phi})} (\delta_{ij} + \sigma_{3ij}) \Lambda^j \left. \right] - \\
& - \frac{\sqrt{2} g (\delta_{ij} + \sigma_{3ij})}{(\phi + \bar{\phi})^{1/2}(b + \bar{b})} \bar{\psi}^i_\mu \sigma^{\mu\nu} \psi_\nu{}^j + \text{h.c.} \} \\
& + \{ -(3i/8) \frac{\bar{\Lambda}^i (\not{\partial} \phi) \Lambda_i}{(\phi + \bar{\phi})^2} - i \frac{\partial_\mu \bar{\phi}}{(\phi + \bar{\phi})} \frac{\bar{\xi}^\alpha \gamma^\mu \xi_\alpha}{(b + \bar{b})} - \\
& - \frac{\delta^i_\alpha \hat{\partial}_\mu b^\alpha{}_j}{(b + \bar{b})} \frac{\bar{\Lambda}^j \gamma^\mu \Lambda_i}{(\phi + \bar{\phi})} \}
\end{aligned}$$

$$\begin{aligned}
& -2 \frac{\bar{\xi}^\alpha (\hat{\partial} b_\alpha^i) \delta_i^\beta \xi_\beta}{(b + \bar{b})^2} \\
& + (3/64) \frac{(\bar{\Lambda}^i \Lambda^j) (\bar{\Lambda}_i \Lambda_j)}{(\phi + \bar{\phi})^2} - (1/2) \frac{(\bar{\xi}^\alpha \Lambda^i) (\bar{\Lambda}_i \xi_\alpha)}{(\phi + \bar{\phi})(b + \bar{b})} \\
& + 2 [(\bar{\xi}^\alpha \xi^\beta) (\bar{\xi}_\alpha \xi_\beta) / (b + \bar{b})^2] \\
& + (\sqrt{2}/4) [- \frac{|\phi|^2 + \phi^2}{(\phi + \bar{\phi})^2} \bar{\Lambda}^i \gamma_{[\mu} \psi_{\nu]}^j \varepsilon_{ij} \\
& + \frac{|\phi|^2 + \phi^2}{(\phi + \bar{\phi})^{3/2}} \bar{\psi}_\mu^i \psi_\nu^j \varepsilon_{ij}] (\tilde{G}^{\mu\nu} - i \hat{F}^{\mu\nu}) \\
& + \frac{\sqrt{2}(\phi + \bar{\phi})^{1/2}}{2(b + \bar{b})} \rho_{\alpha\beta} \bar{\xi}^\alpha \sigma_{\mu\nu\rho}^\beta (\hat{G}^{\mu\nu} - i \hat{F}^{\mu\nu}) \\
& + (1/4) \frac{\bar{\phi}^2}{(\phi + \bar{\phi})^2} e^{-1} \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu^i \psi_\nu^j \varepsilon_{ij} \bar{\psi}_\rho^k \psi_\sigma^m \varepsilon_{km} \\
& + (1/4) \frac{\phi^2 - |\phi|^2}{(\phi + \bar{\phi})^3} e^{-1} \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu^i \gamma_\nu \Lambda^j \varepsilon_{ij} \bar{\psi}_\rho^k \gamma_\rho \Lambda^m \varepsilon_{km} - \\
& - (1/4) \frac{|\phi|^2 - \phi^2 - \bar{\phi}^2}{(\phi + \bar{\phi})^{5/2}} e^{-1} \varepsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu^i \gamma_\nu \Lambda^j \varepsilon_{ij} \bar{\psi}_\rho^k \psi_\sigma^m \varepsilon_{km} \\
& + (1/8) e^{-1} \varepsilon^{\mu\nu\rho\sigma} [\frac{\bar{\phi}^2 + |\phi|^2}{(\phi + \bar{\phi})^2} \bar{\Lambda}_i \gamma_\mu \psi_{\nu j} \varepsilon^{ij}
\end{aligned}$$

$$\begin{aligned}
& - \frac{|\phi|^2 + \bar{\phi}^2}{(\phi + \bar{\phi})^{3/2}} \bar{\psi}_{\mu i} \psi_{\nu j} \epsilon^{ij}] \times \\
& \times \left[\frac{\bar{\Lambda}^k_{\gamma \rho} \psi_{\sigma}^m \epsilon_{km}}{(\phi + \bar{\phi})} + \frac{\bar{\psi}^k_{\rho} \psi_{\sigma}^m \epsilon_{km}}{(\phi + \bar{\phi})^{1/2}} \right] \\
& + e^{-1} \epsilon^{\mu\nu\rho\sigma} \bar{\psi}_{\mu i} \gamma_{\nu} \psi_{\rho j} \times \\
& \times \left[(1/8) \delta_j^i \frac{\partial_{\sigma}(\bar{\phi} - \phi)}{(\phi + \bar{\phi})} + (1/4) \frac{\delta^i_{\alpha} \hat{\partial}_{\mu} b^{\alpha}_j - \delta^{\alpha}_j \hat{\partial}_{\mu} b^i_{\alpha}}{(b + \bar{b})} \right] \\
& + (1/2) \bar{\psi}^i_{\mu} \frac{\not{\partial} \phi \gamma^{\mu} \Lambda_i}{(\phi + \bar{\phi})} + 2 \frac{\bar{\psi}^i_{\mu} \hat{\partial} b^{\alpha}_i \gamma^{\mu} \xi_{\alpha}}{(b + \bar{b})^{3/2}} \\
& - (1/8) (\phi + \bar{\phi})^{-1} \bar{\Lambda}^i_{\gamma} \gamma^{\mu} \gamma^{\nu} \psi_{\mu i} \bar{\psi}^j_{\nu} \Lambda_j \\
& - (1/(b + \bar{b})) \bar{\xi}^{\alpha}_{\gamma} \gamma^{\mu} \gamma^{\nu} \psi_{\mu i} (\bar{\xi}_{\alpha} \psi^i_{\nu} + \rho_{\alpha\beta} \epsilon^{ij} \bar{\xi}^{\beta} \psi_{i\nu}) + \text{h.c.} \}
\end{aligned}$$

From eq.(28), we see that the physical (i.e. positive kinetic energy) domain is

$$\begin{aligned}
(\phi + \bar{\phi}) &> 0 \\
(b + \bar{b}) &> 0
\end{aligned} \tag{29}$$

The scalar potential is absent, as promised, and the mass of the heavy gravitino is

$$m_2 = \frac{2\sqrt{2} \, g}{(\phi + \bar{\phi})^{1/2} (b + \bar{b})} \tag{30}$$

which slides with the value of $\text{Re } \phi$, $\text{Re } b$, which are not fixed by the classical potential.

From eq.(28) it is easy to see that the Lagrangian is invariant under a parity operation similar to the one present in tensor calculus, which on scalars is just complex conjugation. However, all terms in eq.(28) which are linear in the field strengths depend on the particular combinations $G_{\mu\nu} \pm iF_{\mu\nu}$ and this forces the parity assignments of the two spin-1 fields to be opposite, as anticipated. It must be stressed that this parity operation is not connected with the original one of tensor calculus, as it is easily seen from the phases in the field redefinitions, eq.(14,18); the new parity makes no sense for $t^2 \neq 1$ and is just a feature of the $t^2 \rightarrow 1$ limit.

As we have said before, this construction may be generalized in, at least, two directions:

- i) find more general flat models with partial breaking
- ii) find flat models with given m_1/m_2 .

A large class of models with property i) can be constructed by the present methods, since many models with partial breaking in AdS space are known. One can construct also models which, in the limit, have positive definite potential (just a sum of squares). The unfortunate aspect of the situation is the fact that, apart from spectator non-Abelian gauge interactions (which do not enter the above SUSY breaking mechanism) the present method allows only for Abelian gauge interactions in the relevant limit. A non-Abelian case would also have the problem that the θ -angle is of order $(1 - t^2)^{-1}$, and so the limit $t^2 \rightarrow 1$ would be not completely satisfactory. Thus, at the moment we are not able to solve the main problem in N=2 supergravity, namely to generate chiral non-Abelian interaction, at $\Lambda=0$ without generating tachyons.

As for problem ii) a simple example may be obtained just by expanding the field around a different solution of the model of ref.[10]. For instance, if we set

$$z = e^{i\theta} [1 - (1 - t^2)\phi] \quad (31)$$

with $\theta \neq 0$, all other field redefinitions being the same of eqs.(13,15,-,18), we end up with a model with a gravitino mass ratio [10]

$$(m_1/m_2)^2 = (1-\sin\theta)/(1+\sin\theta) \quad (32)$$

Again the Lagrangian turns out to be finite, but now the cancellation is more involved than in eqs.(26,27).

TABLE 1

Covariant and supercovariant derivatives.

$$\hat{\partial}_\mu b^\alpha_i = \partial_\mu b^\alpha_i + ig(A_\mu \sigma_1^\alpha_i - B_\mu \sigma_2^\alpha_i)$$

$$D_\mu b^\alpha_i = \hat{\partial}_\mu b^\alpha_i + (b + \bar{b})^{1/2} [\bar{\xi}^\alpha \psi_{\mu i} + \rho^{\alpha\beta} \epsilon_{ij} \bar{\xi}_\beta \psi^j_\mu]$$

$$D_\mu \phi = \partial_\mu \phi - (1/2)(\phi + \bar{\phi})^{1/2} \bar{\psi}^i_\mu \Lambda_i$$

$$\hat{F}_{\mu\nu} = \hat{F}^+_{\mu\nu} + \hat{F}^-_{\mu\nu} =$$

$$= F_{\mu\nu} + (i/\sqrt{2}) \left[\frac{\bar{\Lambda}^i_\gamma [\mu \psi^j_\nu] \epsilon_{ij}}{(\phi + \bar{\phi})} + \frac{\bar{\psi}^i_\mu \psi^j_\nu \epsilon_{ij}}{(\phi + \bar{\phi})^{1/2}} - \text{h.c.} \right]$$

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \quad G_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$$

$$\hat{G}_{\mu\nu} = \hat{G}^+_{\mu\nu} + \hat{G}^-_{\mu\nu} =$$

$$= G_{\mu\nu} - (1/\sqrt{2}) \left[\frac{\bar{\Lambda}^i_\gamma [\mu \psi^j_\nu] \epsilon_{ij}}{(\phi + \bar{\phi})} + \frac{\bar{\psi}^i_\mu \psi^j_\nu \epsilon_{ij}}{(\phi + \bar{\phi})^{1/2}} + \text{h.c.} \right]$$

$$\hat{\partial}_\mu \psi^i_\mu = \nabla_\mu \psi^i_\mu$$

$$\hat{\partial}_\mu \Lambda_i = (\partial_\mu - (1/2)\omega_{\mu\alpha\beta} \sigma^{\alpha\beta}) \Lambda_i$$

$$\hat{\partial}_\mu \xi^\alpha = (\partial_\mu - (1/2)\omega_{\mu\alpha\beta} \sigma^{\alpha\beta}) \xi^\alpha$$

TABLE II

Auxiliary expressions for the field transformations.

$$\mathcal{A}_\mu = -(i/2)(\partial_\mu(\phi - \bar{\phi})/(\phi + \bar{\phi}) + (i/8)(\bar{\Lambda}^k_{\gamma_\mu} \Lambda_k)/(\phi + \bar{\phi}) \\ - i(\bar{\xi}^\alpha_{\gamma_\mu} \xi_\alpha)/(b + \bar{b})$$

$$\mathcal{V}_{\mu j}^i = [\delta_\alpha^i \partial_\mu b_\alpha^j - (\partial_\mu b_\alpha^i) \delta_\alpha^j]/(b + \bar{b}) \\ - (1/2)[\bar{\Lambda}^i_{\gamma_\mu} \Lambda_j - (1/2)\delta_j^i \bar{\Lambda}^k_{\gamma_\mu} \Lambda_k]/(\phi + \bar{\phi})$$

$$T_{\mu\nu}^- = 2\sqrt{2} (\phi + \bar{\phi})^{-1/2} \{\hat{G}_{\mu\nu}^- + i\hat{F}_{\mu\nu}^-\}$$

$$\eta^i = \frac{\sqrt{2} g (\delta^{ik} + \sigma_3^{ik}) \varepsilon_k}{(b + \bar{b}) (\phi + \bar{\phi})^{1/2}}$$

$$-\gamma^\nu \varepsilon^i \left[\frac{\bar{\Lambda}_j \gamma_\nu \Lambda^i - (1/2)\delta_j^i \bar{\Lambda}_k \gamma_\nu \Lambda^k}{4(\phi + \bar{\phi})} + \delta_j^i \frac{\bar{\xi}^\alpha_{\gamma_\nu} \xi_\alpha}{(b + \bar{b})} \right]$$

$$-p_{\alpha\beta} \sigma^{\nu\sigma} \varepsilon^{ij} \varepsilon_j (\bar{\xi}^\alpha_{\sigma\nu} \xi_\alpha)/(b + \bar{b})$$

$$\Lambda_A = (i/2)(\phi + \bar{\phi})^{-1/2} [\bar{\varepsilon}^i \Lambda_i - \bar{\varepsilon}_i \Lambda^i]$$

$$\Lambda_i^j = (-2(b + \bar{b})^{-1/2}) [\bar{\xi}^\alpha (\delta_\alpha^j \varepsilon_i - (1/2)\delta_i^j \delta_\alpha^k \varepsilon_k)] \\ + (2(b + \bar{b})^{-1/2}) [\bar{\xi}_\alpha (\delta_i^\alpha \varepsilon^j - (1/2)\delta_i^j \delta_\alpha^k \varepsilon^k)]$$

NOTES

(1) Note that in the tensorial framework, we cannot have a nontrivial flat model in the presence of charged hypermultiplets.

(2) In the simplest case, the masses are sliding but the ratio of the two gravitino masses is uniquely fixed.

(3) More generally, they go to the boundary of the positivity domain for the corresponding kinetic terms.

(4) We follow the conventions of the first paper of ref.[2]. In particular, with X_I $I=0,1$ we denote the scalar components of the two vector multiplets.

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